18th European Conference on

Controlled Fusion and Plasma Physics

Berlin, 3-7 June 1991
Editors: P. Bachmann, D.C. Robinson

Contributed Papers
Part I

Published by: European Physical Society

Editor: Prof. K. Bethge, Frankfurt/M.

G. Thomas, Geneva

VOLUME 15 C Part I
Das Brandenburger Tor zu Berlin.
18th European Conference on

Controlled Fusion and Plasma Physics

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The 18th European Conference on Controlled Fusion and Plasma Physics was held in Berlin, Germany, from 3th to 7th June 1991. The Conference has been organized by the Central Institute of Electron Physics (ZIE) on behalf of the Plasma Physics Division of the European Physical Society.

The programme, format and schedule of the Conference were determined by the International Programme Committee appointed by the Plasma Physics Division of the EPS. The programme included 14 invited lectures and 388 contributed papers from which 27 were selected for oral presentation.

This 4-volume publication is published in the European Conference Abstract Series and contains all accepted contributed papers received in due time by the organizers. The 4-page extended abstracts were reproduced photographically using the manuscript submitted by the authors. The invited papers will be published in a special issue of the journal "Plasma Physics and Controlled Fusion" and sent free of charge to each registered participant.

The editors would like to acknowledge the skilful and dedicated support given by colleagues in the Department of Plasma Wall Interactions of the ZIE, in preparing the manuscript for reproduction in these four volumes.

May 1991

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CONTENTS

Part I  A  TOKAMAKS
A1  Tokamaks, General  I - 1
A2  Transport and Modelling  I -177
A3  Turbulence, Waves  I -245
A4  Pellet Injection  I -313
A5  Improved Confinement  I -353

Part II  A6  MHD Phenomena, Sawteeth  II - 1
B  STELLARATORS  II -109
C  ALTERNATIVE CONFINEMENT SCHEMES  II -217

Part III  D  PLASMA EDGE PHYSICS  III - 1
E  PLASMA HEATING AND CURRENT DRIVE  III-249

Part IV  F  GENERAL PLASMA THEORY  IV - 1
G  DIAGNOSTICS  IV -217
PAPER IDENTIFICATION

All contributed papers are listed with their title and the name of the first authors. Each paper may be identified by its classification symbol which starts with a letter that indicates to which topic it belongs followed by a number.
<table>
<thead>
<tr>
<th>Title List</th>
<th>Title</th>
<th>Author(s)</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>A</strong> TOKAMAKS</td>
<td></td>
<td></td>
</tr>
<tr>
<td>A1</td>
<td>Tokamaks, General</td>
<td></td>
</tr>
<tr>
<td></td>
<td>The effects of particle drift orbits on flux deposition profiles at the JET X-point target</td>
<td>Summers, D.D.R., Lesourd, M., Reichle, R., Schulz, J.-P., Zhu, Y.</td>
</tr>
<tr>
<td></td>
<td>Confinement of high performance JET plasmas</td>
<td>Balet, B., Cordey, J.G., Stubberfield, P.M., Thomsen, K.</td>
</tr>
<tr>
<td>A5</td>
<td>Evidence for fine scale density structures on JET under additional heated conditions</td>
<td>Cripwell, P., Costley, A.E.</td>
</tr>
<tr>
<td>A10</td>
<td>The role of the plasma current distribution in L-mode confinement</td>
<td>O'Rourke, J., Balet, B., Challis, C., Cordey, J.G., Gowers, C., et al.</td>
</tr>
</tbody>
</table>
VIII

A11 Current rise studies

A12 JET experiment with 120 keV \(^3\)He and \(^4\)He neutral beam injection

A13 Diffusion of alpha-like MeV ions in TFTR

A14 Observation of strongly localized fast particles ripple losses in TORE SUPRA
Roubin, J.-P., Guilhem, D., Martin, G., Pouri, B., Peysson, Y.

A15 Plasma decontamination during ergodic divertor experiments in TORE SUPRA
Monier-Garbet, P., DeMichelis, C., Evans, T.E., ... Mattioli, M., et al.

A16 Current diffusion and flux consumption in TORE SUPRA

A17 Scaling properties of runaway electrons in TJ-I tokamak
Rodriguez, L., Navarro, A.P.

A18 Effects of electrode polarization and particle deposition profile on TJ-I plasma confinement

A19 Tokamak transport and ohmic confinement
Bartiromo, R.

A20 Pressure anisotropy in ohmic FTU discharges
Alladio, F., Buratti, P., Grolli, M., Marinucci, M., Podda, S., Zerbini, M.

A21 Confinement studies of circular and X-point plasmas in the COMPASS-C tokamak

A22 Theoretical studies of tight aspect ratio tokamaks
Turner, M.F., Gryaznevich, M., Haynes, P.S., Nicolai, A., Sykes, A.

A23 Tight aspect ratio tokamaks
A24 Confinement projections for the burning plasma experiment (BPX)

A25 Isotope dependence of electron particle transport in ASDEX
Gehre, O., Gentle, K.W., ASDEX- and NI-Team

A26 On temperature and density dependence of the ASDEX L-mode confinement

A27 Long term \( Z_{\text{eff}} \) profile behaviour on ASDEX for different heating and wall conditions
Stuever, K.-H., Röhr, H., Dollinger, F., Engelhardt, W., et al., ASDEX-Team

A28 Current ramp experiments on the ASDEX tokamak
Nurmann, H., Stroth, U., ASDEX- and NI-Team

A29 ICRF heating on the burning plasma experiment (BPX)

A30 Radiation asymmetries of ASDEX divertor discharges close to the density limit
Müller, E.R., Hartinger, K.T., Niedermeyer, H., Stäbler, A., ASDEX-Team

A31 (Oral) Experimental observations of ohmic and ECR heated tokamak plasmas in RTP

A32 Diffusion of suprathermal electrons measured by means of ECRH and 2nd harmonic ECE O-mode
Schokker, B.C., Jaspers, R.J.E., Lopes Cardozo, N.J., and RTP team

A33 Simulations of \( \alpha \)-effects in TFTR D-T experiments

A34 The effect of passive stabilizing plates on high \( B \)-low \( q \) disruptions in PBX-M
Okabayashi, M., Bell, R., Chance, M., Fishman, H., Hatcher, R. et al.

A35 Effect of plasma aspect ratio on plasma confinement properties in JIPP T-IIU
Results and update on the TdeV facility  

The influence of hydrogen influx toroidal inhomogeneity on the particle balance analysis in FT-1 tokamak experiments  

The experiments on fast current rampdown in TUMAN-3 device  
Askinasi, L.G., Afanaziev, V.I., Golant, V.E., ... Sakharov, N.V., et al.

Comparison of measured magnetic field pitch profiles on PBX-M with Spitzer and neoclassical theories  
Kaye, S.M., Levinton, F., Hatcher, R., Kaita, R., LeBlanc, B., Paul, S.

Ignition achievement in high field tokamaks  
Airoldi, A., Cenacchi, G., Rulli, M.

Measurement of toroidal and poloidal plasma rotation in TCA  
Duval, B.P., Joye, B., Marchal, B.

Enhanced toroidal rotation in hot-ion mode with nearly-balanced neutral beam injection in JT-60  
Ishida, S., Koide, Y., Hirayama, T.

The initial experiments on the JT-60 Upgrade Tokamak  
JT-60 Team, Funahashi, A.

Comparison of dimensionally similar discharges with similar heat deposition profiles  
DeBoo, J.C., Waltz, R.E., Osborne, T.

Transport and Modelling

The simulation of energy and particle transport, heat and density pulse propagation and H-mode confinement in JET and a reactor  
Boucher, D., Rebut, P.H., Watkins, M.L.

Local transport analysis in L and H regimes  
Taroni, A., Seek, Ch., Springmann, E., Tibone, F.
<table>
<thead>
<tr>
<th>A47 (Oral)</th>
<th>Simulated ash transport experiments in JET using helium neutral beams and charge exchange spectroscopy</th>
</tr>
</thead>
<tbody>
<tr>
<td>A48</td>
<td>Local confinement in neutral beam heated JET discharges</td>
</tr>
<tr>
<td>A49</td>
<td>Interpretation of heat and density pulse propagation in tokamaks</td>
</tr>
<tr>
<td>Sips, A.C.C., Lopes Cardozo, N.J., Costley, A.E., Hogeweij, G.M.D., O'Rourke, J.</td>
<td></td>
</tr>
<tr>
<td>A50</td>
<td>Comparison of the impurity and electron particle transport in JET discharges</td>
</tr>
<tr>
<td>A51</td>
<td>Particle and energy transport properties deduced from the plasma dynamic response</td>
</tr>
<tr>
<td>Noret, J.-M., Bruneau, J.-L., Graud, A., Gil, C., Talvard, M.</td>
<td></td>
</tr>
<tr>
<td>A52</td>
<td>Determination of local transport coefficients from analysis of steady-state profiles and fluxes. Ion loss cones and particle transport in TJ-I tokamak</td>
</tr>
<tr>
<td>Rodriguez-Yunta, A., Pardo, C., Tabares, F., Zurro, B.</td>
<td></td>
</tr>
<tr>
<td>A53</td>
<td>Determination of the electron heat diffusivity from temperature perturbations in FT and FTU tokamaks</td>
</tr>
<tr>
<td>Berton, F., Bruschi, A., Buratti, F., Imperiali, C., Romanelli, F., Tudisco, C.</td>
<td></td>
</tr>
<tr>
<td>A54</td>
<td>A study of the ion species dependence of X by heat pulse propagation</td>
</tr>
<tr>
<td>Giannone, L., Krämer-Flecken, A., Mertens, V., Riedel, K., Wagner, F., Waidmann, G.</td>
<td></td>
</tr>
<tr>
<td>A55</td>
<td>Comparison of anomalous momentum transport with particle and energy transport on ASDEX</td>
</tr>
<tr>
<td>Kallenbach, A., Fussmann, G., Mayer, H.M., Krieger, K., ...et al., ASDEX Team</td>
<td></td>
</tr>
<tr>
<td>A56</td>
<td>Statistical analyses of local transport coefficients in ohmic ASDEX discharges</td>
</tr>
<tr>
<td>A57</td>
<td>Reassessment of the interpretation of sawtooth-induced pulse propagation</td>
</tr>
<tr>
<td>Lopes Cardozo, N.J., Sips, A.C.C.</td>
<td></td>
</tr>
<tr>
<td>A58</td>
<td>Role of boundary plasmas on the energy and particle transport in JT-60</td>
</tr>
</tbody>
</table>

**Notes:**
- The page numbers next to each entry seem to be random and do not correspond to the typical page numbering of a journal or conference proceedings. They may be placeholders or misprints.
- The entries appear to be a list of papers or abstracts, possibly from a conference or journal, with authors, titles, and some additional notes or references.
Local transport analysis of L-mode plasmas in JT-60
Hirayama, T., Kikuchi, M., Shirai, H., Shimizu, K., Yagi, M., Koide, Y., Azumi, M.

Transport simulations using theory-based models
Bateman, G., Singer, C.E., Kinsey, J.

Non-diffusive heat transport during electron cyclotron heating on the DIII-D tokamak
Petty, C.C., Luce, T.C., de Haas, J.C.M., James, R.A., Lohr, J., et al.

A model to evaluate coupling of ion Bernstein waves to tokamak plasmas
Brambilla, M., Cardinali, A., Cesario, R.

Fluctuation measurements by Langmuir probes during LHCD on ASDEX tokamak
Stöckel, J., Söldner, F., Giannone, L., Leuterer, F., ASDEX-Team

Experimental results of runaway electrons and magnetic fluctuations on HL-1 tokamak
Gong Dingfu, Yang Qingwei, Dong Jiafu, Yang Shikun, Shang Zuoyu et al.

Optical visualization of magnetic island structures and comparison with a magnetic turbulence model
Drawin, H.W., Dubois, M.A.

High frequency magnetic modes and particle transport in rotating and locked TOKOLOSHE plasmas

Density fluctuations at the sawtooth crash in TFTR

Long wavelength density turbulence measurements in TFTR beam-heated discharges

Density fluctuations in ohmic, L-mode, and H-mode discharges of ASDEX
A70 Analysis of coupled temperature and density perturbations using Fourier methods

A71 Two phase reduction of microturbulence at the transition into H-mode
measured on DIII-D

A72 ELM precursors on DIII-D

A73 Study of edge electric field and edge microturbulence at the L-H transition
in DIII-D

A74 Framing camera studies of the edge in rotating and locked TOKOSHE plasmas

A75 Comparison of edge fluctuations in toroidal confinement devices
Tsui, H.Y.W., Lin, H., Meier, N., Ritz, C., Wootton, A.J.

A76 Multi-channel Langmuir-probe and $H_e$ -measurements of edge fluctuations on ASDEX

A77 ELM studies on ASDEX
Zohm, H., Wagner, F., Endler, M., Gernhardt, J., Holzhauer, E., Mertens, V.

A78 Detection of coherent structures in the edge of the TEXT tokamak plasma
Filippas, A.V., Ritz, Ch.P., Koniges, A.E., Crottinger, J.A., Diamond, P.H.

A4 Pellet Injection

A79 (Oral) Parallel expansion of the ablation cloud during pellet injection in TORE SUPRA
Pegouri, B., Bruneau, J.L., Picchiottino, J.M.

A80 Density control in TORE SUPRA with ergodic divertor and multi-pellet injection
A81 Energy confinement of high-density pellet-fuelled H-mode plasmas in ASDEX
Mertens, V., Kaufmann, H., Lang, R., Loch, R., Sandmann, W., et al., ASDEX-Team

A82 Impurity pellet experiments in TEXT

A83 Pellet ablation studies in the TCA tokamak
Drakakis, N., Dutch, M., Duval, B.P., Holienstein, Ch., ... Martin, Y., et al.

A84 Studies on fast oscillations and on particle transport during sawtooth crashes in pellet-injected TEXTOR plasmas

A85 Analysis of symmetrization process of density perturbations after deuterium pellet injection in T-10

A86 Plasma cloud near the pellet injected into a tokamak

A87 Plasma perturbation during hydrogen pellet injection on T-10

A88 Behaviour of an ion component during the pellet injection on T-10
Braitsev, S.N., Efremov, S.L., Medvedev, A.A., Pivinsky, A.A.

A5 Improved Confinement

A89 Influence of VB drift direction on H-modes in JET

A90 Control of carbon blooms and the subsequent effects on the H to L mode transition in JET X-point plasmas
Stork, D., Campbell, D.J., Clement, S., Gottardi, N., de Kock, L., et al.

A91 The evolution of $Z_{\text{eff}}$ during H-mode operation in JET
Hot-ion and H-mode plasmas in limiter configuration in JET
Tonga, A., Jones, T.T.C., Lomas, F., Nardone, C., Sartori, R., et al.

ICRH H-modes produced with Be-screen antennas and coupling-resistance position feedback control
Bhatnagar, V.P., Bosia, G., Bures, M., Campbell, D., Fessey, J., et al.

Edge current density in H-mode discharges at JET

Power threshold for L-H mode transition in JET
Nardone, C., Bhatnagar, V.P., Campbell, D., Gottardi, N., Lazzaro, E., et al.

Does the ion confinement improve in ASDEX H-mode discharges?
Gruber, O., Menzler, H.-P., Herrmann, W., Kallenbach, A., Steuer, K.-H.

Long pulse stationary H-mode with ELMS on ASDEX
Volmer, O., Rytie, F., Steuer, K.H., Wagner, F., Zohm, H., ASDEX and NI Teams

Achieving improved Ohmic confinement via impurity injection
Bessenrodt-Weberpals, M., Soiidner, F.X., and ASDEX Team

H-mode studies with microwave reflectometry on ASDEX

Random coefficient H-mode scalings
Riedel, K.S.

Edge plasma behaviour in ohmic H-mode and edge polarization in TUMAN-3

Novel features of H-mode plasmas induced by edge polarization in TEXTOR

Improved confinement regimes in TEXTOR

Canonical profiles transport model for improved confinement regimes in tokamaks
Dnestrovskij, Yu.N., Esipchuk, Yu.V., Lysenko, S.E., Tarasjan, K.N.
A105  On the possibility of reaching the H-mode from initial conditions
Spineanu, F., Vlad, M.

A6 MHD Phenomena, Sawteeth

A106  Influence of ECRH on electron temperature sawtooth oscillations on T-10
Bagdasarov, A.A., Neudatchin, S.V.

A107  Sawtooth oscillations in TFTR

A108  Sawtooth stabilization studies on TFTR

A109  The scaling of sawtooth parameter and the occurrence of single sawteeth in the
start-up phase of TEXTOR
Graffmann, E., Fang, Z.S., Soltwisch, H., Wang, K.

A110  Sawtooth oscillations of various plasma parameters in correlation to periodic
changes of the internal magnetic field structure
Soltwisch, H., Fuchs, G., Koslowski, H.R., Schlüter, J., Waidmann, G.

A111  D simulation of the sawtooth ramp
Gimblett, C.G., Campbell, D.J., Fitzpatrick, R., Hastie, R.J., Martin, T.J.

A112  Latest JET experimental results on the sawtooth
Pearson, D., Campbell, D.J., Edwards, A.W., O’Rourke, J.

A113  Local X-ray emissivity structure in the density limit in the TJ-I tokamak
Vega, J., Navarro, A.P.

A114  Density limit studies on ASDEX

A115  Density limit studies in the TCA tokamak
A116 (Oral)  
Simulation of MHD activity during density limit disruptions in tokamaks  
Bonneson, A., Parker, R., Hugon, M.

A117  
The evolution of the density limit disruption in TEXTOR  
Waidmann, G., Kuang, G.

A118  
Spontaneous appearance of snakes in JET  
Gill, R.D., Edwards, A.W., Pasini, D., Wolfe, S.W.

A119  
MHD studies in JET  

A120  
MHD stability and mode locking in pre-disruptive plasmas on TORE SUPRA  

A121 (Oral)  
Mode-locking and error field studies on COMPASS-C and DIII-D  
Morris, A.W., Fitzpatrick, R., Hender, T.C., Todd, T.N., Bamford, R. et al.

A122  
Statistical properties of intrinsic topological noise in tokamaks  
Evans, T.E.

A123  
Tokamak error fields and locked modes  
Reiman, A., Monticello, D.

A124  
The effects of MHD-activity on the density in the RTP tokamak  

A125  
Excitation of toroidal Alfvén Eigenmodes in TFTR  

A126 (Oral)  
Asymmetric reconnection and stochasticity induced by the m=1 island  
Baty, H., Luciani, J.F., Bussac, M.N.

A127  
Estimation of the major disruption time by energy approach  
Astashkovich, A.M., Kokotkov, V.V., Mineev, A.B.

A128  
a) Systematic investigation of periodic disruptions using a new diagnostic ...  
b) Diagnosis of ion energy distribution based on ejected fast-ions from ...  
Nihar Ranjan Ray, Basu, J., Majumdar, S.K., Mukherjee, S.
<table>
<thead>
<tr>
<th>A129</th>
<th>II-93</th>
</tr>
</thead>
<tbody>
<tr>
<td>Shaping and vertical stability elongated plasmas on the TVD</td>
<td></td>
</tr>
<tr>
<td>Abramov, A.V., Bortnikov, A.V., Brevnov, N.N., Gerasimov, S.N., Polianchik, K.D.</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>A130</th>
<th>II-97</th>
</tr>
</thead>
<tbody>
<tr>
<td>Nonlinear mode mixing in high beta PBX-M discharges</td>
<td></td>
</tr>
<tr>
<td>Sesnic, S., Kaye, S., Okabayashi, M.</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>A131</th>
<th>II-101</th>
</tr>
</thead>
<tbody>
<tr>
<td>Simulation study of stable high-B tokamak plasmas with beam-driven and bootstrap currents</td>
<td></td>
</tr>
<tr>
<td>Murakami, Y., Okano, K., Ogawa, Y., Takase, H., Shinya, K.</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>A132</th>
<th>II-105</th>
</tr>
</thead>
<tbody>
<tr>
<td>Dependence of the DIII-D beta limit on the current profile</td>
<td></td>
</tr>
<tr>
<td>Title</td>
<td>Authors</td>
</tr>
<tr>
<td>----------------------------------------------------------------------</td>
<td>-------------------------------------------------------------------------</td>
</tr>
<tr>
<td>Mercier criterion for stellarator with planar circular axis</td>
<td>Cherepnykh, O.K., Podnebesnyj, A.V., Pustovitov, V.D.</td>
</tr>
<tr>
<td>... in stellarators</td>
<td>Pustovitov, V.D., Pukhov, A.V.</td>
</tr>
<tr>
<td>Charged particle injection and fusion product escape in separatrix</td>
<td>Alladio, F., Batistoni, P., Mancuso, S.</td>
</tr>
<tr>
<td>Optimization of analytic stellarator fields by using mapping methods</td>
<td>Dommaschk, W., Herrnegger, F., Schlüter, A.</td>
</tr>
<tr>
<td>Effects of externally-applied perturbation field on the confinement</td>
<td>Okamura, S., Peranich, L., Matsuoka, K., Iguchi, H., et al.</td>
</tr>
<tr>
<td>Improvement of 1=2 torsatron configuration with additional toroidal</td>
<td>Besedin, N.T., Lesnyakov, G.G., Pankratov, I.M.</td>
</tr>
<tr>
<td>Magnetic fluctuations in Heliotron E</td>
<td>Zushi, H., Harada, N., Osaki, T., Wakatani, M., Obiki, T., and Heliotron E group</td>
</tr>
</tbody>
</table>
B13
Ion confinement and radiation losses in the advanced toroidal facility

B14
Heavy ion beam probe measurements of ECH heated plasma in the advanced toroidal facility

B15
Biasing experiments on the ATF torsatron
Uckan, T., Aceto, S.C., Baylor, L.R., Bell, J.D., Bigelow, T., et al.

B16 (Oral)
Configuration control, fluctuations, and transport in low-collisionality plasmas in the ATF torsatron
Harris, J.H., Murakami, M., Aceto, S., Branas, B., ... Lyon, J.F., et al.

B17
Structure of the magnetic field line diversion in Helias configurations
Strumberger, E.

B18
Neoclassical transport in stellarators - a comparison of conventional stellarator/torsatrons with the advanced stellarator Wendelstein 7X
Beidler, C.D.

B19
On natural islands and the edge structure of the Wendelstein 7-X stellarator

B20
Optimization of coils for divertor experimentation in W7-X
Merkel, P.

B21
Impurity behaviour in W7-AS plasmas under different wall conditions
Brakel, R., Burhenn, R., Elsner, A., et al., W7-AS Team, ECRH Group, NI Group

B22
Particle transport and plasma edge behaviour in the W7-AS stellarator

B23
MHD activity driven by NBI in the W7-AS stellarator
Lazaros, A., Jaenicke, R., Weller, A.

B24
Neutral injection experiments on W7-AS stellarator
Penningsfeld, F.-P., Ott, W., W7-AS Team, ECRH Group, NI Group, Pellet Inj. Group
<table>
<thead>
<tr>
<th>B25</th>
<th>Simulation of the influence of coherent and random density fluctuations on the propagation of ECRH-beams in the W7-AS stellarator</th>
</tr>
</thead>
<tbody>
<tr>
<td>Tutter, M., Erckmann, V., Gasparino, U., W7-AS Team</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>B26</th>
<th>Ion heat conductivity, radial electric fields and CX-losses in the W7-AS stellarator</th>
</tr>
</thead>
<tbody>
<tr>
<td>Afanasjev, V.I., Izvozhikov, A.B., Junker, J., Kick, M., et al., W7-AS Team &amp; NBI</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>B27</th>
<th>Thermal diffusivity from heat wave propagation in Wendelstein 7-AS</th>
</tr>
</thead>
<tbody>
<tr>
<td>Hartfuß, H.J., Erckmann, V., Giannone, L., Maaßberg, H., Tutter, M.</td>
<td></td>
</tr>
<tr>
<td>XXII</td>
<td></td>
</tr>
<tr>
<td>-------------------------------</td>
<td></td>
</tr>
<tr>
<td><strong>C ALTERNATIVE CONFINEMENT SCHEMES</strong></td>
<td></td>
</tr>
<tr>
<td>-------------------------------</td>
<td></td>
</tr>
</tbody>
</table>
| **C1** Experiment on ion beam collider for advanced fuel fusion  
  Tanaka, H., Yuyama, T., Mitchishita, T., Mohri, A. | II-217 |
| **C2** Hot spot formation and emission characteristics of the plasma focus  
  Antsiferov, P., Franz, D., Herold, H., Jakubowski, L., Jonas, A., ... Schmidt, H. | II-221 |
| **C3** Rundown phase of plasma focus in multicharged gas  
  Gursov, A.V. | II-225 |
| **C4** Interaction of plasma with magnetic fields in coaxial discharge  
  Soliman, H.H., Masoud, M.M. | II-229 |
| **C5** Studies of high energy electron beams emitted from PF-type discharges  
  Sadowski, N., Jakubowski, L., Zebrowski, J. | II-233 |
| **C6** X-ray radiation and electron beams in a plasma focus  
| **C7** Interpretation of self sustained wave coupling and microturbulence in  
  SK/CG-1 machine  
  Sinman, S., Sinman, A. | II-241 |
| **C8** Toroidal discharges in an octupole field  
  Hellblom, G., Brunsell, P., Drake, J.R. | II-245 |
| **C9** Experimental detection of locking of vacuum electron drift in curvilinear  
  element of drakon system  
  Perelygin, S.F., Smirnov, V.N. | II-249 |
| **C10** Ion energy measurements in the SPHEREX Spheromak  
  Gibson, K.J., Gee, S.J., Cunningham, G., Rusbridge, M.G., Carolan, P.G. | II-253 |
| **C11** Magnetic equilibria in spheromaks and low aspect ratio tokamaks  
  Browning, P.K., Duck, R., Martin, R., Rusbridge, M.G. | II-257 |
| **C12** Straight axisymmetric trap with average minimum-B  
  Arsenin, V.V., Sergeev, E.B. | II-261 |
Detrapping of hot electrons from magnetic well under ECR heating with parallel HF power launching
Zhil'tsov, V.A., Skvorodka, A.A., Timofeev, A.V., Scherbakov, A.G.

Recent upgraded tandem mirror experiments in GAMMA 10

Observation of various electron velocity distribution shapes using X-ray diagnostics in GAMMA 10
Cho, T., Hirata, M., Takahashi, E., Yamaguchi, N., Ogura, K., Masai, K., et al.

An approach to the fusion neutron source concept based on a mirror with ICRF heating

Electron temperature-gradient instability caused by conducting end-plates in mirror devices
Ryutov, D.D., Tsidulko, Yu.A., Berk, H.L.

On the role of slow compression in self-organization of the reversed field pinch plasma
Sorokin, A.V.

Edge fluctuations and transport in the MST reversed field pinch

Ion and electron temperature measurements on Repute-1 reversed field pinch

Formation, confinement, and stability of compact toroids in LSX

Initial experimental results in the LSX field reversed configuration

Density regimes and dynamo processes in RFP plasmas
Ferrer Roca, Ch., Innocente, P., Martini, S., Paccagnella, R.

Electron and ion temperature studies on ETA-BETA II reversed field pinch
Carraro, L., Costa, S., Martin, P., Puiatti, M.E., Scarin, P., Valisa, N.
C25
Resistive MHD analysis of rotational instabilities in FRC
Santiago, M.A.N., Tsui, K.H., Azevedo, M.T., Sakanaka, P.H.

C26 (Oral)
Particle transport investigation in the HBTXIC reversed field pinch
Walsh, M.J., Carolan, P.G.

C27
Tight neck in quick dense hydrocarbonic Z—pinch at MA current

C28
Modelling of short wavelength m = 0 instabilities in the compressional Z—pinch
Bayley, J.M., Coppins, M., Jaitly, P.

C29
A universal diagram for regimes of Z—pinch stability
Haines, M.G., Coppins, M.

C30
Characteristics of the neutron emission in different plasma focus configurations
One-dimensional model for a description of transitions of a tokamak edge plasma into a strongly radiative state

Tokar', M.Z.

Neutral transport and hydrogen recycling in edge region of HL-1 tokamak plasma


Analysis of a scheme to improve the SOL transport properties by plasma current modulation

Nicolai, A.

Influence of limiter bias and interchange instabilities on the structure of the tokamak plasma edge

Gerhauser, H., Claassen, H.A.

Excitation of an instability by neutral particle ionization induced fluxes in the Tokamak edge plasma

Bachmann, P., Morozov, D.Kh., Sünd, D.

Ion velocity distributions at the tokamak edge

Pitts, R.A.

Experimental verification of simple scaling relations for edge-plasma density in tokamaks

Alexander, K.F., Günther, K., Laux, M.

2-D modelling of the edge plasma with arbitrary high level of impurity concentration

Igitkhanov, Yu.L., Pozharov, V.A.

Edge plasma transport in Grad approach

Rabinski, N.

Instabilities of plasma-collector interaction in imitation experiments

Vizgalov, I.V., Dimitrov, S.K., Kurnaev, V.A., Chernyatjev, Yu.V.

Main characteristics of a high-power full-scale quasi-stationary plasma accelerator QSPA-Kh-50 and some results of preliminary experiments

Kulik, N.V., Nanojlo, V.S., Malikov, V.A., ... Tereshin, V.I., et al.
D12  
Impurity deposition on surface probes during different operation modes at EXTRAP T1  
Gudowska, I., Bergsaker, H., Hellblom, G.

D13  
Analysis of hydrogen and impurity outgassing under carbon, boron and beryllium first wall conditions  

D14  
Sputtering and redeposition of impurities in T15 studied by collector probes  

D15  
Characterization of the wall recycling properties and bulk particle lifetime in TORE SUPRA  
Grisolia, C., Hutter, T., Pourgiez, B.

D16  
Scaling properties for the edge turbulence in the ATF torsatron  
Hidalgo, C., Meier, M.A., Uckan, T., Ritz, Ch.P., Harris, J.H., et al.

D17  
Experimental study on edge electric field in Heliotron-E currentless plasma  

D18  
Effects of energetic electrons on the edge properties of the ETA-BETA II reversed field pinch  
Antoni, V., Bagatin, M., Desideri, D., Martines, E., Yagi, Y.

D19  
Boronisation, recycling and isotope ratio control experiments on COMPASS  

D20  
Experimental determination of the helium pumping by beryllium  
Saibene, G., Clement, S., Ehrenberg, J., Peacock, A., Philipps, V., Sartori, R.

D21  
Deuterium and tritium release on venting the JET torus to air after the beryllium phase  
Coad, J.P., Gibson, A., Haigh, A.D., Kaveney, G., Orchard, J.

D22 (Oral)  
Turbulence studies in the proximity of the velocity shear layer in the TJ-I tokamak  
ohmic and H-mode particle transport in the CCT tokamak edge plasma
Tynan, G.R., Conn, R.W., Doerner, R., Lehmer, R., Schmitz, L.

Edge radial profiles and transport in JET X-point plasmas

Dependence of He retention on X-point plasma parameters in JET

The behaviour of neutral particles in the private region of X-point discharges in JET

Power loading and radiation distribution at the X-point target in JET for normal and reversed toroidal field
Reichle, R., Clement, S., Gottardi, N., Jaeckel, H.J., Lesourd, M., Summers, D.D.R.

Results from edge spectroscopy in JET

Measurement of the radial electric field at the periphery of ASDEX plasmas
Field, A.R., Fussmann, G., Hofmann, J.V.

Ion temperature near the separatrix at ASDEX
Schneider, R., Verbeek, H., Reiter, D., Neuhauser, J., and ASDEX team

Experimental investigation of ExB transport during the transition from attached to detached plasmas in TEXTOR
Bors, D., Fuchs, G., Ivanov, R.S., Samm, U., Van Oost, G.

Temperature profiles of C(6+)-ions in the TEXTOR edge plasma — measured with lithium-beam activated charge-exchange spectroscopy
Schorn, R.P., Claassen, H.A., Hintz, E., Rusbultdt, D., Unterreiter, E.

Flux and energy of neutral Deuterium and radial flux of neutral Boron in Textor
Bergsaker, H., Emmoth, B., Wienhold, P.

Density profile and DC electric field measurements in the TEXTOR boundary for ohmic and ICRF-heated discharges
<table>
<thead>
<tr>
<th>D35</th>
<th>III-137</th>
</tr>
</thead>
<tbody>
<tr>
<td>Suppression of Marfes by plasma position feedback control based on interferometric measurements</td>
<td></td>
</tr>
<tr>
<td>Samm, U., Koslowski, H.R., Soltwisch, H.</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D36</th>
<th>III-141</th>
</tr>
</thead>
<tbody>
<tr>
<td>Wall conditioning by lithium pellet injection on TFTR</td>
<td></td>
</tr>
<tr>
<td>Snipes, J.A., Terry, J.L., Marmar, E.S., Bell, N.G., et al., TFTR Group</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D37</th>
<th>III-145</th>
</tr>
</thead>
<tbody>
<tr>
<td>The effect of density on boundary plasma behaviour in TFTR</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D38</th>
<th>III-149</th>
</tr>
</thead>
<tbody>
<tr>
<td>Power flow thickness and edge density scaling in the scrape-off layer of JET</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D39</th>
<th>III-153</th>
</tr>
</thead>
<tbody>
<tr>
<td>Deuterium depth profiles and impurities on the Be coated carbon belt limiter in JET</td>
<td></td>
</tr>
<tr>
<td>Martinelli, A.P., Hughes, I., Behrisch, R., Peacock, A.T.</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D40 (Oral)</th>
<th>III-157</th>
</tr>
</thead>
<tbody>
<tr>
<td>Radiation cooling with intrinsic and injected impurities in the plasma boundary of a limiter tokamak</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D41</th>
<th>III-161</th>
</tr>
</thead>
<tbody>
<tr>
<td>A model for limiter heat flux protection by local impurity radiation</td>
<td></td>
</tr>
<tr>
<td>Tokar', M.Z., Nedospasov, A.V., Samm, U.</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D42</th>
<th>III-165</th>
</tr>
</thead>
<tbody>
<tr>
<td>Oxygen collection in the limiter shadow of TEXTOR depending on wall conditioning with boron</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D43</th>
<th>III-169</th>
</tr>
</thead>
<tbody>
<tr>
<td>(H_n)-diagnostics of the limiter surrounding in T-15 using a CCD-camera</td>
<td></td>
</tr>
<tr>
<td>Kastelewicz, H., Pigarov, A.</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D44</th>
<th>III-173</th>
</tr>
</thead>
<tbody>
<tr>
<td>Test of a carbonized molybdenum limiter in TEXTOR</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D45</th>
<th>III-177</th>
</tr>
</thead>
<tbody>
<tr>
<td>The preliminary studies of HL-1 plasma with pump limiter</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D46 (Oral)</th>
<th>III-181</th>
</tr>
</thead>
<tbody>
<tr>
<td>Changes in the limiter shadow of T-10 during current and/or field reversal</td>
<td></td>
</tr>
</tbody>
</table>
A sheath model of asymmetric heat flow to a poloidal limiter

Haines, M.G.

Helium and carbon ions flow in the pump limiter channel

Brevnov, N.N., Stepanov, S.B., Khimchenko, L.N.

Magnetic shielding of a limiter


Towards fully authentic modelling of ITER divertor plasmas

Maddison, G.P., Hotston, E.S., Reiter, D., Börner, P., Baelmans, T.

Retention of gaseous impurities in the divertor of DIII-D

Lippmann, S., Mahdavi, A., Roth, J., Krieger, K., Fußmann, G., Janeschitz, G.

Divertor plate biasing experiments on the Tokamak de Varennes


The simulation of the ITER divertor plates erosion in stationary plasma


Influence of kinetic effects on a sheath potential and divertor plasma parameters in ITER

Krasheninnikov, S.I., Soboleva, T.K., Igitkhanov, Yu.L., Runov, A.N.

Carbon radiation in the vicinity to the neutralization plates

Abramov, V.A., Brevnov, N.N., Pistunovich, V.I., Stepanov, S.B., Khimchenko, L.N.

An analytic model for retention of divertor impurities by forced flows

Vlases, G.C., Simonini, R.

Impurity flow at a divertor target

Chodura, R., Zanino, R.

Impurity transport at the DIII-D divertor strike points


Operating conditions of the BPX divertor

<table>
<thead>
<tr>
<th>D60</th>
<th>III-237</th>
<th>Recent gaseous divertor experiments in DIII-D</th>
</tr>
</thead>
</table>

<table>
<thead>
<tr>
<th>D61</th>
<th>III-241</th>
<th>Advanced divertor experiments on DIII-D</th>
</tr>
</thead>
</table>

<table>
<thead>
<tr>
<th>D62</th>
<th>III-245</th>
<th>Dynamic measurements of the hydrogen inventory in graphite exposed to a RF-discharge</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td>Jandl, C., Möller, W., Scherzer, B.</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>D63</th>
<th>IV-369</th>
<th>ALT-II toroidal belt limiter biasing experiments on TEXTOR</th>
</tr>
</thead>
</table>
| E1 | Numerical analysis on neutral-beam current drive with an energy spread of the neutral beam  
    | Okazaki, T., Ohtsuka, M. |
| E2 | Efficiency studies of high frequency current drive  
    | Karttunen, S.J., Pättikangas, T.J.H., Salomaa, R.R.E. |
| E3 | Relativistic electron cyclotron absorption for perpendicular propagation in an inhomogeneous magnetic field  
    | Bornatici, M. |
| E4 | 3D treatment of the effects of radial transport on RF current drive in tokamaks  
    | O’Brien, M.R., Cox, M., McKenzie, J.S., Warrick, C.D. |
| E5 | Three-dimensional simulation of electron cyclotron current drive in tokamak plasma  
| E6 | Spectral properties and absorption of Alfvén waves in toroidal plasmas  
    | Elfimov, A.G., Medvedev, S.Yu., Pestryakova, G.A. |
| E7 | Theory of linear propagation and absorption of the ion Bernstein waves in toroidal geometry  
    | Cardinali, A., Romanelli, F. |
| E8 | Maximizing absorption in ion-cyclotron heating of tokamak plasmas  
    | Bers, A., Fuchs, V., Chow, C.C. |
| E9 | Numerical modelling, analysis, and evaluation of ICRH antennae  
    | Grossmann, W., Riyopoulos, S., Ko, K., Kress, M., Drobot, A.T. |
| E10 | Fast electron transport during lower-hybrid current drive  
    | Kupfer, K., Bers, A., Ram, A.K. |
| E11 | Fast wave helicity current drive in tokamaks  
<pre><code>| Tataronis, J.A., Moroz, P.E., Hershkowitz, N. |
</code></pre>
<p>| E12 | Numerical simulation of current drive by RF fields and helicity injection | Elfimov, A., Churkina, G., Dmitrieva, M.V., Potapenko, I. | III-293 |
| E13 | Quasilinear description of heating and current drive in tokamaks by means of test particle Fokker-Planck equation | Faulconer, D.W., Evrard, M.P. | III-297 |
| E14 | Experimental studies on RF current drive by a standing Alfvén wave | Kirov, A.G., Voytenko, D.A., Sukachev, A.V., Ruchko, L.F. | III-301 |
| E15 | The effect of poloidal phasing of ICRF antennae on wave excitation | Alava, M.J., Heikkinen, J.A. | III-305 |
| E16 | The effect of realistic antenna geometry on plasma loading predictions | Ryan, P.M., Baity, F.W., Batchelor, D.B., Goulding, R.H., Hoffmann, D.J., Tolliver, J. | III-309 |
| E17 | Current drive by EC-waves in stellarators | Castejón, F., Coarasa, J.A., Alejaldre, C. | III-313 |
| E18 | ECRH produced start-up plasmas in RTP | Polman, R.W., van Lammeren, A.C.A.P., Lok, J., et al., RTP-Team | III-317 |
| E19 | Sawtooth stabilization by electron cyclotron heating near $q = 1$ surface in the WT-3 Tokamak | Tanaka, S., Hanada, K., Tanaka, H., Iida, M., Ide, S., Minami, T., et al. | III-321 |
| E20 | High frequency ion Bernstein wave heating experiments with minority and neutral beam heating on JIPP T-IIU tokamak | Kumazawa, R., Ono, M., Seki, T., Yasaka, Y., Watari, T., Shinbo, F. et al. | III-325 |
| E21 | Tuning method for multiple transmission lines with mutually coupled fast wave antennas in JFT-2M | Kazumi, H., Yoshioka, K., Kinoshita, S., Yamamoto, T., Petty, C.C., Saegusa, M. | III-329 |
| E22 | The generation of harmonics and the coupling between MHD activity and fundamental frequency during Alfvén wave heating in TCA | Borg, G.G., Duperrex, P.A., Lister, J.B. | III-333 |</p>
<table>
<thead>
<tr>
<th>E24</th>
<th>Density fluctuations and particle confinement during OH/LHCD on tokamak CASTOR</th>
</tr>
</thead>
<tbody>
<tr>
<td>E25</td>
<td>The effect of lower hybrid current drive on the penetration behaviour of test</td>
</tr>
<tr>
<td></td>
<td>impurities</td>
</tr>
<tr>
<td></td>
<td>Hildebrandt, D., Pursch, H., Weixelbaum, L., Jakubka, K.</td>
</tr>
<tr>
<td>E26</td>
<td>Power deposition profile during lower hybrid current drive in TORE SUPRA</td>
</tr>
<tr>
<td></td>
<td>Pecquet, A.-L., Hubbard, A., Moreau, D., Moret, J.H., Fall, T., et al.</td>
</tr>
<tr>
<td>E27</td>
<td>Lower hybrid wave coupling in TORE SUPRA through multijunction launchers</td>
</tr>
<tr>
<td>E28</td>
<td>Analysis and simulations of lower hybrid current drive in mixed OH-LH</td>
</tr>
<tr>
<td></td>
<td>discharges in TORE SUPRA</td>
</tr>
<tr>
<td>E29</td>
<td>ECCD experiments on T-10</td>
</tr>
<tr>
<td>E30</td>
<td>Fokker-Planck analysis of ECCD experiments in DIII-D</td>
</tr>
<tr>
<td></td>
<td>Giruzzi, G., James, R.A., Lohr, J.</td>
</tr>
<tr>
<td>E31</td>
<td>110 GHz ECH system for DIII-D</td>
</tr>
<tr>
<td>E32</td>
<td>Measurement of the ICRH power absorption from modulation experiments in TEXTOR</td>
</tr>
<tr>
<td></td>
<td>Lebeau, D., Koch, R., Messiaen, A.N., Vandenplas, P.E.</td>
</tr>
<tr>
<td>E33</td>
<td>Neutron yield during ICRH and NBI modulation experiments in TEXTOR</td>
</tr>
<tr>
<td></td>
<td>Lebeau, D., Van Wassenhove, G., Delvigne, T., Hoenen, F., Sauer, M.</td>
</tr>
<tr>
<td>E34</td>
<td>ICRH-mode coupling and heating in BPX and JET</td>
</tr>
<tr>
<td></td>
<td>Scharer, J.E., Lam, N.T., Bettenhausen, M.</td>
</tr>
<tr>
<td>E35 (Oral)</td>
<td>Electron heating in JET by ICRH</td>
</tr>
</tbody>
</table>
E36  Fast electron dynamics during LHCD in JET  

E37  Lower hybrid current drive experiments on JET  

E38 (Oral)  Role of parametric decay instabilities and edge plasma fluctuations on current drive efficiency of lower hybrid waves  
Pericoli-Ridolfini, V., Cesario, R.

E39  Transport analysis of LHCD driven plasmas in ASDEX  
Parail, V.V., Pereverzev, G.V., Soldner, F.X.

E40  Combined operation of lower hybrid and neutral beam injection on ASDEX  

E41  Evidence for nonlinear coupling of the lower hybrid grill in ASDEX  
Leuterer, F., Soldner, F.X., Giannone, L., Schubert, R.

E42  ICRF power deposition and confinement scaling in ASDEX  
Ryter, F., Stroth, U., Brambilla, M., ICRH-Group, ASDEX-Group, NI-Group

E43  Spatial diffusion of fast electrons during the 2.45 GHz experiment on ASDEX  
Barbato, E., Bardiromo, R., Gabellieri, L., Tuccillo, A.A.
<table>
<thead>
<tr>
<th>Page</th>
<th>Title</th>
<th>Authors</th>
</tr>
</thead>
<tbody>
<tr>
<td>F13</td>
<td>Effect of LHCD on tearing mode instability</td>
<td>Sheng, Z.M., Xiang, N., Yu, G.Y.</td>
</tr>
<tr>
<td>F14</td>
<td>Magnetic reconnection in collisionless plasmas</td>
<td>Coppi, B., Detragiache, P.</td>
</tr>
<tr>
<td>F15</td>
<td>Mathematical models displaying the gross features of sawtooth oscillations in tokamaks</td>
<td>Haas, F.A., Thyagaraja, A.</td>
</tr>
<tr>
<td>F16</td>
<td>Sawteeth stabilization by energetic trapped ions</td>
<td>Samain, A., Edery, D., Garbet, X., Roubin, J.-P.</td>
</tr>
<tr>
<td>F17</td>
<td>The effect of the ion bounce resonance on ideal and resistive MHD instabilities</td>
<td>Romanelli, F., Fogaccia, G., Graziadei, S.</td>
</tr>
<tr>
<td>F18</td>
<td>Theoretical calculations and experimental comparisons for high-n toroidal instabilities and quasilinear fluxes</td>
<td>Rewoldt, W., Tang, W.M.</td>
</tr>
<tr>
<td>F19</td>
<td>Effects of current profiles on MHD stability</td>
<td>Lao, L.L., Taylor, T.S., Chu, M.S., Turnbull, A.D., Strait, E.J., Ferron, J.R.</td>
</tr>
<tr>
<td>F20</td>
<td>Interpretation of resonant magnetic perturbation experiments</td>
<td>Hender, T.C., Fitzpatrick, R., Morris, A.W., Haynes, P.S., Jenkins, I. et al.</td>
</tr>
<tr>
<td>F21</td>
<td>Marfe stability</td>
<td>Hender, T.C., Wesson, J.A.</td>
</tr>
<tr>
<td>F22</td>
<td>Free boundary toroidal stability of ideal and resistive internal kinks</td>
<td>Vlad, G., Lutjens, H., Bondeson, A.</td>
</tr>
<tr>
<td>F24</td>
<td>Dispersion relations for global Alfvén modes with m &gt;&gt; 1</td>
<td>Cheremnykh, O.K., Revenchuk, S.M.</td>
</tr>
</tbody>
</table>
Stabilization effects on toroidicity-induced shear Alfvén waves in tokamaks
Berk, H.L., Guo, Z., Lindberg, D.M., Van Dam, J.W., Rosenbluth, M.N.

Stochastic nature of ICRF wave-particle interaction
Helander, P., Lisak, M., Anderson, D.

The influence of electrons heating on the magnetoplasma surface waves propagation at the plasma-metal structure
Azarenkov, N.A., Ostrakov, K.N.

On existence of solitary drift waves in the presence of inhomogeneous electric field
Smirnov, A.P., Sheina, E.A.

Wave propagation in an inhomogeneous relativistic magnetoplasma
Kerkhof, M.J., Kamp, L.P.J., Sluijer, F.W., Weenink, M.P.H.

Influence of plasma density inhomogeneities on the Eigenfrequency of global three-dimensional MHD modes in toroidal plasmas
Cap, F.

Anomalous dispersion of electron-cyclotron-waves on non-Maxwellian, relativistic plasmas
Moser, F., Räuchle, E.

A multiple timescale expansion and anomalous plasma transport
Edenstrasser, J.W.

Electron heat conduction and suprathermal particles
Bakunin, O.G., Krasheninnikov, S.I.

Fast poloidal rotation and improved confinement
Rozhansky, V., Samain, A., Tendler, M.

Net transport equations for a tokamak plasma
Callen, J.D., Hollenberg, J.B.

Fluid-kinetic model of tokamak transport
Feneberg, W., Kerner, W.
<table>
<thead>
<tr>
<th>Number</th>
<th>IV-145</th>
<th>Title</th>
<th>Authors</th>
</tr>
</thead>
<tbody>
<tr>
<td>F38</td>
<td>Neoclassical current and plasma rotation in a helical systems</td>
<td>Nakajima, N., Okamoto, M.</td>
<td></td>
</tr>
<tr>
<td>F39</td>
<td>Numerical solution of the drift-kinetic Vlasov equation</td>
<td>Ghizzo, A., Shoucri, M., Bertrand, P., Feix, M., Fijalkow, E.</td>
<td></td>
</tr>
<tr>
<td>F40</td>
<td>The integrals of drift particles motion with finite Larmor radius</td>
<td>Ilgisonis, V.I.</td>
<td></td>
</tr>
<tr>
<td>F41</td>
<td>Drift effects in the theory of magnetic islands</td>
<td>Smolyakov, A.I.</td>
<td></td>
</tr>
<tr>
<td>F42</td>
<td>Influence of magnetic islands on magnetic field line diffusion</td>
<td>Martins, A.M., Mendonca, J.T.</td>
<td></td>
</tr>
<tr>
<td>F43</td>
<td>Computation of plasma equilibria with semifree and free boundary</td>
<td>Nicolai, A.</td>
<td></td>
</tr>
<tr>
<td>F44</td>
<td>Study of density and temperature profile across the magnetic island</td>
<td>Qingquan Yu, Yuping Huo</td>
<td></td>
</tr>
<tr>
<td>F45</td>
<td>Two-fluid theory of presheath and sheath including magnetic fields</td>
<td>Valentini, H.-B.</td>
<td></td>
</tr>
<tr>
<td>F46</td>
<td>The Grad-Shafranov shift calculated on the basis of magnetic compressive and tensile stresses</td>
<td>Jensen, V.O.</td>
<td></td>
</tr>
<tr>
<td>F47</td>
<td>Profile consistency as a result of coupling between the radial profile functions of pressure and current density</td>
<td>Schüller, F.C., Schram, D.C., Konings, J., van Lammeren, A.C.A.P. et al.</td>
<td></td>
</tr>
<tr>
<td>F48</td>
<td>Impurity control by the radial electric field in a stochastic layer</td>
<td>Nguyen, F., Samain, A., Ghendrih, Ph.</td>
<td></td>
</tr>
<tr>
<td>Session</td>
<td>Title</td>
<td>Authors</td>
<td></td>
</tr>
<tr>
<td>---------</td>
<td>----------------------------------------------------------------------</td>
<td>---------</td>
<td></td>
</tr>
<tr>
<td>F49</td>
<td>Nonlocal dielectric response of a toroidal plasma</td>
<td>Lamalle, P.U.</td>
<td></td>
</tr>
<tr>
<td>F50</td>
<td>State diagrams of tokamaks and scaling laws</td>
<td>Minardi, E.</td>
<td></td>
</tr>
<tr>
<td>F51</td>
<td>Neoclassical poloidal flow bifurcation in the H mode transition</td>
<td>Lazzaro, E., Lucca, F., Nardone, C., Tanga, A.</td>
<td></td>
</tr>
<tr>
<td>F52</td>
<td>On the existence and uniqueness of dissipative plasma equilibria in a toroidal domain</td>
<td>Spada, M., Wobig, H.</td>
<td></td>
</tr>
<tr>
<td>F53</td>
<td>L-H transitions via the Matsuda anomaly</td>
<td>Puri, S.</td>
<td></td>
</tr>
<tr>
<td>F54 (Oral)</td>
<td>The role of Pfirsch-Schlüter currents in plasma equilibrium, stability and transport</td>
<td>Wobig, H.</td>
<td></td>
</tr>
<tr>
<td>F55</td>
<td>Fusion alpha particle transport studies using energy dependent diffusion coefficients</td>
<td>Kamelander, G.</td>
<td></td>
</tr>
</tbody>
</table>
### G DIAGNOSTICS

| G1 | First results of ECRH transmitted power measurements on RTP  
*Smits, F.N.A., Bank, S.L., Bongers, W.A., Oomens, A.A.M., Polman, R.W., RTP-Team* |
| G2 | Measurements of plasma potential in T-10  
| G3 | Twin E-mode reflectometry for magnetic field measurements in Tokamaks  
*Lazzaro, E., Ramponi, G.* |
| G4 | Magnetic field measurements at JET based on the Faraday and motional Stark effects  
| G5 | The correlation of magnetic flux surfaces with soft X-ray iso-emissivity surfaces in COMPASS-C  
*R.D. Durst, Haynes, P.* |
| G6 | Impurity atoms diagnostic by observation of near-resonant Rayleigh scattering  
*Berlizov, A.B., Moskalenko, I.V., Shcheglov, D.A.* |
| G7 | Investigations of light impurities transport in tokamak using small-view optical tomography  
*Kuteev, B.V., Ovsishcher, M.V.* |
| G8 | Measurement of gas injection efficiency for helium, neon and argon impurities in TEXTOR  
| G9 | Neutral density in FT ohmic plasma  
*Bracco, G., Noleti, A., Zanza, V.* |
| G10 | First results from the JET time of flight neutral particle analyser  
*Corti, S., Bracco, G., Noleti, A., Zanza, V.* |
| G11 | Comparison between Rutherford scattering and neutral particle analysis in ohmic discharges in TEXTOR  
G12 A new diagnostic for the tritium phase of JET covering the visible and UV wavelength range

G13 Alpha-particle diagnostics for the D-T phase

G14 Neutron production during deuterium injection into ASDEX
Bomba, B., ASDEX-Team, NI-Team, Feng, Y., Hübner, K., Wolle, B.

G15 Observation of velocity dependence and line-of-sight effects in ion temperature and toroidal rotation velocity measurements at JET
Danielsson, M., Källne, E., Zastrow, K.-D., von Hellermann, M., Mandl, W., et al.

G16 Ti(r) profiles from the JET neutron profile monitor for ohmic discharges

G17 Ion temperature and fuel dilution measurements using neutron spectroscopy

G18 On the opportunity to measure the plasma ion temperature by a photoelectron method
Gott, Yu.V., Shurygin, V.A.

G19 Measurement of ion temperature profiles in the TCA tokamak by collective Thomson scattering

G20 Potential of millimeter-wavelength collective scattering in high field tokamaks
Tartari, U., Lontano, M.

G21 Ion temperature profiles deduced from Doppler broadening of X-ray lines in ASDEX
Chu, C.C., Nolte, R., Fussmann, G., Fahrbach, H.U., et al., ASDEX-Teams

G22 Study of plasma turbulence in the TJ-I Tokamak by a spectroscopic technique
Zurro, B., and TJ-I Team

G23 Simultaneous measurement of 3 fluctuating plasma parameters
Carlson, A., Giannone, L., ASDEX Team
G24  Pulsed radar; a promise for future density profile measurements on thermonuclear plasmas
Heijnen, S.H., Hugenholtz, C.A.J., Pavlo, P.

G25  Correlation reflectometry techniques for TJ-I and ATF

G26  Measurement of fast changes of the edge density profile in TEXTOR
Gunkel, H., Kuszyński, J.O., Pospieszczyk, A., Schwer, B.

G27  Reflectometric diagnostics of plasma density fluctuations in TUMAN-3 tokamak

G28  Plasma turbulence studying on the T-10 by microwave reflectometry
Zhuravlev, V.A., Dreval, V.V.

G29  A fife-camera X-ray tomography system for the RTP Tokamak
da Cruz, D.F., Donn, A.J.H.,

G30  Study of q-profile in LHCD regimes with microwave reflectometry
Silva, A., Manso, M.E., Söldner, F.X., Zohm, H., Serra, F.

G31  Density fluctuation profiles on TORE SUPRA

G32  Electron temperature fluctuations in RTP

G33  TORE-SUPRA X-ray pulse-height analyzer diagnostic

G34  Acceleration of electrons during current increase in the "TUMAN-3" device
Afanes'ev, V.I., Its, E.R., Kiptiliy, V.G., ... Rozhdestvensky, V.V., et al.

G35  ECE measurements using Doppler-shifted observations
Rodriguez, L., Aug, N., Giruzzi, G., Javon, C., Laurent, L., Talvard, M.

G36  First results with the upgraded ECE heterodyne radiometer on JET
Porte, L., Bartlett, D.V., Campbell, D.J., Costley, A.E.
Measurement of $T_e$-profiles in the boundary layer of TEXTOR by means of spectroscopical observation of a thermal helium beam

Schweer, B., Pospieszczyk, A., Hank, G., Samm, U., Brosda, B., Pohlmeyer, B.

Model calculations for a 20 keV neutral lithium diagnostic beam

Unterreiter, E., Aumayr, F., Schorn, R.P., Winter, H.
HIGH THERMONUCLEAR YIELD ON JET BY COMBINING ENHANCED PLASMA PERFORMANCE OF ICRH-HEATED, PELLET-PEAKED DENSITY PROFILES WITH H-MODE CONFINEMENT


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Introduction

The main features of the pellet enhanced performance (PEP) regime in JET have previously been described [JET Team, presented by G.L. Schmidt, Proc. 12th IAEA Conf., Nice 1988]. In these plasmas a strongly peaked density profile was created by injection of deuterium pellets immediately after the current rise. The subsequent additionally heated phase exhibits a strongly increased D-D neutron production rate and improved central confinement, as compared to non-pellet-enhanced discharges. These previous results were obtained in limiter deuterium plasmas with otherwise typical L-mode characteristics. We have now found that the PEP-mode can be combined with H-mode confinement giving higher plasma performance, and particularly, higher thermo-nuclear yield [initial results in B. Tubbing et al., accepted by Nuclear Fusion]. Furthermore, from measurements carried out on these shots a possible explanation for the nature of the PEP features begins to emerge.

Experimental evidence

The PEP H-mode experiments were carried out in a double-null X-point configuration, with plasma currents between 3 and 3.6 MA and toroidal fields between 2.8 and 3.2 T. A typical example is shown in fig. 1. The X-point configuration is formed immediately after the end of the plasma current rise. A 4 mm pellet is injected soon after and well before the onset of sawteeth. The pellet creates a peaked density profile with an initial central value of about $1.6 \times 10^{20} \text{ m}^{-3}$. The pellet injection is followed by additional heating on a level of about 8 to 10 MW of ion cyclotron resonance heating (10-15% hydrogen minority and axial deposition) and 2.5 MW of 80 kV neutral beam (for diagnostic purposes). This leads, in less than 1 second, to equal central electron and ion temperatures of about 11 keV, a central density of about $7 \times 10^{20} \text{ m}^{-3}$ and a central electron pressure up to about 1.2 bar at the time of the maximum D-D neutron rate of $1 \times 10^{15} \text{ s}^{-1}$. 80% of the neutron rate is of thermonuclear origin. This is clearly the highest observed thermonuclear neutron rate on JET for plasmas with $T_e = T_i$. The maximum value of the fusion product $n_e(0) \tau_e(0) T_e(0)$ is

This work has been performed under a collaboration agreement between the JET Joint Undertaking and the US Department of Energy.
7*10^{25} \text{ m}^{-3} \text{s}^{-1} \text{keV} and is one of the highest seen on JET. After about 0.5 seconds an L to H transition takes place, as seen from the typical signature of the edge D_\alpha light and the total plasma energy of 7-8 MJ. The plasma is in PEP H-mode for about 0.5 s, then the PEP-mode terminates and the plasma adopts ordinary H-mode behaviour. The discharge is not saw-toothing before or during the PEP-phase, nor in the subsequent H-mode. In fig. 2 the peak neutron production rate of L- and H-mode plasmas with and without PEP-mode is plotted versus plasma energy, demonstrating that the new PEP H-modes are typically a factor of 5 better than the ordinary H-modes. They extend the trend curve of neutron production rate by a factor of 2 as compared to the limiter PEP pulses. In fig. 3 the normalised plasma energy content is plotted against the loss power for a similar selection of shots, at the time of maximum energy. The lines indicate one or two times Goldston confinement scaling. The new PEP data in this figure contains both discharges with clear PEP H-mode signatures and discharges in which the H-mode signature is less clear (PEP L- or PEP elm/H-modes). The figure shows that the global energy confinement of the good PEP H-modes is comparable to or slightly better than that of ordinary H-modes. The figure further shows (solid triangles) that the confinement of the H-mode that remains after the decay of the PEP is similar to that of ordinary H-modes. - In a number of shots it has been shown that the onset of the H-mode can be significantly advanced in time by an initial boost in power (here, neutral beams) without much influencing the PEP-mode performance; however, these attempts also see the carbon influx due to local overheating of the X-point carbon tiles at earlier times. - Although most of the experiments were performed with central ICRH deposition it was shown that PEP H-modes can also be produced with neutral beam injection. - During PEP-modes several MHD phenomena can occur, which can affect the central parameters and the neutron rate. Hesitations in the slope of the neutron signal have been associated with modes with toroidal mode numbers 2, 3 and 4. In some cases an MHD mode with mode numbers m,n = 1,1 is observed to start coinciding with the neutron decay. Often the PEP-phase is terminated by a fast MHD event but it can also decay more smoothly. In one case (#23100) the post-cursor oscillations following the 1,1 mode corresponds to two 4-8 cm wide non-overlapping rotating island chains with mode numbers m,n = 3,2 and 2,2 and radially located at r = 15 and 30 cm, respectively. The existence of these modes suggests a non-monotonic q-profile with one or even two q = 1 surfaces. The equilibrium code IDENTD, in which the radial position of the q = 1 and q = 1.5 rational surfaces can be imposed, can produce a solution consistent with the measured pressure profile and magnetic fluxes at the vessel. Detailed evaluation
The confinement features found on L-mode and H-mode PEP-pulses alike suggest that the PEP mode is more or less independent of the state of the background plasma it is superimposed upon. During the overlapping period of the PEP- and H-mode, local transport calculations using the FALCON and TRANSPI codes have shown central values (r ≤ 0.4 a) of the electron diffusion coefficient D_e = 0.1 m^2s^{-1} and the effective heat conduction coefficient \chi_{eff} = 0.5 m^2s^{-1}. (These central values are typical also for the old limiter PEP L-mode). Outside this region \chi_{eff} values are found characteristic for H-modes, with typical values of 1-2 m^2s^{-1} as shown in fig. 4. The code calculations also show consistency with the thermonuclear neutron production rate. In many cases it is observed that the neutron rate decays significantly before the termination of the PEP-phase. This can be
explained by three possible causes or their combination: decaying central density, dilution with impurities which accumulate during the good confinement, and in some cases degrading confinement caused by MHD activity. The higher neutron rates in the PEP H-modes, as compared to the old PEP L-modes, may result from the higher ion temperatures. These may be due to the better confinement of the H-mode. However, the PEP H-mode experiments were also conducted with a better ion heating efficiency of the ICRH than the PEP L-mode experiments. This is due to operation at higher minority concentration (below 5% with about 30% of the ICRH power coupled to the ions in the old experiments, and 10-15% with about 50% ion heating in the recent experiments).

Interpretation and conjectures

In this section we attempt to draw a reasonably consistent, but speculative, picture of the PEP-mode and the reasons, why this confinement deteriorates and often ends in a spectacular crash in central electron and ion temperature as well as neutron rate. Earlier work has been performed considering ballooning [Galvao, 1988] or infernal modes [e.g. Charlton, submitted to Nuclear Fusion], due to the high pressure gradients, to be responsible for the observed MHD phenomena. However, in many cases PEP-modes display similar MHD phenomena without having reached high levels of neutron rate or pressure. This suggests that the current density, rather than the pressure, might be the dominant factor responsible for the stability. As mentioned, the existence of a q=1.5 surface inside a q=1 surface has been diagnosed; therefore a region of negative shear (dq/dr<1) exists. This non-monotonic q-profile can be maintained during the PEP by the bootstrap current which is expected due to the steep radial density profile with dn/dr = 5 to 1*10^20 m^-2. Calculations of the bootstrap current densities indicate a value of about 1 MAm^-2 in the region r = a/4, comparable to the ohmic current density. The initial onset of shear inversion may result from the hollow temperature profile created by the central deposition of the pellet, and by subsequent rapid current penetration at low electron temperature. Then, as the temperature reaches higher values due to the additional heating and the improved confinement, a large bootstrap current develops enhancing the q profile inversion, and possibly driving q below 1 off-axis. It has been speculated that the enhanced central confinement is associated with the reversal of the shear. Simulations using the Rebut-Lallia critical temperature gradient model outside the negative shear region and neoclassical transport inside it show qualitative agreement with the experiment, in the time window between the pellet and the onset of MHD phenomena. These simulations were done either with \(\chi_e=\chi_{e,neo}, \chi_i=\chi_{i,neo}\) or with \(\chi_e=\chi_{i,neo}\) inside the negative shear region. Accepting this interpretation of the PEP mode one can devise experiments aimed at increasing its duration and performance. These would be, in the first place experiments with higher axial q during the PEP, to be achieved by operating at higher toroidal field, by non-inductive current drive or by injection of the pellet sooner after the current rise (at lower internal inductance, 1_i). Secondly, one could attempt to delay the development of the q-profile by increasing the electron temperature as soon as possible after the pellet. Thirdly, experiments on control of the bootstrap current, by influencing the density profile (neural injection or pellet deposition control) could be carried out. Much work has still to be done to understand particularly the build-up of the PEP-mode.
Fig. 1: Pulse history of # 22490

Fig. 2: Neutron rate vs plasma energy

Fig. 3: Plasma energy vs power loss

Fig. 4: Heat conductivity vs radius
THE EFFECTS OF PARTICLE DRIFT ORBITS ON FLUX DEPOSITION PROFILES AT THE JET X-POINT TARGET

JET Joint Undertaking, ABINGDON UK.

INTRODUCTION

Surface temperature and particle flux distributions at the X-point target in JET are measured using CCD cameras. Throughout ohmic and L-phases of X-point discharges, the expected two strike zones are observed, corresponding to particle flows in the ion- and electron-drift directions. During H-modes, an additional strike zone is observed at the ion-drift side of the target [1,2]. Evidence is presented indicating that this additional zone is due to a flux of ions with energies of several keV. This interpretation is supported by simulations of ion drift orbits, which also indicate that this flux comes directly from the bulk plasma.

EXPERIMENT

The X-point target in JET consists of 32 bands of tiles. Each band follows the poloidal form of the vessel and in the toroidal direction the tiles are curved so that particle flows from opposite toroidal directions impact on opposite sides of the tile apex. Three bands of tiles at the carbon upper target were studied using two cameras. One of these was fitted with a 1µm wavelength filter to measure temperature distributions, and the other was equipped with a carousel of filters, allowing the distributions of Ha(656.3nm), CI(909.5nm), CII (514.5 nm) and CIII(464.7nm) to be studied. This study was made during 3MA single-null X-pt discharges with NBI heating. Data were taken with the toroidal field, Bt, in the normal direction (ion grad B drift towards the target) and with Bt reversed (ion grad B drift away from the target). The strike zones observed under these conditions are shown schematically in Fig. 1.

RESULTS AND DISCUSSION

The additional strike zone at the ion-drift side (marked B in Fig. 1) was detected during H-modes. The effect was observed with Bt in either direction, although, as Fig 1 shows, the deposition profiles were different. Intense Ha and Carbon spectral emission was observed at A and C, and the sources of CII and CIII emission extended toroidally into the region between bands of target tiles. At B the spectral signals increased rapidly but remained localised until the surface temperature exceeded ~2000°C and persisted after the NBI heating period, implying a thermal source. The effect is illustrated in Fig. 2, which shows radial profiles of CII intensity at the ion-drift side of the target during a discharge with reversed Bt. In the first trace, showing the profile before NBI heating was applied, only the broad strike zone C at the ion-drift side is present, at a mean radius of ~265cm. The next trace was taken 400msec later, shortly after the onset of the H-mode. Although the position of zone C has shifted slightly, the intensity has remained almost constant. At R~255cm, corresponding to zone B, a narrow additional strike zone may be seen, but at this time the intensity is very low. The third trace, taken after a further 400msec shows that the intensity at zone B has increased, and is now higher than at C. Data from the other camera showed rapid surface heating at A and B, with the temperatures rapidly increasing until the termination of the H-mode.
Despite strong recycling and sputtering at C, no surface heating was detected, suggesting a high flux of low-energy particles at this point. Conversely, at B, rapid heating occurred but little apparent sputtering, implying a weaker flux of more energetic particles. At A both sputtering and heating were observed, indicating the presence of both effects.

During pulses in which the X-point to target distance, Dx, was continuously varied during the H-mode period the radial separation between zones B and C, was found to increase as Dx decreased, reaching a maximum when Dx~0. Also, during these discharges, the Ha intensity at zone B was lower (by at least 10x) than at zones A and C. This was taken as further evidence of a low-density, high energy flux at B.

An array of Langmuir probes in another band of target tiles was used to obtain profiles of the electron temperature and ion saturation current at the target with a resolution of ~6cm in the radial direction [3]. Data from a number of discharges showed that although the profiles changed after the transition to H-mode, the electron temperature in the scrape-off layer never exceeded 100ev.

**ION DRIFT ORBIT SIMULATIONS**

The computer code "Orbit" calculates the poloidal drift orbit of an individual ion launched with defined velocity components at any position within the JET vacuum vessel. It uses real JET magnetic equilibrium data, but ignores the effects of collisions and electric fields. For X-point configurations the code includes the grad B drift of ions which approach the poloidal field null. Using data from an H-mode discharge in this code ions were launched parallel and anti-parallel to the plasma current from positions near the midplane outer boundary of the plasma.

The results for ions launched in the ion-drift direction showed that those starting from outside the last closed flux surface (LCFS) strike the target at a position which depends on their energy, with low-energy ions arriving closer to the X-point radius, Rx. The radial dispersion in energy was greatest when the separation between X-point and target was <3cm. For any particular energy, ions launched from larger radii reached the target at a greater distance from Rx.

The orbits of almost all ions launched inside the LCFS were confined within the bulk plasma. However, ions of a few kev energy which were launched very close to the LCFS could cross this surface and strike the target at positions close to Rx. In addition the minimum launch radius, Rmin, from which such ions reach the target was found to be inversely dependent on their energy, so that ions with highest energies reach the target from points deepest in the plasma and arrive closest to Rx.

For ions launched in the electron-drift direction, there is less radial dispersion in energy, so that kev ions launched from just inside the LCFS arrive at similar radii to lower energy ions launched outside, and only a single strike zone is formed. The value of Rmin found for each particular ion energy was not the same for ions launched in this direction as for those in the opposite direction.

Another feature of the code is the possibility to artificially reverse the direction of Bt without changing the poloidal field configuration. Reversing Bt caused the positions of strike zones to be reversed, so that the ion-drift side of the target changed from the inner to the outer strike zone (Fig 1.) It was found that the value of Rmin was lower in each case for ions reaching the target at the ion-drift side. This implies that the energetic ion flux at the outer strike zone is greater for normal Bt direction than for reversed Bt, and that the converse is true for the inboard strike zone. This agrees well with experimental observations [4].
In Fig 3 are plotted ion orbits near the target computed for a reversed Bt discharge, for which the ion-drift side corresponds to radii greater than $R_x$. The orbits shown are of 5kev and 50ev ions launched from the Rmin value appropriate to each. At the ion-drift side a clear separation may be seen between the strike points for the two ion energies, while at the electron-drift side the two ion species arrive at the same point. In this figure the X-point to tile distance is ~18cm, and it is clear that the separation between the orbits is greatest in the region close to the X-point. Even in this region, however, there is a smaller separation at the electron-drift side.

**SUMMARY AND CONCLUSIONS**

CCD cameras have been used to study the additional strike zone which occurs in JET during H-modes. The low hydrogen and impurity spectral emission measured in this region, and the rapid surface heating suggest that it is caused by high-energy particles. From its position relative to the other strike zones and the low measured values of local electron temperature a flux of high-energy ions is inferred. Further support for this conclusion is provided by recent measurements of high edge ion temperature during JET H-modes [5].

Calculations of ion drift orbits indicate that energetic ions can flow to the target directly from the outer regions of the plasma. Consideration of these ion orbits provides an explanation of the increase in power flow to the target during H-modes and for each toroidal field direction the calculations qualitatively agree with the observations. This implies that the total power flow to the target is determined by the ion temperature profile, and that the additional strike zone is another indication of high ion temperatures near the plasma edge. It is worth noting, that since the onset of an H-mode does not cause gross changes in magnetic equilibrium, the additional deposition zone would be observed during L-modes if the edge ion temperature were high enough.

**REFERENCES**

FIG 1.
Schematic view of upper X-point target from below showing the observed strike zone positions during H-modes with forward and reversed B_T.

FIG 2.
CII intensity profiles at the target ion-drift side during reserved B_T discharge (#22266)

FIG 3
Calculated ion orbits for a reversed B_T discharge (#22269). At the inner (electron-drift) side, the orbits of 50ev and 5kev ions are similar, while at the ion-drift side the orbits are widely separated.
CONFINEMENT OF HIGH PERFORMANCE JET PLASMAS

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I. Introduction

The energy confinement properties of different regimes of JET plasmas have been investigated and compared with several theoretical models [1]. Some of these models, such as the Rebut-Lallia critical temperature gradient model and the neoclassical pressure gradient driven turbulence, give reasonable agreement with the radial behaviour of the measured heat flux for the high recycling L-modes in JET. However, they have difficulty in reproducing the better confinement properties of high performance JET plasmas such as the hot ion H-modes or the pellet enhanced phase L and H-modes. These latter regimes are of particular interest for future reactors as the highest $n_D \tau_B T_I$ products as well as record neutron yields are achieved in these circumstances.

The aim of this paper is to construct a gyro-Bohm model which fits the effective thermal conductivity obtained from the $1^{1/2}D$ transport analysis code TRANSP [2]. The measured profiles of temperature, density, $Z_{\text{eff}}$ as well as many other quantities are input data to TRANSP. Checks on the consistency of the data are made by comparing the predicted and measured value of such quantities as diamagnetic energy, fast ion energy (produced by ICRH or NBI), surface loop voltage, total neutron yield and neutron emission profiles. The effective thermal conductivity is defined by: $\chi_{\text{eff}} = (\text{heat flux})/(n_e \nabla T_e + n_i \nabla T_i)$

$\chi_{\text{eff}}$ is known to within 20-30%; however the error increases with radius.

II. Model for $\chi_{\text{eff}}$ and physics constraints

Global scaling laws which both satisfy the Connor-Taylor constraints of plasma physics [3] and give a good fit to the data have been found for both L and H mode operation [4,5]. A similar analysis has also been done for local JET L-mode data [6] to try to derive a local scaling law. All these studies conclude that the constraint imposed by the high $\beta$ - Fokker-Planck equation is satisfied. Here a similar power law model is used:

$$\chi_{\text{eff}} = C \rho v_e \left( \frac{P}{r} \right)^{\chi_1} F \left( \nu_e, \beta_p, \frac{r}{R}, q, \frac{T_i}{T_e}, \frac{L}{R} \right)$$

where $C$ is a constant, $\rho$ is the poloidal electron Larmor radius and $v_e$ is the electron thermal velocity.
The high $\beta$-Fokker Planck constraint is satisfied by that expression for $\chi_{\text{eff}}$. Short wavelength turbulent fluctuations correspond to $x_1 = 1$ and long wavelength turbulent fluctuations to $x_1 = 0$.

The function $F$ whose arguments are dimensionless depends on the ion-electron collisionality $\nu_e$, beta poloidal $\beta_p$, local inverse aspect ratio $r/R$, safety factor $q$, ratio $T_i/T_e$ and $L/R$ where $L$ is a profile scale length for density temperature, shear etc. defined by $L_n = -n/V_n e$ etc.

$$F = \nu_e^{x_2} \beta_p^{x_3} \left( \frac{r}{R} \right)^{x_4} q^{x_5} \left( \frac{T_i}{T_e} \right)^{x_6} \left( \frac{L}{R} \right)^{x_7} \ldots$$

III The different JET regimes

Several pulses are included in this study and their confinement properties have previously been analysed [1,7]: high recycling L-mode, hot ion L and H-modes, hot electron plasmas and sawtooth free discharges achieved with ICRH. More recent results are also included: the pellet enhanced phase of L and H-modes where a very peaked density profile, with $n_e(0) \sim 10^{20} \text{m}^{-3}$ formed by pellet injection is heated using either ICRH alone or in combination with NBI (the PEP mode) [8]. Thus a local database covers a wide range of plasma parameters and profile features: from very flat density profiles for the H-modes to very peaked profiles for the PEP mode; very peaked ion temperature profiles, with $T_i(0) \sim 24 \text{ keV}$ achieved with high power D$^+$ beam injection at 140 keV.

The characteristic time for each shot is chosen to be representative of the steady state of a regime: the rate of ion (or electron) stored energy $\dot{w}$ is small compared with the conduction losses. The inner region of the plasma ($r/a \leq 1/3$) is excluded from the analysis if the uncertainties on the power deposition profile are too great.

The database is also limited in the outer region of the plasma ($r/a \geq 0.8$) due to the lack of experimental measurements. Each radial profile consists of 8 - 10 radial points all treated as independent observations. The data set comprises 51 observations (L-mode) and 39 observations (H-mode).

IV Analysis of the results

The best fits to the different groups of data (all pulses, L-mode alone, H-modes alone) are obtained with (1) and the density scale length $L_{ne} = -n_e/V_n e$. The other scale lengths dependences such as $L_T/R$ or $L_s/R$ play no significant role. The parameter values and their standard errors are shown in Table I for the different cases. If L and H modes are treated together, the fit to the data is poor ($\sigma = 37.7\%$) and the standard error of each parameter is significant. If a parameter is introduced to distinguish L from H modes and the data is combined into one scaling law, then $\chi_{\text{eff}}$ for the L-modes is larger than $\chi_{\text{eff}}$ for the H-modes by a factor 1.5 - 1.9; however the statistics remain poor ($\sigma = 34.3\%$). The L and H mode data have not so far been reconciled into a unique scaling law; this is also the case when global data is analysed.
Fig. 1 and Fig. 2 show the fits obtained with L-modes alone and H-modes alone respectively. The statistics are now better than in the previous cases (σ = 26.2% for L-modes, σ = 16.2% for H-modes). The radial behaviour of χ_{eff} for hot ion L and H-modes (Fig. 3) and for PEP L and H-modes (Fig. 4) is well reproduced with these fits.

<table>
<thead>
<tr>
<th>TABLE I</th>
<th>σ(%)</th>
<th>log C</th>
<th>ρ/r</th>
<th>ν</th>
<th>β_p</th>
<th>r/R</th>
<th>q</th>
<th>T_i/T_e</th>
<th>L_{ne}/R</th>
</tr>
</thead>
<tbody>
<tr>
<td>L and H modes</td>
<td>37.7</td>
<td>-3.65</td>
<td>0.33</td>
<td>0.49</td>
<td>-0.72</td>
<td>-0.2</td>
<td>-0.41</td>
<td>-0.46</td>
<td>-0.05</td>
</tr>
<tr>
<td>L and H (1) modes</td>
<td>34.3</td>
<td>-3.76</td>
<td>0.11</td>
<td>0.43</td>
<td>-0.37</td>
<td>0.6</td>
<td>-0.52</td>
<td>-0.17</td>
<td>0</td>
</tr>
<tr>
<td>L modes</td>
<td>26.2</td>
<td>-19.3</td>
<td>-1.7</td>
<td>-0.03</td>
<td>0.35</td>
<td>0.39</td>
<td>0.97</td>
<td>0.49</td>
<td>1.0</td>
</tr>
<tr>
<td>H modes</td>
<td>16.2</td>
<td>-38.1</td>
<td>-3.4</td>
<td>0.58</td>
<td>0.41</td>
<td>-4.87</td>
<td>0.73</td>
<td>0.15</td>
<td>0.16</td>
</tr>
</tbody>
</table>

(1) In case L and H-mode data is combined into one scaling law but with different multipliers C we obtain χ_{eff} (L) = (1.5 - 1.9) χ_{eff} (H)

V Conclusion

A unique power law model for χ_{eff} with the constraint imposed by the high β Fokker-Planck equation cannot so far reconcile the different confinement properties of L and H-modes. However separate laws can well reproduce the radial behaviour of χ_{eff} for L-modes or H-modes alone if the ratio T_i/T_e and the density scale length L_{ne} are taken into account.

References

FIG. 1. $X_{\text{eff}}$ versus fit for L-mode pulses

FIG. 2. $X_{\text{eff}}$ versus fit for H-mode pulses

FIG. 3. Radial profile of $X_{\text{eff}}$ from the fits and that derived by TRANSP for the hot ion L- and H-mode pulses

FIG. 4. Radial profile of $X_{\text{eff}}$ from the fits and that derived by TRANSP for the PEP L- and H-mode pulses
THE PERFORMANCE OF HIGH CURRENT BELT LIMITER
PLASMAS IN JET

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1. Introduction.

Experiments have been performed for plasma currents 5-7MA using the
beryllium belt as a limiter and with total input powers up to 30MW. The objective
of these experiments was the optimisation of fusion performance and central
temperatures at maximum L mode confinement. This paper describes experiments
on sawtooth suppression, electron heating at low density, ion heating in sawtooth
free discharges. Complementary experiments are described in a density regime more
appropriate to a reactor with emphasis on the density limit.

2. 7MA plasmas.

Operation at 7MA has been demonstrated for the first time on the beryllium
belt limiter. A number of other improvements have been made to the 7MA plasma
described in [1]. Firstly, strengthening of the vacuum vessel has permitted larger
shaping fields which have been optimised to make the plasma more elongated and
D shaped as shown in fig.1, and therefore raise \( q_\phi \) above 3. Secondly, improved
toroidal field coil cooling has made it possible to ramp the plasma current at
constant toroidal field. A strong aperture expansion scenario was developed such
that \( q_\phi = 4 \) could be passed early in the current ramp without instabilities arising
from unstable current profiles leading to disruption [2]. Thereafter \( q_\phi \) was kept
between 3 and 4 by programming the minor radius, elongation and triangularity,
and this allowed a faster 1MA/sec rise to 7MA as shown in figure 2. This scenario
was more economical in Volt-sec consumption and, even after the 3sec flat top
shown, 8 Volt-sec remained for further extension of the flat top duration. So far
JET has withstood 2 disruptions at 7MA caused by unexpected contact with inconel
components. In these disruptions an halo current up to 1MA was transferred from
plasma to vessel giving rise to a vertical force of 350 Tonnes.

3. Electron Heating and Sawtooth Suppression

In the discharge illustrated in fig.2, 6MW of ICRF applied during the rise gave
\( T_e(0) \sim 9keV \) with long sawtooth free periods sawteeth until the RF switched off at
6MA. The sawtooth period then decreased as the current continued to ramp to
7MA. In cases where the RF was kept on longer, continuous m = 1 modes were observed in the plasma centre suggesting only marginal sawtooth stabilisation at currents above 6MA. Sawtooth free plasmas were routinely achieved at 5MA when a similar scenario was employed. Examination of this data shows that the electron temperature increases approximately linearly with total power per particle up to $P_{tot}/n_e \sim 3 \times 10^{-19} \text{MWm}^3$. This is consistent with a simple prediction based on the Rebut-Lallia-Watkins model. However, for $P_{tot}/n_e$ in the range 3 – 6 $\times$ $10^{-19} \text{MWm}^3$ $T_e$ remains in the range 10-12 keV. This 'apparent' saturation is explained by a reduction in the central heating power density [3].

4. Ion heating

In these RF heated plasmas the ion temperature was low and so neutral beam heating was applied in order to heat the ions. At 5MA, high ion temperatures in the range 10-13 keV were obtained at low density, $n_e \leq 4 \times 10^{19} \text{m}^{-3}$, as illustrated in fig.3. The values of $n_{D_0}T_e \tau_e$ were low because of poor fuel concentration. In the example shown, the beryllium source increased abruptly following a sawtooth crash 0.3sec after the NBI was switched on. The D-D reaction rate fell despite high $T_e$ and $T_i$. At this same time the limiter viewing camera showed a 'blob' of material detaching from the belt limiter tiles.

At 6.5MA, 6MW of ICRF combined with 8MW of NBI produced a sawtooth free period 0.8 sec long with $T_{i0} \sim 7.5 \text{keV}$, $T_{e0} \sim 8.0 \text{keV}$, $n_{e0} \sim 5 \times 10^{19} \text{m}^{-3}$ and $\tau_e \sim 0.65 \text{sec}$. In this case the central value of fuel concentration was high $n_{D_0}/n_e \sim 0.88$ giving $n_{D_0}T_e \tau_{e0} \sim 2.1 \times 10^{20} \text{m}^{-3} \text{sec keV}$. The confinement time supports the hypothesis of a linear scaling with plasma current, but data at higher current is as yet limited in additional heating power.

5. High Densities

High density gas fuelled 5MA plasmas were explored for heating powers up to 20MW. The temperature fell with increasing density such that $T_e(0) \sim 2.5 \text{keV}$ at $n_e(0) \sim 1.3 \times 10^{20} \text{m}^{-3}$ at 20MW. The plasmas were very clean with $Z_{eff}$ in the range 1.2-1.4. As at lower current [4], the maximum density reached with gas fuelling was set by the appearance of MARFEs, and the dominant radiation was from Chlorine. Fig.4 shows that higher densities can be reached at higher input power because the MARFE is postponed to higher density. At 20MW the density reaches $n_{e0} \sim 1.45 \times 10^{20} \text{m}^{-3}$ with a flat profile. Since the density profile is flat it is appropriate to compare the data with the scaling of edge density at MARFE onset with power deduced from 3MA data. The 5MA data shows a modest improvement over the 3MA scaling.

6. Summary

7MA discharges have been established at higher $q_0$, and sawteeth have been suppressed up to 6MA with ICRF and up to 6.5MA with combined heating. The improved central parameters and high L mode confinement have realised a promising improvement in the fusion performance of JET limiter plasmas. Whilst
plasma purity is a major concern in such regimes, it has been shown that, at high density, it is possible to maintain clean plasmas even with high power heating and high plasma current. The density limit shows a modest improvement at high current.

References

Fig 1. Equilibrium for 7MA pulse 22416.

Fig 2. Various time traces for 7MA pulse 22416. the pulse illustrated in fig 1.
Fig 3. Various time traces for high temperature pulse 23256.

Fig 4. Central density versus input power for 5-7MA limiter data with gas fuelling.
- MARFE onset
- 5MA sawtoothing
- 5MA sawtooth free
- 6-7MA data
- 5MA carbon limiter
EVIDENCE FOR FINE SCALE DENSITY STRUCTURES ON JET UNDER ADDITIONAL HEATED CONDITIONS

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Introduction: Measurements of density fluctuations are important in experimental efforts to identify the mechanisms which give rise to anomalous particle and energy transport in tokamak plasmas. One technique used for studying density fluctuations is microwave reflectometry [1,2,3]. We have developed a version of this technique, called correlation reflectometry, which has the potential to provide detailed information on the fluctuations [4]. In this paper we describe a new correlation reflectometer and present measurements made under a wide range of plasma conditions on JET. The results show that there are fine scale density structures present in the plasma and that the size of the structures depends on plasma conditions.

Correlation Reflectometry: In correlation reflectometry two or more independent reflectometers operate along the same line of sight at different fixed frequencies, they thus probe reflecting layers at different radial positions. The signals from the reflectometers are analysed using spectral analysis techniques. From the analysis, the crosspower spectrum, G(\omega), the crossphase spectrum, \Theta(\omega), and the coherence spectrum, \gamma(\omega), for any two signals are derived. The coherence has a maximum value of 1 (identical signals) to a lower level of significance \gamma . The latter is the coherence between two random signals and is in general greater than zero due to the digital nature of the analysis [5]. By analysing measurements from combinations of reflectometers with different frequency separations, corresponding to different separations between the reflecting layers \Delta x, it is possible to estimate the radial scale size and the radial wavenumber of density fluctuations. The former is obtained from the coherence spectrum and the latter is obtained from the crossphase spectrum.

A four channel correlation reflectometer has been constructed and operated on JET, figure [1]. The probing frequencies are 75.5 GHz, 75.6 GHz, 75.75 GHz and 76.2 GHz. The reflectometer operates in the extraordinary mode utilising the upper cutoff and probes the plasma along a major radius. For these probing frequencies, the corresponding separations between the reflecting layers are typically in the range 2 mm - 20 mm. Signal detection is by heterodyne techniques which give a good channel separation (> 40 dB) and a high signal to noise ratio (\approx 30 dB).

Measurements: Results have been obtained for three different plasma regimes: ohmic, L - mode, H - mode.

Ohmic: At the shortest interlayer distance of 2 mm, the coherence level is in the range 0.4 - 0.5 across the entire frequency band in the recorded signal (0 - 60 kHz). In such conditions, the crossphase has a constant value of 0 but with a large amount of scatter in the measurement. As the channel separation is increased, the level of coherence falls rapidly to the random level for a channel separation of \approx 4 mm, figure [2]. This behaviour has been found at all radii examined in the range of normalised
Figure 1: Schematic of four channel correlation reflectometer.

minor radii 0.15 - 0.75.

L-Mode: The coherence between channels increases as the amount of additional power is increased, figure [3]. When the additional power exceeds 10 MW, the level of coherence between all channels is high (> 0.65). When the coherence is high, the crossphase = 0 across the entire frequency band, figure [4]. This result has been found in both neutral beam heated and combined neutral beam/radio frequency heated plasmas. These results have been found for all the measurements in the range 0.15 - 0.9 of the normalised minor radius.

H-Mode: The levels of coherence is similar to that observed in ohmic plasmas, i.e. for small channel separations the level of coherence is significant (≈ 0.4 - 0.5). However, as the interlayer distance is increased the coherence falls rapidly to the random level. For significant coherence levels the crossphase is constant at zero. The range of normalised minor radii examined in this case is 0.55 - 0.95.

Interpretation: The results suggest that the density structures investigated have a radial scale size in the range 2 mm to ≈ 20 mm depending on plasma conditions. Moreover the measurement of a zero crossphase shows that the fluctuations are propagating transverse to the reflectometer beam (ie toroidally and/or poloidally) and not radially. It is known that under additionally heated conditions the plasma rotates toroidally. A possible interpretation of the reflectometer data therefore is that within the plasma there are fine scale density structures which are convected past the reflectometer antenna by the motion of the plasma. Under this assumption, and
from a knowledge of the plasma rotational velocity (obtained from charge exchange spectroscopy), it is possible to derive the effective toroidal wavenumber spectrum of the density structures. An example is shown in figure [5]. We use the crosspower spectrum because it is a measure of the relative amplitude of the density structures that exist between the two reflecting layers.

In order to obtain an estimate of the magnitude of the density perturbations associated with the structures we use a one dimensional model of the measurement. We assume that the perturbations have a gaussian shape in the radial direction:

$$\Delta n(r) = \sum \Delta n_0 \exp \left\{ -\frac{(r_c - r)^2}{l^2} \right\}$$

where $l$ is the correlation length and $r_c$ the central position of the perturbations. The model predicts the change of coherence with interlayer distance for given values of $\Delta n$ and $l$. The values used in the simulation are varied to obtain the best fit with the measured data. In the simulation the equilibrium upper cutoff profile is calculated from experimentally diagnosed values (LIDAR Thomson scattering and magnetics). For a typical example, the correlation length of the density perturbations in a beam heated plasmas ($P = 6$ MW) at 3.40 m is found to be $\approx 12$ mm. The amplitude of the perturbations is estimated to be $\Delta n/n \approx 1\%$.

Conclusions: Measurements with a four channel correlation reflectometer have shown the existence of fine scale density structures in JET plasmas. The structures exist throughout the observed range of normalised radius ($\rho = 0.15 - 0.95$). They have a radial scale size $\approx 2$ mm under ohmic and H - mode conditions and $\approx 20$ mm under L - mode conditions. The radial scale size of the structures appears to scale with the additional heating power. Estimates using a one dimensional simulation model suggest that the amplitude of the associated density perturbations is typically a few percent. The implications of these results for energy and particle is not clear at this stage but it is interesting to note that the structures are large when the particle confinement is low.

References:
Figure 2: The change of coherence with interlayer distance for different plasma conditions. Ohmic = +, H-mode = 0, L-mode (8 MW) = x and L-mode (15 MW) = n.

Figure 3: Variation of the coherence with neutral beam power in L-mode conditions. The interlayer distance $\approx 10$ mm.

Figure 4: The crossphase spectrum measured in an L-mode plasma.

Figure 5: Relative amplitude of the density fluctuations as a function of the toroidal wavenumber $k_\phi$. 
TRITON BURNUP IN JET - PROFILE EFFECTS

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Abstract
Measurements of the 14 MeV neutron emission from triton burnup show that the 14 MeV emission profile shadows closely the 2.5 MeV profile but after a delay corresponding to the triton slowing down time. The slightly greater width of the 14 MeV neutron profile is a consequence of the finite Larmor radius of the tritons. It has not so far been possible to identify unambiguously any effects on the triton burnup that are attributable to sawtooth crashes. Finally, the time dependence of the triton profile indicates that the triton diffusion coefficient is very small (<<0.1 m²/s).

Introduction
The study of the burnup of tritons released from d-d reactions is of special interest because the 1.0 MeV tritons possess similar kinematic properties to the 3.5 MeV alpha-particles released from d-t reactions. The birth profiles of both particles are indicated by the emission of the associated neutrons but, in contrast to the alpha-particles, the tritons also provide a measure of their radial movement whilst slowing down because their burnup (t-d) reactions are most probable just before they become thermalized (see fig 1). Previous work has indicated [1] that the confinement and slowing down of the fast tritons is in good accord with classical expectations. However, the measurements were global, representing an integral effect for the whole plasma volume. In the present paper, we report the first spatially resolved measurements of 14 MeV neutrons from triton burnup obtained at JET.

Experimental
The spatial profile of the 2.5 MeV neutron emission from JET plasmas is routinely measured using the Neutron Profile Monitor [2]. This diagnostic comprises two massive cameras with fan-shaped arrays of collimators and associated neutron detectors. The neutron detectors are NE213 liquid scintillators coupled to photomultipliers from which the anode signals are processed by Pulse-Shape Discriminator (PSD) units which permit the desired neutron-induced events to be separated from the accompanying gamma-rays. Events are only processed when the signal amplitudes correspond to neutron energies in the range 2.0 to 3.5 MeV, so as to reject neutrons that have scattered off the vacuum chamber wall.

The PSD units also provide an output signal corresponding to neutrons with energies above the upper discrimination level, normally set to 3.5 MeV. This permits the 14 MeV neutron-induced events to be recorded with an efficiency comparable with that for 2.5 MeV events. However, the typical fraction of tritons undergoing burnup reactions is less than 2% so that, in order to obtain useful statistics in the 14 MeV data channels, it proves...
necessary to confine the data analysis to those discharges with the highest 2.5 MeV neutron emission rates.

The 3.5 MeV discrimination level discussed above is rather low if pulse pile-up effects due to operation at very high count-rates are to be avoided; moreover, it is certainly too low when ICRF heating is employed, since RF heating is known to generate an extra component of neutron emission with energies up to 6 MeV or more. Unfortunately, when the discrimination level was raised to an effective 9 MeV, by reducing the photomultiplier voltages by 200 V, the PSD’s proved extremely difficult to set up and the deduced detection efficiencies for both channels were totally unreliable. Thus, profile information was obtained only with the 3.5 MeV discrimination level.

Results

Only a few discharges have so far provided optimal conditions for the triton profile measurements. Out of these, we select number 22517, a 3.16 MA, 3.44 T discharge with 12 MW of D⁰ Neutral Beam Heating (7 MW at 80 keV and 5 MW at 140 keV). As indicated in fig 2, the instantaneous neutron emission rate rose to 1.3 × 10¹⁵ n/s, falling by a factor of ten when the applied heating power was reduced to 2 MW. Ramping down the power in this way has the beneficial effect of maintaining the electron temperature at a high level and providing a conveniently long slowing down time for the tritons. The global 14 MeV neutron emission rate, as measured with a silicon diode [1], clearly shows that the average slowing down time is over one second (fig 2).

Inspection of the time-traces for the 14 MeV neutron count-rates in the individual channels of the profile monitor provides interesting information. Somewhat unwelcome is the discovery that the end channels of the vertical camera are sensitive to down-scattered 14 MeV neutrons that reach the detectors through the rather thin ends of the concrete shielding block; the 2.5 MeV neutron signals in these channels are similarly affected. All other channels appear to be satisfactory.

An immediate observation concerns the time delays between the 2.5 MeV and the 14 MeV neutron signal strengths. For 22517, the time delays are greatest for the central channels and least for the peripheral channels. This discharge is a typical H-mode discharge with a nearly flat electron density profile and a broad electron temperature profile that falls only a factor of two between the centre and the periphery as viewed by the cameras. These variations explain the observed time dependencies, at least qualitatively.

Because of the non-optimal settings of the PSD units, analysis is restricted to the period following 15.5 seconds, where the 14 MeV emission has reached its peak and the 2.5 MeV emission has begun its steep decline. It was necessary to correct the 14 MeV neutron signal for pulse pile-up of the 2.5 MeV neutron signals for the first few time intervals (until the 2.5 MeV emission had fallen below 2×10¹⁵ n/s), using an effective pulse pair resolution time for the PSD of 60 ns.

The emission profiles have been analyzed with the predictive code Orion [3], which assumes the neutron emission to be uniform around the magnetic flux surfaces. It fits the measured line integrals of neutron emission with a parabolic form (1−ψ^α)² for the radial variation of specific emissivity, where ψ is the flux surface label and α is the peaking factor to be optimized through a least squares procedure. The variations of the peaking factors for both 2.5 MeV and 14 MeV neutron emission profiles are shown in fig 3. It should be noted that, for discharge 22517, the flux surfaces do not change significantly between 14.5 and 16.5 s although the peaking of the electron density and temperature profiles increases when the H to L mode transition
takes place (at about 15.5 s). The moments at which sawtooth collapses occur are indicated in fig. 3. The 2.5 MeV profile maintains a peaking factor of about 5.8 until the sawtooth at 15.2 s, which briefly broadens the profile; then at 15.5 s the main NB heating ceases and the profile peaks ($\alpha = 7.5$). Another sawtooth crash at 16.4 s apparently broadens the profile suddenly, after which it peaks up again slowly.

The first reliable measurement of the 14 MeV peaking factor is at 15.9 s, giving $\alpha = 4.6$, which is considerably less than the value of 7.5 for the 2.5 MeV neutrons at the same time (i.e. the fwhm of the 14 MeV profile is 17 cm broader than the 2.5 MeV profile). However, it is not appropriate to compare profiles at the same time because the 14 MeV profile is defined by the 2.5 MeV profile at a time earlier by approximately the slowing down time for the 1.01 MeV tritons (about 1.4 secs for the conditions applying between 14.5 s and 16.5 s, decreasing to 0.8 s at 57.0 s). Thus, the comparison should be with the 2.5 MeV peaking factor after 14.4 s, i.e. $\alpha = 5.8$. In this case the the 14 MeV profile would have a fwhm broader by about 9 cm. This is explicable in terms of trapped orbit effects; orbit code calculations for a representative population of tritons indicate that a spreading of the fwhm by about 11 cm can be expected. (We note that the 14 MeV profile is displaced to larger radius than the 2.5 MeV profile by about 4 cm). A further broadening can be expected as a result of the weak sawtooth crash at 15.2 secs; it is generally understood that the effect of a sawtooth crash is to displace fast particles from within the central region of the plasma (within the electron temperature inversion radius) to just outside that region. However, a rearrangement of the innermost region would not have a strong effect on the width of a relatively broad triton distribution. Certainly, we have never observed changes in 14 MeV neutron emission at the time of a sawtooth crash, indicating that triton redistribution is not a major effect.

The fact that the two peaking factors are very similar towards the end of the discharge is purely coincidental. The abrupt rise in $\alpha$ for the 14 MeV neutrons begins at 16.5 s and corresponds to the rise in the 2.5 MeV peaking factor that starts at 15.3 s (for which the expected delay would be about 1.4 s, as observed).

The results presented in fig 3 also show that radial diffusion of tritons is not an important feature of discharge 22517, since there is no obvious tendency for the 14 MeV neutron emission profile to flatten out with time. Numerical analysis of the time-dependence of the 14 MeV neutron signal obtained from the silicon diode for an unusual discharge (number 20729), for which 14 MeV neutron profile data are not available, was presented earlier [4]; for 20729, the triton slowing down time was exceptionally long, indicating the need for a triton loss term representing either charge exchange or triton diffusion (the diffusion coefficient was estimated as being about 0.1 m$^2$/s). Unfortunately, for discharge 22517, the ancillary data needed for the full calculation of triton burnup are not of the required quality. Nevertheless, by comparison with 20729, the clear lack of broadening with time demands that the diffusion coefficient be appreciably smaller than 0.1 m$^2$/s. From this, we infer that losses such as those due to charge exchange are important for discharges with exceptionally long triton slowing down times.

Conclusions

The study of triton burnup is greatly facilitated by the availability of 14 MeV neutron emission profiles, since these permit a distinction to be made between charge exchange and particle diffusion losses. The present work indicates that the triton diffusion coefficient is very small ($\ll 0.1$ m$^2$/s).
Fig 1 Showing the time taken for a 1.01 MeV triton to slow down to the energy corresponding to its peak reactivity, for a plasma with the following conditions:

\[
\begin{align*}
  n_o &= 3.0 \times 10^{19} \text{m}^{-3} \\
  T_o &= T_i = 8.0 \text{ keV} \\
  Z_{\text{eff}} &= 4 \\
\end{align*}
\]

Fig 2 Time-traces for discharge 22517:

(i) - 2.5 MeV neutron emission (×10^6 n/s)
(ii) - 14 MeV neutron emission (×10^14 n/s)
(iii) - Total beam power (MW)
(iv) - Soft X-Ray signal, to identify sawtooth crashes.

Fig 3 Profile peaking factors for both 2.5 and 14 MeV neutrons as obtained with the neutron profile monitor. Times at which sawtooth crashes occur are indicated. The effective slowing down time for 1.01 MeV tritons is about 1.4 s for times between 14.5 and 16.5 s.
LONG PULSE HIGH POWER HEATING OF JET PLASMAS


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I. The observed rapid contamination of deuterium plasmas in JET during high power heating is attributed to evaporation of limiter and X-point targets, due to overheating of edges of target tiles. A specimen deuterium plasma, shot #22948, limited by beryllium outer belt-limiters, with $I_d=3$MA, \( <n_e> = 1.6 \times 10^{19} \text{m}^{-3} \), \( <n_D> = 1.3 \times 10^{19} \text{m}^{-3} \), and \( <n_{Be}> = 5 \times 10^{17} \text{m}^{-3} \), was sustained by deuterium and beryllium influxes \( \Phi_e \approx 10^{22} \text{atoms/s} \) and \( \Phi_{Be} \approx 4 \times 10^{20} \text{ atoms/s} \), measured using the $D_\alpha$ (6561Å) and $Be\ II$ (4361Å) emissions, giving $Z_{eff} = 1.5$. This was heated with $P_{CRF} = 8$MW, and $P_{NE} = 2$MW, giving $T_e(0) \approx 9$keV, $T_i(0) \approx 5.5$ keV, and D-D fusion rate $R_{DD} \approx 1.5 \times 10^{15} \text{s}^{-1}$ due to RF heating alone. Deuterium and beryllium influxes $\Phi_e \approx 10^{22} \text{atoms/s}$ and $\Phi_{Be} \approx 3 \times 10^{21} \text{atoms/s}$ accompanied the heating, giving $<n_e> = 2 \times 10^{19} \text{m}^{-3}$, $<n_D> = 1.2 \times 10^{19} \text{m}^{-3}$, and $<n_{Be}> = 1.5 \times 10^{18} \text{m}^{-3}$ and $Z_{eff} = 2.4$. This level of contamination could be sustained for $\approx 1.5 \text{s}$, during which $\Phi_{Be}$ continuously increased and $R_{DD}$ rapidly declined to insignificance.

II. The limiter surface area in contact with the plasma was $\approx 7 \text{m}^2$. Thermographic measurements of the beryllium limiter tiles showed that the temperature of the edges increased by $\approx 980 \text{K}$ to $\approx 1550 \text{K}$ within $1.5 \text{s}$, implying that the tile edges intercepted $\approx 18 \text{MW/m}^2$. The total tile edge area that was thus heated is estimated to be $\approx 5 \times 10^{-2} \text{m}^2$. Computations have been made of beryllium vapor pressure above the hot tile surface and the ablation rate. This is shown in fig.1, where $P_s$ is the saturated vapor pressure, and $\Gamma_{Be}$ is the flux of ablated beryllium atoms leaving the surface. At $1500 \text{K}$, the tile edges alone would contribute $\Phi_{Be} \approx 5 \times 10^{21} \text{atoms/s}$. This evaporation flux, plus that due to sputtering, with yield $\gamma_{SPUT} = \Gamma_{Be}/P_s \approx 0.1$, exceeds the measured $\Phi_{Be}$ given in sec.I. The discrepancy is most likely due to inaccuracies in measurement of beryllium influx, as discussed further in sec.III below. Beryllium ablation is a strong function of surface temperature $T_s$, $\Delta \Gamma_{Be}/\Gamma_{Be} \approx 72/100 \text{K}$ at $T_s \approx 1000 \text{K}$. In order to maintain the total Be influx tolerably low, $\Phi_{Be} \approx 10^{21} \text{atoms/s}$, then, even for perfectly uniform power loading, the limiter surface temperature must be kept low, $T_s < 1000 \text{K}$.

III. Temperature of beryllium atoms entering the plasma was deduced from measurements of Doppler width of BeI emission at 4407Å. In analyzing the spectrum the three Zeeman components of the measured line were assumed to be identically broadened. A spectrometer with an instrument width of $\approx 0.8 \text{Å}$ was used to measure a typical Doppler width of $\approx 0.3 \text{Å}$. Fig.2 shows evolution of the neutral beryllium temperature, $T_{Be}(\text{eV})$, so deduced. The statistical error is indicated, but larger systematic errors, $\approx \pm 4 \text{eV}$, are possible. Notice that in going from Ohmic heating, when sputtering is expected to dominate, to strong additional heating, when evaporation is thought to dominate, $T_{Be}$ is nearly unchanged. Evolution of $T_{Be}$ shown in fig.2 is contrary to expectation that $T_{Be}$ (sputtered) $\approx 100$ - $T_{Be}$ (evaporated). The measurement is expected to be able to register such a large change in $T_{Be}$. We conclude that the measured $T_{Be}$ relates to atoms of sputtering origin, and that the observed limiter surface is not representative of the whole. Fig.3 shows evolution of BeI and BeII emissions from the same plasma volume. BeI emission is very localized, whereas the BeII emission could be due to ions convected into the observation volume from a point of birth.
outside the view. Strong divergence between emissions of BeI and BeII shown in fig.3 indicates such a process. We must conclude that the measured BeII intensity is only an approximate measure of the total beryllium influx.

IV. In shot #22957, with identical OH heated target plasma as #22948 previously, deuterium gas was injected at 80mbl/s at the same time as ICRF heating, giving \( n_D = 4 \times 10^{19} \text{m}^{-3}, \) \( n_F = 3.4 \times 10^{17} \text{m}^{-3} \), \( n_{\text{Be}} = 2 \times 10^{16} \text{m}^{-3} \), and \( Z_{\text{eff}} = 2.4 \), sustained by \( \Phi_D = 2.5 \times 10^{22} \text{s} \) and \( \Phi_{\text{Be}} = 3 \times 10^{22} \text{s} \). \( \Phi_{\text{ICRF}} = 19 \text{MW} \) gave \( T_e(0) = 9 \text{keV} \), \( T_i(0) = 4 \text{keV} \), and \( R_{\text{BD}} = 2 \times 10^{15} \text{s} \). This level of \( R_{\text{BD}} \) could be sustained for several seconds. Although \( \Phi_{\text{Be}} \) increased by \( \geq 10 \) compared to that in #22948 during ICRF heating, runaway contamination by beryllium was arrested by gas injection without reducing \( R_{\text{BD}} \), as observed earlier with carbon limiters[1]. This technique for impurity control has been exploited widely during high power heating in JET, as well as in the X-point configuration.

V. Investigations are continuing of why injection of a small amount of gas, \( \Phi_D(\text{gas injection}) = 0.1 \Phi_D(\text{recycling}) \), simultaneously with high power heating reduces the contamination efficiency of edge impurities by a large factor. Conventional impurity transport through the SOL and edge plasma[2] has been modelled employing the 2-D LIM Code[3]. Flux of impurities in the bulk plasma is modelled by \( \Gamma_z(r) = -D(r) \nabla z(r) + n_z(r) V(r) \), where \( D = 0.54 \text{m}^2/\text{s} \) and \( V = -2Dr/\text{a}^2 \) are used, on the basis of previous measurement. Consequences of reducing the impurity source energy are computed. \( <E> \) is the deduced average energy of the source atoms, the relative contamination efficiency is expressed as impurity confinement time \( \tau_z \); for constant impurity influx contamination efficiency is proportional to \( \tau_z \). Results for sputtering \( (<E> = 10.5 \text{eV}) \) and evaporation \( (<E> = 0.1 \text{eV}) \) sources are tabulated.

<table>
<thead>
<tr>
<th>Source</th>
<th>(&lt;E&gt;(\text{eV}))</th>
<th>(\tau_z(\text{ms}))</th>
</tr>
</thead>
<tbody>
<tr>
<td>D+ sputtering</td>
<td>10.5</td>
<td>33</td>
</tr>
<tr>
<td>D+ sputtering plus</td>
<td>13.2</td>
<td>38</td>
</tr>
<tr>
<td>self-sputtering</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Evaporation</td>
<td>0.1</td>
<td>4.6</td>
</tr>
<tr>
<td>Evaporation plus</td>
<td>3.2</td>
<td>11.5</td>
</tr>
<tr>
<td>self-sputtering</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

Table I for #22948 shows that sputtered beryllium atoms are more efficient at contaminating the bulk plasma than an equal number of evaporated atoms by a factor \( \approx 3.3 \). This confirms the observation in sec.I where, going from OH to ICRF heating, the increase in \( \Phi_{\text{Be}} \) due to evaporation is \( \times 2.5 \) greater than the corresponding increase in \( n_{\text{Be}} \). For the gas injection conditions of #22957 the LIM calculations showed only a small reduction in \( \tau_z \) compared to #22948. Therefore, in terms of the conventional model[2] for impurity contamination, the observed reduction in \( \tau_z \) by a factor \( \approx 10 \) during gas injection, as witnessed by comparing sec.I and sec.IV, would require adjustment of D(r) and V(r) in LIM during gas injection. No justification has been found yet in JET for taking such a step[1].
VI. However, further experiments to measure modifications of impurity transport due to gas injection have been performed using impurity injection by laser blow-off[4]. The measured 1/e decay time, $\tau_{\text{NiXXVII}}$, of NiXXVII emission after a non-perturbing number of Ni atoms was injected into the plasma from the periphery is shown in fig. 4. $\tau_{\text{NiXXVII}}$ is plotted against $dn_e/dt$ during gas injection, with different heating powers in fig. 4(a). In fig. 4(b) $\tau_{\text{NiXXVII}}$ is plotted against gas influx $F_g$. We conclude from fig. 4 that $\tau_{\text{NiXXVII}}$ does not vary with $dn_e/dt$ or $F_g$. This however is inconclusive. The peak of NiXXVII emissivity occurs at $T_e \approx 8 \text{keV}$, thus near the center of the plasma. Measured $\tau_{\text{NiXXVII}}$ is thus representative of the plasma core. The contamination efficiency of edge produced impurities is controlled by transport coefficients in the plasma edge. These have not been accessed by the above measurements, for lack of suitable materials with observable low ionization spectral lines.

VII. To summarize, gas injection during high power heating is widely employed in JET to delay runaway contamination of the plasma by target materials, in limiter and X-point configurations. In limiter operation, measurements and analysis confirm that runaway contamination occurs due to evaporation of the Be targets. Consideration of vapour pressure and surface ablation show that surface temperatures $\geq 1000 \text{K}$ for Be targets in JET are intolerable. The experiments support calculations that the contamination efficiency of evaporated beryllium is smaller by a factor $\approx 3$ than that of sputtered beryllium. Injection of deuterium gas during heating increases shielding by a factor $\approx 10$, without excessive increase in $\langle n_e \rangle$ or reduction in $R_{\text{DD}}$. This large increase in shielding due to gas injection can not be understood in terms of conventional models[2,3].

We have conjectured previously that the observed increase in impurity screening due to gas injection may be associated with modification of flows in the plasma edge and SOL[1]. Whereas strong screening is feasible due to flows in the SOL in divertor configurations[5], it is not clear how flow related screening can occur in limiter configurations. However, progress has been made in this direction. A model of thermal transport in tokamaks due to micromagnetic stochasticity proposed earlier[6] has been extended to considerations of density transport[7]. Here, anomalous radial transport is a manifestation of classical parallel transport wherein the field lines of $B \phi$ are slightly tangled and caused to migrate radially due to the presence of magnetic islands. Injection of deuterium gas in such a plasma will give rise to average outward deuteron flow along field lines which will oppose by friction impurities flowing inward, also along field lines. Frictional impurity screening effects like those reported in this paper have been qualitatively reproduced[7].

ACKNOWLEDGEMENTS

We thank P.R. Thomas for encouraging this investigation. Valuable assistance of J. How, R. König, and R. Martin-Solis is acknowledged.

REFERENCES

Fig. 1: Saturated vapour pressure $P_s$ and ablation rate $P_{Be}$ of a beryllium surface of temperature $T_s$.

Fig. 2: Evolution of beryllium atom temperature $T_{Be}$ during Ohmic and ICRF heating, using BeI emission at 4407A.

Fig. 3: Evolution of total beryllium influx $\Phi_{Be}$ during Ohmic and ICRF heating, viewed in BeI and BeII emission.

Fig. 4: Variation of confinement time of nickel, $\tau_{Ni}$, measured from NiXXVII emission, with (a) the rate of change of $\langle n_e \rangle$, and (b) deuterium gas injection rate $F_D$. 

$\Phi_{Be} (\text{atoms/s})$ vs. time (s) for BeI and BeII emission.

$T_{Be} (\text{eV})$ vs. time (s) for Ohmic and 19 MW ICRF heating.

$\tau_{Ni}$ (ms) vs. $\langle n_e \rangle (10^{19} \text{m}^{-3}/\text{s})$ and $F_D (\text{mb}^{-1}/\text{s})$ for Ohmic, 15-17 MW ICRF, 7 MW ICRF + 3 MW NBI heating. $1.6 < n_e (10^{19} \text{m}^{-3}) < 2.5$. 

Graph showing confinement time $\tau_{Ni}$ against $\langle n_e \rangle$ and $F_D$. 

Legend: •: Ohmic, *: 15-17 MW ICRF, #: 7 MW ICRF + 3 MW NBI.
Abstract

The $^3$He–D fusion yield has been further optimized in a new series of experiments using the ICRF heating system to generate a high energy $^3$He minority tail. In excess of 15 MW has been coupled to the plasma in the $^3$He minority regime. The newly commissioned $^3$He neutral beam injection system was used to provide a central source of $^3$He in order to counteract the relatively strong pumping of He by Be that has been observed on previous occasions.

Best results were obtained with 3.5 MA discharges in the double null configuration with the wall power loading being shared between X-point dump plates, the belt limiter and antenna protection tiles.

The generated $^3$He–D fusion power increased to a record level of 140 kW (corresponding to a reaction rate of $4.6 \times 10^5$ reactions/s). The fusion multiplication factor $Q_{\text{rf}} = P_{\text{rf}}/P_{\text{rf}}$ improved from just below 1% as obtained in the Be gettering phase (1989) to 1.25%. However, as observed in previous campaigns, a saturation of $Q_{\text{rf}}$ seems to occur at around 10 MW of coupled RF power. The beneficial effect of central $^3$He deposition by the NBI system could not be exploited to its full extent due to excessive $^3$He concentrations already present in the plasma caused by the outgassing of the inner wall components.

Time resolved measurements show a clear correlation between generated fusion power and energy stored in the fast $^3$He ions.

Introduction

The fusion power generated by $^3$He–D reactions in JET plasmas has steadily increased since the first successful observation in 1986, when 10 kW was reported [1]. In 1988, the first systematic approach to optimize the fusion power resulted in 60 kW of generated power [2]. One of the limitations in these experiments was the rather low deuterium fuel concentration ($n_d/n_e=0.3$) in the plasma. With the introduction of Be (evaporation) into the machine, the plasma fuel concentration increased - leading to an improved yield of 100 kW [3].

The introduction of Be into the JET machine soon led to the realization that the presence of Be impurity ions, even in minute quantities, can give rise to the emission of large amounts of neutrons and gamma-rays from reactions of fast particles with the Be ions, which interfere with neutron and gamma-ray measurements [4,5]. Indeed, these reactions have been observed to contribute up to 5 kW of released nuclear power. Plasma purity, therefore, is not only vital from the plasma performance point of view but also plays a dominant role in the nuclear diagnostic capabilities.

In this paper, we describe a new series of experiments performed after the Be belt limiters and Be antenna screens were installed. The new, redesigned, antenna screens reduce the RF induced impurity influx to a negligible level.
Experimental

In order to keep the influx of impurities to a minimum, the plasma was set up in a double null configuration with the wall loading being shared between X-point dump plates (upper C, lower Be), belt-limiter (Be) and antenna protection tiles (C). Care was taken to avoid any excessive Be contamination, often at the expense of an increased influx of carbon from the upper dump plates. (The data from another series of experiments conducted in 5 MA limiter discharges had to be disregarded because of excessive interference from Be $\gamma$-rays). The configuration described above allowed good density control (at a value close to the optimum of $3 \times 10^{19}$ m$^{-3}$) and the coupling up to 15 MW of RF power. Fig 1 gives the evolution of the general plasma parameters for a typical discharge. In general, these discharges were free of sawtooth until the occurrence of a monster crash towards the end of the heating period (5-7 sec of sawtooth-free periods could be routinely observed). The effective ionic charge $Z_{\text{eff}}$ was estimated from charge exchange recombination spectroscopy to lie in the range $Z_{\text{eff}} = 2$ to 4.

The fusion power was monitored by measuring the $\gamma$-rays from the weak secondary branch $^3\text{He} + \text{D} \rightarrow ^2\text{Li} + \gamma$ (16.6 MeV); the detection system consists of a 125 mm x 125 mm long NaI(Tl) scintillator and a 75 x 75 mm BGO crystal. The detailed experimental setup has been described previously [1,2]

Results

Fig 2 shows the fusion power obtained as a function of coupled RF power. As can be seen, the highest yields achieved approach 140 kW, which is well above the previous record of 100 kW obtained during the Be gettering phase. All the data points, except one, have been obtained with 3.5 MA discharges in the double null configuration as described above. The exception was a 5 MA limiter discharge for which the relatively modest performance can be explained by a low deuterium content caused by a large influx of Be. All 5 MA limiter discharges and a few double null X-point discharges suffered from this problem. Fusion yields well above 140 kW have almost certainly been achieved but any quantitative analysis was rendered impossible due to a contamination by high energy $\gamma$-rays ($>12$ MeV) from $^3\text{He}$-Be reactions.

The most successful performance was obtained with the RF antennas in the monopole configuration with the $^3\text{He}$ minority gas being introduced by gas puffing. The dipole configuration was also used; there seems to be little difference between the two configurations as far as fusion power generated per MW coupled to the plasma is concerned [fig 2]. However, the maximum power coupled in this mode was only 4.5 MW since lack of time prevented the coupling from being optimized.

In order to counteract the strong pumping of He by Be as observed in previous campaigns, the newly commissioned $^3\text{He}$ neutral beam injection system was used to provide a central source of $^3\text{He}$. As can be seen [fig 2], all data points obtained this way lie significantly below those obtained when the minority gas was introduced by gas puffing. Excessive $^3\text{He}$ concentrations already present in the plasma due to outgassing of inner wall components after extensive $^3\text{He}$ NBI injection are thought to be the cause of this poorer performance. Shorter beam bursts and/or the use of the RF system in the dipole configuration to avoid mode conversion at these high levels of minority concentration could improve the situation.

An interesting observation is the correlation between fusion yield and energy content of the fast $^3\text{He}$ ion population as deduced from measurements with the diamagnetic loop. This can be clearly seen in fig 3, where the time evolution of the 16 MeV $\gamma$-ray yield is compared to that of the stored energy for a typical discharge. Also shown are the time traces for RF power and the
central electron temperature.

Using the full formulation for the RF driven minority distribution function as given in ref 6, it can be shown that the maximum fusion yield obtainable at optimum minority concentration scales linearly with the perpendicular energy content of the fast ions [7]:\[ P_{\text{rus}}(\text{MW}) = 0.2 \frac{n_d}{n_e} W_{\text{fast}}(\text{MJ}) \]\nprovided the tail temperature exceeds the critical energy \( E_{\text{crit}} = 300 \text{ keV} \) in present experiments.

In fig 4 the measured fusion power is plotted against the energy content of the fast particles for the present series of discharges. Also shown is the calculated maximum fusion yield obtainable as a function of stored energy for assumed fuel concentrations, \( n_d/n_e = 0.5 \) and \( n_d/n_e = 0.8 \). It can be seen that only for a few discharges has the theoretical maximum compatible with a clean plasma been achieved.

First comparisons with simulations using the PION [8] code show reasonable agreement, although some experimental observations are not fully reproduced (doubtless due to the lack of reliable diagnostic information on \( n_d/n_e \) and the \(^3\)He concentration).

Conclusions

Up to 140 kW of fusion power in the form of charged particles from RF driven He minority ions reacting with a deuterium background plasma has been generated in a recent series of experiments, where up to 15 MW of RF power has been coupled to the plasma in the \(^3\)He minority regime.

The power multiplication factor \( Q_{\text{rt}} \) reached a maximum of 1.25% at 10 MW of coupled power.

The clear correlation of generated fusion power with fast minority ion stored energy that is predicted from theory has been observed.

The improved fusion performance cannot be explained by higher \( n_d/n_e \) values alone. Part of the improvement has to be ascribed to the better confinement properties of the discharges used in these experiments (3.5 MA as opposed to 3 MA and double null X-point configuration instead of pure limiter plasmas) leading to higher plasma temperatures and therefore a higher fast ion energy content at equal coupled power.

As fast ion energy contents of up to 3 MJ have already been obtained in JET it should be possible to increase the fusion power further up to 300 kW, provided the deuterium fuel concentration \( n_d/n_e \) can be kept high (\( >> 0.5 \)).

It is worthwhile noticing that this power is almost exclusively released by charged particles, i.e. 14.7 MeV protons and 3.7 MeV \( \alpha \) particles and is, therefore, equivalent to 0.7 MW of D-T total fusion power for the same heat deposition; it should be possible to increase this to 1.4 MW in the near future.

References

5) Loughlin M.J. et al, This conference.
Fig 1: Evolution of main plasma parameters during a typical high fusion yield discharge; \(^{3}\)He beams for fueling and D\(^{3}\) beams as a diagnostic probe.

Fig 2: Variation of generated fusion power as a function of coupled RF power:
- \(^{3}\)He gas puff, monopole, 3.5 MA
- \(^{3}\)He NBI, monopole, 3.5 MA
- \(^{3}\)He puff and/or \(^{3}\)He NBI, dipole, 3.5 MA
- \(^{3}\)He puff, monopole, 5 MA Be limiter

Fig 3: Showing correlation between D\(^{-}\)\(^{3}\)He \(\gamma\)-ray signal and fast ion energy content.

Fig 4: Generated fusion power vs fast ion energy content; experimental points and theoretical maximum achievable power for \(n_d/n_e = 0.5\) and 0.8.
Introduction - The injection of metallic impurities by laser blow-off has been used on JET to study impurity transport. Fig.1 shows a comparison of the Ni XXVI brightnesses time evolution following Ni injection into OH, L-mode and H-mode plasmas. These results are typical of many other injection experiments in plasmas of similar conditions. The L-mode data clearly show a reduction of the impurity confinement time, $\tau_{\text{imp}}$, with additional heating as indicated by the faster decay of the Ni XXVI signal (~200-250 ms) compared to the OH case (~300-400 ms). The higher values of $\tau_{\text{imp}}$ are observed at higher electron densities. In H-mode plasmas the impurity confinement was dramatically improved with values for $\tau_{\text{imp}}$ more than ten times the OH values. This is observed both in ICRF and NBI heated H-mode. The radial dependencies of the transport parameters D and V have been obtained by simulating the soft X-ray profiles and line brightnesses time evolution following injection. To illustrate the technique the results of an H-mode shot will be discussed first. Similar data obtained in OH and L-mode plasmas have already been presented in Ref.1. Impurity transport parameters determined for OH, L-mode and H-mode will then be compared and the role of sawteeth on transport will be discussed.

Ni injection into H-mode - Typical traces of a 3 MA, double null X-point, neutral beam heated H-mode shot are shown in Fig.2. The $n_e^\alpha$ trace shows the H-phase to start at 10.2 s, 200 ms after the beginning of beam heating, and to last up to 12 s when the beam power is stepped down from 9 MW to 4 MW. The continuous rise of $n_e$ is due to both beam fuelling and the improved particle confinement during the H-phase. Ni was injected near the middle of the H-mode at 10.985 s. The subsequent evolution of the Ni XXV and Ni XXVI signals is given with an expanded time scale in Fig.2. The intensity step seen 100 ms after the initial rise on the Ni XXV signal is due to a small additional burst of Ni impurities. Fig.3 shows the time evolution of the soft X-ray emissivity distribution (background emission subtracted) along a horizontal central chord. Since only a fixed background level was subtracted corresponding to the emission immediately before injection, the emissivity of the last two profiles could be overestimated by up to 15%. The experimental data have been simulated using an impurity transport code where the radial flux density $\Gamma_z$ was described with the sum of a diffusive and a convective term. Both the diffusion coefficient D and the inward convection velocity V were taken to be a function of the radial coordinates and were specified, independently, to reproduce best the experimental data, namely, the absolute X-ray emissivity profiles and the spectroscopic line brightnesses time evolution. The shape of
the source function was taken so as to reproduce the time evolution of peripheral Ni lines (FWHM ~ 20 ms). The calculated soft X-ray emissivity profiles are shown in Fig.3 (dashed lines) for comparison with the measured profiles, together with the corresponding Ni density profiles also calculated with the same code (lower part of Fig.3). The transport parameter curves which have been used are shown in Fig.4 along with the corresponding curves used to describe the transport behaviour of OH and ICRF heated L-mode plasmas (see Ref.1).

Fig.1 Evolution of the Ni XXVI brightnesses divided by the electron density following Ni injection into OH, L-mode and H-mode plasmas.

Comparison between OH, L-mode and H-mode transport - The results of the experiments, indicate in all cases, that the impurity diffusion coefficient was much lower in the central region of the plasma than on the outside and that the transition between these two regions was rather abrupt. This explain the behaviour of the soft X-ray profiles, shown in Fig.3, which remain hollow up to the sawtooth crash. It is worth mentioning that such a structure in the diffusion coefficient was also observed in discharges where q was above 1 on axis. The values of D over the central region where the confinement is good (low D) are of similar magnitude in the three cases. However, this region is more extended in the H-mode case and the transition region is more gradual. These values are 0.03, 0.15 and 0.1 m²/s in the OH, L-mode and H-mode case, respectively, which compare well with the corresponding neoclassical values.
which are, taking into account collisions with the deuterium ions and background impurities, 0.04, 0.07, and 0.1 m$^2$/s, respectively. In the external region, the values of D are 1, 3 and 0.4 m$^2$/s assuming that D is constant up to the plasma edge. They are much larger than the corresponding neoclassical values which are 0.06, 0.06 and 0.08 m$^2$/s. There are some hints that the value of D may be reduced near the plasma edge. However this is difficult to ascertain because of the large experimental uncertainties in that plasma region.

Fig. 3 (top) Time evolution of the tomographically reconstructed X-ray emissivity distribution (background emission subtracted) along a horizontal central chord after Ni injection (shot #22122). The dashed lines show the result of a simulation. (bottom) Calculated Ni density profiles at the corresponding times.

Fig. 4 Radial profile of D (logarithmic scale on the left) and of V (linear scale on the right) determined for OH (#17661), L-mode (18112) and H-mode plasmas (#22122).
To simulate the H-mode data a convective transport barrier can be set at the plasma edge to retain impurities as already discussed in Refs. 2-3. It may be explained by the localised edge perpendicular electric field which has been shown to appear at the L-H transition. In the H-mode case discussed here the inward convection velocity was increased sharply in a thin region at the separatrix as shown in Fig. 4. The thickness $\Delta r$ of the zone where a high inward velocity is needed was determined from the behaviour of the Ni XXV line and of lines from lower ionization stages which are very sensitive to this parameter. The behaviour of the Ni XXV and Ni XXVI lines whose intensities remained almost constant up to the end of the H-mode (see Fig. 2) clearly show that impurities are very well confined during the H-mode. In our simulation D was assumed to have the same value at the plasma edge that it had in the plasma interior. If D is smaller near the edge the value of $V$ would also be smaller since the critical parameter to simulate the H-mode behaviour is $V\Delta r/D$. In one L-mode case which was near the H-mode threshold and where D was 5 m/s a strong edge barrier was also required to explain the data. This suggests that the transition from the L- to H-mode transport behaviour may be a gradual one where first a convective barrier builds up at the edge and then the diffusion gets gradually reduced in the plasma interior. In OH plasmas and in ordinary L-mode there is no indication of a transport barrier. However, some inward convection is still required to explain the data. The peaking parameter is of the order 1 for the impurities in comparison with 0.5 for the electrons. This is consistent with the slightly peaked profiles of $Z_{eff}$ often inferred from the X-ray emissivity profiles.

Sawteeth effects - During a sawtooth crash (~100 µs) the soft X-ray data show that impurity transport is greatly enhanced over the central region of the plasma up to a mixing radius roughly 30 % larger than the sawtooth inversion radius. In the case of laser impurity injection, when the impurity distribution is hollow during the inflow phase, a sawtooth crash allows for a rapid inward movement of the impurities which can quickly fill the central plasma region. This effect can be seen in Fig. 3 where the X-ray profiles just before and immediately after a sawtooth crash are shown together with the corresponding impurity density profiles on the lower part of the figure. In the simulation the sawtooth discontinuity was simulated as in Ref. 1 by taking ad-hoc enhanced transport parameters which reproduce the X-ray emissivity profile variation and the density jump visible only on the Ni XXVI line (Fig. 2). Once the impurities have peaked in the plasma center the reversed effect can be seen during the outflow phase, i.e. the impurities are expelled from the central region at each sawtooth crash (See Ref 1). Sawteeth therefore allow for much more rapid movement of impurities between the central region and the outside and effectively short-circuit these two distinct transport regions. When the the sawtooth period is much shorter than the impurity confinement time the transport of impurities is effectively controlled by the value of D in the anomalous transport region.

References


Acknowledgements: Contributions and valuable discussions with K Lawson and D Boucher are gratefully acknowledged.
THE ROLE OF THE PLASMA CURRENT DISTRIBUTION IN L-MODE CONFINEMENT

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Introduction

In JET steady-state L-mode discharges, the energy confinement time scales with the total plasma current [1]. This paper reports on two preliminary sets of experiments, conducted in discharges with non-stationary current distributions in order to elucidate the functional dependences of the local heat diffusivity, $\chi_{\text{eff}}$, which gives rise to this scaling. In sections 2 we describe the experiments. In section 3 we model the results. We find that a dependence of the form $\chi_{\text{eff}} \sim 1/B_p$ is not sufficient to reproduce the experimentally observed time evolution of $\tau_{\text{Ee}}$ and that an additional dependence on the magnetic shear length $L_s \equiv q/\nabla q$, such as the one proposed in [2], is required to reproduce the experimental observations.

Experimental observations

Current ramp experiments

In the first series of experiments, ICRF power ($D(He^3)$, 4 - 6 MW) is applied in discharges in which the plasma current is ramped from 2 to 3 MA, or down from 3 to 2 MA, at a rate of $\sim 0.5$ MA/sec. The toroidal field, electron density, $Z_{\text{eff}}$, auxiliary power, and radiated power are nearly constant. The auxiliary power is 3 to 4 times larger than the ohmic power input. It should be noted that the different magnetic and kinetic measurements of the stored energy have a somewhat different time evolution. In order to avoid ambiguities related to the fast ion energy content (which is not negligible in these discharges) as well as possible parasitic dependences of the magnetic measurements of the stored energy on plasma current, we shall use kinetic measurements to infer the confinement properties of these discharges.

Modest ($\sim 50\%$) variations in the internal inductance, $l$, are achieved with these ramp rates. The discharges are sawtoothing, and, although the sawtooth period is affected by the current ramp, the position of the $q = 1$ surface (determined from polarimetric measurements and from the electron temperature inversion radius) does not change significantly ($< 5\text{cm.}$). Field diffusion calculations using the TRANSP code and assuming neo-classical conductivity reproduce the measured evolution of the internal inductance and the $q = 1$ surface, indicating that no anomalous current penetration is taking place.

In the ramp-up case (see figure 1), the electron energy content, $W_e$, begins to increase almost immediately, although not proportionally to the plasma current. The ramp-down case, (figure 2) has a greater variation in $l$, and exhibits
a delay before $W_e$ begins to decrease. This result is similar to that obtained in current ramp experiments carried out in the TFTR tokamak [3] in which large variations in $l_e$ ($X < 3$) were produced and which also exhibited a delayed (>1 sec.) response of the confinement time to the plasma current. Note that the confinement deteriorates markedly only at the end of the current ramp, when $l_e$ and the inverse magnetic shear length, $Vq/q$ have stopped rising.

**Slow current penetration experiments**

In the second series of experiments, sawtooth-free periods of up to 7 sec. at constant plasma current (3.6 MA) are produced by the application of 11 MW of ICRF at the start of the current plateau. The resulting high electron temperatures (10-12 KeV) retard the peaking of the current distribution. $l_e$ varies from 0.7 to 1.1 and $q_e$, inferred from Faraday rotation measurements and calculated in TRANSP, descends from ~1.6 to ~1. The electron energy content increases approximately linearly with $l_e$ during the sawtooth-free interval. See figure 3.

![Figure 1](image1.png)  
Figure 1 (above, left) Time evolution of the plasma current, ICRF power, central electron temperature, electron energy content, self inductance and shear (at r/a=0.6) in pulse 23412.

![Figure 2](image2.png)  
Figure 2 (above, right) Time evolution of the plasma current, ICRF power, central electron temperature, electron energy content, self inductance and shear (at r/a=0.6) in pulse 23410.

![Figure 3](image3.png)  
Figure 3 (across, left) Time evolution of the plasma current, total power, central electron temperature, electron energy content, self inductance and axial safety factor (solid line--measured; dashed line--TRANSP calculation) in pulse 23400.
Simulations

The above experiments are simulated with a 1D, 1 fluid predictive transport code, which simultaneously integrates the energy, particle and current diffusion equations. Since the confinement time in JET is known to obey the constraints of short-wavelength gyro-kinetic turbulence, [4], we take

\[ \chi_{\text{eff}} = \rho^2 \frac{V_{th}}{L} \cdot F(q, \frac{L_S}{L}, \frac{L_T}{L}, \frac{L_n}{L}, \nu L/V_{th}, \beta, \varepsilon, \kappa) \]

where \( \rho \) is a Larmor radius, \( V_{th} \) is a thermal velocity, \( L \) is a plasma dimension, \( q \) is the local safety factor, \( L_T \) is a thermal scale length, \( L_n \) is a density scale length, \( \nu \) is a collision frequency, \( \beta \) is a local beta, \( \varepsilon \) is the inverse aspect ratio and \( \kappa \) is the elongation. The leading, “gyro-Bohm”, term in this expression scales as \( T_{\parallel}^2/B_0^2 \) and gives rise to the observed L-mode dependence of \( \tau_E \) on plasma current. Therefore, we shall try to determine what additional dependence, introduced by \( F \), can reproduce the observed time evolution of the electron energy during current transients.

We begin with the current ramp-down case (pulse 23410). As noted above the shear is strongly increased during the current ramp (53 - 55 sec), so that a dependence of the form \( F \sim L_n/L \) introduces a transient improvement in the confinement relative to the \( F \sim \) constant case. However this improvement is sustained for too long following the end of the current ramp. To obtain a better match to the data we take \( F \sim (qL_n/L)^\alpha \). The simulations, (normalized at \( t = 53 \) sec.) are compared with the measured electron energy for \( \alpha = 0, 1, 2 \) in figure 4. With no dependence of \( \chi_{\text{eff}} \) on shear, \( (\alpha = 0) \), the energy content decays too rapidly in the current ramp. This result is confirmed for different radial profiles of \( \chi_{\text{eff}} \).

Simulations of the current ramp-up (pulse 23412) exhibit a similar pattern, although the variation with \( \alpha \) is less marked corresponding to the fact that the shear is less drastically altered in this case. Nevertheless, with \( \alpha = 0 \), the energy content rises too quickly at the start of the current ramp and saturates too soon whereas simulations with \( \alpha > 0 \) show a more linear time evolution, in conformity with the data.

Figure 5 shows the evolution of the electron energy content in a discharge with slow current penetration (pulse 23400). Again it is necessary to invoke a dependence of \( \chi_{\text{eff}} \) on the shear in order to account for the variation in \( W_e \).

The simulations presented here take the “gyro-Bohm” scaling as a point of departure, in order to ensure a dimensionally correct scaling. One difficulty with this approach is that any dependence of \( F \) on the temperature (eg. through \( \beta \)) will affect the dependence of \( \tau_E \) on \( L_n \). Therefore the dependence of \( \chi_{\text{eff}} \) on \( L_n \) can not be determined unless its dependence on the temperature is also known. It should be noted that a scaling of the form \( \chi_{\text{eff}} \sim q^2/N_q \) (as proposed in [2]) also reproduces the qualitative features of the data.

Conclusions

The preliminary experiments conducted at JET to investigate the role of the current distribution in L-mode confinement indicate that \( \chi \sim L_n \) or \( L_\parallel^2 \). This
dependence accounts for the observed evolution of the electron energy in both current ramp and current penetration experiments. In order to confirm this dependence it will be necessary to conduct experiments with faster ramp rates and with higher levels of auxiliary power. In addition, it would be preferable to repeat the experiments with Neutral Beam heating in order to eliminate ambiguities related to ICRF-generated fast particles and the concomitant pressure anisotropy. Use of NBI would also permit a measurement of the ion energy confinement.

References

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Figure 4 Measured and modelled evolution of the electron energy content in pulse 23410 (current ramp-down). Solid line--data; open circles--$\alpha = 0$; squares--$\alpha = 1$; closed circles--$\alpha = 2$.

Figure 5 Measured and modelled evolution of the electron energy content in pulse 23400 (slow current penetration). Solid line--data; open circles--$\alpha = 0$; squares--$\alpha = 1$; closed circles--$\alpha = 2$. 
CURRENT RISE STUDIES


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1. Introduction

During current rise at ramp rates of order 1 MA/s, discharges evolve through a phase of MHD instabilities, as the field line safety factor $q_\psi$ at the plasma edge is falling. A quiescent regime is reached, when $q_\psi$ is held constant. The MHD behaviour in the early phase has been monitored by the magnetic diagnostic and the soft X-ray cameras with a fast sampling rate. Particular attention has been paid to the crossing of the $q_\psi=4$ surface, which is an important parameter for the 7 MA limiter discharge scenario [1]. The measurements are discussed and their implications for the $q$-profile are compared with the predictions of the 1 1/2 D LARS code, which evolves the field diffusion/transport equations and monitors toroidal tearing mode stability [2]. The current penetration during the later MHD stable phase has been studied using the 1 1/2 D analysis code TRANSP [3] to determine which form of resistivity gives the best data consistency. The effects on the delay of the onset of sawteeth by electron heating by ICRF are also studied.

2. Early MHD phase

a. Experimental data

Fig. 1 shows the temporal evolution of the plasma current $I_p$, $q_\psi$ and signals proportional to the amplitude of rotating magnetic modes with toroidal mode numbers $n=1, 2$ and $3$. The dominant poloidal mode number of these modes determined from pick-up coil data analysis is marked in the figure, when there is no ambiguity. As the magnetic modes are observed to be destabilized, when $q_\psi$ is larger than the appropriate low-order rational value, they are resistive, since their resonant surfaces lie within the plasma.

When crossing $q_\psi=4$, two cases can be distinguished:

- At low inductance, the radiation is low with high central electron temperature and low density. The magnetic diagnostic measurements indicate the presence of a dominant $m=4, n=1$ mode at the plasma edge and give a radial width around 10 cm for a magnetic island on the $q=4$ surface. When the soft X-ray emission in the central region of the plasma does not show any sign of mode activity, the current rise is successful. In the opposite case, a dominant $m=4, n=1$ mode is measured by the soft X-ray cameras between about 20 and 60 cm from the plasma centre indicating that there are at least two $q=4$ surfaces in the plasma. Soon thereafter, there is a persistent locked
mode, which can terminate the discharge by a disruption.

At moderate inductance, the radiation is high with low central electron temperature and high density. Sawtoothing starts about 2s earlier than in the low inductance case, showing that the current penetration is faster. A dominant \( m=2, n=1 \) mode is measured by the magnetic pick-up coils. A tomographic reconstruction of the soft X-ray emission shows an \( m=2, n=1 \) island 8 to 10 cm wide radially (see Fig.2). The island width and the position of the \( q=2 \) surface agree with the magnetic measurements and the q-profile computed with a magnetic equilibrium code.

b. Current diffusion and mode stability predictions

The current ramp phase is simulated using the LARS diffusion and transport code [2]. This code solves the equations derived from a high order expansion of the full toroidal 1 1/2 D model generalised to incorporate self-consistent shaping of the magnetic surfaces and permit time programming of the minor and major radii and plasma current. The simulation starts 0.5s after the beginning of the discharge, because the experimental data are not reliable before. The current and the major and minor radii all follow the experimental values; this gives a good fit to the measured \( q(t) \) traces. Neo-classical resistivity is used together with a constant thermal conductivity. The computed \( I_1 \) versus \( q_\psi \) diagram for the low inductance case is in reasonable agreement with the experiment.

Fig.3 exhibits a time sequence of radial q-profiles in the low inductance case: they are non monotonic. At \( t=1s \) (1.5s from the start of the discharge), \( q_\psi=4 \) is about to be crossed and there is only one \( q=4 \) surface inside the plasma. However, the off-axis maximum of \( q \) depends sensitively on the initial conditions and can be larger than 4; this gives three \( q=4 \) surfaces in the plasma. In practice, the q-profile in the middle of the plasma might be reconnected by a double tearing mode. Note that, in the experiment, the observation of the \( (4,1) \) mode in the plasma centre critically depends on the previous history of the discharge (including MHD behaviour) before crossing \( q_\psi=4 \). In the moderate inductance case, the q-profiles are also non-monotonic in \( r \), but with \( q(0) \) falling quicker to 1, in agreement with the experiment.

The programme is also coupled to a toroidal stability code so that stability of the evolving q-profiles can be monitored. Considering the internal MHD modes for integral q-values, the sequence of events met in Fig.3 is the following. A resonance enters at the axis as the central \( q(0) \) drops; the \( m=4 \) mode is initially stable, when the resonance emerges from the axis in this way, while \( m=3 \) and 2 are initially unstable. A triple resonance is encountered as the off-axis minimum in \( q \) drops in the outer region of the plasma. Because this condition is approached from above, the current gradient is stabilising, when it is negative. Note that the triple resonance cannot be handled as yet by the stability code, but presumably the tearing mode associated with at least one of the resonances is unstable. Finally, the triple resonance is lost as the off-axis maximum drops in the inner region of the plasma; in this case, the current gradient is destabilising, when it is negative. At this point, the \( m=4 \) and 3 modes, which remain in the outer region of the plasma, are unstable.
3. Quiescent phase

The magnetic field diffusion during the later MHD stable phase is studied using the 1.5D transport analysis code TRANSPI [3]. The fast current rise with a constant \( q_\psi \) maintained by controlling the plasma aperture expansion is followed from 2.5 (2.5s) to 5MA (5s). 9MW of ICRF is applied at 3s.

The calculations start from an initial state prescribed by experimental data and only the profile shape of the toroidal electric field is assumed. During the field diffusion, data on electron/ion temperatures, density, \( Z_{ eff} \) and the plasma boundary shape are used with a model for resistivity. Consistency checks are made between measured and predicted values on loop voltage, total energy, internal inductance and line of sight integrals from the interferometer and polarimeter.

These checks demonstrate that a neoclassical resistivity model yields the best overall agreement with the experimental data. This confirms the result already obtained for an ohmic discharge with the same current rise phase [4].

In the ohmic pulse, sawteeth start to be observed at 3s. The safety factor on axis \( q(0) \) computed by TRANSPI with a neoclassical resistivity reaches 1 at about the same time, whereas it remains above 1 throughout the calculation using a Spitzer resistivity. In the ICRF heated discharge, sawteeth are suppressed for a long period. However, \( q(0) \) continues to decrease below 1 as shown in Fig.4, which compares the values computed by TRANSPI with the measurements obtained with the polarimeter. Thus, these calculations, which are fully consistent with a large set of experimental data, show that the suppression of sawteeth does not arise from the reduction of resistivity caused by ICRF heating.

4. Conclusion

The fast current rises developed for the 7MA limiter discharge scenario are successful in a limited range of operating conditions only. Crossing of \( q_\psi = 4 \) can cause resistive \((4,1)\) tearing modes at low inductance or a \((2,1)\) island at moderate inductance, both of which can lead to a disruption. During the subsequent quiescent phase, current penetration is well described by neoclassical resistivity and electron heating by ICRF does not delay the appearance of \( q(0) = 1 \).

References
Fig. 1: Plasma current $I_p$, safety factor at the edge $q_\psi$ and amplitudes of rotating magnetic modes with toroidal mode numbers $n=1, 2$ and $3$ versus time.

Fig. 2: Tomographic reconstruction of the soft X-ray emission showing a $(2,1)$ island, when crossing $q_\psi=4$ in the moderate inductance case.

Fig. 3: Time evolution of $q$-profiles in the low inductance case. The computation starts $0.5s$ after the beginning of the discharge.

Fig. 4: Time variation of the safety factor on axis $q(0)$: TRANSP calculation (dash-dotted line); polarimeter data (full line).
JET EXPERIMENTS WITH 120 keV $^3$He and $^4$He NEUTRAL BEAM INJECTION


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ABSTRACT

Multi-megawatt, long-pulse NBI of $^3$He and $^4$He into JET has produced good ion heating and similar global and local energy confinement to results obtained with D NBI, in both limiter and H-mode regimes.

INTRODUCTION

Neutral beam injection (NBI) of helium (He) is of interest for a variety of reasons, as proposed in [1]. In particular, it should be possible to attain very high central ion temperatures, whilst limiting the neutron emission (and hence the induced radioactivity of the machine). The employment of monoenergetic He NBI as a diagnostic probe and for heating greatly facilitates the study of plasma conditions. He offers better beam penetration than with deuterium (D) beams, higher energy, and the possibility of diagnostics from D-$^3$He fusion. Accordingly, initial experiments have been carried out in JET with He NBI of up to 5 MW of $^3$He or 7 MW of $^4$He at 120 keV for 3 sec using Argon frost pumping. These are the first NBI experiments in which He has been used for heating a tokamak, and a record power of $^4$He has been injected.

Results have been obtained for diagnostic comparisons of ion temperatures [2], particle transport studies [3], ICRF heating with NBI [4], and edge diagnostics for He [5]. In what follows, we concentrate on ion temperature measurements from the neutron generation and global and local power balance assessments in plasmas with He NBI in comparison to similar plasmas employing D NBI.

GLOBAL POWER BALANCE AND LOCAL TRANSPORT ANALYSIS

To investigate the relative heating efficiency of He and D NBI, we compare two discharges which are similar except for the heating species. Both are 3.6 MA, 2.5T, double-null D plasmas with 2.7 MW of 80 keV D NBI for 3.5 s. Discharge #22975 has an additional 3.7-5.1 MW of 120 keV $^4$He
injected for 0.5 s, and #22976 has an additional 4.3 MW of 80 keV D. The stored energies, central ion temperatures, and NBI heating for both discharges are compared in Fig. 1. These discharges enter the H-mode regime after 0.15 s of He and 0.25 s of D additional heating. The incremental confinement time of the plasma is 0.6 s for the He injection and 0.5 s for the D injection. The ion temperature is substantially higher in the He injection discharge. Both discharges are sawtoothed.

A concern with He injection is that He neutrals with electrons in a metastable state are formed during beam neutralization; they ionize at the plasma edge and cause localized limiter or dump plate heating, leading to impurity generation and enhanced plasma radiation. The most recent calculations indicate that about one tenth of the beam would be metastables. A 3He beam at 4-5 MW for 3 s does not produce a significant increase in plasma radiated power, as seen in Fig. 2 for the 5 MA limiter discharge #23252. Localized limiter heating from direct impact of metastables ionized in the edge is below the detection level of cameras viewing the limiter impact zone for ionized metastables. The global energy confinement time of 0.53 s is similar to those for D NBI in 5 MA belt limiter plasmas. However, surface probe measurements indicate that some metastables are trapped at the plasma boundary.

A time-slice analysis of sawtooth-averaged local transport for limiter discharge #23252, which had steady state conditions, derives an effective thermal conductivity (regarding the plasma as a single fluid) \(\chi_{eff}\) between 0.5 and 1.0 m²/s in the radial range 0.3<r/a<0.8. This \(\chi_{eff}\) is similar within error bars to that corresponding to comparable discharges with D NBI.

Detailed particle transport studies with CXRS use a 0.5 s He NBI pulse to provide an axial source. A rapid decay (0.5 s) in the presence of sawteeth is seen in L-mode plasmas. In H-mode, the central decay is less marked, and absent in some ELM free cases.

NEUTRON AND CHARGE EXCHANGE SPECTROSCOPY (CXRS) MEASUREMENTS

He NBI eliminates the beam-beam and beam-plasma reactions occurring with deuterium beams, resulting in a reduction of total neutron generation while still providing efficient plasma heating. The neutron rate shown in Fig. 2 from 5 MW of He is comparable to the rate from 1.3 MW of D NBI at 80 keV.

The elimination of beam produced reactions also allows an improved measurement of the ion temperature with neutron spectrometers, since the broadening of the D-D neutron spectrum is due only to thermal ion motion. Measured temperatures agree with CXRS measurements taken by adding a low-power D beam at the end of the He heating pulse. In discharge #23252, the axial ion temperature at the end of the He NBI is about 3 keV from CXRS. This temperature agrees with the time-of-flight neutron spectrometer measurement of 2.9 keV and with the electron temperature (due to high plasma density).

The ion temperature is also deduced from the neutron emissivity and compared to other measurements, as follows: By averaging the neutron profile monitor data from #23252 over 8 to 11 s to obtain good statistics, the 2-D emissivity profile is derived by tomographic analysis and shown in Fig. 3. The axial neutron emissivity is 10⁵ n/m²s. The profile was averaged over many sawteeth (see Fig. 2), and therefore the result is subject to the problems of ion redistribution by sawtooth crashes and restrictions discussed in [10].
Both the neutron line-of-sight integrals and the resulting emissivity are peaked on axis, although there is some off-axis structure due to the effects of sawtooth crashes. The Full Width Half Maximum is 0.35 m, compared to the wider and hollow emissivity profile from tomography for the 5 MA ohmic discharge #15119 [10].

The local neutron emissivity depends only on ion temperature and deuterium ion density. The measured Zerr from both CXRS and visible Bremsstrahlung is 1.5±0.1 and the time-averaged axial electron density is 6x10^{19} m^{-3}. The CXRS measurements indicate that the axial ratio of deuterium to electrons is 0.8 to 0.95. The lower value is consistent with the amount of He injected during 3 s, which could provide 20% of the electrons when volume averaged. If the axial deuterium density is taken to be 4.8x10^{19} m^{-3}, then the 3 s averaged axial ion temperature based on axial neutron emissivity is 2.6 keV, within the error bars of the other measurements.

HOT-ION H-MODES

He and D NBI together are able to produce a hot-ion H-mode in double-null divertor plasmas. When 4.7 MW of 120 keV 3He is combined with 10.5 MW of 80 keV D NBI into a low density deuterium plasma, a hot-ion H-mode is produced with axial ion temperatures of up to 25 keV. The H-mode regime is entered 0.6 s after NBI begins. The NBI powers and axial temperatures are shown in Fig. 4 for discharge #23275. These results are comparable to those with pure D NBI. Similarly, the temperatures eventually decrease due to impurity influx. Using the arguments in [1] related to thermalizing injected beam energies, the maximum ion temperature is limited to ~30 keV by the low (species-averaged) directed-beam energy (~60 keV) of the higher (factor of ~5) particle current D beam.

At somewhat higher plasma densities, for example in discharge #23272, diamagnetic loop measurements give stored energy and energy confinement time values of 8 MJ and 1 s.

CONCLUSIONS AND DISCUSSION

He NBI is shown to produce both H-modes and (with D NBI) hot-ion H-modes with reduced neutron yields. The global and local energy confinement is similar (within error bars) to those obtained with D NBI. The central ion temperature is higher with He NBI than for equal power D NBI. Diagnostic possibilities have been enhanced, especially for neutron spectrometer and neutron profile measurements. Increased ion temperatures are expected in future experiments with ~15 MW of 3He NBI into the plasma at energies around 155 keV.

REFERENCES

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Fig. 1 Stored energies, central ion temperatures, and NBI heating waveforms for discharges 22975 (with He) and 22976 (with D).

Fig. 2 Global neutron emission, radiated power, and D and He NBI waveforms for discharge 23252.

Fig. 3 2-D neutron emissivity profile derived by tomographic analysis from 8 s to 11 s for discharge 23252.

Fig. 4 Ion and electron temperatures and NBI waveforms for discharge 23275.
DIFFUSION OF ALPHA-LIKE MeV IONS IN TFTR

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Introduction:
Single particle confinement of alpha particles is of crucial importance in reactor-grade tokamaks like BPX and ITER. Besides the well-known process of first-orbit losses, mechanisms that could lead to significant loss of alpha particles are turbulence-induced diffusion and toroidal field ripple stochastic diffusion. These two mechanisms have been separately studied in TFTR using two different detectors (one at the bottom of the machine and the other near the outer midplane) which can detect escaping charged fusion products, namely the 1 MeV triton and the 3 MeV proton in D-D plasmas (and also the 3.5 MeV alpha in D-T). The main difficulty in this type of experiment lies in the necessity of distinguishing the diffusion process from the always-present first-orbit loss process. In this paper, we show how these two processes can be distinguished using the pitch-angle discrimination of the detectors. The pitch-angle is defined here as the angle of the particle trajectory with respect to the toroidal direction and so is a measure of the ion magnetic moment $\mu$. Results obtained at the midplane would be the first reported evidence of TF ripple diffusion in a tokamak.

Experimental results:

a) Turbulence-induced diffusion of passing MeV ions:
Because of their relatively small shift from the flux surfaces most of the passing particles are confined on their first-orbit. Any cross-field radial diffusion due to the background turbulence can bring passing particles further out in minor radius until they mirror at the higher magnetic field on the inner side of their orbit and become trapped. Confined counter-going passing particles can thus be lost to the wall because of the much larger shift of the trapped particle orbit (co-going particles would simply be reflected to a smaller minor radius and would not be lost). This extra loss of MeV ions, if present, would then be perceptible at the bottom detector as an additional
loss precisely at the pitch-angle $\chi_{fb}$ corresponding to the passing-trapped boundary (the fattest-banana orbit). Meanwhile, first-orbit loss would occur over a range of pitch-angles between this pitch angle and $\sim 90^\circ$.

The number of confined passing particles increases substantially with plasma current (i.e. from 0.6 to 1.8 MA) and so is the number of particles “available” for radial diffusion in the central part of the plasma. In figure 1, we show the two extreme cases in first-orbit confinement, 0.6 and 1.8 MA (with a major radius of 2.60 m and a toroidal field of 3.8 T). At 0.6 MA, the observed pitch-angle distribution agrees with the numerically calculated distribution of first-orbit losses, as expected. At 1.8 MA, however, the ratio of possible diffusive loss to first-orbit can be as high as $2.5 \pm 1$ in the limit of $D \to \infty$, thus the shape of the calculated pitch-angle distribution varies considerably with the amount of anomalous losses, e.g. for $D \to \infty$ the shape approaches an experimentally broadened delta function at $\chi_{fb}$. Models using different diffusion coefficients are plotted along with the observed pitch-angle distribution. The calculations also include the finite resolution of the detector of $\pm 6^\circ$ (FWHM). The good agreement of the experimentally measured distribution with the $D=0$ case in figure 1b) indicates a very small diffusion of energetic ions, with an upper bound at about 0.1 m$^2$/sec.

b) TF stochastic ripple diffusion of trapped MeV ions:

In the case of trapped particles, diffusion would be dominated by the toroidal field ripple stochastic diffusion [1]. In the stochastic region, which would be typically a region of ripple strength > 0.1%, the bounce point of energetic ions diffuses vertically (in conserving $\mu$ and $E$) until the ion hits the wall. In TFTR, the amplitude of the vertical drift of the bounce tip is typically on the order of a centimeter near the plasma edge. In this case, the previously confined alpha-like ions will tend to impact the wall just below the midplane. This loss mechanism is relatively fast, since particles are expected to escape the plasma in typically less than 50 ms (for the 1 MeV triton, and less for the 3.5 MeV alpha and the 3 MeV proton), which is an order of magnitude smaller than their slowing down time.

At the end of 1989, a new charged fusion product detector [2] was installed in TFTR just below the midplane to look specifically for TF ripple stochastic losses of alpha-like particles. Since the detector is pitch-angle resolved, we can distinguish particles lost on their first-orbit (with the maximum at $\chi_{fb}$ described above) from ripple losses which originate mainly from trapped particles with their bounce points below the magnetic
axis and $\chi > \chi_{fb}$. At 0.6 MA all trapped particles are lost on their first orbit, whereas at 1.8 MA a substantial number are confined and thus subject to TF ripple stochastic diffusion. The results reported here are the first indications of TF ripple diffusion and the numerical simulations should be considered as preliminary.

In figure 2, we show the pitch-angle distribution of escaping ions as measured by the midplane detector for the same 0.6 and 1.8 MA discharges as above. The dashed lines correspond to numerical calculations from a mapping code which uses a guiding-center code for calculating the first poloidal transit of particles and a random walk code in which bounce points are diffusing until they reach the wall. In the 0.6 MA case, the measured distribution is in rough agreement with first-orbit calculations. The discrepancy is believed to originate from the simple profiles used in the simulations (i.e. for the plasma current, magnetic field, particle source and ripple). In the 1.8 MA case the first-orbit losses are expected to peak near $52^\circ$, in good agreement with the measured distribution. However, at higher pitch-angle (near $62^\circ$) we observed a second peak which is consistent with calculated TF ripple losses and which cannot be explained by just first-orbit losses. Further improvements in the numerical simulations are underway and will include more realistic profiles and a more detailed first wall geometry.

Summary:

Experiments performed on TFTR on the diffusion of passing MeV ions indicate a small turbulence-induced diffusion rate (less than 0.1 m$^2$/sec), consistent with the reduced diffusion rate expected from orbit-averaged diffusion theory for energetic particles [3]. New measurements at the outer midplane are the first evidence of the presence of TF ripple-induced losses as predicted in ref. [1]. Further work is actually being performed in order to simulate more adequately the experimental results by using more realistic profiles and tokamak geometry.


Acknowledgements:

This work was supported by the United States Department of Energy, under Contract No. DE-AE02-76-CHO3073.
fig. 1 Measured and calculated pitch-angle distributions for 0.6 MA and 1.8 MA discharges for the bottom detector. For the 0.6 MA case the measured distribution is compared with two calculated first-orbit distributions (in dashed lines) with different current profiles. For the 1.8 MA case, the measured distribution is compared with first-orbit losses only (D=0), with D=0.03 m²/sec and with D=0.3 m²/sec calculated distributions.

fig. 2 Observed and calculated pitch-angle distributions for the same discharges but for the midplane detector. The 0.6 MA case exhibits first-orbit losses only whereas the 1.8 MA case exhibits the presence of ripple losses at higher pitch angle consistent with numerical calculations (dashed line).
OBSERVATION OF STRONGLY LOCALIZED FAST PARTICLES RIPPLE LOSSES IN TORE-SUPRA

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INTRODUCTION

During additional heating in tokamaks, hot spots are commonly observed on the first wall components and are generally related to misalignments and/or limiter edge load (e.g. [1,2]), though more specific cases due to fast particles or toroidal field (TF) ripple have also been reported [3,4]. Recently, during long pulses experiments with LHCD in TORE-SUPRA, strongly localized hot spots have been observed deeply inside each vertical ports. In this paper, an interpretation of this unusual feature is given on the basis of a comprehensive analysis of the ripple trapped particle dynamics.

HOT SPOTS AND FAST TRAPPED ELECTRONS

The most striking feature of the hot spots is that they always occur far below the first wall components (≈ 30 cm) at the same location on a pumping tube inside each open vertical ports, on the electron drift side (Fig.1). At these points, the port angular aperture is small (3°). The phenomenon is particularly strong during low density operation and a thermal load of 100 W/cm² localized on about 5 cm² is deduced from infrared thermography.

Fig.1: Localization of the hot spots.
Though the associated total power loss is evaluated to be only of order of 10 kW, the local temperature elevation can be a severe limitation for long pulse operation. An accurate understanding of this "unexpected" phenomenon can be given by a model based on the interaction of trapped particles with the TF ripple perturbation. Besides the fast circulating electron tail generated by the LH wave, angular scattering creates a small suprathermal trapped component (banana orbits). As soon as the ripple $\delta$ is high enough to get $B_0/(NB\delta)<1$, these electrons can be trapped in the magnetic ripple wells (depth $\delta^*$) that appear along a field line [5]. $N$ is the number of TF coils and $B_R$ the horizontal component of the magnetic field. This condition is satisfied in a large peripheral zone in TORE SUPRA. The trapped electrons are lost to the wall if collisions are unable to scatter them out of the well during their drift. It can be shown that this condition is fulfilled for electrons above 50 keV. In the following we will consider that such electrons are collisionless for the ripple related phenomena of interest here.

**ADIABATIC INVARIANCE FOR RIPPLE TRAPPED PARTICLES**

We may consider that a collisionless ripple trapped particle is confined in a magnetic mirror whose depth and length vary in time because of the particle vertical drift across the plasma. For these electrons, there is a strong ordering between the rate of change of the mirror characteristics and the ripple bounce frequency ($v_{\text{drift}}/a_\delta << \omega_{\text{rb}} = (\delta/2)^{1/2}Nv/R$). As a consequence, their dynamics is adequately described by an adiabatic invariant given by an action variable $I_\parallel = (1/mv) \int p_\parallel dq$ associated to the ripple bounce motion. As indicated on fig. 2, the particle always enters the ripple well on a marginal orbit that determines the value of the invariant (curve 1) and then drifts along a $B=C^d$ surface towards higher ripple zones and enters more deeply into the well (curve 2). During the drift the well becomes more and more symmetric as the ripple increases.

Fig. 2: Evolution of the particle position in the ripple well during its drift along an iso-B surface.
This invariant links the bounce angles \((\phi_{b1}, \phi_{b2})\) to the magnetic well shape during the drift motion along the iso-B surface, independently of the particle mass or energy. For non-symmetric wells (for example near the trapping point), \(J_{\parallel}\) must be computed numerically. However, the main feature of the evolution of the bounce angles is given by the case of symmetric wells (high \(\delta\)), where the action integral is:

\[
J_{\parallel} = 8 R \sqrt{\delta^*} \left\{ E(k)-(1-k^2)K(k) \right\}/N
\]

where \(E\) an \(K\) are the complete elliptic integrals, \(k = \sin(N\phi_b)\) and \(\phi_b\) the toroidal bounce amplitude. For harmonic oscillations (deeply trapped particles), this expression reduces to \(J_{\parallel} = (\pi/2) N R \sqrt{\delta^*} \phi_b^2\). The bounce angle \(\phi_b\) decreases as \((\delta^*)^{-1/4}\) during the drift towards the vacuum vessel. Consequently electrons that are first marginally trapped on mirrors with \(\delta^* = 0\) and \(\phi_{b} \leq 10^\circ\) can get in the vertical port where \(\delta^* > 10\%\) and \(\phi_{b} < 1.5^\circ\).

**INTERPRETATION OF THE HOT SPOTS LOCALIZATION**

The plasma zones where an iso-B surface is tangent to an iso-ripple surface play a major role in the observed phenomena. They correspond to large portions of the plasma, where all the banana fast electrons are ripple-trapped with the same bounce invariant \(J_{\parallel}\), and have exactly the same dynamics during their vertical drift along the considered iso-B curve (fig.3). If these electrons get in the vertical port, they hit the port structure at the point where their bounce angle matches the port width. Consequently, a hot point can be associated in the port to each iso-B curve and all these points form a line, somewhat analogous to a rainbow or a caustic. The heat load along this line is reinforced if an obstacle (such as the vertical pumping tube represented on fig.1) intercepts more directly the particles (horizontally drifting in this area because of the proximity of the coils). The key point that validates the model is the following: the invariant \(J_{\parallel}\) computed at the hot spot effectively corresponds to particles originating from the above mentioned tangency zone. Guiding center trajectory numerical computations with the exact magnetic field and edge geometry fully confirms this analytical approach (fig.4).

Besides this result, the \(J_{\parallel}\) invariance allows a detailed description of the fast ion ripple losses inhomogeneity during NBI or ICRH, where the banana orbits distribution are driven by different mechanisms. Faraday collectors disposed in order to measure not only the poloidal angle but also the bounce amplitude of the escaping ripple trapped particles will allow to determine the location of their trapping point inside the plasma and constitute a potential method of ripple losses measurements with spatial resolution.

**REFERENCES**

Fig. 3: Topology of the iso-ripple (thick lines) and iso-B contours (thin lines) and of the tangency zone that corresponds to trapped particle with similar dynamics.

Fig. 4: Projections of the guiding center motion of a fast particle hitting the hot spot. The banana trapping point lays in a tangency zone iso-ripple, iso-B.
PLASMA DECONTAMINATION DURING ERGODIC DIVERTOR EXPERIMENTS IN TORE SUPRA


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I. Introduction

In Tore Supra an ergodic divertor (ED) has been integrated in the machine design and successfully operated, as already reported (1, 2). This paper analyses the decontamination effect resulting from the creation of an ergodic boundary zone. Two plasma geometrical configurations (outboard and inboard) are studied, the plasma being limited respectively either, on the low field side (lfs), by an outboard limiter (3 to 5 cm ahead of the ED modules) or, on the high field side (hfs), by the graphite inner wall. Strong decontamination effects have already been reported for the first configuration (3) by observing line emission of the intrinsic (carbon and oxygen) and purposely injected (nitrogen) impurities. When limited by the inner wall, the plasma is several centimetres farther from the ED modules than in the lfs configuration. The magnetic perturbation is then greatly reduced, and much smaller decontamination effects should be expected. In this paper, the hfs configuration data is compared with that from the lfs configuration. Preliminary experiments combining lower hybrid current drive and ED operation in the hfs configuration are also reported.

II. Ergodic divertor operation in ohmic plasmas

In this section we describe ohmic deuterium plasmas having an impurity content dominated by light impurities (carbon and oxygen), carbon being more abundant except at the largest electron density values. Light impurities are completely ionized in the plasma center, where therefore they are not accessible to spectroscopic line observation; however, as discussed in (3), the Lyman α line brightnesses for both carbon and oxygen, when divided by \( \bar{n}_e^2 \) (\( \bar{n}_e \) being the volume averaged electron density) are a reliable indication of the concentrations of the two elements. Both are plotted in figure 1 as a function of \( \bar{n}_e \). The dotted lines schematically show the concentration values without ED. This dependence on \( \bar{n}_e \) is observed in most tokamak devices. The solid lines are the concentrations obtained when the ED is activated with a maximum value of the magnetic perturbation (that is, in the lfs configuration, with the maximum current (\( I_{\text{ED}} = 45 \text{kA} \)) in the ergodic divertor coils). It has to be recalled that in this configuration, activation of the ED reduces the electron density (typically \( \Delta \bar{n}_e = 0.5 \times 10^{19} \text{ m}^{-3} \)). Consequently, when analysing quantitatively decontamination, one should compare discharges with the same \( \bar{n}_e \) value; since carbon is more abundant than oxygen, the carbon concentration reduction is dominant with respect to the oxygen concentration increase, and purification occurs (\( \Delta Z_{\text{eff}} \) being about 1 at \( \bar{n}_e = 1.5 \times 10^{19} \text{ m}^{-3} \)). This purification is due both to a change in the impurity sources (reduced edge temperature and different wall connection) and to impurity screening, as it has been proven by nitrogen injection experiments. Figure 2 (left) shows the
results obtained when injecting nitrogen in an outboard configuration plasma (solid and dotted lines, with and without ED, respectively); given from top to bottom, the $\bar{n}_e$ value, the normalized Ly$\alpha$ brightnesses of nitrogen and of the intrinsic carbon, along with the nitrogen feeding (black). Two puffs are injected: the first one before, and the second one during ED. The intrinsic carbon content is reduced in agreement with figure 1, whereas a large screening effect is seen on the nitrogen signal corresponding to the second puff, the N$^6+$ ion density being reduced by a factor of 2.5 during the ED. This screening is due to the presence of a peripheral ergodic layer reducing the impurity confinement time at the very edge of the plasma. However, because of the better knowledge of the impurity source, the nitrogen experiment clearly shows the screening effect without the complications introduced by the poor knowledge of the intrinsic impurity sources (and specially of their modifications induced by the ED perturbation).

The same injection experiment has been performed in the hfs configuration (figure 2 right). Note that in this experiment the wall conditioning resulted in a decreasing $n_e$ when the ED was applied (not the usual case for the hfs configuration). The intrinsic carbon content is still reduced, but the injected nitrogen is now practically unscreened. The reduction of the carbon content is now interpreted as a large modification of the carbon source, which is supported by two observations: the edge electron temperature drops to a very low value, and a highly radiating zone appears on the plasma high field side. However, there is no screening effect because, in the hfs configuration, the ED magnetic perturbation is very low (the plasma is farther from the ED modules). So the magnetic perturbation is large enough to lower the peripheral electron temperature (through a modification of the heat diffusion coefficient and/or through increased radiative losses), but not to affect the particle transport.

III. Combined ergodic divertor (ED) and lower hybrid current drive (LHCD) experiments

A series of combined ED and LHCD experiments were carried out in D$_2$ hfs plasmas with LH power levels between 0.5 and 2.5 MW. Results from this series have confirmed observations from earlier experiments of this kind (4) in which the divertor produced significant effects on the reflection coefficients of individual LH waveguides and on the impurity content of the plasma during the LHCD pulse. The timing sequence used in these experiments is exemplified in fig. 3 for shot number 5669. During the first 2 seconds of the ohmic phase $I_p$ (not shown) is ramped to a plateau value of 1.6 MA with $B_T = 3.3$ T. The loop voltage stabilizes at $V_1 = 1.3$ Volt by time $t = 3.0$ sec. The LH power is initiated at $t = 3.5$ sec, and reaches $P_{LH} = 2.5$ MW by $t = 3.8$ sec. The ergodic divertor is pulsed at $t = 5.0$ sec. with a current $I_{ED} = 36$ kA. The diverted LHCD phase lasts until $t = 6.5$ sec, and is followed by a diverted ohmic phase. At $t = 8.0$ sec, the divertor is switched off and $I_p$ is ramped down. The signal responses shown in fig. 3 indicate that the duration of each phase of the experiment is sufficient to approach a steady state condition and that the final diverted phase reaches a steady state in roughly half the time needed for either the LH or diverted LH phases.

It is noted that $V_1$ begins to decrease immediately with the application of the LH pulse while $\bar{n}_e$ first shows a small decrease followed by 50% increase during the LH phase. The C$^5+$ density shows a gradual reduction during this phase with a rate of change which is significantly less than that of $\bar{n}_e$. $V_1$ reaches a minimum at approximately the same time as $\bar{n}_e$ begins to increase and drifts slightly higher with $\bar{n}_e$. When the divertor is applied $\bar{n}_e$ initially decreases and $V_1$ takes a short upward transient excursion. The internal plasma inductance is increasing during this period as the resonant magnetic flux surfaces are being rearranged. At approximately the peak in this $V_1$ transient, the C$^5+$ density begins to decrease at a more rapid rate reaching a stable value about half way up the second, divertor induced, density rise. During this time $V_1$ has decreased to a value which is approximately the same as it had just before the divertor was applied. $V_1$ then remains constant as $\bar{n}_e$ increases during the remainder of the LH pulse. At the end of the LH pulse $V_1$ recovers its original ohmic phase value and $\bar{n}_e$ decays. The C$^2+$ density increases slightly during this final diverted phase but remains relatively low compared to the initial ohmic phase.
Data from fig. 3 is used to determine the C impurity concentration in each of the four discharge phases described above and is plotted in fig. 1 (● OH, ♦ LH, ▲ LH+ED, ■ ED) for comparisons with diverted and undiverted ohmic IFS plasmas. Diverted IFS plasmas have lower C concentrations than undiverted plasmas but only to the extent of the \( \bar{n}_e \) increase. Plasma density increases are typically observed when the divertor is pulsed in this configuration \((4,5)\) as opposed to a \( \bar{n}_e \) reduction during the divertor phase with IFS plasmas. The divertor and the LH both produce C concentration reductions which follow the IFS ohmic concentrations when \( \bar{n}_e \) is increased with gas injection. The divertor produces larger \( \bar{n}_e \) increases during the LHCD phase and lower C concentrations than in undiverted LHCD plasmas. In addition, both the divertor and LHCD, when used individually, have approximately the same effect on the C concentration for this value of \( I_{ED} \) and \( P_{LH} \). Again, impurity screening effects, such as those observed in the low-field limited configuration, apparently do not occur under these conditions. This is thought to be due to the lower effective magnetic perturbation levels in IFS plasmas. Fig. 4 shows the change in \( H_\alpha \) recycling and carbon density from the ohmic value (the zero value thus referring to the ohmic case) as a function of LH power. The boundary layer \( H_\alpha \) recycling continuously increases, whereas the C density shows a decrease only at the largest powers; however, the magnetic perturbation effect is apparently unaffected by the LH power increase. It should be noted that in this limiter configuration \( H_\alpha \) emissions from within the magnetically perturbed region of the boundary layer typically increase in \( D_2 \) plasmas \((4)\) while in He plasmas with the same conditions these emissions are either unchanged or show a slight decrease \((5)\).

IV. Conclusions

The ergodic divertor has been shown to have an important impurity decontamination effect in the IFS configuration (where the magnetic perturbation is stronger), due to both a change in the source terms and a screening by the ergodic layer. The screening effect is not observed in the hfs configuration, probably due to the much lower value of the magnetic perturbation. Combined ergodic divertor and lower hybrid current drive experiments have demonstrated that relatively large C density reductions can be maintained throughout the ergodic divertor pulse and that this reduction is larger with the LH than without it. For these hfs limited plasmas, it appears that this is due to a \( \bar{n}_e \) increase rather than a screening enhancement as seen in the IFS case. Because of a non zero loop voltage, the analysis of the divertor effect on LHCD efficiency has not yet been completed. These experiments confirmed earlier results \((5)\) demonstrating that the LH phase does not appreciably change the different divertor impurity responses in the IFS and hfs configurations.

V. Acknowledgements

This work has been performed as part of the US/France Collaboration on Fusion Research and is partially supported by the U. S. DoE under Contract No. DE-AC03-89ER51114.

VI. References

fig. 1: C and O concentrations versus $n_e$. Symbols refer to LH + ED plasma of fig. 3.

fig. 2: Identical nitrogen injections (black) in outboard and inboard plasma configurations.

fig. 3: Time evolution of a few parameters for LH (2.5 MW) + ED plasma.

fig. 4
CURRENT DIFFUSION AND FLUX CONSUMPTION IN TORE SUPRA


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1. CURRENT DIFFUSION EXPERIMENTS- TORE SUPRA has been designed to study long pulse plasmas (t ≥ 30 s) at high plasma current (I_p ≤ 2 MA) associated with high additional power (P_{add} ≤ 20 MW). Current diffusion studies are essentially based on the analysis of the plasma discharge paths in the plane (q_{w}, l_i) where l_i represents the internal inductance of the plasma and q_{w}, the safety factor at the edge of the plasma. The current diffusion rate during the current rise phase is analysed with a numerical code using plasma resistivity profiles from Tef profiles measured by the ECE diagnostic.

On TORE SUPRA, the current build-up between plasma ionization and the current plateau is established in two phases. A first phase, referred to as rapid, corresponds to the initial current rise with a rate between 3 and 5 MA/s. This phase is mainly dependent on the discharge breakdown conditions (electric field, filling pressure, leakage field and state of the walls) and is very often characterized by intense MHD activity and rapid current penetration. The dynamics of the current profile has been studied in the second phase, referred to as slow, where the plasma is well established and where it is possible to vary the current ramp rate (series 1).

The same analysis was carried out during a second current ramp up between two stationary phases (series 2).

The current ramp rate is the parameter which has the greatest effect on the variation of the inductance, l_i, in relation to q_{w}(a) (fig. 1). After a short time, during which the current profile is reorganized, the internal inductance l_i varies quasi linearly with q_{w}(a). The lowest internal inductance corresponds to the faster current ramp rate. MHD activity, which is not observed for cases with slow current ramp-up, appears in the vicinity of rational values of q_{w}(a) at higher ramp rate and increases with the ramp rate, until a major disruption occurs.

In order to study the variation of the current profile in the (q_{w}(a), l_i) space we have characterized each discharge by its initial inductance (after the fast current rise), l_i(0), by the variation of l_i with respect to q_{w}(a), dl_i/dq_{w}(a), and by its inductance during the equilibrium phase, l_i(t).

Figure 1
The quantity \( \frac{dl_i}{dq_y}(a) \) is practically proportional to \( \frac{dl}{dp} \) as long as there is no MHD activity (fig 2). As soon as MHD activity is observable, \( \frac{dl_i}{dq_y}(a) \) is quickly saturated with respect to \( \frac{dl}{dp} \). This saturation indicates an acceleration of the current diffusion rate.

For the other series, S2, a second current ramp, with different ramp rate, was applied after a first plateau was obtained. In that case, the initial \( l_i(1) \) values are different and so does the values obtained for \( \frac{dl_i}{dq_y}(a) \) for the same value of \( \frac{dl}{dp} \). However the linear dependence on \( \frac{dl}{dp} \), followed by saturation at high ramp rate are still observed. As for the previous series, MHD activity is detected shortly before saturation.

A numerical 1-D code computing the cylindrical current density profile from Michelson interferometry has been developed (fig.3). Peaked, flat or even hollow \( T_e \) current profiles have been observed, depending on the current ramp rate. We constructed \( T_e(\rho) \) profiles from \( T_e(\rho) \), where \( \rho \) labels the most likely magnetic surface radius. The poloidal magnetic field diffusion equation is solved iteratively using classical Spitzer resistivity as a first try. A \( Z_{eff} \) value independent of the radius is used, and neoclassical correction is neglected. An anomaly factor, \( A \), is then introduced in order to fit the plasma surface voltage \( V_s \). The current density profiles obtained with this code, have been successfully tested with a cylindrical \( \lambda' \) code used to study the tearing modes stability [1].

In the two series S1 and S2, this anomaly factor is smaller or of the order of 2. This is even the case at large \( \frac{dl}{dp} \), with large MHD activity and saturated \( \frac{dl_i}{dq_y}(a) \). The acceleration of the current diffusion indicated by the saturation of \( \frac{dl_i}{dq_y}(a) \) is due to an increase of the resistivity allowing the current to penetrate more rapidly.

![Figure 3](image-url)
It is clear that the same \( \frac{dI_p}{dt} \), applied to two different current profiles characterized by two different initial internal inductances, \( l_i(i) \), should not evolve towards the same equilibrium, characterized by \( l_i(f) \), following an identical path. This was tested by applying the same second current ramp with different time delays after the first current plateau. It was found that the derivative of the internal inductance with respect to \( q_W(a) \) depends practically linearly on its "initial value", \( l_i(i) \), as long as there is no MHD activity. For a given ramp rate, a value, \( l_i(i) = l_{io} \), can be found which results in a stationary current profile during the current build-up phase. The \( l_{io} \) value depends on the disturbance applied, \( \frac{dI_p}{dt} \), such as,

\[
l_{io} = l_i(i) + C \frac{dI_p}{dt} \quad \text{where} \quad C = 0.6 \text{ s/MA}.
\]

The duration of the TORE SUPRA ohmic discharges, typically 12 s between plasma breakdown and the current plateau, is long enough to reach steady state. To a given \( q_W(a) \) steady state OH discharge, corresponds a given current profile characterized by its specific inductance value \( l_i(f) = f(q_W(a)) \). Values higher than this equilibrium value result, during the current plateau phase, in major disruptions (figure 4).

The limits for stable operation of TORE SUPRA are shown in figure 4. Within that domain, the internal inductance may be modelled as

\[
l_i = (l_i(f) - C \frac{dI_p}{dt} - l_i(i)) \left(1 - e^{-t/\tau_D}\right) + l_i(i)
\]

Thus, given an initial (breakdown) value of \( l_i \), the current profile and plasma stability can be controlled by proper adjustment of the ramp-up rate.

2. FLUX CONSUMPTION - Owing to the fact that the quantity of magnetic flux available in a tokamak is limited, perfect knowledge is required of the various components of the flux consumed in order to minimize consumption and to be able to define a suitable transformer size for future high current tokamak projects.

Flux consumption in TORE SUPRA was analyzed using the Poynting's vector approach:

\[
\frac{1}{\mu_0} \iint_S \left( \vec{E}_\phi \times \vec{B}_\alpha \right) \cdot d\vec{S} = \int_v \frac{\partial}{\partial t} \left( \frac{B_\alpha^2}{2\mu_0} \right) dV + \int_v E_{\phi} j_{\phi} dV
\]

where \( \vec{E}_\phi \) is the toroidal electric field, \( \vec{B}_\alpha \) is the poloidal magnetic field and \( j_{\phi} \) is the toroidal current density. In order to define a purely inductive flux which is cancelled out with the plasma current at the end of discharge, and following Sugihara, we combined resistive losses with the loss term due to current diffusion. We carried out a systematic flux consumption study for various current ramp-up rates. These range from the slow scenario in which the current is built up within a period of time comparable to the current diffusion time to the rapid scenario in which a high level of MHD activity subsists throughout the whole build-up. In Figure 5, we show numerous results of flux consumption at the core for all current ramp-up rates to reach the current plateau in relation to that plateau current value. A very simple model allows an accurate estimation of the flux consumed at the plasma surface to achieve a current plateau, \( I_p \):

\[
\Delta \Phi_s = \Delta \Phi_D + \Delta \Phi_I = (\Phi_A + V_D \cdot t_{r,u}) + l_i I_p
\]

where \( \Phi_A \) is the flux dissipated on breakdown, \( V_D \) is the dissipative voltage and \( t_{r,u} \) is the plasma current ramp-up time.
As shown in dashed line in the figure 5, the flux consumed at core has a lower limit proportional to the plasma current:

$$\Delta \Phi_\text{n} = \Delta \Phi_\text{s} + \Delta \Phi_\text{ext} = 3.52 I_p = 1.18 \mu_0 R I_p$$

The corresponding limit for the flux at the plasma surface is

$$\Delta \Phi_\text{s} / \mu_0 R I_p = 0.82 \text{(Wb/MA)}$$

with for the inductive component,

$$\Delta \Phi_\text{i} / \mu_0 R I_p = 0.45 \text{(Wb/MA)}$$

and for the dissipative component,

$$\Delta \Phi_\text{D} / \mu_0 R I_p = 0.38 \text{(Wb/MA)}.$$  

In fact, the flux consumption calculated in such a way is usually not sufficient to reach steady state. The inductive flux continues to increase until the internal inductance reaches its equilibrium value, $L_i(f) = L_i(q_{\psi}(a))$. 

In Figure 6, we show, as a function of $dl_p/dt$, the consumption of inductive flux, $\Phi_\text{i}$, dissipated flux, $\Phi_\text{D}$ and we compare the flux at the plasma surface, $\Phi_{\text{plateau}}$, when the plasma discharge reaches the current plateau and the flux at the plasma surface, $\Phi_{\text{equi.}}$, at the moment when the plasma discharge reaches a state of equilibrium. The following conclusion may be made:

- The inductive flux consumed to reach the current plateau becomes smaller as the current ramp-up rate increases, in that the internal inductance, $l_i$ is that much smaller.
- The dissipative component is a quasi-linear function of time and the flux dissipated to reach the current plateau is increasingly large as the current ramp-up rate is reduced.
- However, the flux consumed to reach equilibrium is practically independent of the current ramp rate because the time required to reach steady state varies very little ($\approx 3$ s on TORE SUPRA) and the inductive flux consumed is determined by the equilibrium internal inductance $l_i(f)$ which is mainly a function of the $q_{\psi}(a)$ flat top value. The flux consumption at the plasma surface to reach a steady state is accurately estimated by the formula:

$$\Delta \Phi_{\text{Se}} = (\Phi_\text{A} + \Phi_{\text{D}}) + L_i(f) I_p$$

where $t_e$ is the time required to reach the equilibrium and $L_i(f) = L_i(q_{\psi}(a))$ is the equilibrium internal inductance. The flux consumption to reach a stationary current profile is not a direct function of the flat top current value but depends upon the current diffusion time $\tau_D$ ($\approx 1$ s on TORE SUPRA) and the $q_{\psi}(a)$ value. This flux can be two times larger than the minimum flux to reach the flat top current. Thus even if ramping the current faster results in a longer current plateau, the steady state phase is not extended.

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SCALING PROPERTIES OF RUNAWAY ELECTRONS IN TJ-I TOKAMAK
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INTRODUCTION

Runaway confinement time in ohmic and additionally heated tokamak plasmas, presents an anomalous behaviour when compared with theoretical predictions based on neoclassical models.

In this paper, a one-dimensional numerical model [1] including generation and loss effects for runaway electrons is used to deduce the runaway energy $\varepsilon_r$ dependence on the runaway confinement time. The simulation results are presented in the form of a scaling law for $\varepsilon_r$ on plasma parameters. The scaling of $\varepsilon_r$ and therefore the runaway confinement time, has been studied in the TJ-I tokamak ($R_0=0.3$ m, $a=0.1$ m, $b=0.12$ m, $B_T<1.5$ T, $I_p<65$ kA), by measuring hard-X-ray spectra under different experimental conditions. Runaway confinement time has been deduced for the range of plasma parameters available in the TJ-I. A tentative explanation for the scaling of the obtained data, based on effects from magnetic turbulence will be presented [2].

SCALING LAW FOR RUNAWAY ENERGY FROM 1-D MODEL

Hard-X-ray spectra have been simulated by calculating the bremsstrahlung radiation caused by runaway collisions with the background plasma. To this aim, electron distribution functions are obtained numerically following a simple scheme of deformation of a maxwellian distribution function: The source of fast electrons is created by those ones that overcome the critical velocity and run away from it; the gap left by these electrons is immediately (into the equilibration time of the bulk, assumed as the collision time $\tau_{ee}$) refilled; this process is repeated during all the acceleration time. Simultaneously losses are considered: The probability that a runaway born at $t_0$ is still in the plasma at $t$ is given by $\exp\left[\left(t-t_0\right)/\tau_r\right]$, where $\tau_r$ is a runaway confinement time for these runaways in the plasma without any previous consideration for driven mechanism for losses generation; when $\tau_r>>\tau_{ee}$ we obtain a slide away distribution function. The kinetic model that is applied to obtain the runaway energy has been described in [1], and considers electron acceleration in a parallel electric field $E_{||}$ and drag forces due to ion-electron and electron-electron collisions and impurities. Input parameters are $n_e$, $T_e$, $Z_{eff}$, $E_{||}$ for the source and $\tau_r$ for the losses. The output value is the inverse slope of the simulated hard-X-spectra, that we called "runaway
average energy" $\varepsilon_r$. To obtain a stable distribution function the simulated acceleration time is
taken $t = 10 \tau_r$. Using more than 100 different simulations the following scaling law is
obtained for $\varepsilon_r$:

$$\varepsilon_r = (n_e^{0.16 \pm 0.001} T_e^{0.07 \pm 0.01} Z_{\text{eff}}^{-0.02 \pm 0.01}) V_L^{+1.04 \pm 0.01} \tau_r^{+1.1 \pm 0.01}$$

(1)

where $\varepsilon_r$ is in KeV, $n_e$ is in cm$^{-3}$, $T_e$ in keV, $V_L$ in V and $\tau_r$ in ms.

In figure 1 numerical results are compared with the obtained scaling law and with the
energy obtained from free fall acceleration in the electric field in a time $\tau_r$. In our model
electrons seem to be accelerated longer ($\tau_r \sim 1$) than in the free fall case. As in free fall case
runaway energy is almost proportional to $V_L$. The expression in parenthesis plays the role of
a constant, since its value does not change in spite of the variation range of $n_e$ (0.1-1.5x 10$^{13}$
cm$^{-3}$), $Z_{\text{eff}}$ (1-6), $T_e$ (0.2-0.7 keV) used for these simulations (see figure 2).

**EXPERIMENTAL SCALING**

Using the expression (1) we can deduce experimentally the value of $\tau_r$, as a global
measurement, assuming that the inverse slope of the hard-X-ray spectrum is the mean energy
of runaway electrons in the plasma. Data have been obtained under different operational and
plasma conditions, $B_T$ (0.8-1.4 T), $I_p$ (16-44 kA), $V_L$ (2.8-5 V), $n_e$ (0.25-2.5 x 10$^{13}$ cm$^{-3}$),
$T_e$ (0.1-0.7 keV), $Z_{\text{eff}}$ (1.75-7.5) in a total of 35 spectra. The range variation of the results
for $\varepsilon_r$ is 60 keV to 900 keV, that gives variation in the confinement time from 0.1 to 2.5 ms.
Figure 3 shows the evolution of the confinement time with the density in TJ-I (bottom
curve), together with the values of the expression in parenthesis in (1). It can be check that
this term is almost constant in all this experimental range in accordance with the results from
the simulation.

Experimentally the runaway confinement dependences on other parameters that have
not being included in our model, may show what kind of phenomena are governing losses.
Electron energy confinement time and runaway confinement time are plotted in figure 4 and 5
as function of safety factor at $a=9.5$ cm (last magnetic surface). As a general tendency from
figure 4 and 5, $\tau_r$ and $\tau_{\text{EC}}$ seem to be decorrelated since $\tau_r$ increases with $q(a)$ while $\tau_{\text{EC}}$
decreases. But if we choose discharges with similar densities, we can see that for $n_e > 10^{13}$
cm$^{-3}$ (open circles), both $\tau_r$ and $\tau_{\text{EC}}$ decrease with $q(a)$, and for $n_e \leq 10^{13}$ cm$^{-3}$ (black circles)
$\tau_{\text{EC}}$ is almost constant while $\tau_r$ increases. $\tau_r/\tau_{\text{EC}}$ is plotted in figure 8 and its ratio becomes
below 1 for densities higher than $10^{13}$ cm$^{-3}$. These results could suggest in accordance with
the OH studies at the ASDEX tokamak [2] that runaway losses are governed by a turbulent
process depending on electron density value. The evolution of runaway drifts calculated for
the value of $\varepsilon_r$ on the last magnetic surface is plotted depending on the density (figure 6) and in figure 7 the experimental evolution of $\tau_r$ with the drift that shows a clear improvement in confinement with the drift. These facts are in good agreement with ASDEX scaling studies that seem to be explained by resistive balloning modes [2], as well as by Alfvénic microturbulence. Figure 9 compares runaway confinement time in TJ-I with the scaling given in [2].

**CONCLUSION**

A good agreement between experimental values for runaway confinement time and ASDEX results for magnetic turbulence has been obtained in TJ-I tokamak. Magnetic turbulent regime seems to be present only for low and intermediate density cases, and for higher densities the drift effects tend to vanish. More measurements are in progress in order to obtain at the same time the value of $b_0$, the radial correlation length and the spectra of hard-X-ray fluctuations.

Figure 4: Energy confinement dependence on q(a) for thermal electrons.

Figure 5: Runaway confinement time dependence on q(a).

Figure 6: Calculated runaway electron drift as function of electron density

Figure 7: Runaway confinement time behaviour with drift.

Figure 8: Variation of $\tau_{nr}/\tau_{Te}$ on electron density

Figure 9: Fitting of $\tau_r$ with the inverse of $b_0$ given by Kwon [2]
EFFECTS OF ELECTRODE POLARIZATION AND PARTICLE DEPOSITION PROFILE ON TJ-I PLASMA CONFINEMENT

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1. Introduction. The role of self-created radial electric field on particle confinement in TJ-I plasmas was addressed in Ref. (1), using plasma rotation data in conjunction with particle confinement times measured by laser ablation (2). In this paper following the pioneer work of Taylor (3), we have started to study the influence of a polarized electrode inserted into the plasma on particle confinement and plasma rotation in this ohmically heated tokamak. To have a supportive frame of reference, the confinement time of background particles and their transport into the plasma, without electrode, has been studied by measuring with space-time resolution the \( H_\alpha \) emission on varying plasma conditions. These experiments have been carried out in ohmically heated discharges of the TJ-I tokamak \((R_0 = 30 \text{ cm}, a = 10 \text{ cm})\) which was operated with plasma currents between 20 and 45 kA and a toroidal field ranging from 0.8 to 1.5 T. In this paper, firstly the experimental plasma and specific diagnostics are described, secondly, the parametric dependence of the particle confinement time and radial transport of background plasma is presented and finally, the influence of polarizing an inserted electrode on a particular discharge is given and discussed in the context of other polarization experiments.

2. Particle confinement time in TJ-I. Although several models have been proposed to predict the particle confinement properties of fusion plasmas in terms of the particle deposition profile (4), experimental information about the possible correlation is scarce. Since one of the relevant parameters is the ratio of the mean ionization radius for neutrals, \( \lambda_{iz} \), to the plasma minor radius, these models should be tested in machines of different sizes. Although particle confinement (5), (2) and poloidal distribution of the source term (6) have been investigated in some detail in the TJ-I small tokamak, the possible correlation between particle deposition profiles and particle confinement has not been studied yet. In the present work, the dependence of the global particle confinement time on electron density and toroidal field and the \( H_\alpha \) radial emission profiles and their fitting to simple model calculations for some of the cases studied are reported.

Relative particle fluxes were evaluated from the line integrated \( H_\alpha \) emission intensity perpendicular to the plasma in the equatorial plane at two different toroidal positions and electron densities were measured with a 2 mm interferometer located at the central chord. A set of two optical fibers, interference filters and detectors were used in reproducible discharges to measure the \( H_\alpha \) emission profile from a top window, shot to shot. The line of sight used in
the outer side was slightly tilted in order to avoid the wall gap existing at that position. A spatial resolution of 1 cm at the equatorial plane was used for the experiments here reported. Figure 1 shows the evolution of the global particle confinement time, $\tau_p$, with line average electron density for three values of the toroidal field, $B_T = 0.8$, 1, and 1.4 T, at a constant value of the plasma current of 35 kA. The data were taken during the density plateau, at $t = 8$ ms. The volume average proton densities required for the calculation of $\tau_p$ were evaluated from the interferometer data by correcting them with the relative contribution from impurities (oxygen) to the electron density and the evolution of the density profile with line average electron density when available (7). The recycling coefficient was evaluated from the ratio of $H_\alpha$ signal to gas puffing level (5). A constant value was found for the conditions of the experiments here reported, indicating that saturation of the part of the chamber acting as toroidal limiter in TJ-I takes place in the discharge time for the fluxes here involved ($\approx 10^{18}$ cm$^{-2}$ s$^{-1}$). The interaction area was assumed to be the same for all the conditions (6). As it can be seen in the figure, a different scaling of $\tau_p$ with electron density is observed depending on the toroidal field value. For low densities, a $B_T^{-2}$ scaling is found, in agreement with previous determinations (2) whereas at higher densities the dependence is almost linear. The fueling efficiency was evaluated from the dependence of the electron density at the plateau with the amount of injected gas. For the experiments at 1 T and 1.4 T the same value was found for the whole range of electron densities here studied.

The radial profiles of $H_\alpha$ emission, integrated along a vertical chord, for $\tilde{n}e = 1.2$ and $2 \times 10^{13}$ cm$^{-3}$ and $B_T = 1$ T are shown in figure 2. As it can be seen, aside from the comparatively stronger interaction with the inner wall in the case of high density, in agreement with previous
determinations (6), the emission profiles for the two values of electron density are very similar. In order to evaluate the average value of the poloidally averaged source term from the emission profiles reported, a simple model was developed. The distribution of neutrals inside the plasma is modeled as the product of a radial function, depending on electron density profile and neutral attenuation, and a poloidally asymmetric source term. Due to the small size of the TJ-I tokamak, a "mirror-like" effect, that concentrates particles from the outer wall, was included. The effect of ion temperature profile in the contribution of the density of CX neutrals to the source term was found to be small due to the compensation effect of penetration and velocity. Neither the S.O.L nor the effect of molecules at the edge was included in this simple model, the atoms being produced at the plasma periphery with a typical Frank-Condon energy of few eV. As it can be seen in the figure (dashed lines) good fitting to the experimental data can be obtained. The electron density profile used to fit the data at higher density is narrower than the one required for lower density, in agreement Thomson scattering data in TJ-I (7), and the ratio of central to edge electron densities used agree with previous probe measurements (6). No significant difference in the radial position of the average source term was found for the two cases studied (r = 6.8 cm). Single shot Hα profile studies at 1.5 T, where a strong evolution of confinement time with electron density is observed, are in progress.

3. Electrode influence on plasma confinement. To elucidate the role of radial electric field on TJ-I particle confinement, an isolated electrode has been introduced up to a reduced radius of 0.7. Its influence on plasma rotation and particle confinement, when polarized between ±450 V with respect to the vacuum chamber, was studied in a low current (I_p = 20 kA) and rather moderate density discharge n_e = 1.5x10^{13} cm^{-3}. The electrode polarization is achieved with a power supply made of electrolytic capacitors which is triggered at a prefixed discharge time by means of a thyristor, a maximum radial current of 40 A has been driven. This type of discharge was chosen to avoid a rapid electrode deterioration due to runaway electrons or to an excessive plasma thermal load. Since the TJ-I tokamak does not have a position feedback system the preprogrammed vertical field and primary OH voltages must be slightly readjusted when the electrode is inserted into the plasma to keep constant the current and plasma position.

The Langmuir characteristic of the electrode was studied within the voltage range quoted above, which allows us to estimate the plasma potential and electron temperature at the electrode tip position. The values of these parameters so deduced support, along with data from a CCD camera, that the electrode itself is not limiting the plasma size. The maximum density achieved as a function of the electrode polarization, keeping constant the background hydrogen pressure and the amount of gas puffed into the discharge, was studied and is plotted in Fig. 2. Higher density values were achieved in this particular discharge for electrode polarizations in the range of -350 to +50 V. This line average density increase seen by the microwave interferometer is also observed by a Langmuir probe located at the plasma edge.
The rise in density is not due to impurity production as evidenced by the behaviour of visible continuum and pyroelectric signatures. Therefore, it is attributed to a direct influence of the applied electric field. The observed confinement improvement does not correlate with maximum radial current but rather well with a particular range of negative voltages with respect to the natural plasma potential of the plasma at the electrode tip position.

The influence of electrode polarization on poloidal plasma rotation was monitored measuring the relative change in the spectral shift of the OV (2781 Å) line along a line of sight with a chordal radius of 6.5 cm. When these data shift were converted to absolute rotation by reference to the plasma central chord, the deduced poloidal rotation as a function of the electrode polarization obtained is shown in Fig. 2. It must be noticed that the improvement in confinement correlates with this reduction on poloidal rotation which must be due to a reduction of the radial electric field contribution to this flow, produced at determined electrode polarization. The experiment is being extended to more general discharges and for this aim an upgrading of the electrode power supply by a factor 2.5 has been already performed.

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Heat transport in tokamaks is empirically characterized by the existence of different regimes. In ohmically heated discharges the energy confinement time $t_E$ first increases linearly with density (the LOC regime) and then saturates at high density (the SOC regime). In additionally heated discharges $t_E$ degrades with the heating power $P$ and improves with the plasma current $I_p$, irrespective of the heating method. This is the L-mode confinement.

The aim of this paper is to derive a scaling which describes the LOC, SOC and the L-mode. L-mode confinement is now observed with the most different heating methods, including discharges with non-inductive current drive; it would therefore be surprising that the basic mechanism underlying its behaviour is not active in ohmic conditions. Moreover it was observed in FT that some of the scaling proposed for the L-mode can give a satisfactorily account of the ohmic confinement in a wide range of plasma parameters [1].

When inspecting ohmic data, it is generally recognized that the LOC regime is characterized by a decrease of $Z_{eff}$ and of the electron temperature $T_e$ with increasing plasma density $n$ and that SOC is reached when $Z_{eff}$ becomes constant and $T_i = T_e$ [2]. This suggests that LOC confinement could be caused by the same mechanism responsible for power degradation in the L-mode, for example by a thermoconductivity increasing with temperature. Therefore we will first concentrate on the SOC regime where $Z_{eff}$ variations are less important. For this regime scaling have been proposed by Shimomura et al [3]. Both these scaling however fail to fit FT and Asdex data and, more importantly, do not respect the dimensionality constraints of Connor and Taylor [5].

These constraints are a very useful tool because they allow to determine completely the size dependence of the scaling once the density and magnetic field dependence are known. Therefore, in order to check their validity we have selected, from the ohmic data set of FT and Asdex, a subset of shots with the same value of $B_t R_1^{1.25}$ and we have compared the behaviour of the normalized confinement $t_E B_t$ as a function of the normalized density $n/B_t^{1.6}$. The result is given in Fig.1 which shows that dimensionally equivalent shots have similar values of the normalized confinement as enforced by dimensionality constraints.

The two mentioned scaling reproduce the isotope effect as observed in FT and Asdex too, as well as the weak dependence of $t_E$ in SOC upon the plasma density and the edge safety factor $q_a$. Accurate $B_t$ scaling turns
out to be impossible with data from a single device because of the interference of the density limit. We have therefore collected in the published literature data from different devices and we have found the following scaling (units used throughout this paper are MW, MA, Tesla, 10^{13} \text{cm}^{-3} \text{ and meters)}

\begin{equation}
\tau_E^{SOC} = 6.0 \times 10^{-2} a R_0^{0.875} B_t^{0.5} k A_i^{0.5}
\end{equation}

which has to be considered approximate because we have neglected the weak n and qa dependence and because the data set is biased by the higher Z_{eff} of larger devices probably caused by the use of carbon walls.

This result can be used to analyse the LOC regime when the link of the plasma temperature T_e with density and magnetic field, imposed by the ohmic heating constraint, is released. With a little algebra we obtain

\begin{equation}
\tau_E = \tau_E^{SOC} (1 - 0.6 \alpha) n^{0.6 \alpha} k_0^{0.6 \alpha} R_1^{1.2 \alpha} B_t^{-0.2 \alpha} I_p^{\alpha} P^{-\alpha}
\end{equation}

where the dependence upon Z_{eff} has been neglected as usually done for all L-mode scaling.

At this point the scaling is completely determined once the exponent of the power term \alpha is known. Published results usually give 0.5 < \alpha < 0.6 when fast particle contribution is not corrected for and \alpha = 0.75 when this is done. We have also derived \alpha for a set of lower-hybrid heated discharges at low density in Asdex where the increase in plasma energy occurs only in the electrons which are well diagnosed by kinetic measurements and we have found \alpha = 0.75. This scaling has been applied to a data set of ohmic discharges from FT, Asdex and JET, covering the LOC and
SOC regimes. The comparison with the experiment shows a systematic underestimate of high confinement data irrespective of the value of $\alpha$. This is imputed to the approximations made in the derivation of $\tau_E^{\text{SOC}}$. Regression of the errors gives an additional factor $(n^2)^{0.1}(B_t R^{1.25})^{0.1}$ and, with $\alpha=0.75$, we finally obtain

$$\tau_E = 1.92 \times 10^{-2} n^{0.55} A_t^{0.3} k a^{0.55} R^{1.7} B_t^{0.2} I_p^{0.75} p^{-0.75}$$

which when compared with the ohmic data set, Fig. 2, gives a good agreement and a mean square root deviation of 16%.

Fig. 2

The experimental confinement time of ohmic discharges in FT, Asdex (with and without boronized walls) and JET is plotted versus the scaling given by equation (1). Also shown are LH heated discharges of Asdex.

Except for the power exponent, only ohmic observations have been used to derive this scaling. Its prediction capability has been tested for Asdex LH discharges also shown in Fig. 2. A more stringent test has been done by using the Kaye dataset for neutral beam heated discharges. In order to minimize the contribution of fast ions, the comparison has been restricted only to data with $n R / B_t > 2$. The result is shown in Fig. 3 and gives a mean square root deviation of 19%.

Finally when we convert the scaling to an ohmicon one we obtain

$$\tau_E^{\text{OH}} = 2.8 \times 10^{-2} n^{0.2} A_t^{0.5} k a R^{1.45} B_t^{0.7} Z_{\text{eff}}^{-0.5}$$

which compares satisfactorily with published data of saturated confinement.

In conclusion we have shown that by means of simple arguments and imposing dimensionality constraints it is possible to construct a scaling which is able to predict tokamak confinement in the LOC, SOC and L-mode.
This should possibly indicate that there exist a main transport mechanism which is responsible for the observed confinement in the three regimes.

![Graph](image)

**Fig. 3**

The experimental confinement time for neutral beam heated shots in the Kaye dataset (1370 data points) is plotted versus the scaling given by equation (1)

**ACKNOWLEDGMENTS**

Thanks are due to Dr.s G. Bracco, P. Buratti, S.M. Kaye, A. Tanga, F. Wagner for providing the data used in this work.

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The diamagnetic measurements of the toroidal magnetic flux provides on tokamaks a direct evaluation of the perpendicular beta poloidal of the plasma $\beta_p^\perp$. The diamagnetic measurement is performed on FTU by compensated diamagnetic loops that are mounted on the inside of the toroidal field magnet. The signal of the main loop that surrounds the plasma is compensated by the difference between the signals of two auxiliary loops (one external and the other internal to the main one) that just measures the vacuum toroidal flux. The most careful zeroing of the compensated signal in absence of the plasma has been performed by trimerizing all the electronics that makes the analogue signal processing; however, due to the time evolution of the spatial ripple of the toroidal field, such zeroing does not produce a zero voltage signal. At the best regulation one has obtained a reproducible signal for a given toroidal field current waveform (in absence of any other machine current). The reproducibility of the signal was perfectly constant during months within the arbitrary addition of an offset and of a linear ramp both due to the minimal thermal drifts of the analogue electronics. This has allowed to obtain the real diamagnetic signal by a simple subtraction and allowing for an additional offset and ramp (see Fig. 1b).

This operation was performed on two independent sets of compensated diamagnetic loops, one sitting on the minimum and the other on the maximum of the toroidal field ripple.

The two remaining relevant systematic error sources are: the spurious vertical field pick-up (due to slight planar misalignments of any of the loops) and the plasma

![Fig. 1 a) Plasma current for FTU shot 1794. b) Rough compensated diamagnetic signals for the two independent set of loops where the arbitrary offset and ramp is shown. c) $\beta_p^\perp$ values from the two set of loops and kinetic $\beta_p^k$ value](image)
current pick-up (due to slight nonplanarities of any of the loops). The former problem
was solved by subtracting from the compensated signal a quantity proportional to the
poloidal flux function difference measured by the equilibrium system across the
vacuum vessel with a proportionality constant determined by a pure vertical field
shot. The latter one was dealt with by subtracting the plasma current with a
proportionality constant determined by reversing the toroidal field direction without
changing (as detectable) any other plasma parameter. The amount of both corrections
is less than 1 mWb for a plasma current $I_p = 1$ MA, where the maximum true
diamagnetic signal is about $\Delta \phi = 10$ mWb at a toroidal field $B_T = 6T$.

The well known cylindrical relation $\beta_p^* = 1 - 8 \pi n (B_T \Delta \phi) / (\mu_0 I_p^2)$ has been used; it has been extensively checked by a predictive equilibrium code that in the case of
FTU the finite aspect ratio correction to such a simple formula are negligible. Figure
1b shows the rough compensated diamagnetic signals from the two sets of coils and
Fig. 1c shows the derived $\beta^*$ values; along them the kinetic $\beta_{ik}$ is shown as obtained
from electron temperature profiles as measured by ECE and line averaged density as
measured by DCN interferometer assuming parabolic density profiles and $T_i = T_e$.

Through the analysis of the poloidal fields and fluxes distribution around the plasma one can obtain a measurement of $1/2 \beta_p^* + 1/2 \beta_{ik} + 1/2 \beta_1$ [1] where $\beta_p$ is the parallel poloidal beta and $\beta_{ik}$ is the energy defined internal inductance. In FTU a set of
16 poloidal field coils, 16 voltage loops and 16 saddle coils are used for the equilibrium measurements; their main systematic errors are due to the spurious toroidal field pick-up and to the non zero sum of the saddle coils. The former has
been cured by subtracting the signals of a pure toroidal field reference shot. The latter turned out to be very small, up-down symmetric and proportional to the
applied loop voltage because of non perfectly symmetric up-down mounting of the twisks of the saddle coils and it has been properly distributed on the saddles.

An independent measurement of the poloidal beta is then possible from the
equilibrium analysis if the two quantities $1/2 (\beta_p^* + \beta_{ik})$ and $1/2 \beta_1$ can be evaluated separately. The parametric analysis performed by the predictive calculation of a
variety of plasma equilibria in the FTU configuration showed exhaustively that, whereas in perfectly circular plasma conditions a sum of the two quantities exceeding
1.2 allows such separate evaluation, a slight elongation (1.03-1.05) of the plasma
cross section allows the separation even at values of the sum like 0.7+0.8 [2]. A full
confirmation of this fact results from a thorough analysis of hundreds of FTU shots
which exhibit the degeneracy removal due to the elongation and a small triangularity
of the plasma shape (see Fig. 3), even at low values of poloidal beta. The agreement
of the equilibrium beta with the measurement of $\beta_{ik}$ from the diamagnetic coils and
with the $\beta_{ik}$ from kinetic measurements is generally good (see Fig. 2).

A very large discrepancy appears (as expected) in runaway shots where the parallel component of the electron poloidal beta becomes dominant; however this
situation will not be discussed in this paper.

A more remarkable discrepancy is observed as a transient behaviour in a
number of FTU shots. Figure 2b shows such an event in a 350 kA shot during the time
interval 0.15 $< t < 0.33$ s. The value of $\beta_p^*$ exceeds $\beta_{ik}$ up to the impressive value of
0.7 (the latter being consistent with $\beta_{ik}$). The event is distinctly visible on the
behaviour of two chords of the DCN Interferometer: the line averaged density at -18
cm with respect to the geometrical axis of the FTU vacuum vessel decreases sharply
and stays quite low until the event terminates; on the other hand the line averaged
density through the center starts climbing, then stays constant during the decrease of
the beta anisotropy and finally declines (Fig. 2c). The values of $\beta_{ik}/2$ is consistent (at
least after the current rise) with the values of $I_i/2$ as deduced from ECE
temperature profiles assuming Spitzer resistivity (Fig. 2d). The sum $\beta_{ik}^p + I_i/2$
(Fig. 2e) shows that during the event some energy is missing from the kinetic part
whereas it is present on the $1/2 \beta_p^* + 1/2 \beta_{ik} + 1/2 \beta_1$ of the equilibrium. This extra
Fig. 2 a) Plasma current for FTU shot 1531

b) $\beta_p^{\perp}$ from diamagnetic, $1/2 \beta_p^{\parallel} + 1/2 \beta_p^{\perp}$ from equilibrium and $\beta^{\perp k}_p$ from kinetic data
c) Line averaged electron density on the central and the $r=\pm18 \text{ cm}$ chords
d) $l_{1/2}$ from equilibrium and $l_{1/2}^{k}$ from kinetic data
e) $1/2 \beta_p^{\parallel} + 1/2 \beta_p^{\perp} + 1/2 l_1^{k}$ from equilibrium and $\beta^{\perp k}_p + 1/2 l_1^{k}$ from kinetic data
f) $W_{\text{dia}}$-diamagnetic isotropic plasma energy and $W_{\text{tot}}$-total plasma energy from equilibrium and diamagnetism data.
g) $R_{\text{axis}}$ equilibrium magnetic axis and $R_{\text{axis}}^k$ ECE plasma axis (without plasma fields corrections)
h) neutron emission
i) Dynamic confinement time of the isotropic part of the plasma energy
energy is calculated to be roughly 10 kJ (Fig. 2f). A further evidence of the high value of $\beta_p^{\perp}$ as indicated from the equilibrium is obtained by comparing the magnetic axis position as determined by the equilibrium analysis with the weight centre of the electron temperature profile as measured by ECE (Fig. 2g).

The neutron emission is remarkably enhanced during the event reaching its maximum at the maximum of the density (Fig. 2h). The dynamic energy confinement time as deduced from the diamagnetic and magnetic measurements is strongly enhanced during the transient phenomenon (Fig. 2i).

The onset and the evolution of this phenomenon is connected with the plasma position: for its onset it seems necessary that the plasma is in contact with the external limiter being removed from the internal one to such an extent that no lines of force connect the two limiters (Fig. 3 b,c); when, as a matter of fact, during the event the plasma reapproaches the internal limiter the discrepancy starts to decline (Fig. 3d). This transient phenomenon has been detected up to a plasma current of 700 kA and at any operating density. However it becomes more relevant and detectable at low plasma current and density, as the amount of extra energy measured by the equilibrium seems limited to 10-15 kJ.

A pressure anisotropy in the electrons seems to be ruled out by the fact that no signs of non-maxwellian distribution function occur on the ECE spectra. A pressure anisotropy in the ion distribution function seems ruled out by the behaviour of the neutron emission that follows the increase of the density and the reapproching of the ion temperature to the electron one at higher densities. So a possible tentative explanation is that the extra energy is due to the bulk plasma toroidal rotation (with a velocity of the order of the ion thermal velocity) induced by the plasma becoming non-neutral with the contemporary charging of the limiters that are both electrically floating. This picture could be supported by the critical role that is played in the event by the magnetic field line connection in between the two limiters and even the decrease of the internal chord density could be explained by the centrifugal force decoupling the plasma density from magnetic surfaces. Further experimental work is underway to get a more accurate and reliable picture of the phenomenon.

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Confinement Studies of Circular and X-Point Plasmas in the COMPASS-C Tokamak


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Introduction
A wide variety of plasma conditions in COMPASS-C is investigated to examine how the observed confinement compares with scaling laws and whether the separatrix configurations available can lead to conditions of improved confinement with auxiliary heating. Stable inboard X-point configurations [1] have been obtained with \( I_p = 70 - 100\, \text{kA}, B_T = 1.04 - 1.4\, \text{T} \) and \( \bar{n}_e = 0.5 - 2 \times 10^{19}\, \text{m}^{-3} \). Current rampdown experiments (150kA to 100kA in 20ms), preceding, during and following X-point formation, were investigated. Comparisons are made with limiter bounded plasmas. Additional heating is provided by 60GHz, second harmonic ECRH (X-mode, low field launch). Results are reported on the confinement of energy and particles in hydrogen, deuterium and helium plasmas and with a boronized vacuum vessel [2].

Confinement in ohmic heated plasmas
Diamagnetic measurements of the energy confinement time are compared in Fig. 1 with those expected from neo-Alcator scaling. The \( \tau_B \) is obtained from \( \tau_B = W/\dot{P}_{\text{oh}} \) and at times when \( \dot{W}/\dot{dt} \) is small (i.e. \( \dot{W}/\dot{dt} < < W/\tau_B \)). For centred circular cross-section discharges, the minor radius is given by the limiter radius \( a = 19.6\, \text{cm} \). In the case of inboard X-point discharges [1] (elongation \( \sim 0.9 \)), we take the equivalent-area radius \( a = 16.5\, \text{cm} \) and a geometric correction to the diamagnetic beta, \( \delta \beta_{\text{pol}} = +0.06 \), as obtained from flux contour reconstructions. The \( \tau_B \) in the unsaturated regime is close to the neo-Alcator value for both configurations. At the maximum current of \( \sim 150\, \text{kA} \), the confinement saturates at \( \bar{n}_e \approx 8 \times 10^{19}\, \text{m}^{-3} \), with \( \tau_B \approx 8.5\, \text{ms} \) and \( \sim 10\, \text{ms} \) for He and deuterium plasmas, respectively. Laser ablation of aluminium in deuterium plasmas suggests that \( \tau_p \approx 5\, \text{ms} \) for \( I_p \approx 100\, \text{kA} \) and \( \bar{n}_e \approx 1.2 \times 10^{19}\, \text{m}^{-3} \) for both circular and X-point ohmic discharges.

Confinement in ECRH plasmas
In ECRH experiments, the RF power (\( \leq 760\, \text{kW} \)) is launched into the plasma from the low field side. The second harmonic resonance position is generally close to the axis (\( \tau_{\text{res}}/a \leq 0.17 \)) for this series of plasmas. The absorbed power, \( P_{\text{ech}} \), is assumed to be 90% of the gyrotron power. The heating power exceeds the ohmic power of the target plasmas by a factor of \( 1 \rightarrow 6 \).

The dependence of the plasma energy on the total power, \( P_{\text{tot}} = P_{\text{oh}} + P_{\text{ech}} \), is shown in Fig. 2 for deuterium plasmas. Ohmic X-point plasmas have \( P_{\text{oh}} \approx 100\, \text{kW} \), central ion temperature \( T_i(0) = 140\, \text{eV} \) and central electron temperature \( T_e(0) = 550\, \text{eV} \). The diamagnetic plasma energy is in good agreement (\( \pm 10\% \)) with that obtained from the temperature and density measurements. At \( P_{\text{ech}} = 400\, \text{kW} \), the electron temperature reaches \( T_e(0) \approx 1.1\, \text{keV} \).
and the ohmic power drops to \( P_{\text{oh}} \sim 60kW \). Only small changes (\( \leq 5\% \)) to the ion temperature are observed. The increase in plasma energy with the total input power shows an offset-linear dependence, with a corresponding incremental energy confinement time of \( \tau_{\text{inc}} \sim 0.6ms \). The \( \tau_{\text{E}} \) decreases by a factor of \( \sim 2.2 \) when the total power rises from the ohmic level of \( \sim 100kW \) to \( \sim 470kW \). These observations are also consistent with the L-mode-like scaling of \( \tau_{\text{E}} \propto P_{\text{tot}}^{0.6} \). However, the values given by Kaye–Goldston[3] and Goldston[4] scalings overestimate the results by a factor of 2.5 and 3.2, respectively (cf. Fig. 2). The offset-linear scaling of Rebult–Lallia[5] provides a lower \( \tau_{\text{inc}}(\sim 0.23ms) \) than observed but furnishes the better absolute value agreement, cf. Fig. 2. Laser ablation of aluminium gives \( \tau_{\text{p}} \sim 3ms \) for \( I_{\text{p}} \sim 90kA \) and \( n_e \sim 1.2 \times 10^{19}m^{-3} \) for both circular and X-point ECRH discharges.

When ECRH starts before the simultaneous rampdown and plasma shaping, with the X-point detached from the inside limiter (1 \( \rightarrow \) 2cm) for \( t \geq 200ms \), \( \tau_{\text{E}} \) is only slightly less (\( \sim 10\% \)) than in the corresponding ohmic cases. However, when the ECRH is applied only after the X-point formation, the discharges have a similar \( \tau_{\text{E}} \), irrespective of whether or not current rampdown is employed. (cf. Fig. 2).

Results from slightly off-axis heating (\( \tau_{\text{cas}}/a = 0.3 \)) in circular deuterium plasmas are also shown in Fig. 2. The \( \tau_{\text{inc}} \) is similar to the X-point plasmas (\( \tau_{\text{inc}} = 0.44ms \)) and the \( \tau_{\text{E}} \) degrades by a factor of \( \sim 4 \) when \( P_{\text{tot}} = 96 \rightarrow 634kW \) leading to \( \tau_{\text{E}} \propto P_{\text{tot}}^{-0.7} \). In these cases, the Kaye–Goldston scaling overestimates the experimental results by a factor of \( \sim 1.5 \). The Rebult–Lallia scaling underestimates the observed \( \tau_{\text{inc}} \) by a factor of \( \sim 1.7 \).

The results of ECRH in low density hydrogen X-point plasmas are shown in Fig. 3, in which the total power per particle is more than twice those displayed in Fig. 2. Applying ECRH generally increases the density from \( n_e \sim 0.5 \times 10^{19}m^{-3} \) to \( \sim 0.8 \times 10^{19}m^{-3} \) in this case. The low initial ohmic \( \tau_{\text{E}} \) may account in part for the small change in \( \tau_{\text{E}} \) with increasing power.

In none of the various plasma conditions examined, for X-point configurations, was there any of the usual signatures associated with a transition to H-mode, or improved confinement, (e.g. reduction in \( H_\alpha \) light intensity, appearance of ELMs or a change in rotation, as seen by a high resolution visible spectrometer, with \( \sim 20\mu s \) time resolution and \( \sim 1km \) velocity resolution, which monitored boron\(^{3+} \) at 2821Å).

**Probe measurements in ECRH X-Point plasmas.**

Initial experiments with inboard X-point configurations, in hydrogen plasmas [1], had a twin reciprocating double Langmuir probes that measured local electron temperature and density as well as the floating potential, \( V_f \), and the saturated ion current, \( I_{\text{sat}} \). At the separatrix there is a change in the electron density gradient scale–length, from \( \geq 12cm \) to \( \leq 3cm \), but no obvious corresponding change in the electron temperature gradient. In this respect, the behaviour of the edge parameters is similar to that observed in the JFT-2M tokamak [6] during L-mode operation.

Localised cross–field particle fluxes can be estimated from \( \Gamma_r = (\tilde{n}_e \tilde{E}_\perp)/B \), where \( \tilde{E}_\perp \) is the fluctuating electric field, along a flux surface, and perpendicular to \( B \). A second estimate, \( \Gamma_r^* \), of the flux is obtained from the cross–spectrum of the density–potential fluctuations, \( P_{n\phi}(f) \), and the average wavenumber, \( \tilde{k}_\phi(f) \), of the potential fluctuations, \( \Phi \). This gives, in the electrostatic limit, \( \Gamma_r^* = 2Re \int_0^\infty \tilde{k}(f)P_{n\phi}(f)df/B \). The values obtained for \( \Gamma_r^* \) are typically a factor of two less than those for \( \Gamma_r \).

In the vicinity of the separatrix, the particle replacement time, \( \tau_{\text{p}} \), corresponding to the estimated fluxes, is \( \sim 3ms \) and comparable to the \( \tau_{\text{p}} \) obtained from laser ablation of aluminium.
Conclusions
The energy confinement of ohmic plasmas is close to neo-Alcator scaling for both circular and X-point configurations. During ECRH, at $n_e \sim 10^{19} m^{-3}$, the energy confinement degrades with power, with an offset–linear relationship, but at lower densities ($n_e \sim 0.5 \times 10^{19} m^{-3}$) there is little confinement degradation. The Rebut–Lallia scaling provides reasonable estimates of $\tau_e$ but underestimates the observed $\tau_{inc}$ by a factor of $\geq 1.7$. The fluctuation–driven convection yields particle confinement times comparable to those obtained from laser ablation. In the many different plasma conditions studied, we have found no evidence of improved confinement associated with this class of separatrix configuration using ECRH for auxiliary heating.

References

Figure 1. COMPASS-C ohmic confinement v.s. neo-Alcator scaling. The plasma parameters are:
$I_p = 75 - 147 kA, B_T = 0.66 - 1.32 T, q_{cyl}(a) = 2.3 - 5.6, n_e = (0.5 - 12.1) \times 10^{19} m^{-3}$.
Symbols: Circular plasma: $\triangle - H_2, \circ - D_2, \square - He$. X-point plasma: $\times - H_2, \bullet - D_2$. 
Figure 2. Power dependence of the plasma energy for deuterium discharges. The plasma conditions for circular cross-section are:

- \( I_p = 75 - 78 \text{ kA}, B_T = 1.20T \),
- \( n_e = (1.1 - 1.5) \times 10^{19} \text{m}^{-3} \).

For X-point plasma:

- \( I_p = 83 - 88 \text{ kA}, B_T = 1.07T \),
- \( n_e = (0.9 - 1.3) \times 10^{19} \text{m}^{-3} \).

Symbols:
- \( \triangle - \text{neo-Alcator (circular)} \),
- \( \triangledown - \text{Kaye-Goldston (circular)} \),
- \( \square - \text{neo-Alcator (X-point)} \),
- \( \odot - \text{kinetic energy} \),
- \( n_e = n_i \sim 1 - (r/a)^2 \),
- \( T_e,i \sim 1 - (r/a)^2 \),
- X-point + - with, \( x \) - without current ramp-down,
- \( \Box - \text{Rebut-Lallia (X-point)} \),
- \( \Diamond - \text{Rebut-Lallia (circular)} \).

Figure 3. Power dependence of the plasma energy for low density hydrogen X-point discharges. The plasma conditions are:

- \( I_p = 75 - 78 \text{ kA}, B_T = 1.06 - 1.10T \),
- \( n_e = (0.5 - 0.8) \times 10^{18} \text{m}^{-3} \).

Symbols:
- \( \Diamond - \text{Kaye-Goldston} \)
- \( \square - \text{neo-Alcator} \)
- \( \Box - \text{Rebut-Lallia} \).
THEORETICAL STUDIES OF TIGHT ASPECT RATIO TOKAMAKS

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Abstract
Computer simulations of the start-up and compression sequences of the Culham START Experiment have been carried out using a number of different codes. Results from TOPEOL and PROTEUS equilibrium codes were used to assist the initial design of START. These are discussed together with more recent simulations of the compression sequence from ASTRA (Kurchatov Institute), TORUS II (Julich) and TSC (Princeton). The choice of transport model is shown to have considerable importance for the prediction of the final state of plasma.

Introduction
The growing interest in Tight Aspect Ratio devices (ie. those with \( A = \frac{R}{a} < 2 \)) and the possible advantages to be gained by operating in such a regime has inspired the building of such a device at Culham Laboratory. START (Small Tight Aspect Ratio Tokamak) is designed to investigate the physics of this largely unexplored region and to discover whether there is any validity in extrapolation of conventional scaling laws. The basic features of the experiment are described elsewhere [1]. The novel design feature whereby the plasma is expected to form with large aspect ratio at an initial major radius of about 60-65cm and then to be pushed in, by increasing the vertical field, to its final tight aspect ratio state at about 20cm major radius has presented interesting problems for computer simulation both in the initial design stage and in prediction of the quality of plasma which may be achieved.

Design Work
A free boundary equilibrium code (TOPEOL) was used for the initial theoretical studies of the equilibrium characteristics of tight aspect ratio plasmas [2]. This code was not however totally suitable for the design of poloidal field (PF) coil system and calculation of the required PF power supply current because it could not accurately represent the effect of currents induced in the conducting wall of the vacuum vessel \((r \approx 300ms)\) on the timescale of the discharge \((\approx 10ms)\). In the TOPEOL model the presence of the vacuum vessel was either ignored or was represented as an infinitely conducting wall which could be simulated by a set of current carrying 'image' filaments, appropriately positioned and weighted.

The equilibrium code PROTEUS [3] was used to improve the modelling of the PF system. The version of PROTEUS used solved the equilibrium problem together with the circuit equations for the coils and conducting vessel. It did not include a self consistent transport model for the diffusion of the plasma current density; this profile could be held fixed or arbitrarily changed in time.
The required PF currents predicted by PROTEUS agreed well with those of TOPEOL in both the case of no vacuum vessel and where the wall was taken to be infinitely conducting. With realistic estimates of the vessel conductivity, PROTEUS was used to model the inward compression of the plasma from $R = 65\, \text{cm}$, given $I_P = 30\, \text{kA}$ (Fig.1A) to $R = 20\, \text{cm}$ with given $I_P$ increasing to $200\, \text{kA}$ (Fig.1B). The results were used to optimise coil positions and current waveforms.

Simulations with ASTRA Code

The ASTRA transport code [4] solves the plasma equilibrium plus heat, particle and field transport. Ohmic heating and electron-ion equilibration only have been included in the heat balance for the START simulations and Plasma dimensions, current, density and toroidal field have been prescribed. Ion heat transport is described by the cylindrical neo-classical model, including contributions from the Plateau, Banana & Pfirsch-Schluter regimes. Several different models for electron heat transport were used eg. Alcator, Neo-Alcator, Merezshkin (T-11), ITER89 & T-10 scalings. The values of $\chi_e$ could vary by a factor of 50 over the different models.

Table I shows typical results from the ASTRA simulations for the initial plasma state and after its inward compression to small major radius. During the compression $I_P$ & $n_e$ were increased whilst $B_T$ was reduced in order to achieve the prescribed shape and parameters.

In the initial state the plasma parameters show low dependence on transport model but this becomes more important at tight aspect ratio. However it should be noted that the cylindrical model of neoclassical ion transport as used in ASTRA can be expected to give a pessimistic result for ion temperature at tight aspect ratio. For example the Artsimovich formula [5] for the ion temperature predicts $T_i = 74\, \text{eV}$ for hydrogen ions at the conditions prescribed here for the final plasma state.

<table>
<thead>
<tr>
<th>Prescribed Parameters</th>
<th>$I_P$ (kA)</th>
<th>$B_T$ (T)</th>
<th>$q(95)$</th>
<th>$\bar{n}_e$ ($10^{19}, \text{m}^{-3}$)</th>
<th>$T_{e0}$ (eV)</th>
<th>$T_i$ (eV)</th>
<th>$\tau_T$ (ms)</th>
<th>$\tau_e$ (ms)</th>
<th>$\tau_i$ (ms)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Initial State ($R = 0.57, \text{m}, a = 0.2, \text{m}$)</td>
<td>32.0</td>
<td>0.17</td>
<td>2.0</td>
<td>1.0</td>
<td>100-200</td>
<td>50</td>
<td>1.1-2.5</td>
<td>1.0-3.5</td>
<td>1.5-1.6</td>
</tr>
<tr>
<td>After Major Radius Compression ($R = 0.21, \text{m}, a = 0.16, \text{m}, e = 1.8, \delta = 0.4$)</td>
<td>150.0</td>
<td>0.28</td>
<td>5.0</td>
<td>2.8</td>
<td>240-1100</td>
<td>25-40</td>
<td>0.5-5.5</td>
<td>0.5-7.5</td>
<td>0.5</td>
</tr>
</tbody>
</table>

Table I Results from ASTRA (with $Z_{eff} = 2$)

Free Boundary Simulations (TSC & TORUS II)

The TSC Tokamak Simulation Code [6] provides for self consistent evolution of transport and free boundary equilibrium processes. Following some earlier work by M.Peng et al at ORNL, TSC has been used at Culham to simulate the START compression sequence but without the presence of the vacuum vessel. The plasma was initialised at a radius of 58cm with a toroidal current of 40kA, a vacuum toroidal field $B_\phi = (0.1/R)T$ and $n_{e0} = 2.0 \times 10^{19}\, \text{m}^{-3}$ with parabolic profile. The plasma was allowed to evolve in situ for 1ms and then the currents in the PF coils were increased over the next 4.5ms causing the inward movement of the plasma; the simulation was then continued for a further 1.5ms while the coil currents again remained fixed. During the plasma compression $B_\phi$
was reduced to \((0.033/R)T\) and \(n_{e0}\) allowed to rise to \(3.2 \times 10^{19} \text{m}^{-3}\) but the profile shape was kept fixed. Four separate simulations were carried out with various different transport models and values of \(Z_{\text{eff}}\). The conditions for each run are summarised in Table II. In all cases \(\chi_i = \chi_e\). Despite the different transport models the energy confinement times at the initial large aspect ratio were similar for each case \((\tau_E \approx 2 \text{ms} \pm 0.5 \text{ms})\) but at the compressed tight aspect ratio state the variation was by factor > 6 as indicated in Table II. These results are in broad agreement with those from ASTRA even though the ion transport model was completely different. Fig 2 shows variation in toroidal current [A], electron and ion temperatures [B], major radius [C] and aspect ratio [D] during the course of the simulation. This again illustrates the greater dependence upon transport model of the state of the compressed plasma as compared to the state of the initial plasma.

<table>
<thead>
<tr>
<th>Case</th>
<th>Transport Model</th>
<th>Include Degradation with Power?</th>
<th>(Z_{\text{eff}})</th>
<th>Final (\tau_E) (ms)</th>
</tr>
</thead>
<tbody>
<tr>
<td>(a)</td>
<td>ALCATOR ((n_{e} \chi_e = 10^{19} \text{m}^{2} \text{s}^{-1}))</td>
<td>Yes</td>
<td>3</td>
<td>3.2</td>
</tr>
<tr>
<td>(b)</td>
<td>ALCATOR ((n_{e} \chi_e = 3.5 \times 10^{19} \text{m}^{2} \text{s}^{-1}))</td>
<td>Yes</td>
<td>5</td>
<td>1.8</td>
</tr>
<tr>
<td>(c)</td>
<td>Coppi/Tang</td>
<td>Yes</td>
<td>3</td>
<td>0.2</td>
</tr>
<tr>
<td>(d)</td>
<td>Coppi/Tang</td>
<td>No</td>
<td>3</td>
<td>0.5</td>
</tr>
</tbody>
</table>

Table II Transport Models & \(Z_{\text{eff}}\) for TSC Simulations

Simulations with the TORUS II code [7] have indicated the importance of including the vacuum vessel in the model showing then the need for higher currents in the PF coils resulting in the more rapid inward movement of the plasma (3ms) and the shrinkage of the plasma which occurs subsequent to contact with the inner limiter. TORUS II simulations have also shown that, if plasma resistivity is not taken into account then \(I_P\) may be expected to rise to \(150 - 200 \text{kA}\) during its inward movement; however, when plasma resistivity is included \(I_P\) is restricted to \(\approx 100 \text{kA}\), in broad agreement with the TSC simulations.

Figures
1A: PROTEUS Simulation of Plasma before Compression \((I_P = 30 \text{kA}, R = 65 \text{cm})\)
1B: PROTEUS Simulation of Plasma after Compression \((I_P = 200 \text{kA}, R = 20 \text{cm})\)
2: Variation of Plasma Parameters with time during TSC Simulations

References
TIGHT ASPECT RATIO TOKAMAKS

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Abstract

The physics of tight aspect ratio plasmas is being studied on two experiments in the UK: on the SPHEX spheromak (with central rod) at UMIST; and on the START (Small Tight Aspect Ratio Tokamak) device at Culham.

Introduction

Although large proposed devices such as ITER should approach or achieve ignition, the ultimate goal of economic fusion power may be harder to attain. The very successful JET results indicate good behaviour at the tight aspect ratio of 2.4, but present reactor designs based on a toroidal containment vessel are limited to aspect ratios of at least 4 by requirements of centre column shielding and breeding blanket thickness. However the ultimate limit is the ‘Spherical Torus’ [1] having a spherical containment vessel with a simple rod up the centre to provide the toroidal field. This has aspect ratio of 1.5 or less, and may offer a simpler, more compact (and hence cheaper) reactor than conventional designs.

Many changes are expected in the physics of plasmas at tight aspect ratio; most are predicted to be favourable (eg beta limit; density limit; high ohmic heating) others are unknown (eg plasma stability sawtooth behaviour) and for one important parameter, the confinement time scaling, the predictions of the various empirical models vary by a large factor when applied in this unfamiliar and unexplored region [2].

Experimental evidence is urgently needed to verify or discount these predictions, and also to improve empirical scaling models by widening the available database. Useful initial information, especially on equilibrium properties, has been obtained from the SPHEX spheromak at UMIST which has installed a current-carrying rod down the axis; and the START experiment which has now become operational at Culham.
SPHEX

SPHEX [3] has been modified by the temporary addition of an axial copper rod carrying up to 60kA (Fig.1), so creating an additional toroidal magnetic field of $\sim 50$ mT in the region of the magnetic axis ($R \sim 23$ cm). It is found that the relaxation process which creates and sustains the toroidal current is not inhibited by this imposed field, and indeed the toroidal current rises to over 3 times its value in the spheromak [4]. This increase allows the plasma to remain near the Taylor relaxed state ($\mu = J/B \sim$ constant), and since $q_0 = 2/(\mu_0 R)$, $q_0$ also increases rather slowly with $I_{\text{ROD}}$ (Fig.2). The increased helicity content is reflected in a rise in the voltage on the Marshall gun from $\sim 400V$ to $\sim 700V$, increasing the helicity input rate $\dot{k} = 2\dot{\psi}V$. The input power is thus also increased, causing a rise in $T_e$ from $\sim 12eV$ to $\sim 25eV$ even though $\tau_{E\phi}$, which is believed to be radiation limited, remains unchanged.

The plasma magnetic fields continue to show large fluctuations, and although the detailed structure is altered, some features remain. Radial shows amplitude decreasing from the geometric to the magnetic axis, and phase at the magnetic axis trailing that at the edge (Fig.3). Further work to quantify the power carried by these fluctuations should give a better understanding of the toroidal current drive mechanism.

START

Construction of START has been achieved on schedule (Oct 1990) and first plasmas were obtained in January 1991. The main features of the assembly are shown in Fig.4; the 2m diameter aluminium vacuum tank is inherited from the 1970 TORSO experiment. START is designed to induce an initial plasma of about 30kA in an octupole null at major radius $\sim 65cm$, and then compress this plasma in towards the centre rod [6]. During this process the current is predicted to increase to over 100kA. The design central rod current is 500kA, giving a toroidal field of 0.17T at 65cm.

In practice however it has been found easier to break down the initial plasma in the quadrupole null at a major radius of $\sim 45cm$ formed by the two inner poloidal field coils, supported by a small vertical field. The breakdown loop voltage is $\geq 30$ volts. Initial results from START are encouraging, and show that hot compressed plasmas of up to 70kA have already been obtained. Fig.5 shows typical results. In this discharge, the plasma moves in to contact the central column graphite limiter at 23.3msec, and maintains an aspect ratio of 1.5 for a further 2msec. The high value of light emission is produced by ionised gas surrounding the plasma; highspeed cine film indicates a relatively hot, clean plasma. Future operation at higher plasma current, featuring major radius compression coupled with a simultaneous reduction in toroidal field, should enable the properties of plasmas at tight aspect ratio to be investigated.

Conclusion

Tight aspect ratio tokamak-like equilibria have been generated on the SPHEX ‘rodomak’ and show improved plasma properties compared to the spheromak. Hot tokamak plasmas at aspect ratios $\sim 1.5$ have been obtained on START.
References

1. J B Hicks and Y-K M Peng, Proc. 16th Symposium on Fusion Technology, R70
2. M F Turner et al. this conference
4. P K Browning, this conference
5. R T C Smith et al. Proc. 16th Symposium on Fusion Technology, p. 62

Fig 1. SPHEX geometry

Fig 2. $q_0$ vs $I_{ROD}$

Fig 3. $\vec{B}$ for the dominant 20kHz mode, spheromak and $I_{ROD} = 20$KA.

FLUX CONSERVER
ELECTRODES
SOLENOID
GAS VALVES

AXIAL ROD
QUARTZ SLEEVE

96 cm
Fig 4. START geometry

Fig 5. START Results  
○ plasma centre from CCD diagnostic  
● plasma centre by equilibrium reconstruction
The mission of the Burning Plasma Experiment (BPX, formerly CIT) is to study the physics of self-heated fusion plasmas (Q = 5 to ignition), and to demonstrate the production of substantial amounts of fusion power (P_{fus} = 100 to 500 MW). Confinement projections for BPX have been made on the basis of 1) dimensional extrapolation, 2) theory-based modeling calibrated to experiment, and 3) statistical scaling from the available empirical database. The results of all three approaches, discussed below, roughly coincide. We presently view the third approach, statistical scaling, as the most reliable means for projecting the confinement performance of BPX, and especially for assessing the uncertainty in the projection.

The dimensional scaling approach\cite{1} is based on the observations that the key dimensionless parameters governing tokamak transport behavior are \( v* = v_{90}/\omega_p \), \( \beta \), and \( \rho^* = \rho_i/a \), and that all present theories of transport fall close to one of two extreme confinement scalings, Bohm (\( \omega_c\tau_E \sim \rho^* \)) or "gyro-Bohm" (\( \omega_c\tau_E \sim \rho^*^{-3} \)) when \( v^* \) and \( \beta \) are held fixed. Thus it is reasonable to assume that plasmas in present devices which achieve values of \( v^* \) and \( \beta \) which are accessible in BPX can be used as bases for Bohm or gyro-Bohm scaling of confinement to BPX conditions, as illustrated in Table 1. [Aspect ratio effects are factored out by using the experimental result\cite{2} that \( nT \) depends on \( I_p \) and \( R/a \) in the combination \( (I_p/R/a) \)]. The design parameters for BPX are \( R = 2.59m, a = 0.8m, \kappa_{95} \sim 2, \delta_{95} \sim 0.35, I_p = 10.6MA, B_T = 8.1T, q_{95} \geq 3.2, P_{ICRF} = 20 MW \). The device and facilities can accommodate upgrades in tokamak power supplies such that \( I_p = 11.8MA \) and \( B_T = 9T \) can be provided, and in heating power up to 30 MW of ICRF, or 50 MW of ICRF + ECH. Extrapolation based on dimensional scaling gives Q's in the required range, even with the more pessimistic Bohm scaling, at plasma currents below the nominal 10.6MA operating point. The "starting point" discharges in DIII-D and JET have not been optimized for this purpose. Developing optimal starting point conditions, and discriminating experimentally between Bohm and gyro-Bohm scaling, are two key elements of the BPX Physics R&D Plan\cite{3}.

Confinement projections have also been made using simulations based on a Multi-Mode model\cite{4} which includes transport due to trapped electron modes, ion...
temperature gradient modes, and resistive ballooning modes. This model has been calibrated against experimental data from TFTR, DIII-D, JET, PDX, and ASDEX. A preliminary edge stabilization model to simulate H-mode confinement effects\[6\] is included. The density profile is simulated in an ad-hoc manner to provide a very flat $n_e(r)$. Q~7 is predicted for $I_p = 10.6\text{MA}, B_T = 8.1\text{T}$. The transport coefficients of the Rebut-Lalliam model\[6\] however, are very optimistic in comparison with the Multi-Mode model, giving both $\chi_e$ and $\chi_i$ 2 - 3 times lower in the confinement zone ($\chi_e \sim 0.3\text{ m}^2/\text{sec vs.} 0.8\text{ m}^2/\text{sec}$), and so implying a much greater predicted Q.

**TABLE I**

<table>
<thead>
<tr>
<th>Device</th>
<th>DIII-D</th>
<th>BPX</th>
<th>JET</th>
<th>BPX</th>
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<tr>
<td>$a$ (m)</td>
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<td>0.79</td>
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<td>1.9</td>
<td>1.9</td>
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<tr>
<td>$I_p$ (MA)</td>
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<td>9.2</td>
<td>4.2</td>
<td>10.0</td>
</tr>
<tr>
<td>$B_T$ (T)</td>
<td>2.1</td>
<td>9</td>
<td>2.8</td>
<td>9</td>
</tr>
<tr>
<td>$T_i/T_e$</td>
<td>1.29</td>
<td>0.56</td>
<td>1.09</td>
<td>2.9</td>
</tr>
<tr>
<td>$n_{D0}$ ($10^{20}/\text{m}^3$)</td>
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<td>7.0</td>
<td>0.56</td>
<td>1.24 - 2.23</td>
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<tr>
<td>$\tau_E$ (sec)</td>
<td>0.21</td>
<td>0.51 - 1.64</td>
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</tr>
<tr>
<td>$\beta$</td>
<td>4.1%</td>
<td>4.1%</td>
<td>2.4%</td>
<td>2.4%</td>
</tr>
<tr>
<td>$n_e$</td>
<td>0.083</td>
<td>0.083</td>
<td>0.027</td>
<td>0.027</td>
</tr>
<tr>
<td>$\rho^*$ ratio</td>
<td>0.31</td>
<td>0.56</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$P_{\alpha} + P_{aux}$ (MW)</td>
<td>12</td>
<td>190 - 58</td>
<td>7.2</td>
<td>36 - 21</td>
</tr>
<tr>
<td>$Q$</td>
<td>7 - $\infty$</td>
<td></td>
<td>25 - $\infty$</td>
<td></td>
</tr>
</tbody>
</table>

**TABLE II**

<table>
<thead>
<tr>
<th>Device</th>
<th>(\tau_{\text{E}}/\tau_{\text{E}}^{\text{ITER89-P}})</th>
<th>(\tau_{\text{E}}^{\text{MHD}}/\tau_{\text{E}}^{\text{ITER89-P}})</th>
</tr>
</thead>
<tbody>
<tr>
<td>JET</td>
<td>2.10 ± 0.38</td>
<td>1.86 ± 0.35</td>
</tr>
<tr>
<td>DIII-D</td>
<td>1.70 ± 0.13</td>
<td>1.70 ± 0.21</td>
</tr>
<tr>
<td>ASDEX</td>
<td>2.23 ± 0.22</td>
<td>(2.73 ± 0.30)</td>
</tr>
<tr>
<td>PBX-M</td>
<td>-</td>
<td>2.05 ± 0.26</td>
</tr>
<tr>
<td>PDX</td>
<td>-</td>
<td>1.56 ± 0.33</td>
</tr>
<tr>
<td>JFT-2M</td>
<td>1.51 ± 0.16</td>
<td>1.79 ± 0.21</td>
</tr>
<tr>
<td></td>
<td>1.88 ± 0.34</td>
<td>1.79 ± 0.18</td>
</tr>
</tbody>
</table>

At present we judge that the most reliable method for projecting the performance of BPX is based on analysis of the recently-developed ITER H-mode database\[7\]. Table II shows the mean ratio of the H-mode confinement data to the
ITER89-P scaling relation for each machine in the database, sorted by data type. The data analyzed is the "standard" dataset defined by the authors (ELM'ing and ELM-free) further constrained by dW/dt < 0.2P_{\text{heat}} and q_{95} < 5.0. The ranges in the body of the Table indicate shot-to-shot scatter, while the ranges in the averages indicate machine-to-machine scatter. Constraining to the diamagnetic data where available, in order to minimize systematic uncertainties and fast ion effects, and taking the MHD data for PDX and PBX-M which had perpendicular and mixed parallel and perpendicular injection respectively, one obtains a machine-averaged L-mode enhancement factor of 1.85 ± 0.31 (figure 1).

![Figure 1. ITER H-mode database for B\tau_E plotted vs. ITER89-P L-mode scaling. Data type (MHD vs. diamagnetic) selected to minimize fast ion contribution.](image)

Scaling relations have also been developed for H-mode data, but we view these as too preliminary for use as the basis for extrapolation. For typical Q ~ 10 conditions in BPX (P_{aux} + P_{RF} - P_{\text{brems}} = 50 \text{ MW}, \bar{n}_e = 2.5 \times 10^{20}/\text{m}^3) we project \tau_E = 1.85 \times \text{ITER89-P} = 1.01 \text{ sec}. The DIII-D - JET scaling\[^8\] gives 0.961 sec, the ITER90-H ELM-free scaling gives 1.26 sec, the ITER90-H direct regression fit to the full standard dataset gives 1.90 sec, and the random coefficients model\[^6\] gives 1.25 sec, with an uncertainty of ± 26%. This uncertainty is consistent with a 2-step uncertainty estimate based on combining the 17% uncertainty in the H/L ratio in quadrature with the 17% uncertainty in L-mode extrapolation\[^10\] for machine-averaged performance. Thus we take as the "standard" confinement extrapolation for BPX 1.85 \times \text{ITER89-P}, with an uncertainty of ± 25%. Combining this result with the Monte-Carlo approach to uncertainties\[^11\] in Z_{eff} and n_e(r) gives the projected range in Q shown in figure 2.
We conclude that the BPX device has adequate performance to achieve its mission of determining the confinement physics, operational limits, and $\alpha$-particle dynamics of DT plasmas with $\alpha$ power greater than auxiliary heating power, while producing more than 100 MW of fusion power. The upgrade capabilities of the device provide assurance that this mission can be achieved even in the case of unfavorable plasma performance.

(Work supported by U.S. DOE Contract No. DE-AC02-CHO3073)


[9] Riedel, K.S., this conference


ISOTOPE DEPENDENCE OF ELECTRON PARTICLE TRANSPORT IN ASDEX

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Introduction

The isotope dependence of electron particle transport in the ASDEX bulk plasma was investigated using gas oscillation techniques /1/. Ohmic discharges in hydrogen and deuterium were evaluated for densities up to the density limit, and some experiments were carried out in helium. In addition, neutral beam heated L-type plasmas at ohmic target densities corresponding to saturated confinement could be analyzed by this method for both hydrogen isotopes.

Method

Small density perturbations about equilibrium, induced by sinusoidal modulation of the gas valve, were analyzed for the different channels of the ASDEX HCN-laser-interferometer. The measured pattern of amplitudes and phase shifts is compared to solutions of the particle conservation equation, \( \frac{\partial n}{\partial t} = - \nabla \cdot \Gamma + P \). A transport law with a diffusive and a convective flux component \( \Gamma = -D \nabla n - V n \) is assumed. The coefficients \( D(r) \) and \( V(r) \) are determined with a crude radial resolution, using a spatial transport model with constant inner and peripheral values of \( D \) and \( V/r \) and a linear interpolation in between. The physical justification for this model, which seems adequate to the limited number of experimental chords, is that the central and outer plasma regions normally have greatly different parameters and often different transport characteristics. The absolute numbers of \( D \) and \( V \) for each plasma zone are gained from a numerical fitting procedure.

Results for ohmic plasmas

Using the gas oscillation technique, density scans including the range of linear ohmic confinement (LOC), where the energy confinement time \( \tau_E \) linearly increases with density, to saturated ohmic confinement (SOC) were performed for standard values of plasma current and main field (320 kA / 2.17 T) and standard wall conditions (stainless steel or carbonized). A clear isotope effect in both diffusion \( D \) and inward convection \( V \) was found in all plasma regions, except for \( V \)
in the outer zone, which was comparable for both isotopes. As an example, fig. 1 shows the diffusion $D_p$ in the confinement region versus average density for both hydrogen isotopes.

Particle confinement is superior for deuterium as compared to hydrogen for all densities, except possibly the lowest values around $1 \times 10^{13} \text{ cm}^{-3}$, where the effect is less pronounced. Quantitatively, the diffusion coefficient roughly follows a law $\propto A_i^{-1/2}$ in LOC and $\propto A_i^{-1}$ in SOC, where $A_i$ is the isotope mass. This is demonstrated in fig. 2, which shows the ratio of diffusion $D_{\text{hyd.}} / D_{\text{deut.}}$ for the central and the confinement region as a function of average density.

Results with neutral beam heating

While the quiescent H-phase in ASDEX due to its non-stationary character did not allow gas oscillation experiments, this technique could be applied to stationary L-phases. Hydrogen beams up to 1.4 MW heating power were injected into either hydrogen or deuterium target plasmas at comparable parameters. The results indicate a continuation of the isotope effect into the L-mode, with an $A_i^{-1}$ dependence similar to the ohmic phase for target densities corresponding to SOC conditions. This is illustrated in fig. 3, which shows $D$ versus normalized radius for hydrogen and deuterium cases at different injection powers.

In addition, during the neutral beam phase inward convection $V$ in the confinement region was clearly larger for deuterium, in contrast to the ohmic case, where no difference was found. This behaviour, which for deuterium further improves the particle confinement, is documented in fig. 4, showing the inward convection for the same shots as in fig. 3.

Comparison with helium

During the later experimental phase with boronized ASDEX vessel gas oscillation experiments were performed in ohmic deuterium and helium plasmas, allowing a direct comparison of transport in both filling gases. As compared to the ohmic results with stainless steel or carbonized wall, the absolute values of $D$ were generally lower with the new wall conditioning, close to the values evaluated for improved ohmic confinement (IOC) /2/. No clear difference was found between both gases, although for helium an $A_i^{-1}$ dependence, expected at the chosen density, should lead to another factor of two improvement in transport with respect to deuterium. The values of $D$, plotted for both gases versus normalized radius, are shown in fig. 5.

As both $A_i$ and the charge $Z$ are increased by a factor of 2 for helium, these
results seem to indicate, that the isotope effect in D cannot be described by a pure mass dependence, but rather follows a pattern $\alpha Z \cdot A_i^{-1}$.

**Summary and Conclusions**

ASDEX gas oscillation results have shown, that particle transport is often closely associated with thermal energy confinement /3/. Particle diffusion and thermal transport share a common density dependence for the transition from the LOC- to the SOC- or IOC-regime. Similarly, a degradation in both energy and particle confinement is found during the neutral beam L-phase. The comparison between transport in hydrogen and deuterium plasmas, presented in this paper, indicates, that the favourable energy confinement of deuterium with respect to hydrogen also finds its concomitant in particle confinement. Particle diffusion roughly scales with the inverse of the isotope mass in the SOC- and L-regime and the results for helium suggest an additional Z-dependence.

**References**

Fig. 2:
Ratio of diffusion in hydrogen versus deuterium as a function of average density, shown for the central and the confinement region. (Broken line sketched as a guide)

Fig. 3:
Particle diffusion versus normalized radius during Nl- L-phases in hydrogen (hatched lines) and deuterium (full lines) at different Nl-power levels.

Fig. 4:
Inward convection for the same cases as in fig. 3, deuterium full lines, hydrogen hatched lines. (Average densities 3 - 4 \cdot 10^{13} \text{cm}^{-3})

Fig. 5:
Diffusion versus normalized radius for ohmic deuterium (full line) and helium plasmas (hatched line) at an average density of 5 \cdot 10^{13} \text{cm}^{-3}. 
ON TEMPERATURE AND DENSITY DEPENDENCE OF THE ASDEX L-MODE CONFINEMENT

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1. Introduction

The global energy confinement time of additionally heated tokamak plasmas is empirically well described by a power law scaling. The physical processes, however, which make transport in a tokamak behave according to a simple law are not understood. From comparison of parameter dependences of theoretical transport coefficients with scaling laws a working hypothesis can be deduced. According to this ansatz, dissipative trapped electron modes (DTEM) could drive energy transport in the plasma confinement region while sawteeth dominate the centre and electromagnetic modes are responsible for enhanced transport at the edge of the plasma.

Beside the isotope effect, which is not predicted correctly by almost all theories, the absence of the plasma current as a confining agent and a too strong dependence on toroidal field strength $B_{tor}$ are the most important failures of DTEM theory. Strongly enhanced transport outside the $q=2$ surface had to be evoked, introducing a favourable current and curing a too strong $B_{tor}$ dependence. Tearing or resistive ballooning modes (RBM) have been discussed as candidates for the transport in this region. Ion energy transport is considered to be enhanced above the prediction from neoclassical theory due to ion-temperature-gradient (ITG) driven modes.

The validity of this working hypothesis cannot be unambiguously clarified by comparison to global energy confinement times. Transport calculations have to be performed and theories must prove themselves in describing the dependences of local transport coefficients on the plasma parameters correctly. The following questions are addressed here: (1) Does the power deposition profile shift to unfavourable regions when density is increased and therefore conceal a density dependence of the transport coefficient, as predicted from DTEM theory? (2) Is the temperature dependence of the DTEM transport coefficient the reason for power degradation of $\tau_E$ and are RBM a possible candidate for enhanced transport in the plasma edge? And finally (3), is the transport strongly enhanced outside the $q=2$ surface?

In the following, we try to extract answers to these questions from transport analyses of well diagnosed ASDEX discharges using the TRANSP transport code.

2. Low Density L-Mode and Density Dependence of the Power Deposition Profile

Three series of L-mode discharges at low densities ($\bar{n}_e=1.3$, $2.0$, and $2.6\times10^{19}$m$^{-3}$) have been analyzed and in Fig.1 their kinetic confinement times are compared to those of OH plasmas. If we correct for the power dependence according to $\tau_E\sim(P_{lum}/P_{OH})^{-0.58}$ we obtain the open squares in Fig. 1. It is interesting that these corrected values fall in a line with the high density SOC discharges. From this consideration we could conclude that confinement is density independent and that an additional effect causes $\tau_E$ to decrease with decreasing density in the LOC regime.

Analyses of LOC plasmas show that a process contributes to the increase of $\tau_E$ with $\bar{n}_e$ in the LOC regime which is absent in the L-mode: The ion temperature only slowly approaches the electron temperature as density is increased. In the SOC regime and the L-mode, ion
and electron energy content are practically identical. On the other hand, in the L-mode a density scan at low $\bar{n}_e$ is simultaneously a power scan (see Fig. 2). At the lowest density, shine through and secondary charge exchange losses lower the absorbed power. However, correcting for the two differences cannot entirely reconcile the $\bar{n}_e$ dependence of LOC and the lack of it in L-mode confinement. Also the shape of the power deposition profile does not change much in the density range of the LOC regime. No indication can be deduced from Fig. 2 for the hypothesis, that a broadening of the deposition could conceal a confinement improvement with density.

In Fig. 3a the diffusivities of the discharge with $\bar{n}_e=2.0 \times 10^{19}$ m$^{-3}$ are shown. They are compared with predictions from DTEM [1] and RBM [2] theories. In all discharges there exists a radial region (0.25\leq r/a \leq 0.75) where electron and ion transport coefficients are comparable in size. In this region also $\chi_e$ from DTEM has its maximum values. The absolute values are not reproduced by DTEM theory, which predicts a reduction in $\chi_e$ with density whereas the experimental values are independent on it or slightly increasing (see Fig. 6b). The discharge at $\bar{n}_e=1.4 \times 10^{19}$ m$^{-3}$ has to be taken with caution, because it is already touching a regime which is dominated by the hot beam ion population and characterized by more peaked density profiles with the consequence of impurity accumulation.

At the plasma edge (0.75\leq r/a \leq 1), $\chi_e$ from RBM shows the right radial dependence and is of the order of the experimental values if it is enhanced by a neoclassical multiplier (7/e$^{1/2}$). The dependence with density is opposite to that of the experimental results (see Fig. 6c).

Ion heat conductivity is a factor of 3 higher than predicted from neoclassical theory. The $\eta_i$ value ($\eta_i = \partial_i \log T_i / \partial_i \log n_i$) is clearly below the critical value $\eta_i^{\text{crit}}$ for the onset of the ITG modes, which is taken in its density gradient dependent form with a lowest value of 1.5 [3]. Hence, ITG driven transport cannot explain the enhanced ion energy transport at low $\bar{n}_e$.

3. Medium Density L-Mode and Comparison to DTEM Theory

DTEM theory predicts $\chi_e$ to depend strongly on electron temperature and density: $\chi_e^{\text{DTEM}} \sim T_{e}^{7/2}/n_e$. Whether this is the cause of the power degradation of the L-mode confinement can be investigated in discharges which realize the increase in plasma energy content with neutral beam heating either at constant $\bar{n}_e$ and increasing $T_e$ or with increasing $\bar{n}_e$ at constant $T_e$. We analyzed power scans on pairs of this kind of discharges. The energy confinement times of both types of discharges obey the common L-mode scaling law. Discharges with higher density have slightly higher confinement times according to $\tau_E \sim n_e^{0.26}$. The absorbed power is practically identical in each pair of discharges. Although DTEM theory predicts important changes in transport of up to a factor of 6 the experimental confinement times differ only by a few % (see Fig. 4). The confinement degradation with increasing $T_e$ as predicted from DTEM theories is not found in discharges where $T_e$ has been changed at constant heating power, consequently also the degradation of $\tau_E$ with power cannot be attributed to enhanced DTEM energy transport.

In Fig. 3b the transport coefficients of a discharge with $P_{NI}=1$ MW and $\bar{n}_e=4.4 \times 10^{19}$ m$^{-3}$ are plotted. Fig. 3a illustrates the trends in transport if $\bar{n}_e$ is increased: Ion energy transport becomes the leading loss channel in the inner half of the plasma and $\chi_e$ decreases. But the electrons are still responsible for the transport in the edge. $\chi_e$ from DTEM theory follows the trend of the experimental values in the region where electrons contribute only little to the total transport. At the edge RBM induced transport is of the right order to describe the experimental $\chi_e$.

The ITG driven mode is destabilized and can qualitatively account for the enhanced transport in the ion channel. Taking into account the influence of ITG driven transport on the
electron channel [1], yields $\chi_e$ values which lie factors 5 to 10 above the experimental ones.

4. Energy Transport at the Plasma Edge
The enhanced electron energy transport at the plasma edge does not depend on the location of the $q=2$ surface. This is depicted in Fig.5 where the experimental $\chi_e$ at two current values ($I=420$ and $280\text{kA}$) of an L-mode discharge are plotted. The $q=2$ surface is moved from $r=35$ to $27\text{cm}$ without changing the radial shape of $\chi_e$ which is strongly enhanced outside $r=35\text{cm}$.

5. Summary
In Fig.6, the transport calculations of the density scan and a pair of discharges from the comparison with DTEM theory are summarized. In the discussion it must be kept in mind, that these discharges do not belong to identical heating power, as it is indicated in Fig.6b. We compare radially averaged quantities and the "error bars" indicate the variation of the functions in the range we used for averaging.

In the inner half of the plasma, ion energy transport contributes more as density is increased and becomes the dominant loss channel at high densities. The enhancement factor above neoclassical theory increases with density in the same way as the experimental $\eta_i$ approaches the critical value to de-stabilize the ITG driven modes (see Fig. 6a). As in OH plasmas, ITG modes can describe qualitatively the increasing importance of the ions for the energy transport in this region.

At low densities, in an intermediate range, $\chi_e$ and $\chi_i$ are of comparable level (see Fig.6b). There ITG driven modes should be stable and cannot explain the enhanced ion energy transport. This happens in the same range, where DTEM theories predict large transport coefficients. It is questionable, however, whether DTEM give rise to the same level of ion and electron transport, since neither the absolute values nor the parameter dependences are in agreement with the experimental results (Fig.6b). Only at higher densities and in regions where electron transport loses importance, DTEM induced transport can explain absolute values and parameter dependences of the experimental $\chi_e$.

Furthermore, we have shown that the lack of a density dependence of L-mode confinement cannot be explained by an unfavourable shift of the power deposition profile with increasing density. The absence of a plasma current scaling of DTEM cannot be reconciled by the assumption, that strongly enhanced transport particularly outside the $q=2$ surface introduces an effective $q$-dependent plasma radius. RBM theories (enhanced by a neoclassical multiplier) predict transport at the edge of the right magnitude but with the wrong dependence on density.

In conclusion, energy transport in tokamak plasmas is probably governed by a variety of transport processes with changing importance as parameters change.

References
Fig. 1: OH and L-mode energy confinement times. Open squares: Confinement time of a hypothetical OH plasma out of which the L plasma would emerge if power degradation is valid at low densities.

Fig. 2: NI power deposition profiles from Monte-Carlo calculation for L-mode plasmas at three densities.

Fig. 3: Experimental and theoretical electron and ion heat diffusivities at different densities.

Fig. 4: Comparison of L-mode discharges where the energy has been built up by density (I) or by temperature (II): $T^4_E(I)/T^4_E(II)$ and $(n_0T^{3.5}_I)/n_0T^{3.5}_E(II)$.

Fig. 5: Experimental electron diffusivities at different plasma currents and otherwise identical plasma parameters. The $q=2$ surfaces are indicated.

Fig. 6: Radially averaged transport coefficients. (a): enhancement factor of experimental $\chi_e$ above the neoclassical prediction and the relative difference of experimental and critical $\chi_e$ values. (b): Experimental electron and ion diffusivities compared to the prediction from DTEM theory. (c): Electron diffusivity at the plasma edge compared to the prediction from RBM theory.
LONG TERM $Z_{\text{eff}}$ PROFILE BEHAVIOUR ON ASDEX FOR DIFFERENT HEATING AND WALL CONDITIONS

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Introduction

$Z_{\text{eff}}$ is an overall measure of the impurity content in a plasma and can be directly derived from bremsstrahlung emission. It should ideally be close to 1 to minimize both fuel dilution and radiation losses. Obviously, the type and amount of impurities depend on the wall material and also on the long term wall conditioning. In this paper we refer to ASDEX divertor discharges where the main chamber was covered with graphite tiles; we investigate the influence of uncoated, carbonized and boronized walls on $Z_{\text{eff}}$ profiles.

$Z_{\text{eff}}$ diagnostics

The detection system of the Nd:YAG laser scattering device is used to measure the bremsstrahlung along 16 chords in different broad wavelength bands in the near infrared where line radiation can generally be neglected \cite{1}. Using the simultaneously measured $n_e$ and $T_e$ profiles we obtain radial $Z_{\text{eff}}$ profiles with a spatial resolution of 4 cm. The statistical errors are mainly caused by the density measurements and vary in the range between 10\% on axis and 30\% in the plasma edge.

Carbon and oxygen density profiles and their $Z_{\text{eff}}$ contributions are also absolutely measured by charge exchange recombination spectroscopy (CXRS) during beam injection and are compared with the results from bremsstrahlung measurements.

Experimental results and discussion

$Z_{\text{eff}}$ profiles are analyzed for different wall coatings and heating methods. The influence of the different wall conditions on $Z_{\text{eff}}$ on axis is clearly demonstrated in Fig.1 for standard ohmic discharges ($\overline{n}_e = 3 \times 10^{19}$ m$^{-3}$, $I_p = 0.32$ MA, $B_t = 2.17$ T, $q_a = 3.3$ ). There is a strong reduction in $Z_{\text{eff}}$ from the uncoated to the boronized vacuum vessel. With each new boronization the central $Z_{\text{eff}}$ was reduced and reached 1 after the last boronization (bor 7) which was done in a gas mixture of 90\% He and 10\% B$_2$H$_6$. The effectiveness of the boron layer is maintained over a long operation period; after about 1000 plasma discharges $Z_{\text{eff}}$ increases to about 1.3. The increase of $Z_{\text{eff}}$ is faster when strong additional heating is frequently applied (bor 5).

The reduction of $Z_{\text{eff}}$ due to boronization is observed over the whole plasma radius. Close to the separatrix $Z_{\text{eff}}$ is about 3 in standard ohmic discharges which is considerably less than for non-boronized walls. At higher densities ($\overline{n}_e > 4.5 \times 10^{18}$ m$^{-3}$) the boundary $Z_{\text{eff}}$ is even less than 2. The improvement in $Z_{\text{eff}}$ correlates well with the rise of the density limit. In the presence of boronized walls low $Z_{\text{eff}}$ values over the whole plasma cross section were also observed with ICRH, LH and neutral beam injection, both in the L- and H-mode. In long pulse (up to 3.5 s) ELMy H-mode plasmas, stationary $Z_{\text{eff}}$ profiles with values close to 2 have been measured \cite{2}.

As an example of the improvement of $Z_{\text{eff}}$ in an auxiliary heated plasma Fig. 2 shows the power dependence of the central $Z_{\text{eff}}$ in a LH-heated plasma ($\overline{n}_e = 3 \times 10^{19}$ m$^{-3}$) before and after boronization. During LH operation the increase of $Z_{\text{eff}}$ is halved from values between 3 and 6 (non-boronized) to 1.5 and 3 (boronized), resulting in improved current drive efficiency.
Impurity profiles

For uncoated graphite walls, a relatively high impurity level is found leading to a significant dilution of the hydrogen plasma. Under these conditions the main metallic impurity is Cu which is sputtered from the neutralizer plates in the divertor. It turns out, however, that for medium and high densities ($n_e > 2.5 \times 10^{19} \text{ m}^{-3}$) the contribution of Cu to $Z_{\text{eff}}$ can be neglected. $Z_{\text{eff}}$ is determined by low Z impurities (O and C). Typical impurity mixes for uncoated and boronized wall conditions of neutral beam heated plasmas (L-mode) are shown in Fig. 3. Densities of fully stripped oxygen and carbon ions were obtained by CXRS with an absolutely calibrated detection system /3/.

Under uncoated conditions oxygen is the main impurity ion having about twice the $Z_{\text{eff}}$ contribution of carbon. The $Z_{\text{eff}}$ profile obtained from the bremsstrahlung measurement and the profile constructed from the (fully ionized) carbon and oxygen contributions agree very well.

With boronization, the oxygen content is reduced by a factor of 7. The carbon concentration is also reduced by a factor slightly greater than 2. The $Z_{\text{eff}}$ profile obtained from the contributions of $O^{8+}$ and $C^{6+}$ has now smaller values in comparison with the bremsstrahlung values. The gap may originate from the contribution of boron to $Z_{\text{eff}}$, which could not be measured directly.

Impurity accumulation

Impurity accumulation in high confinement regimes with peaked electron density profiles (pellet injection and counter NI - heating) is observed under different wall conditions. In non-boronized cases it has been shown that metallic impurities accumulate whereas low Z impurities which mainly determine $Z_{\text{eff}}$ do not exhibit such a clear tendency /5/. $Z_{\text{eff}}$ profiles show only slight peaking towards the plasma centre. With boronization, however, a more pronounced peaking of $Z_{\text{eff}}$ is observed. Fig. 4 shows for example impurity accumulation after repetitive pellet injection. Both, $Z_{\text{eff}}$ on axis as compared to the half radius value and the soft x-ray signal give clear evidence for accumulation. The proton density profile which can be derived from the measured $Z_{\text{eff}}$ profile /4/ is still peaked. For counter beam injection (Fig. 5) the peaking of $Z_{\text{eff}}$ is even stronger leading to a strong dilution of hydrogen in the plasma centre and a nearly complete flattening of the proton profile /4/. To understand the differences in the accumulation behaviour in presence of non-boronized and boronized walls one has to realize that with boronization metallic impurities are suppressed and oxygen is reduced. In this case we find an accumulation of light impurities which was possibly prevented by the accumulation of metallic impurities under non-boronized conditions. A further hint for this explanation is provided by CXRS showing an even stronger peaking of oxygen as compared to carbon in the accumulation phase. If neon (Z = 10) is puffed into the discharge the peaking of the oxygen is reduced. This behaviour suggests an interpretation in terms of multispecies interaction in such a manner that heavier elements tend to expel the light ones from the central region /5/.

References

/2/ F. Wagner et al. IAEA - CN - 53 / A - 4 - 2 (Washington 1990)
Fig. 1: Long-term dependence of $Z_{\text{eff}}$ on ASDEX in standard ohmic discharges ($n_e = 3 \times 10^{19} \text{ m}^{-3}$) with uncoated, carbonized (carb.) and boronized (bor 1-7) walls.

Fig. 2: Improvement of $Z_{\text{eff}}$ due to boronization in LH-heated plasmas.

Fig. 3: Radial $Z_{\text{eff}}$ profiles of L-mode discharges from bremsstrahlung measurements and CXRS for uncoated and boronized wall conditions. ($P_{\text{tot}} = 1.3 \text{ MW}, n_e = 3 \times 10^{19} \text{ m}^{-3}, I_p = 0.4 \text{ MA}$.)
Fig. 4: Time evolution of soft x-ray signal and $Z_{\text{eff}}$ on axis and half radius with repetitive pellet injection.

Fig. 5: Time evolution of $Z_{\text{eff}}$ on axis and half radius for counter neutral beam injection.
CURRENT RAMP EXPERIMENTS ON THE ASDEX TOKAMAK

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Abstract

The confinement behaviour of an L-mode plasma in ASDEX is studied under transient conditions, when the total plasma current is changed on a time scale comparable to the energy confinement time. The investigation reveals a decoupling of the conductivity - from the current density profile and shows that the scaling laws derived for steady state conditions can not be applied. The analysis of the total energy content of the plasma constituents reveals similar temporal behaviour. While the current density diffuses on a time scale determined by the plasma resistivity, the loss of thermal energy is delayed until the central confinement region is affected within the q = 1 flux surface; this is correlated with a sudden change in the sawtooth frequency and the shape of the electron temperature profiles. When the current is ramped down, the total energy on the other hand begins to drop slightly earlier; it is also remarkable that the electron energy exchange across the plasma radius due to sawtooth action occurs almost simultaneously on a time scale much shorter than skin time or energy confinement time. The observed delay of the electron energy change indicated by the temperature profiles depends on whether the current is ramped down or up. In the latter case the response of the central electron temperature occurs much faster although the central q-value has not altered yet. The current ramp experiments exhibit an interesting tool for clarifying the different properties and confinement physics of the plasma boundary as compared to the hot core.

Introduction

The transport of thermal energy across the magnetic fieldlines is influenced by the shape of the current density profile. Thus the profiles of the electron and ion temperature and density are controlled by the safety factor q_A. An additional constraint must exist to explain the resilience of the electron temperature profile and its insensitivity to the change of plasma parameters other than the ratio of toroidal magnetic field over plasma current. This has been studied previously /1/ by showing that the current density profile cannot be proportional to \( T_e^{3/2} (r) \) in a transient case where the total current is changing in time. A peculiar effect is the fact, that the change of the total energy of the plasma is delayed considerably against the current (or q_A) change, when the current is ramped down. The delay is shorter however in the opposite case, when the current is ramped up. This can be observed on all traces of the electron and ion temperature vs. time (see fig.2); the onset of the \( T_e \) drop (left picture) moves inward with time from the plasma boundary (lower trace) to the center \( r = 0 \) according to the diffusion process of the current density shown in fig 4 (top right). The \( T_e \) profiles remain almost constant in shape but have the same tendency as the \( j(r) \) profiles: at the beginning of the current ramp there is a small \( T_e \) disturbance that diffuses inward with time; the amplitude of this disturbance is very small, however, and seems to be leveled out by some superior constraint. In the opposite case, when the current is ramped up, the \( T_e \) profiles react much quicker, as seen in fig 2 (right). The development of the profiles is masked by the activity of sawteeth, however, which propagate much faster across the plasma and cannot be resolved by the 16 ms time resolution of the laser system. The mechanism, responsible for sawtooth propagation, is obviously of a different nature and not restricted to diffusion processes. A possible explanation for this different behavior during ramp up - or down, is the development of the current density in these two cases especially in the plasma boundary region, leading to a very different deposition of ohmic power and heating close to the separatrix thus causing the different delay times described above. These early experimental
data do not allow a clear separation of the electrons from the ions. For this reason these current ramp experiments have been repeated with more diagnostic equipment to be able to distinguish between the energy transport phenomena through the electron and the ion loss channel.

The current ramp experiments

The purpose of the experiment was to establish a stationary beam heated plasma in the L-regime with 2 different plasma current levels, separated by a short transition phase in which the current is altered. Discharges with a negative and positive current ramp are shown in fig. 1 to 3. The electron temperature and density profiles are measured with Nd:YAG Thomson scattering every 16.7 ms, the ion temperature profiles were measured by active and passive charge exchange diagnostic equipment /2/ before, during and after the current ramp, when the plasma has resumed a stationary state again. The line averaged density has been kept constant during the discharge. The thermal energy of the plasma was calculated from the measured \( T_e \), \( n_e \) and \( T_i \) profiles (see fig 2,3,4). An analysis has been carried out with the TRANSP code, revealing a slight difference between the temporal behaviour of the ion and electron thermal energy content shown in fig. 5. There is a small discrepancy between the calculated total energy \( E_{tot} = E_e + E_i \) and the diamagnetic signal \( E_{dia} \), but one can see the tendency that the ion energy \( E_i \) begins to drop earlier than the electron energy \( E_e \). This indicates that the loss of confinement is initiated by the ions - the electrons react with a delay typically of the order of skin time, about 0.1 sec for the L-regime in ASDEX. Ramping the current down, reduces the ion temperature by a constant additive value at all radii, while the electron temperatures are reduced by a constant factor, leaving the profile shape unchanged (fig 4). Thus the ions in the outer plasma region seem to be responsible for the energy loss at the beginning of the transition phase, while the electrons begin to contribute not before the plasma centre is affected by the current transition. The change of the central ion temperature derived from the neutron diagnostic /3/ is also delayed by typically one resistive time constant. A consequence of the delay of the total energy content against the plasma current is a overshooting of the \( \beta_p \) trace vs. time immediately after the ramp down phase. In the case of a positive current ramp this overshooting is much smaller indicating a shorter delay time between total energy and plasma current.

Summary

- the conductivity profile is decoupled from the electron temperature profile; the shape of \( T_e(r) \) has the tendency to follow \( j(r) \) but deviates only little from its normal shape. A small \( T_e \) disturbance diffuses radially inward on a resistive time scale;

- the change of energy begins in the outer plasma region with the ions; the electrons contribute little at the beginning of a transient phase;

- the change of confinement of the bulk plasma occurs synchronously with that of the inner plasma core after a considerable delay against the current change.

References

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/2/ Herrmann, W., Fahrbach, H.U. IPP Garching
/3/ Bosch, St., IPP Garching
Fig. 1.
(a) line density
(b) gas valve sig.
total current:
neutral beam
power: 1 MW
soft X-signal
loop voltage

Fig. 2.
T_e(t) from
laser scattering
during current
ramp:
top to bottom:
r = 0, 4,
8, 12,...
...40 cm

Fig. 3.
T_i(t) from the
pass. charge
exchange
measurement
at 3 radii

Fig. 1-3: basic plasma parameters versus time in a discharge in which the current is
ramped down (0.42 MA -> 0.28 MA, left) and up (0.28 MA -> 0.42 MA, right)
within 150 ms. Current ramp from t = 1.2 to t = 1.35 sec.
Fig 4. top: development of the electron temperature and current density as calculated from TRANSP during the negative current ramp 0.42 MA → 0.28 MA with time as parameter; bottom: ion temperature from the active neutral charge exchange diagnostics before (1.1 s) immediately after (1.45 s) and in the steady state again (1.8 s) with the current ramped down (left) and up (right).

Fig 5. plasma energy content from TRANSP code (plasma current vs. time on top).

$E_{\text{tot}}$: integrated energy of the ions and the electrons incl. beam contribution;

$E_{\text{dia}}$: diamagnetic signal;

$E_{\text{e}}$: electron energy from laser scattering;

$E_{\text{i}}$: ion energy from the charge exchange diagnostics;

$E_{\text{beam}}$: fraction of the beam energy
ICRF HEATING ON THE BURNING PLASMA EXPERIMENT (BPX)*


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Introduction

RF power in the ion cyclotron range of frequencies (ICRF) has been chosen as the primary heating technique for BPX. This decision is based on the wide success of ICRF heating in existing experiments (JET, TFTR, JT-60), the capability of ion cyclotron waves to penetrate the high-density plasmas of BPX, the ability to concentrate ICRF power deposition near the plasma center, and the ready availability of high-power sources at the appropriate frequency. The primary task of the ICRF system is to heat the plasma to ignition. However, other important roles are envisaged; these include the stabilization of sawteeth, preheating of the plasma during current ramp-up, and possible control of the plasma current profile by means of fast-wave current drive. We give a brief overview of the RF system, describe the operating scenarios planned for BPX, and discuss some of the antenna design issues for BPX.

Overview of the RF System

The RF system will provide at least 20 MW coupled to the plasma throughout the frequency range 50–100 MHz for pulses at least 15 s long. This is to be provided by 16 power sources having a power capability of at least 2 MW per unit at 82 MHz. Quick frequency change operation will be provided in two frequency bands centered at 82 MHz and 54.7 MHz, with operation at any frequency in the tuning range available upon reconfiguration of the tuning system.

The antenna will be of a 'hybrid' design with the current straps and feeders inserted through a port, while the Faraday shield assembly will be mounted on the vacuum vessel wall from inside the device (Figure 1). The rationale for this is to maximize the radiating area of the antenna by permitting the electrically active, current strap components to fill the port area. This optimizes the $k_{\parallel}$ spectrum within the limitations of port size but allows removal of the electrical assembly through the port for ease of remote maintenance. There are four current straps (a $2 \times 2$ array) in each antenna unit and a total of four antenna units. Each current strap is grounded at the center and driven at top and bottom by coaxial transmission lines. The Faraday shield is a single-layer, open design (33% optical transparency) which is passively cooled. The material is high-conductivity carbon-carbon tiles, mechanically attached to Inconel rod. For an estimated 'worst-case' specific antenna loading resistance of 8.3 $\Omega$/m (discussed in section 4) the maximum RF voltage in the system for 20 MW of coupled power is 33 kV.

**ICRF Heating Scenarios**

The standard operating scenario for bulk heating at full toroidal field ($B_0 = 8.1$ T) is to use $^3$He minority heating in a bulk plasma consisting of 50-50 deuterium-tritium (D-T) at $f_{RF} \sim 82$ MHz. As the bulk plasma beta rises, direct second harmonic heating of the T component, which coincides with the $^3$He resonance, takes over. At the high densities employed on BPX, and because of the enhanced collisionality of $^3$He$^{2+}$, strong tail formation is not expected. This will result in bulk heating of the background ions rather than electrons. With the high plasma current of BPX, confinement of tail populations is not expected to be an issue in any case, but it will certainly not be important for $^3$He. Because the wave power propagates essentially radially inward from the antenna and since density gradients provide an additional refractive focusing, the wave energy and subsequent absorption are concentrated near the plasma center. Radial profiles of power deposition in each species as calculated by the ORION 2D full-wave code$^1$ show that the heating is indeed centrally peaked with most of the power deposited in the $^3$He and T components. The single-pass absorption/reflection/transmission/mode-conversion has been calculated using a 1D full wave-code.$^2$ Even at the modest density and temperature of start-up conditions ($n_e(0) = 1.8 \times 10^{14}$ cm$^{-3}$, $T_e = T_i = 4$ keV), the single-pass absorption is $\sim 60\%$ for $\eta_{^3\text{He}} = 1.5\%$. At 2/3 field ($B_0 = 5.4$ T), two heating modes are available, minority $^3$He at $f_{RF} \sim 54.7$ MHz and minority H at $f_{RF} \sim 82$ MHz.

Additional RF scenarios will be employed for special purposes such as heating during the pre-activation phase of BPX using $^3$He in H or $^4$He majority plasmas at $f_{RF} \sim 54.7$ MHz, off-axis heating during the $B_T$ ramp, fundamental D heating in DT plasmas at $f_{RF} \sim 61.5$ MHz, high-frequency minority H heating at $f_{RF} \sim 123$ MHz, and fast-wave current drive at low field.

**Antenna Design Issues**

Plasma physics has an impact on ICRF antenna design in two primary areas: antenna loading resistance and Faraday shield design. For a given antenna design with specified voltage limits, the power that can be coupled to the plasma per antenna is proportional to the antenna load resistance. This quantity can vary over a wide range depending on antenna design and is particularly sensitive to details of the plasma density profile. Therefore, considerable effort has been invested in using the Oak Ridge
recessed antenna modeling code RANT\textsuperscript{3} to determine the ‘worst-case’ density profiles, determining loading for these profiles, and investigating the sensitivity of loading to variations around these difficult profiles. It is known that ICRF loading typically drops during the transition from L-mode to H-mode and that H-mode density profiles tend to be very flat with steep gradients just inside the separatrix and with low plasma density in the scrape-off region. To model this we use the piecewise linear model shown in Figure 2. At relevant densities, loading is a decreasing function of volume-average density, a decreasing function of profile steepness, $\Delta \rightarrow 0$, and an increasing function of scrape-off length, $\lambda_s$.

Figure 3 shows loading versus $\Delta$ for a high-density case near the Greenwald limit $\langle n_e \rangle = 4.4 \times 10^{20}$ m\textsuperscript{-3} for various values of separatrix location $d_s$ and $\lambda_s$. In all cases the density at the separatrix is taken as 10% of the peak density. The current straps are driven out of phase (dipole phasing). It is clear that the critical issue is loading at the highest volume-average density, with steep profiles ($\Delta \sim 0$) and with very little plasma in the scrape-off layer. Therefore, the reference profile for specification of the RF system is taken as the ‘worst case’ of a flat density profile, $\Delta = 0$, with no scrape-off plasma, $\lambda_s = 0$. We see that at the nominal separatrix location, $d_s = 4$ cm, the minimum loading is $8.3 \, \Omega/m$. This value must be reduced by an end effect factor of 0.92 as obtained by scaling from measurements on the TFTR antenna.

Another important consideration in antenna design is minimization of ICRF-related impurities. Theoretical modeling and correlation of calculations with experiment have indicated that local ICRF impurity generation is dominated by energetic ion impacts caused by enhanced, RF-driven sheaths.\textsuperscript{4} The BPX Faraday shield design makes use of the principles of impurity minimization indicated by this work. In particular, the front surfaces of the shield blades are made of carbon, a low-Z material with self-sputtering yield $< 1$ at all energies with normal incidence; the shield blades are tilted and contoured to minimize the RF magnetic flux linkage of field lines that intercept the shield structure; limiters will be strategically placed to minimize the plasma density near the Faraday shield; and the system has sufficient loading resistance to permit full coupling power with dipole phasing of the current straps.

Conclusion

The heating scenarios planned for BPX are those commonly used on existing experiments ($^3$He and H minority). Adequate single-pass absorption is expected with quite small minority concentrations ($\eta_{^3\text{He}} \leq 1.0\%$) for minimal fuel dilution. The RF system design is conservative from a voltage/loading standpoint and makes use of the latest understanding of impurity influx physics.

References

Fig. 1. Perspective view of the BPX Hybrid Antenna.

Fig. 2. Piecewise linear density profile model for H-mode.

Fig. 3. Loading resistance ($\Omega$/m) for the BPX Hybrid Antenna for a density at the Greenwald limit.
RADIATION ASYMMETRIES OF ASDEX DIVERTOR DISCHARGES CLOSE TO THE DENSITY LIMIT

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Introduction
The understanding and control of plasma edge radiation is important for the realization and stabilization of a cold layer at the plasma edge for fusion reactors. Plasma edge radiation and its asymmetries such as the phenomenon of Marfe's also influence the density limit (DL) of tokamak discharges /1, 2, 3/. This paper presents investigations of radiation asymmetries of ohmically heated D$_2$ discharges in ASDEX close to the DL. This includes discharges in the original DV-I and the new DV-II divertor configuration /4/, with uncoated and with boronized vessel walls.

2-dimensional bolometer tomography
In order to study radiation asymmetries, the already existing horizontally viewing 19-bolometer array on ASDEX has been extended by further arrays with 24 more bolometers under different viewing angles in the poloidal plane (Fig. 1). With the new diagnostic equipment, local radiation emission can be determined two-dimensionally as a function of radius and poloidal angle. For data analysis, the one-dimensional treatment by Abel-inversion has been replaced by a two-dimensional tomographic method. The soft X-ray tomography code /5/ had to be adapted to the specific requirements of bolometry, because bolometry deals with radiation profiles which are peaked at the plasma edge and exhibit asymmetries at the edge.

Density Limit and Marfe Limit
In ASDEX, the density limit (DL) is determined in discharges where the line-integrated density $n_e$ is linearly ramped up by feedback controlled gas puffing (Fig. 2). The DL is defined as the maximum obtainable $n_e$ value before the discharge disrupts. The Hugill diagram in Fig. 3 shows that, with regard to the DL in ASDEX, we have to distinguish between low-$q_a$ and high-$q_a$ discharges. At $q_a < 2.6$, the DL is determined by MHD instabilities and we get a DL minimum at $q_a = 2.1$ (with boronization). At $q_a > 2.6$, the DL is a limit to edge density /6/ and we observe the expected linear increase of the DL with $1/q_a$.

In ASDEX, at $q_a > 3$ the DL is preceded by the formation of a Marfe. A Marfe is a cold cloud of edge plasma at the plasma inside boundary with locally enhanced plasma density and radiation emission /7/. Usually, in ASDEX (the discharge with $q_a = 3.0$ in Fig. 2 is a typical example) the Marfe lasts very shortly and leads to an unstable discharge. Therefore, it makes sense to define the Marfe Limit (MaL) as that critical density $n_e$ where the Marfe starts to develop, i.e. when main chamber radiation RAD begins to rise steeply. The Hugill diagram of Fig. 3 compares the MaL and the DL of ohmic D$_2$ discharges in DV-II configuration with boronized walls.

The MaL is reached when the CIII-line emission (977 Å), which is related to the divertor plasma temperature, after a continuous fall during the density built-up has dropped to almost zero intensity and when the D$_{\alpha}$ signal in the divertor has saturated (Fig. 2). This means that a cold and dense plasma has established in the divertor chambers. Prior to the MaL, $n_e$ and RAD start to increase stronger than linearly. Close to the MaL, RAD is typically about 40 % of the heating power. The major fraction of the plasma energy losses is still guided into the divertor chambers, as can be seen from the divertor radiation signal DIV.
Radiation asymmetries at the MaL

Bolometer tomography reveals that discharges at $n_e$ values below the MaL develop significant radiation asymmetries (Fig. 4). The angular distribution of radiation in the poloidal plane exhibits an $m=1$ component with radiation enhancement at the outer plasma edge and an $m=2$ component with radiation enhancement at the top and bottom of the plasma edge close to the X-points. Since the bottom half of the plasma radiates about 15% more than the top half, there is a pronounced down/up radiation asymmetry.

Even at $n_e$ levels far below the MaL, the degree of down/up asymmetry grows approximately linearly with $n_e$ (Fig. 5). But, when $n_e$ approaches the MaL the increase of the asymmetry degree with $n_e$ roughly doubles. Then, the bolometric chord-intensity at 32 cm plasma radius through the bottom plasma becomes 70% larger than the corresponding upper one. Bolometer tomography locates the local radiation enhancement at the bottom of the plasma inside the separatrix at a = 40 cm and around the entire toroidal circumference. Though it is observed in the plasma interior, the asymmetry is associated with the X-points, because the asymmetry vanishes when the plasma is shifted vertically away from the lower X-point and vice versa.

Marfe

Only with boronized walls and at extremely high $q_a$ values of about 6.0, ASDEX discharges develop quasi-stationary Marfes (Fig. 6). During the Marfe, $n_e$ continues to grow over more than 1 s and reaches much higher $n_e$ values (at the DL) than in the pre-Marfe phase (at the MaL). The transition into the Marfe phase occurs suddenly and is characterized by a steep rise of main chamber radiation RAD and a sharp drop of the energy flow into the divertor, as indicated by the divertor radiation signal DIV.

At the MaL of ohmic D$_2$ discharges, RAD does usually not exceed 45% of the heating power input $P_{OH}$. During the Marfe and at the DL, RAD attains values up to 75%. This value is obtained by bolometer tomography and thus corrected for poloidal radiation asymmetries (Fig. 7). Radiation losses during the Marfe are not 100% as reported from other tokamaks/8/. During the Marfe, divertor radiation DIV is more than halved, but remains at about 10% of $P_{OH}$.

Conclusions

At $q_a > 3$, the density limit (DL) is usually preceded by a Marfe. At the onset of the Marfe, the Marfe limit (MaL), main chamber radiation RAD is about 40% of the heating power. During the Marfe, RAD reaches about 75% and there is still a significant energy flow into the divertor. With boronized walls and only at very high $q_a$ ($q_a = 6.0$), the Marfe can be quasi-stationary.

At the MaL, RAD is enhanced over the linear scaling with $n_e$ which is observed at $n_e$ values far below the MaL. The over-proportional rise of RAD can be correlated quantitatively to a local radiation enhancement at the bottom of the plasma close to the lower X-point.

The DL of discharges with Marfe is in agreement with our understanding of the DL as an edge density limit /6/. Enhanced radiation power losses in the main chamber due to the Marfe reduce the heat flow into the scrape off layer. This results in a decrease of plasma edge density /9/ and a peaking of the $n_e$ profile /10/. Hence, during the Marfe (at the DL) higher $n_e$ values are reached than before the Marfe (at the MaL).

References


Fig. 1: Arrangement of viewing chords at ASDEX for 2-dimensional bolometer tomography.

Fig. 3 (right): Marke limit (MaL) and density limit (DL) for ohmic D$_2$ discharges in DV-II divertor configuration with boronized walls.
Fig. 4: Contour plot of radiation distribution in the poloidal plane immediately prior to the MFL. 
\(D_2, I_p = 308 \, \text{kA}, q_a = 3.6, \bar{n}_e = 5.4 \times 10^{13} \, \text{cm}^{-3}\)

Fig. 5: Variation of down/up-asymmetry with \(\bar{n}_e\). The degree of asymmetry is characterized by the ratio of a lower and an upper chord-intensity. Data from discharge intervals with \(\bar{n}_e\) far from and at the MFL are compared.

Fig. 6: Evolution of plasma parameters during a discharge forming a Marfe at 1.18 s 
\(D_2, I_p = 227 \, \text{kA}, q_a = 5.9, \text{boronized walls}\). During the Marfe, diagnostic signals may be false if diagnostics view the Marfe region.

Fig. 7: Radiation distribution in the poloidal plane prior (1.10 s) and during (1.70 s) a Marfe at the inside boundary (same discharge as in Fig. 6).
Experimental Observations of Ohmic and ECR Heated Tokamak Plasmas in RTP


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Introduction

The Rijnhuizen Tokamak Project RTP is dedicated to the study of transport processes in tokamaks. In addition to the examination of steady state conditions of plasmas with ohmic and intense additional heating, a large effort will be made to study the evolution of relevant plasma parameters in time and space during transient states. Theoretical transport models will be compared to experimental findings. In this paper we present an overview of the main results obtained in the first year of physics operation with ECR heating from a single gyrotron. The diagnostics presently operational are a movable, single-point Thomson scattering system [1], a 19-chord interferometer [2], two ECE detection systems (radiometer and polychromator) [3, 4], SXR pulse height analysis, a neutral particle analyser, HXR- and visible light monitors, magnetics, and a 9-point measurement of the single-pass transmitted ECR power [5].

Experimental conditions

The tokamak has a stainless steel vessel with dimensions $R_0 = 0.72$ m, $r_0 = 0.23$ m and a poloidal carbon limiter, usually at a radius of 0.164 m. The maximum toroidal field is 2.4 T, $I_p \leq 250$ kA and the discharge duration is of the order of 250 ms. Conditioning of the first wall is done by prolonged baking ($T \leq 150$ °C) and standard glow- and pulse discharge cleaning techniques yielding $Z_{\text{eff}} = 2 - 3$. Some of the results have been obtained with $Z_{\text{eff}} = 1 - 2$ after boronization of the vacuum vessel with B(CH₃)₃. Reproducible breakdown is achieved by using a short (up to 10 ms) ECR pulse. The results reported in this paper have been obtained in hydrogen plasmas with $<n_e> = (0.5 - 3).10^{19}$ m⁻³, $I_p = 40 - 120$ kA and in the ohmic phase $<T_e> = (0.4 - 0.6)$ keV. The electron drift direction is anti-parallel to the toroidal field.
A single Varian pulse-type gyrotron (60 GHz, 100 ms) is used as a microwave source. It has been operated at various levels of output power $P_{\text{ECR}}$. The transmission line is connected to a launcher, an open circular wave guide with inner diameter of 63.5 mm. The $\text{TE}_{11}$ $O$-mode is launched in the mid-plane perpendicular to the toroidal field from the low field side. Opposite the launcher is a smooth roof-shaped mirror at an angle of 8° from perpendicular. The mode purity is $\approx 90\%$, the overall power loss in the transmission line is $\approx 15\%$, and the vacuum FWHM is $\approx 3$ cm in the centre of the vessel, resulting in a power density $\leq 40$ MW/m$^3$.

Results

The single-pass transmitted power has been measured with a decreasing toroidal field during the ECR pulse for several values of $n_e$ and $I_p$ during the flat top of the discharge. This has revealed a pronounced minimum in transmission, when the resonance is located near the magnetic axis and a smaller minimum with resonance off axis [5]. The plasma behaviour for central resonance has been investigated in more detail. One of the most striking characteristics is the high value (up to 4 keV) of $T_e$ on axis. As shown in Fig. 1, $T_e$ decreases with increasing $n_e$ from the lowest values obtained to cut-off. It is found that $T_e(0)$ is nearly independent from $P_{\text{ECR}}$ over the range 60-180 kW. The $T_e$-values are confirmed by the SXR-PHA diagnostic. By comparing $T_e$-profiles at different levels of $P_{\text{ECR}}$ it is seen that the profiles are almost identical (Fig. 2). The profiles are sharply peaked with gradients up to 50 keV/m inside $r/a = 0.25$, while further out the increase is less but still significant. The large spread in data points, which is limited to the central region, is also seen by the central channel of the radiometer. Thomson measurements of $T_e(0)$ as a function of time reveal that the peak values are reached within 1 ms for $P_{\text{ECR}} = 180$ kW and within 5 ms for $P_{\text{ECR}} = 60$ kW. Off axis ($r = 4.5$ cm) the increase is much slower. In contrast herewith ECE measurements suggest that the central heating occurs in less than 50 μs, which is confirmed by the time behaviour of the transmitted power.

The typical time evolution of the $n_e$-profile is illustrated in Fig. 3. It is seen that during the ECR pulse the profile becomes flat (often even hollow) in the centre and much broader. Compared to the ohmic phase the total number of particles increases. The density pump-out in the centre decreases with increasing $n_e$ and $P_{\text{OH}}$.

Confinement analysis

A power balance analysis has been performed using an adapted version of the SNAP transport code from Princeton. The measured input parameters are $I_p$, $V_1$, $P_{\text{ECR}}$, $T_e(r)$ from Thomson scattering, and $n_e(r)$ from Thomson scattering and interferometry. It is assumed that $\partial q/\partial t = 0$ and $T_i = 0.7 T_e$ in the ohmic phase (which is confirmed by preliminary NPA
measurements). The code calculates the ohmic power density $P_{OH}(r)$ based on neo-classical resistivity, and the classical ion-electron term $P_{Cl}(r)$. At present a gaussian $P_{ECR}(r)$ is assumed. The energy confinement time $\tau_E$ is calculated from the total stored kinetic energy $W$ and the total input power, and the effective electron thermal diffusivity is calculated by

$$\chi_e = - q_e (n_e V T_e)^{-1}.$$ 

This yields typical values of $\chi_e \approx 2 \text{ m}^2/\text{s}$ (ohmic) and $\chi_e \approx 4 \text{ m}^2/\text{s}$ (ECRH phase). In ohmically heated plasmas $\tau_E$ is proportional to $n_e$ and is consistent with the Goldston ohmic scaling law (Fig. 4). For ECR heated discharges with $<n_e> = (0.8 - 1.5).10^{19}$ m$^{-3}$ and $I_p = 60$ kA, $\tau_E$ improves by $\geq 50\%$ if 60 kW ECR is applied. It is interesting to note that although the total input power is almost constant, a much larger $W$ is confined. Increasing $P_{ECR}$ up to 180 kW does not change the value of $W$ appreciably. The calculated value for $\tau_E$ is confirmed by measuring $W(t)$ after switching off the ECR power.

**Conclusions**

Effective ECR heating is found for power deposition inside $r/a = 0.15$. In general, our observations on the heated plasmas are in agreement with those reported earlier by a.o. TFR, PDX, and TEXT [6]. New to our knowledge are the observations: - the modification in $T_e(r)$ and in the total stored energy content are independent from $P_{ECR}$; - replacing part of the $P_{OH}$ by $P_{ECR}$ without changing the total power improves the confinement.

**Acknowledgements**

The authors are indebted to M.C. Zamstorff (Princeton) for making the SNAP code available and to H.G. Esser (KFA-Jülich) for the advice and assistance rendered with the boronization. This work was performed under the Euratom-FOM association agreement with financial support from NWO and Euratom.

**References**

Fig. 1. $T_e(0)$ versus $n_e(0)$ as measured by Thomson scattering for various $P_{ECR}$, 40 ms after switch on.

Fig. 2. $T_e(r)$ for $P_{ECR} = 0, 60, 120, and 180$ kW. The profiles for 60 and 120 kW are at a slightly lower $\langle n_e \rangle$.

Fig. 3. Typical $n_e$-profiles before (dotted line) and during (solid line) the ECR-pulse.

Fig. 4. The observed $\tau_E$ versus the Goldston scaling. The black symbols are for ohmically and the open symbols represent ECR-heated discharges. The values indicated with a $\Delta$ are obtained before and with a $O$ after boronization.
DIFFUSION OF SUPRATHERMAL ELECTRONS MEASURED BY MEANS OF ECRH AND 2ND HARMONIC ECE O-MODE.

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Introduction
In the study of anomalous transport in thermonuclear plasmas, the diffusion of suprathermal electrons deserves special attention. From certain energies onward electrons are effectively collisionless, and therefore follow the field lines. Thus, they can be used to probe the stochasticity of the magnetic field structure. For high energy, electrons are eventually insensitive to magnetic stochasticity as their curvature B-drift becomes larger than the radial correlation length of the turbulence. Hence, by studying the confinement of collisionless electrons in different energy ranges, both the level of magnetic turbulence and the radial correlation length can be established.

A study of the confinement of suprathermal electrons has been reported by Kwon et al, who used measurements of hard X-ray in ASDEX[1]. This study focussed on runaway electrons in the MeV-range, created in the start-up phase of the discharge.

In this paper, we concentrate on the transport of suprathermal electrons with an energy of a few times $T_e$. The advantages of this approach are that a) the curvature B-drift of these electrons is small, so that the transport is sensitive to small scale magnetic turbulence, and b) as we shall show, a local study of the diffusion of these electrons can be made using ECE spectroscopy.

We describe experiments performed in the RTP tokamak, in which ECRH O-mode was launched from the low-field side. In this way, a population of suprathermals in the center of the plasma is almost instantaneously raised in perpendicular energy. This population is diagnosed by ECE with a grating polychromator in the optically thin 2nd harmonic O-mode.

Experiment
The experiments were done in the RTP tokamak ($R_o= 0.72$ m, $a= 0.165$ m, pulse length $> 250$ ms) with $B_T= 2.14$ T, $60$ kA $< I_p < 90$ kA, $n_e(0)= 1.7\pm0.3 \times 10^{19}$ m$^{-3}$. ECR power was launched in O-mode from the low field side (60 GHz, 180 kW, $< 100$ ms pulse length). The ECR power is deposited in the center of the plasma, and single pass absorption was near 100% [2]. ECE spectroscopy was performed using a 6-channel grating polychromator for 2nd harmonic O-mode, and a 20-channel heterodyne receiver for 2nd harmonic X-mode. Other diagnostics included a 19-chord interferometer, Thomson scattering, soft X-ray spectroscopy, transmitted ECR-power measurements and magnetics. In the ohmic phase $T_e(0)= 0.7$ keV, in the ECR-heated phase the central temperature rises to $T_e(0)= 2$ keV (see Fig. 1).
Fig. 1
Te-profile measured by Thomson scattering with and without ECRH. \( B_T = 2.14 \, T \), \( I_p = 60 \, kA \), \( n_e(0) = 1.6 \times 10^{19} \, m^{-3} \), and \( P_{ECRH} = 180 \, kW \).

Fig. 2
ECE-intensity of three O-mode channels and one X-mode heterodyne channel at respectively higher frequencies from top to bottom. The dashed line indicates the leveling of the X-mode intensity.
Fig. 2 shows the evolution of the ECE intensities at frequencies > 120 GHz, i.e. corresponding to positions at the high field side of the magnetic axis. In the ohmic phase of the discharge the O-mode emission shows a constant intensity, attributed to scrambled X-mode radiation. Upon switch-on of the ECRH, the X-mode emission increases rapidly (within <4 ms) in all channels, to typically 2 times the ohmic level, due to the plasma heating. The fast heating is corroborated by Thomson scattering studies. The O-mode emission shows a similar fast increase in the channel which is tuned to the resonance layer. For higher frequencies it also shows an initial fast rise, but it continues to increase at a lower rate for several milliseconds after the X-mode has reached a constant level. The interesting feature of this second phase is that the rise time is a function of the frequency: the higher the frequency, the slower the increase.

**Interpretation and modelling**

Whereas the initial fast rise of the O-mode intensity can be explained by the increase of the X-mode emission, the slow second rise phase must have a different origin. A likely explanation is that the emission is due to suprathermal electrons which have acquired a large perpendicular velocity from the ECRH in the center of the plasma. The slower rise for higher frequency, corresponding to a larger distance from the resonance zone, is then attributed to a diffusion process. Within this interpretation, a diffusion coefficient for the suprathermal electrons can be derived from the time lag of the increase between different measuring channels.

The energy of the suprathermal population can be estimated roughly from the consideration that the strong O-mode emission must be induced by the absorption of ECRH radiation. This puts an upper limit to the energy of roughly ten times $T_e$. For this type of discharge, in the ECR heated phase, the soft X-ray spectrum shows a non-thermal tail at energies above 15 keV, in agreement with this estimation (see Fig. 3). A lower limit is given by the slowing down time: only electrons with sufficient energy can be observed in the O-mode and can retain their energy for the typical diffusion time of a few ms observed.

For this class of electrons, the curvature B-drift is $\approx 0.1$ cm, so that the diffusion of these electrons is sensitive to small scale magnetic turbulence. The frequency down-shift for these energies (=1-2 cm) is of the order of the frequency resolution of the polychromator so that the localization of the emission is not much affected.

A simple model assuming a constant source of suprathermals on the axis, and solving the diffusion equation in cylindrical geometry with zero population at the limiter, qualitatively reproduces the trends in the data. Fig. 4 shows a fit to the data obtained with this model using a diffusion coefficient $D = 0.5$ m$^2$/s for the suprathermal electrons, comparable to the diffusion coefficient for thermal electrons. Although a reasonable agreement with the measurements is obtained, we stress that this is a preliminary result, and that more measurements and modelling is required to come to a reliable estimate of $D$.

**Discussion**

With the method proposed in this paper it is in principle possible to study the diffusion of electrons in an interesting energy range for transport studies. Through the use of ECRH and the possibility to diagnose the transport of the suprathermals locally, well defined experiments can be set up. The interpretation will require much more sophisticated modelling of the
transport in both real space and velocity space. Efforts in this direction are presently being undertaken. Another effect not discussed here is that during ECRH both the critical field and the actual electric field are changed, which can affect the production of suprathermal electrons. This implies that the diffusive model discussed above gives an overestimate of $D$.

Acknowledgements
We are indebted to M. Verreck, M.J. van der Pol, and D. Ferreira da Cruz for making available the radiometer and soft-X-ray data, and to E. Westerhof for enlightening discussions. This work was performed under the Euratom-FOM association agreement with financial support from NWO and Euratom.

References
SIMULATIONS OF $\alpha$ EFFECTS IN TFTR D-T EXPERIMENTS

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Experiments with DT are planned for TFTR to study plasma conditions near breakeven and to study effects of fusion $\alpha$ particles. The discharges which should show the largest effects from the $\alpha$'s are those with a high D-T fusion rate (the $\alpha$ production rate), with a long slowing-down time for the $\alpha$'s, and with a high plasma current to confine the $\alpha$'s as they thermalize.

Simulations of D-T plasmas in TFTR have been performed with the TRANSP transport code. The Monte-Carlo beam-ion calculation was recently extended to include fusion products. The $\alpha$ sources, orbits, slowing down, and heating of the thermal plasma can be calculated. Not included in the TRANSP $\alpha$ simulation are charge exchange, accumulation of the thermalized $\alpha$'s, or collisional coupling with other fast ion populations (D$^+$ and T$^+$).

The simulations considered here are derived from a supershot in deuterium (#55851). This supershot was selected since it had a current high enough (1.6 MA) for good $\alpha$ confinement, a very high central $T_e$ (12.5 keV) for low electron drag, maintained a high neutron rate for a long duration, and extensive diagnostic measurements. The major and minor radii were $R=2.45$ m and $a=0.8$ m, with a toroidal field of 5.1 T, and $q(\psi) = 5.5$.

Two simulations are discussed. The first is a D-T equivalent of the actual deuterium supershot and the second is an extrapolation to higher neutral beam injection (NBI) power. For both simulations, the D:T ratio of the initial plasma and of the D$^0$ and T$^0$ NBI were assumed to be 50:50. For the D-T equivalent, the measured $n_e$, $T_e$, $T_i$, and rotation profiles were taken directly from the shot #55851, along with the measured $Z_{eff}$, NBI powers, voltages, and full and half energy fractions. For the extrapolation, the NBI power was increased from the actual 24.6 MW to 35 MW with beam voltages increased to 120 keV. The electron temperature was...
scaled up 10%, the electron density and ion temperature by 20%, the plasma current by 25%, and $Z_{\text{eff}}$ was scaled down as $1/n_e$.

Results from the D-T equivalent simulation are summarized in the Figures. The fusion yields were $Q_{\text{DT}} = 0.23$ for the D-T equivalent and 0.40 for the extrapolation. The D-T equivalent simulation is conservative in several respects. Since actual measured profiles were used, the simulation does not include mechanisms which might increase the yield. If there is an improvement of confinement in D-T plasmas due, say, to species scaling, then the increased temperatures or ion densities would improve the yield. Also $\alpha$ heating will increase $T_e$. Dilution of the D and T may be overestimated in the simulation due to the assumption that there is only one impurity species.

The extrapolation should be obtainable when the D-T experiments are done in TFTR, due to the accumulation of more operating experience, better heating (increasing the beam powers or voltages, or adding ICRH), and better wall conditioning to control $Z_{\text{eff}}$.

Even in the extrapolated simulation, the heating from the $\alpha$'s is predicted to be only marginally observable. The peak values of $\beta_\alpha$ are 0.3% for the equivalent and 0.9% for the extrapolation. These are large enough that $\alpha$ collective effects such as the toroidal Alfvén eigenmode and high-$n$ ballooning instabilities may be observable.

References

Figures
Figure 1 Volume-integrated heating powers versus time. The $\alpha$ heating is dominated by the beam heating until 4.1 sec, and then by the Ohmic heating.
Figure 2 Components of $\beta(0)$ versus time.
Figure 3 D-T fusion reaction rates versus time.
Figure 4 Profile of the D - T neutron emissivities. The profiles are plotted versus $\phi = \sqrt{\phi}$ the square root of the toroidal flux normalized to the value at the plasma boundary. The value of this variable is very close of the normalized minor radius $r/a$.
Figure 5 Measured profile of $n_e$, and the computed profiles of $n_{\text{beam}}$, $n_D$, $n_T$, and
\( n_{\text{Imp}} \) (the impurity was assumed to be carbon).

Figure 6 Profile shapes of the \( \alpha \) production, density, heating, and \( \beta_\alpha \).

Figure 7 Ratio of the volume-integrated \( \alpha \) heating power to the volume-integrated power loss rate for electrons and ions separately, and for both. Central ignition would be indicated when the ratio for the thermal plasma is greater than 1.

Figure 8 Profiles of the Alfven speed and the Alfven frequency, defined as the Alfven speed / 2 \( \omega_p / (R+r) \).

Figure 9 Profile of the ratio of the initial \( \alpha \) speed and the Alfven speed. This ratio indicates the occurrence of toroidal Alfven eigenmode instabilities.

Figure 10 Profiles of the \( \alpha \) slowing down time and the average \( \alpha \) energy.

**Volume integrated heating power**

**\( \beta_{\text{tor}}(0) \)**

**Neutron emission rates**

**D-T neutron emission profiles**
Density profiles at 3.9 sec

Ratio of $\alpha$ heating powers and losses

Profiles of the Alfvén frequency and speed

Ratio of the initial $\alpha$ and Alfvén speeds

$\alpha$ slowing down time and average energy
THE EFFECT OF PASSIVE STABILIZING PLATES ON HIGH $\beta$ - LOW $q$
DISRUPTIONS IN PBX-M


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Introduction

More than 70% of the plasma boundary in PBX-M is surrounded by an aluminum conducting shell with a thickness of 2.5 cm to stabilize surface MHD modes, especially external kinks. In this paper, we describe the effect of the passive stabilizing plates on high $\beta$ - low $q$ disruptions by measuring the eddy current patterns on the plates with Rogowskki coils mounted on the passive stabilizer electrical junctions. The study includes (1) an evaluation of the plate effectiveness by comparing circular plasmas "loosely coupled" to the plates to "tightly coupled" bean shaped plasmas and (2) an evaluation of the effectiveness of the distance from the passive stabilizer plate and the plasma surface by comparing the MHD behavior for two configurations of the passive stabilizer: "limited poloidal coverage" with L/R time constant = 20 ms and "extended poloidal coverage" with L/R time constant = 40 ms. The degree of conformity of the plates relative to the plasma surface should affect the growth of surface MHD modes.

These studies indicate that the duration of the precursor for oscillatory modes prior to the high $\beta$ low $q$ disruption was prolonged to $\approx$ 30 -50 ms (the order of the L/R time constant) and the duration for locked modes was increased $\approx$ 10-20 ms. However, simultaneous excitation of $n=1$ and $n=2$ modes occurred 1-2 ms before the thermal collapse and led to the total disruption. This final phase is still under investigation.

Passive stabilizing plates

PBX-M passive stabilizers are poloidally segmented into 5 elements, as shown in Fig. 1 for a high $\beta$ equilibrium configuration, with one break in the toroidal direction[1, 2]. A 40 cm separation exists between the top and bottom "hockey stick plates" (the plates closest to the mid-plane on the outboard side near R = 1.8 m) to provide physical access for NBI and horizontal midplane diagnostics. In order to allow helical eddy current flow from the top plate to the bottom plate, 10 electrical connection bars are located between the hockey sticks, every 36 degrees in the toroidal direction. Induced eddy currents are measured by Rogowski coils mounted on these connection bars.

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The characteristic skin time constant of the plates is = 1.5 ms. In the “extended poloidal coverage” configuration, electrical jumpers were added between the outermost hockey sticks plates and their poloidally adjacent plates (“outer cone plates” located at R = 1.6 m). This arrangement increases the poloidal coverage for the helical current by over a factor of two, and increases the L/R time from 20 ms to 30 - 50 ms. Other inner passive plates near R = 1.2 m, 0.75 m and 1.0 m had their top and bottom sections connected in a saddle coil arrangement to stabilize vertical plasma motion.

Plasma conditions for studying high $\beta$ - low q disruptions

When the bean-shaped configuration is produced, the plasma surface geometrically conforms to the passive stabilizer plates. This tight coupling to the passive stabilizers should affect the onset of surface kink modes. In the present experiments the bean shaped configuration was produced by flattening the current profile with a large Ip ramp (dIp/dt = 2 MA/s). The plasma current of 450 - 600 kA with a toroidal field of 11 - 13 kG produced low q (3-4) - high $\beta$ discharges. The typical plasma parameters were $B/Bc = 3.5$, $B = 4.5$ - 6.8%, $\beta_p = 1 - 1.5$, $Ip = 450 - 600$ kA, elongation = 2, and indentation = 28%[3]. The plasma loosely coupled to the passive plates were only slightly elongated and bounded by the carbon limiters. The elongation was 1.2 - 1.3 without any indentation. These discharges were near the q=2 limit. The plasma conditions were $B/Bc = 3-3.5$, $B = 3\%$ $\beta_p = 1.5$, $Ip = 250$ kA with a modest Ip ramp at a toroidal field of 10 - 13 kG.

Effect of distance between the passive plates and the plasma surface

Examples of the toroidal pattern of the induced eddy current measured by Rogowski coils are shown in Fig. 2(a,b) for both cases. Most disruption precursors in the loosely coupled configuration were oscillatory and propagated in the toroidal co-direction relative to NBI with a dominant growth time of 100 - 200 $\mu$s. The duration time of the precursor was comparable to the growth time.

On the other hand, in tightly coupled cases, both oscillatory and locked modes were observed with almost equal probability, and grew on a time scale of 10 - 15 ms, which is one order of magnitude longer than in the loosely coupled cases. The detailed time-behavior is shown in Fig. 2(c,d); shown are the x-ray and vertical position monitor flux signals and the amplitude and phase of the n=1,2 components. For loosely coupled elongated plasmas, the n=1 amplitude grew and the phase velocity decreased near the time of the thermal disruption, while the n=2 amplitude grew very rapidly in the final phase. The n=1 phase velocity decreased and reversed during the thermal collapse phase (the global behavior is seen in Fig. 2-a). The locked mode in the tightly coupled case showed preferential lock locations near 250 - 330 degrees as shown in Fig 2-d; the reason for this is not yet understood. The n=2 amplitude increased 1-2 ms before the thermal disruption. This behavior of simultaneous n=1, 2 excitation seems similar to the loosely coupled case. The vertical plasma position remained constant during the thermal collapse until the electron temperature was reduced to nearly zero.

Figure 3 summarizes comparisons of the duration period for the weakly and strongly coupled passive plates vs. the n=1 amplitude at the time just before the simultaneous excitation of n=1 and 2 modes. In this figure, the “+” symbol corresponds to the nearly circular loosely coupled discharges, and the remaining points are for the tightly coupled bean shaped configuration. The loosely coupled configuration exhibited a two orders magnitude shorter duration period compared with tightly coupled cases.

Effect of extended coverage of n=1 passive stabilizing plates

Figure 3 also includes a comparison of the duration periods for the limited and extended coverage configurations. The data marked in solid symbols are for the “extended coverage” configuration with oscillatory precursors (circles) and locked mode (squares). The data marked in open symbols are for the limited coverage configuration with oscillatory precursors.
Fig. 2 Comparison of eddy current behavior between loosely and tightly coupled configurations.
(circles) and locked modes (squares). The duration time for the extended coverage geometry is longer by a factor of two compared with the limited coverage configuration for both oscillatory and locked modes. Furthermore, the oscillatory modes persist for times which are longer by a factor of two compared with locked modes. These results suggest that the growth time was increased with the increase of the L/R time and that the decreased skin-penetration time due to the oscillations improved the stabilizing effect. It is to be noted that even locked mode stabilization took place on the L/R time-scale, which is much longer than the skin time of 1-2 ms.

Summary

The analysis of eddy currents in the passive plates indicates that the plates function as a means for slowing down the growth time in a way consistent with passive stabilization concepts. This prolonged duration may provide time interval long enough to actively control profiles and to avoid the strong crash associated with the final disruption. However, during the final phase just before thermal collapse, the modes grow on a much faster time scale. These fast growth times may be related to the location of non-linear phenomena, which might be taking place further inside the plasma where the passive plates are not as effective.

Acknowledgement

We would like to acknowledge the technical contributions of L. Gereg and the technical staff headed by J. Semler. The PBX-M project is supported by the U.S. DoE under contract No. DE-AC02-76-CHO-3073.

References


Fig. 3 Comparison of the duration period of precursors between the loosely coupled elongated configuration, tightly coupled bean shaped configuration, and the limited and extended coverage configurations.
EFFECT OF PLASMA ASPECT RATIO
ON PLASMA CONFINEMENT PROPERTIES IN JIPP T-IIU

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1. Introduction

Studies of plasma confinement are intensively carried out in many tokamaks. The transport mechanism is still unclear, but the plasma transport may be enhanced by plasma toroidicity. L-mode scaling of energy confinement time $[1]$ suggests this possibility, because it has fairly strong dependence on plasma major radius ($R$) and weak dependence on plasma minor radius ($a$). On the other hand, a lot of plasma turbulence theories discuss the effect of plasma toroidicity. For instance, the diffusivity is proportional to $\varepsilon^{1.5}$ for the dissipative trapped-electron mode $[2]$, and to $\varepsilon^{0.5}$ for the collisionless trapped-electron mode $[3]$, where $\varepsilon$ is the local inverse aspect ratio $r/R$.

In JIPP T-IIU$[4]$, we study the effect of plasma toroidicity on the transport of neutral beam heated plasmas, changing the minor radius.

2. Experimental Conditions

The minor radius is varied from 24 cm to 12 cm by a rail type carbon limiter moved vertically. Here, we fix the major radius ($R=93$ cm), plasma current ($I_p=88$ or 110 kA) and toroidal magnetic field ($B_t=2.9$ T), because there is the strong $I_p$- and $R$-dependence and no $q(a)$-dependence in L-mode scaling. The aspect ratio $R/a$ is varied from $\approx 4$ to $\approx 8$. Hydrogen beam with the energy $E=36$ keV is injected.
tangentially, and total deposited power is kept constant for various minor radii: \( P_{NB} = 350 \text{ kW} \) at \( I_p = 110 \text{ kA} \). The ratio of \( P_{NB} \) to ohmic heating power is in the range of 1.4-3.5. Plasma stored energy is measured by a diamagnetic loop, and is also estimated from the measurements of electron temperature, density and ion temperature profiles. Two Mirnov probes (MP) and four electrostatic single probes (LP) are mounted on this movable limiter, to detect low frequency fluctuations \( f \leq 250 \text{ kHz} \) of poloidal magnetic field and electron density or floating potential behind the limiter, where each probe position is \( r/a = 1.12-1.25 \) for MP and \( 1.05-1.09 \) for LP.

3. Experimental Results

Figure 1 shows energy confinement time as a function of the aspect ratio \( R/a \). The confinement time is almost independent of \( R/a \). Note that, in these discharges, line averaged electron density \( \langle n_e \rangle \) increases linearly with the decrease in the minor radius (from 24 cm to 13 cm) during beam injection. The data in Fig.1 will reflect the \( R/a \)-effect rather than the difference in \( \langle n_e \rangle \), because the density dependence in L-mode confinement is fairly weak. The data in Fig.1 suggest the effective thermal diffusivity \( (\chi_{\text{eff}} = a^2/\tau_E) \) becomes small for the larger \( R/a \). The radial profiles of electron and ion temperatures are shrunk with the decrease in \( a \), keeping the values at the plasma center (\( T_e(0)=800 \text{ eV}, T_i(0)=600 \text{ eV} \)), while \( n_e \) increases from \( 3.3 \times 10^{13} \text{ cm}^{-3} \) to \( 8.5 \times 10^{13} \text{ cm}^{-3} \) with the decrease in \( a \). The radial profile of beam power density calculated from NFREYA-code[5] changes from fairly peaked profile such as \( [1-(r/a)^2]^2 \) to parabolic one \( [1 - (r/a)^2] \) with the decrease in \( a \), raising the power density at the center by a factor of two. We calculate the following simplified local thermal diffusivity: \( \chi_{\text{flux}}(r) = -(Q_e+Q_i)/(n_e\nabla T_e+n_i\nabla T_i) \),
where \( Q_e \) and \( Q_i \) are electron and ion heat fluxes flowing out of a magnetic surface of minor radius \( r \) [6], because collisional power flow from electron to ion is considerably large in these discharges. We assume \( n_e = n_i \), taking account of relatively low \( Z_{\text{eff}} \)-value (\( \leq 1.5 \)). Moreover, we neglect radiation, charge-exchange and convection losses. This \( X_{\text{flux}}(r) \) will give the upper bound near the plasma edge.

Figure 2 shows the \( X_{\text{flux}}(r) \) in the range \( r/a \leq 0.8 \) for various limiter radii as a function of minor radius \( r \). The abscissa can be regarded as local inverse aspect ratio \( r/R \), where \( R \) is fixed at about 93 cm. The confinement of the plasma with higher aspect ratio is almost equivalent to the central good confinement of the plasma with lower aspect ratio. This clearly indicates the plasma toroidicity plays an important role in tokamak transport. The results are similar to the data obtained from \( R- \) and \( a- \) scan experiments in TFTR [7].

Fluctuations detected in scrape-off layer plasma sometimes reflect the transport of a core plasma. We measure the relative amplitude of incoherent fluctuations of poloidal magnetic field and ion saturation current. Figure 3 shows the dependence of the relative amplitudes on \( R/a \). The density and temperature fluctuations predicted from the ion saturation current decrease

![Fig.2 Radial profile of the local thermal diffusivity defined in the text for various limiter radii, where \( L_p = 110 \) kA and \( R=93 \) cm.](image)

![Fig.3 R/a-dependence of relative amplitudes of incoherent fluctuations of ion saturation current and poloidal magnetic field detected behind the movable limiter, where \( f=90-110 \) kHz.](image)
with the increase in \( R/a \), and this feature agrees well with the behavior of \( \chi_{\text{flux}}(r) \). On the other hand, the fluctuations of poloidal magnetic field exhibit the reverse dependence on \( R/a \). Electrostatic modes such as drift-wave mode may be closely correlated to the transport of these discharges, rather than electromagnetic modes with enhanced magnetic fluctuations.

4. Conclusion

We study the effect of plasma toroidicity on transport properties of L-mode plasmas changing only minor radius with movable limiter at fixed major radius. The global energy confinement time exhibits less dependence on the minor radius. Local thermal diffusivity derived from the simple transport equation is reduced with the decrease in the plasma toroidicity. This dependence is resemble to that of relative amplitude of ion saturation current detected behind the limiter, while it has inverse tendency for that of poloidal magnetic field fluctuations.

RESULTS AND UPDATE ON THE TdeV FACILITY


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INTRODUCTION

TdeV (Tokamak de Varennes) was operated in 1988 in a limiter mode [1] and upgraded in 1989-90 to run in a double or single null poloidal divertor configuration with neutralizing plates in closed chambers behind an Inconel first wall liner (R=0.87 m, a=0.28 m and I_p up to 300 kA). Graphite limiters and divertor plates are electrically insulated from the liner and support structures for plasma biasing and current injection experiments. A new pair of internal fast horizontal positioning coils is also available. A comprehensive set of diagnostics permits detailed impurity and current transport experiments as well as studies of scrape-off layer physics and plasma-wall interactions under various conditions. A 1 MW, 3.7 Ghz lower-hybrid-current-drive system is currently under construction for long pulse plasma experiments (~30 sec) scheduled to begin in 1993.

Impurity transport, production and control is the main long term objective of the TdeV program. Accordingly, the machine can be operated either in a divertor or limiter mode with full added capabilities for plasma biasing experiments under several configurations. New plasma facing material testing and studies are also emphasized (boronized graphite, TiC, composites,...). Plasma current transport is also studied using either perturbation techniques or fast plasma current rampup or rampdown; this work takes advantage of the very accurate and fast plasma position control on TdeV. The addition of the long pulse (~30 sec), high power (1 MW/m^3) lower hybrid current drive system will add significant capabilities for continued studies on plasma current transport, impurity control and material testing.

This paper summarizes significant features now available on TdeV and reports briefly on interesting results on divertor plasma biasing experiments.

TdeV FEATURES AND PARAMETERS

A cross-sectional view of the machine is shown in Fig.1. Its main feature is the double null divertor geometry with closed chambers. The divertor triplet coil configuration

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* Supported by Atomic Energy of Canada Limited, Hydro-Québec and the Institut national de la recherche scientifique.
Figure 1 Cross-sectional view of TdeV

provides very good control of the separatrix X-point at the divertor throat. The graphite neutralizing plates are mounted on water-cooled copper heat sinks for the forthcoming long pulse regime; in addition, separatrix sweeping on the outer plates further reduces the thermal load. An Inconel liner surrounds the plasma, except for the diagnostic and auxiliary ports and a large fraction of the inside wall is covered with graphite tiles used as a guard or main limiter. Movable graphite poloidal limiters are located at the top, bottom and on the outboard equator; they can be heated to 500 °C or cooled to a fixed temperature during machine operation. All three limiters and the divertor neutralizing plates are electrically insulated from the chamber and can be connected in various combinations to three different power supplies (<600 V, several hundred Amps) for biasing or current injection. The neutralizing plates are separated toroidally into two insulated halves. Currents or floating potentials on all these components are monitored together with the current flowing to the two halves of the chamber.

A pair of coils (±6 kA, few ms response time) inside the vacuum chamber supplements the external equilibrium field coils for good horizontal position control (±1 mm with a= 0.27 m) during perturbative experiments or rapid plasma parameter changes (vertical position is stable). For added flexibility, control is performed through gated digital waveforms and a 0.3 msec digital feedback loop on all significant parameters (l_p, \bar{n}_e, R_p, Z_p, V_{\text{blunt}}, I_{\text{blunt}}, I_{\text{DP}}, I_{\text{FP}}, I_{\text{HP}}, I_{\text{TP}}, fueling valves ...). System-wide timing and sequencing is provided by coded markers on a 1 MHz fiber optic highway clock with in-crate encoders/decoders available to all diagnostics and subsystems.

At present, the machine is routinely operated with ~1 s pulses every few minutes (l_p = 250 kA, q_{cyf}= 3.8, OH flux= 2.5 V-s, R_p = 0.87 m, a= 0.27 m and B_T=1.5 T). Water cooling for the coils and power supplies is already adequate for long pulse operation and pumping
Table 1.  

TdeV diagnostic set (some are not yet fully commissioned). Emphasis is on diagnostics for impurity and current transport studies, as well as edge and divertor physics.

<table>
<thead>
<tr>
<th>Diagnostic Set</th>
<th>Description and Measurement Parameters</th>
</tr>
</thead>
<tbody>
<tr>
<td>Flux coils</td>
<td>$R_p(t), I_p(t)$ fast digital feedback (0.3 ms), magnetic measurements</td>
</tr>
<tr>
<td>$\mu$-wave interferometer</td>
<td>$n_u(t)$ vertical line densities for feedback</td>
</tr>
<tr>
<td>SMM interferometer</td>
<td>$n_s(r,t)$ 7 channels at 393.6 $\mu$m, 10 kHz profile sampling</td>
</tr>
<tr>
<td>Thomson scattering</td>
<td>$T_e(r,t)$ multipulse glass (50Hz), 4 spatial points</td>
</tr>
<tr>
<td>Elect. cycl. emission</td>
<td>$T_e(r,t)$ 2 channels, fixed or fast sweeping frequency</td>
</tr>
<tr>
<td>PHA soft X-rays</td>
<td>$T_e, n_{e-high}(r,t)$ 4 independent ch., .5-30 kev, fast radial scan</td>
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<td>Hard X rays</td>
<td>suprathermal 4 energy channels</td>
</tr>
<tr>
<td>VUV I spectroscopy</td>
<td>$T_i(r,t)$ Doppler broadening, multi-ch.($\lambda, r$), MCP/anode</td>
</tr>
<tr>
<td>VUV II spectroscopy</td>
<td>$n_{i-low}$ MCP/reticon ($\lambda, r$), multi-spectral lines</td>
</tr>
<tr>
<td>Charge exchange</td>
<td>$T_{\theta}(t)$ escaping neutrals</td>
</tr>
<tr>
<td>Bolometry</td>
<td>$P_{rad}(r,t)$ 16 channels, horizontally viewing array</td>
</tr>
<tr>
<td>Z$_{eff}$ meter</td>
<td>$Z_{eff}(r,t)$ 8 channels, visible bremsstrahlung</td>
</tr>
<tr>
<td>X-ray tomography</td>
<td>MHD, ($T_e, n_e$) 5 filtered arrays, full 2-D reconstruction 100 kHz</td>
</tr>
<tr>
<td>Mirnov coils</td>
<td>$B_0$ 14 poloidal + 4 toroidal coils, $\leq$250 kHz</td>
</tr>
<tr>
<td>CO$_2$ laser collect. scat.</td>
<td>$n_e$ $6 &lt; k &lt; 180$ cm$^{-1}$, $f &lt; 3$ MHz</td>
</tr>
<tr>
<td>SMM polarimetry</td>
<td>$J(r,t)$ 6 channels at 393.6 $\mu$m</td>
</tr>
<tr>
<td>$H_\alpha$ emission</td>
<td>$H_{recycling}$ 15 ch. poloidal array + filtered TV cameras</td>
</tr>
<tr>
<td>Edge visible spectrosc.</td>
<td>$n_{i, v_{\theta}}$ (edge) filtered diodes, spectro/OMA, spectro/camera</td>
</tr>
<tr>
<td>Edge Li/C laser ablat.</td>
<td>$n_{e, T_e}$ (edge) filtered reticon</td>
</tr>
<tr>
<td>Langmuir probes</td>
<td>$n_{e, T_e}$ (edge) equatorial, fixed and fast scanning</td>
</tr>
<tr>
<td>Plate/limiter probes</td>
<td>$n_{e, T_e}$ (plate) 100 flushmounted, separatrix position on plate</td>
</tr>
<tr>
<td>Fast mass spectrometry pressures</td>
<td>partial/total, main/divertor chamber</td>
</tr>
<tr>
<td>Sniffer probe</td>
<td></td>
</tr>
<tr>
<td>Surface station</td>
<td>samples</td>
</tr>
<tr>
<td>Wall neutral flux</td>
<td></td>
</tr>
<tr>
<td>TV cameras</td>
<td></td>
</tr>
<tr>
<td>under development:</td>
<td></td>
</tr>
<tr>
<td>Core visible spectrosc.</td>
<td>$T_P, Z_{eff}(r,t)$ CID camera, $2000 &lt; \lambda &lt; 7000$ $\AA$</td>
</tr>
<tr>
<td>Neutral beam</td>
<td>$n_{i-low}, J(r,t)$ 40 kV, 1A, 200 mA/cm$^2$, $H/H_\alpha$</td>
</tr>
<tr>
<td>Edge neutral beam</td>
<td>$n_{e, \phi}$ (edge) $\sim$40 kV, thallium beam, electrostatic analyzer</td>
</tr>
</tbody>
</table>

will soon be available in the divertor chambers to control pressure buildup during operation. Preliminary results indicate that the divertor provides a significant impurity reduction with typically a 25% reduction in loop voltage with respect to poloidal limiter operation. Quite interestingly, a hybrid mode with limiters and energized divertor coils also yields good impurity control, pointing out the importance of plasma on stray field lines several centimeters away from the last closed flux surface. Table 1 lists the diagnostic set currently installed on TdeV with some under development.

**PLASMA BIASING EXPERIMENTS**

Plasma biasing [2] is a good example of TdeV's capabilities and flexibility.
approach is very promising for controlling the impurity influx to the plasma and affecting the electric field inside the plasma, the latter being possibly related to H-mode physics. An electrical schematic of the setup is shown in Fig. 2a. Biasing between -200 and +150 V was applied to the upper outer divertor plate. Other plates, electrically floating, are used to ascertain that a uniform potential is indeed applied to the separatrix with respect to the chamber walls.

Partial results are shown in Fig. 1b. Global plasma parameters such as the electron temperature and density profile do not appear to be affected, although gas puffing was required during biasing to maintain the density. Negative biasing is expected to decrease ion sputtering from the walls and reduce the impurity influx. This appears to be the case in Fig. 1b with a 15-20% reduction in loop voltage and a significant reduction of soft X-ray emission on all spatial channels. Part of this improvement must be credited to a significant increase in the divertor pumping efficiency as can be seen with the pressure buildup in the upper divertor chamber (unpumped) during biasing. Since the lower chamber pressure was not affected, an expected $E_{\text{ex}}B_{\text{b}}^2$ drift in the scrape-off layer can be assumed to be responsible for the pressure increase. Transport could also have been affected inside the separatrix as suggested by the microturbulence $n^2$ ($\text{CO}_2$ laser scattering) in Fig. 2b. Poloidal rotation, observed via spectroscopy on a CII line, is strongly affected by biasing [2]. Velocity shear could then be responsible for microturbulence stabilization. Velocity measurements also indicate that the electric field inside the separatrix is indeed modified and is directed oppositely to the field applied between the separatrix and the walls.

**CONCLUSIONS**

TdeV, a medium size tokamak with a double null divertor and a full diagnostic set, is very well equipped for impurity control and current transport studies. The high power LHCD system now under construction will extend its capability to long pulses with good current profile control. Plasma biasing with a closed divertor has been shown to be a promising technique to reduce impurity influx and to study electric field effects both just outside and inside the separatrix.

**References:**
THE INFLUENCE OF HYDROGEN INFLUX TOROIDAL INHOMOGENEITY ON THE PARTICLE BALANCE ANALYSIS IN FT-1 TOKAMAK EXPERIMENTS.

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In the study of particle and energy balance of a hot tokamak plasma, one must take into account the role of neutral hydrogen (deuterium), especially near the wall, where its density may be rather high, up to $N_H\sim (10^{10} - 10^{11})$ cm$^{-3}$. Being the source of charged particles, the atomic hydrogen at the same time due to the charge-exchange processes changes the ionization state of impurity ions, increasing strongly the radiation losses [1]. These well known processes are used for energy dissipation in divertor volumes of large tokamaks [2].

In FT-1 tokamak experiment ($R=62.5$ cm, $a=15$ cm) the role of poloidal limiters and gas puffing valve in particle recycling processes was studied. The measurement of radial and toroidal distribution of atomic hydrogen density was carried out by resonance fluorescence and passive spectral diagnostics [3].

In the calculations the ionization flux of hydrogen atoms $Q_i^j(r)$ was presented as a sum of toroidally averaged fluxes: wall flux $Q_w^j$, limiter flux $Q_l^j$ and gas puffing flux $Q_p^j$.

From our data it is clear that accounting of the additional hydrogen entering plasma from the limiter and gas puffing valve increases the flow $Q_i^j$ in comparison with the flow from chamber wall $Q_w^j$ 1.7 times at $n_e(0)=1\times10^{12}$ cm$^{-3}$ and 2.3 times at $n_e(0)=2\times10^{13}$ cm$^{-3}$ [4].

The accurate calculation of the radial diffusion coefficient of charged particles distribution is connected with defining all electron sources, i.e. the impurity contents are measured. In the experiment on tokamak FT-1 the local ionization state was determined spectroscopically with the help of scanning.
relatively calibrated vacuum monochromator VM (500Å-3000Å).
The absolute calibration was done comparing the Abel-inverted
distribution of Lα line with calculated Lα intensity based on
the resonance fluorescence measurements of atomic hydrogen
density in plasma. These data and the radial radiation prof-
iles of the oxygen, carbon and nitrogen lines permitted to get
the radial distribution of impurity ions.

The density of "k" ion of "i" charge of light impurities is
evaluated comparing the experimentally determined intensity and
the calculated quantity:

$$E_{ps} = \frac{\Omega}{4\pi} \int \left( \frac{n_e (r)n_i (r)q_{ps}}{n_k} + n_e (r)n_{i+1}q_{ps} + n_{k+1} N_H q_{ps} \right) dV,$$

(1)

The integration is done over volume of the cone of VM ob-
ervation V_k. The q_{ps}, \tilde{q}_{ps}, and \tilde{E}_{ps} are collision, recombination and
charge-exchange excitation functions of ps transition of n_k ion; E_{ps} is the energy of this transition.

Excitation by electron collisions is the most important in
a hot plasma [5]. But in the plasma periphery (T_e < 50 eV, n_e <
5 x 10^{19} cm^{-3}, n_0 = 5.10 cm^{-3}) recombination and charge exchange excita-
tion are important also, especially in the case of nonresonant
transitions.

VUV monochromator was used for measurements of emission of
some spectral lines from similar plasma volumes surrounding the
limiter and remote from it. The resonant lines Lα (1216 Å),
OVI (1032 Å) and nonresonant one OVII (1632 Å) were studied.
The scheme of experiment is shown at Fig. 1. The movable aluminium mirror in a limiter shadow in front of the entrance slit
of a monochromator permits to observe the optical emission of
plasma along chords in the equatorial plane of a torus, from
both sides of a diagnostic cross-section. As Fig.1 shows, the
observation area is 12x3 cm in the plane of a limiter. The com-
parison of signals F. obtained in ohmic discharge of FT-1, from
the side of a limiter and from the opposite side, in dependence
on the mirror positions angle at Fig 2. The position of limiter
edges is also shown there.

The signals of spectral lines Lα, OVI, OVII are present-
The increased influx of a hydrogen due to recycling at the limiter results in the increased emission of \( L^\alpha \) from the limiter side. These results depend on vacuum conditions and the level of the heat flow on the wall and limiter. For example, the differential signal \( \Delta F_{L^\alpha} \) was increased during the LH-experiment (the dotted line on Fig. 2). The increased emission of hydrogen line from the side of a limiter indicates the reinforcement of hydrogen recycling here due to fast particles [4].

The emission of O\( \text{VII} \) and O\( \text{VI} \) spectral lines is also increased in the direction of limiter side. As a far as the concentration of \( O^+ \) and \( O^+ \) ions is constant around the torus this increased emission of oxygen lines indicates the reinforcement of charge exchange processes of oxygen ions near-by the limiter. The increased O\( \text{VII} \) and O\( \text{VI} \) line intensity is explained by reaction of charge-transfer excitation of impurity ions [6]. There are the mathematical complications in evaluation of the cross-sections of these processes. Nevertheless, the velocity functions of charge-exchange excitation \( \bar{\sigma}_{\text{ps}} \) can to present through ratio of F signals and well-known velocity function \( \tilde{q}_{\text{H}} \) of hydrogen atoms from eq. (1). For example, it's easy to obtain for O\( \text{VII} \) expression:

\[
n_7^+ \tilde{\sigma}_{52} = (\Delta F_{\text{OVII}}/\Delta F_{L^\alpha})(E_L/E_{52})q_{L^\alpha}.
\]

Thus, the ionisation state of impurity ions in a tokamak plasma can depend on the elevated atomic hydrogen density in some specific volumes of a discharge, such as limiter and gas puffing cross sections. This phenomenon must be taken into account in analysis of the ionisation balance of impurity ions and in the study of energy losses by radiation.

REFERENCES
Fig. 1
The scheme of experiment

Fig. 2
The signals $F$ from the side of limiter and from the opposite side in dependence on mirror positions angle $\theta$. The dotted line is $\Lambda_\theta$ signal during LH-hitting.
THE EXPERIMENTS ON FAST CURRENT RAMPDOWN IN TUMAN-3 DEVICE


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INTRODUCTION. This paper describes the evolution of the plasma parameters in the TUMAN-3 tokamak (R = 52 cm, a = 22 cm, B_T = 5.7 kG, n_e = 1-1.5 \times 10^{13} \text{cm}^{-3}) in the regimes with fast current rampdown. In the described experiments plasma current decreased in 2 times from initial steady state value I_p = 85 kA down to 42 kA during the period \Delta t = 5-7.5 ms. The time of current decrease was a factor of 4-6 shorter than the characteristic time for the diffusion of magnetic field \tau_m. The typical energy confinement time \tau_E in steady state before the rampdown did not exceed the values \tau_E = 1-1.5 ms. This means, that the proportion \tau_E << \Delta t << \tau_m took place. An approximate estimations show that more than a half of the total value of I_p flows inside plasma radius r = a/2 before the rampdown phase. The rapid decrease of the plasma current must produce an appreciable narrowing of the current density profile and thus help to understand the dependence of the plasma heating and the confinement of energy and particles on the j(r) distribution in ohmic regime.

Similar experiments performed in TFTR [1] in neutral beam heated plasma showed that global confinement in L-mode was insensitive to the current outside r = a/2 but depended on the current in the core of the plasma. Hence, for fixed I_p, \tau_E can be substantial increased relative to L-mode scaling [2] by peaking the current profile. The rapid I_p-decrease in TdeV [3] in ohmically heated plasma showed that the peak temperatures had a tendency to remain approximately constant although I_p decreased more than a factor of 5
EXPERIMENTAL RESULTS. The time evolution of the main plasma parameters in TUMAN-3 current rampdown experiments is shown in fig.1. The $I_p$ decrease was followed by the negative inductive spike in loop voltage $U_p$ (fig.1a). If one neglects the change in plasma inductance $L_p$ during rampdown phase, then the amplitude of negative spike in $U_p$ due to the variation of the term $L_p \frac{dI_p}{dU_p}$ only should be a factor of 2 larger than the one observed experimentally. The rude analysis shows, that the behaviour of $U_p$ in rampdown phase is described by the approximately 2-fold increase in plasma internal inductance $l_1$.

The average electron density, the plasma energy $W_\perp$ (fig.1b), measured by diamagnetic loop, and the electron temperature in plasma core $T_{e0}$ (fig.1c), defined from the analysis of the soft X-ray emission did not change significantly during $I_p$ decrease. The evolution of the ion temperature $T_{i0}$ in the vicinity of plasma center, measured by the 5-channel neutral particle analyser (fig.1d), was the same. An insignificant change in plasma energy resulted in appreciable, approximately a factor of 3, rise in $\beta_p$ (fig.1b). After the end of $I_p$ rampdown a slow decrease of energy and temperatures with a characteristic time about 10 ms was observed.

The comparison of the plasma energy in the steady state conditions and just after the end of the $I_p$ decrease (fig.2) shows, that in the last case the values of $W_\perp$ were 40-50% higher than for the plasma current plateau 42 kA.

Fig.3 shows the radial distributions of electron density, observed by the multichannel 2-mm interferometer during current plateau before the rampdown (1) and at the moments 7.5 ms (2) and 12 ms (3) after the start of the $I_p$ decrease. The external gas puffing remained constant during all period of observation. One can see, that the decrease in plasma current was followed by the narrowing of density profiles. The shape of profile (3) corresponds to the density distribution in the steady state plasma with initial $I_p = 42$ kA. Fig.1e shows the
time evolution of the intensity of $D_\alpha$ emission. The appreciable decrease in $D_\alpha$ emission as well as the narrowing of electron density profiles indicates that the fast current rampdown results in some improvement of particle confinement.

**CONCLUSION.** Thus, an essential, a factor of 2, fast current rampdown did not changed significantly main plasma parameters - energy, temperatures in plasma center and density. Apparently, in ohmically heated plasma, as well as in the regimes with neutral beam injection [1], the heating and the global plasma confinement are determined by the plasma current in the core of the plasma. The time of the decrease of energy and temperatures after the end of the rampdown was comparable with the characteristic time for the diffusion of magnetic field.

![Fig.1 Time evolution Of the plasma parameters in the experiment with $I_p$ decrease from 85 kA to 42 kA.](image)
Fig. 2 Plasma energy $W_\perp$ (from the diamagnetic loop) versus average density $\bar{n}_e$.
- $\triangle$ - steady state conditions with $I_p = 42$ kA
- $\square$ - just after the end of $I_p$ decrease from 85 kA to 42 kA.

Fig. 3 Measured radial profiles of electron density.
1 - current plateau before the rampdown, 2 - 7.5 ms after the start of the $I_p$ decrease, 3 - 12 ms after the start of the $I_p$ decrease.

References
Introduction

A fundamental question in tokamak physics is to what extent can plasma current penetration be explained by either Spitzer or neoclassical resistivity. Knowing the current profile is crucial to determining the plasma equilibrium and stability. Recent experiments have shown results indicating better agreement with the neoclassical model [1,2], although these results were generally based on indirect measurements (loop voltage, location of the \( q=1 \) surface). Measured and calculated magnetic field line pitch profiles from different types of plasma discharges in PBX-M [3] have been compared for the purpose of addressing this issue. The magnetic field line pitch is defined as \( \alpha = \tan^{-1}\frac{B_P}{B_t} \), and it is measured on PBX-M by the Motional Stark Effect (MSE) diagnostic [4]. Model field line pitch profiles were computed by solving the fully time dependent poloidal field diffusion equation in the TRANSP code using measured \( T_e \), \( n_e \), and \( Z_{\text{eff}} \) profiles and assuming either Spitzer or neoclassical resistivity. For cases in which these two models could be distinguished, the measured field line pitch profiles were in best agreement with the Spitzer model, especially near the discharge center, although the results of this analysis clearly indicate that aspects of the current diffusion may be anomalous.

Data and Analysis

Discharges used for this analysis were steady-state (\( \frac{dI_p}{dt}=0 \) for 300 msec), and had discharge parameters of \( R=165 \text{ cm}, \theta_{\text{mid}}=28 \text{ cm}, I_p=350 \text{ kA}, B_t=1.35 \text{ T}, \) elongation=1.8, indentation=15\%, \( n_e=6 \times 10^{13} \text{ cm}^{-3}, \ T_e(0)\leq 2 \text{ keV}, \ T_i(0)\leq 4.5 \text{ keV} \), and had \( P_{\text{inj}}=2.5-5 \text{ MW of D}^0 \text{ injection into D}^+ \text{ plasmas in either the co- or counter-injection direction. The co-injection plasmas studied were H-mode plasmas with associated flat density profiles and confinement times up to 2.5t}^{-1}\text{ITER-89P} [5]. No H-mode transitions were observed in the counter-injection plasmas even with injection powers up to 5 MW. In these counter-injection plasmas, radiation losses steadily increased during the neutral beam injection phase, reaching levels of 70% of the input power towards the end of the auxiliary heating phase.

Fig. 1 is an example of a plasma discharge in this operational regime with 5 MW of co-injection. The onset of injection was at \( t=250 \text{ msec} \). The H-mode transition occurred at 380...
msec, and the good confinement lasted until 490 msec, at which time a Giant ELM occurred and caused a sudden loss of plasma energy. The plasma energy was further degraded by the presence of large amplitude continuous MHD mode activity and subsequent mode locking, as is typical of PBX-M "β-collapse" discharges. Four other types of discharges were used in this study; these had only two beams injected, in either the tangential (R_{tan}=135 cm) or perpendicular (R_{tan}=36 cm) direction. This injection scheme was performed in the co- and counter-directions.

The Spitzer and neoclassical model magnetic field line pitch profiles used for the comparison were calculated by solving the time-dependent poloidal field diffusion equation in the TRANSP code. Inputs to the calculation included a series of 56-point Thomson scattering T_e and n_e profiles obtained at different times from a number of similar discharges, and Z_{eff}(r,t) as measured by the tangential visible bremsstrahlung array. The calculation includes the driven currents (e.g., beam and bootstrap) as well as the ohmic current.

Of considerable importance to the calculation, especially in the relatively short-lived PBX-M plasma discharges, is the initial value (at the start of neutral beam injection) of q_0. The calculations indicate current diffusion times, and consequently a relaxation to a steady-state condition, on the order of several hundred msec. Therefore, the calculated field line pitch profiles, especially those early in the neutral injection phase, are very sensitive to the initial value of q_0 representative of the ohmic phase. In addition, because of this finite current penetration time, differences between the calculated Spitzer and neoclassical profiles do not emerge until later in the neutral injection phase. Early in the neutral injection phase, the current profile is still close to that in the ohmic phase. For these reasons, we will show results only from measurements taken and calculations done for the period late in the neutral injection phase. The initial q_0 is taken to be 0.72; this value was obtained from measurements during the counter-injection experimental run. However, since the ohmic target plasmas were similar during the co- and counter-injection runs, and since sawteeth were observed during the ohmic phase of the co-injection discharges, we have assumed an initial q_0 value of 0.72 in the co-injection case as well.

Two examples of the comparison between the measured and model magnetic field line pitch profiles are shown in Figs. 2 and 3. The first example is taken from a two beam counter-injection case approximately 250 msec into the neutral beam injection phase. Plotted are the MSE measurements (stars), and model pitch profiles assuming the Spitzer model with no bootstrap current (solid line), and the neoclassical model with bootstrap current (dashed line) as functions of major radius. The absolute value of the field line pitch is plotted; actually, B_p, and thus α, reverse sign across the current center (α=0). For these
comparisons, the measured and model pitch profiles were shifted with respect to each other in order to bring the current centers of the MSE measurement ($\alpha_{\text{MSE}}=0$) and the model ($R(T_{e,\text{max}})$) profiles into agreement. The shift required was typically of order 0.5-2.5 cm, which is within the combined position uncertainties of the MSE and Thomson scattering systems. Also note that the uncertainties in the MSE measurement are about twice as large at the edge than near the center. As can be seen in Fig. 2, the measured pitch values on average are considerably lower than those of the neoclassical model, and, in fact, are in reasonably good agreement with the Spitzer model at all major radii, for this case. The standard deviation of the measured points from the Spitzer and neoclassical values are 0.64° and 1.10° respectively. While, admittedly, both values are small, the comparison clearly indicates nearly twice as good agreement with the Spitzer as with the neoclassical model.

The second comparison is from the late phase of a four beam co-injection discharge, and is shown in Fig. 3. In this discharge, the beams were sequentially turned on in 50 msec intervals starting at $t=250$ msec. The comparison was made at $t=450$ msec, 50 msec after the last beam was activated. The H-mode transition for this discharge occurred at $t=370$ msec. The discrepancy between the measured and model values in this case is clearly greater than that in the previous case for either model. As far as the overall fit goes, the standard deviations are 0.75° and 1.15° from the Spitzer and neoclassical models respectively. Here again, the overall agreement with Spitzer is better than that with neoclassical; however, significant departures ($\geq 10°$) from even the Spitzer model are apparent in the middle part of the plasma ($R=175$ cm). Near the axis, the agreement with the models, especially with Spitzer, is better.

As is evident in Fig. 3, the discrepancy between measurement and model is such that the measured field line pitch generally falls below the model value, indicating lower local poloidal fields. This is not necessarily a common trend for cases in which there are discrepancies between the model and the

![Fig. 2) Two-perp counter magnetic field line pitch profiles](image1)

![Fig. 3) 5 MW co magnetic field line pitch profiles](image2)
measurements. In some cases the measured pitch values are lower, but in others they are higher, than the model values. However, if there is a discrepancy near the center of the plasma, it is usually the case that the measured values are lower than the model values, indicating slower current penetration than would be expected from the models.

Summary

In summary, radial profiles of magnetic field line pitch measured on PBX-M have been used to assess the main two models believed to characterize the diffusion of current in tokamak discharges. In many cases, the discharge parameters are such that a clear distinction between the two models, Spitzer and neoclassical, can be made. In these cases, the data agree better with the Spitzer model, indicating an effective particle detrapping mechanism, although there is evidence that current diffusion is anomalous at least in some of the cases, as is evidenced by the discrepancy between the measured and model values. These results suggest caution in assuming the validity of either model in characterizing the current profile in tokamaks.

The PBX-M project is supported by the U.S. Dept. of Energy Contract No. DE-AC02-76-CHO-3073.

References

IGNITION ACHIEVEMENT IN HIGH FIELD TOKAMAKS

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Introduction - The possibility of producing record confinement parameters in ohmic heating regimes by using high magnetic field tokamaks is well known. The merit fusion factor, $nT_{e}\nu_e$ that, according to recent suggestions (Waltz et al., 1990) could be $\propto B_r^3/\alpha^{6/2}$, stresses the importance of the high field line. In this frame the Ignitor experiment conceived and upgraded by Coppi during several years (Coppi, 1977 and Coppi et al., 1990) combines its promising physics prospects with a well assessed engineering feasibility. The reference data here considered ($B_r = 13$ T at $R_0 = 1.3$ m) represent an upgrade of the project with respect to that already studied (Airoldi and Cenacchi, 1991a), so as to assure a plasma current flattop ($I_p = 12$ MA) lasting for 4 s. The fusion prospects of Ignitor are contrasted with those of another high field project: Omitron ($B_r = 18$ T at $R_0 = 1.5$ m and $I_p = 13.5$ MA), supposed to operate with a current flattop of 2 s. All simulations, carried out using a 1 1/2-dimensional transport code, begin at full current, magnetic field and density, and terminate at ignition. Sawtooth oscillations, if not inhibited, control temperature and density profiles in the center of the plasma, where most of the fusion power is produced, thus affecting ignition. This work analyzes the influence on the fusion prospects of the model chosen for the central region.

The models - The transport module of the simulation code solves the diffusion equations of the electron and ion energy densities, primary ion and impurity densities, and toroidal current density across the magnetic surfaces determined by the equilibrium module. All simulations have the magnetic equilibrium configuration relevant to the toroidal current flattop as initial condition. The plasma evolution is followed only until ignition, or the end of flattop, is reached. The anomalous thermal diffusion, derived from the profile consistent drift wave model initially proposed by Tang (1986) and furtherly assessed by Tang and coworkers, is adopted in the version published by Redi and Bateman (1990). The particle transport and the other modelling assumptions are described in our above mentioned paper. Two impurity species, carbon and oxygen, are explicitly treated; their content is such that the effective charge, $<Z_{\text{eff}}>$, results to be $\approx 1.2$. The initial peak temperatures are assumed to be 4 keV.

The following different hypotheses are adopted to study the influence on ignition of the central region modelling:

a) current density profile flattened by enhancing the electrical resistivity in the region defined by $q<0.7$. In the same region the electron and ion thermal diffusivities are increased. This model was used all over our previous simulations.
b) Kadomtsev-like reconnection flattening both temperature and density profiles in the region defined by $q < q_{\text{min}}$ with preset repetition period $\Delta t_r$ (\( \Delta t_r = 0.2, \ldots, 0.45 \)). The choice for $q_{\text{min}}$ spans from 0.8 to 1.

c) Reconnection model as in case b), but assigning the minimum value allowed for the safety factor by the condition $q(0)_{\text{min}} = q_{\text{min}} - \Delta q$. In the cases considered: $0.8 < q_{\text{min}} < 1$ and $\Delta q = 0.2$.

Table I summarizes the design parameters here considered for the two machines and some calculated values, specifically the Murakami density limit evaluated according to: $<n_{\text{Mur}}> = 1.5 B_t/R_o q_{\text{cyl}}$, and the Greenwald limit given by $\bar{n}_{\text{gr}} = l_p/(\pi a^2)$. The $\beta$ limit, commonly referred as Troyon limit, has the expression $\beta_{\text{max}} = g l_p/a B_t$, where $g$ assumes the conservative value of 2.5.

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Results - By adopting model a), we get for each machine an ignition curve in the plane (average density, ignition time) that represents the time required, from the beginning of flattop, to reach ignition vs. the volume averaged density. Both curves, obtained under identical transport assumptions, are plotted in Fig.1. The ignition is defined by the balance between the power delivered to the plasma by the produced $a$ particles, $P_a$, and all plasma losses. The minimum of each curve points out, in the hypothesized scenario, the privileged value for the density ($<n> = 5.4 \times 10^{14} \text{cm}^{-3}$ for Ignitor and $<n> = 4.8 \times 10^{14} \text{cm}^{-3}$ for Omitron). From previous simulations (Airoldi and Cenacchi, 1991b) it results that the trend of the ignition curve, when degrading hypotheses are assumed, is a shrinking in the density range together with a shifting to higher values of the ignition time. For this reason it is interesting to consider the ignition time measured by the flattop time for each machine, as in Fig.2. Note that for Ignitor the ignition times, albeit higher than those for Omitron, are less than 50% of the flattop time. In all cases the energy replacement time, defined by $\tau_r = W_{\text{tot}}( \equiv W_e + W_i)/(P_a + P_{\alpha} - \partial W/\partial t)$, lies between the values obtained from the scaling expressions known as Goldston and Kaye-All-Complex (Stotler and
The explicit values for $\tau_F$, relevant to the cases in Fig. 1, span from $0.40s$ to $0.46s$ for Ignitor and from $0.36s$ to $0.42s$ for Omitron.

By choosing the best density values found, calculations were performed by model b). The results are summarized in Fig. 3, where the maximum $Q$ reached ($Q = P_{\text{in}}/P_\Omega$) is plotted against the repetition period. The points for which $Q > 10$ correspond to cases of ignition attained. The better performance of Ignitor when the reconnection region is smaller (label 1 in Fig. 3) is due to the delay in the onset of sawtooth oscillations. In fact the initial equilibrium condition exhibits $q(0) = 0.93$, while Omitron has the initial $q(0)$ as low as $0.75$, and in any case the first reconnection is switched on when $q(0) = 0.7$.

Results obtained using model c) are not substantially different from those given by model b). In fact they confirm that ignition may be attained provided the repetition period is greater than $0.4s$ and the reconnection region does not exceed $40\%$ of the plasma poloidal cross-section. Note that the promising outcomes of model a) depend on the low $q_{\text{min}}$ assumption.

The results confirm that a large width of the mixing region is critical for the performance of both machines. As these instabilities can prevent ignition even in a very high magnetic field machine as Omitron, the requirement of suppressing or, at least, delaying the onset of sawtooth oscillations becomes mandatory. To solve this problem different authors already suggested pellet injection and localized plasma heating (Bateman, 1986) or a careful planning of the current and density rise (Coppi et al., 1990). On the other hand, the possibility of obtaining
long sawtooth-free periods is confirmed by experimental evidence on JET (JET Team, 1990).

Fig. 3 - Maximum Q reached vs $\Delta t$. Full lines refer to Ignitor and dashed lines to Omitron. The reconnection region extends up to $q = 0.8$ (label 1), $q = 0.9$ (label 2) and $q = 1$ (label 3).

The engineering feasibility of both machines is not discussed here, but in our opinion the greater constructive problems of Omitron are not compensated by much better ignition prospects.

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MEASUREMENT OF TOROIDAL AND POLOIDAL PLASMA ROTATION IN TCA

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Introduction

The stimulus for measuring the toroidal and poloidal rotation velocities on TCA (R=0.615m, a=0.18m, Ip=130kA), was the large density rise observed with high power Alfvén Wave Heating (AWH) [1]. From the measured velocities, we can deduce the radial electric field in the plasma, a change in which could be responsible for the density rise. Quartz fibres with quartz windows allowed observations of the CV (227nm) and CIII (230nm) spectral lines in second order which are strong in TCA, and are emitted from r=-2a/3 and in the plasma edge respectively. A 1m Czerny-Turner visible spectrometer with a multichannel OMA detector equipped with a 24001/mm grating recorded a spectral line profile every 2ms through the 200ms discharge. On TCA, these spectral features are completely unpolluted by other spectral lines which simplifies data analysis and improves the data quality. For the toroidal measurements, opposing tangential windows permitted measurements with and against the plasma current, and a long slot window on top of TCA permitted poloidal measurements along near vertical chords on both the low and high field side of the plasma. When access via a quartz window was impossible, helium was injected into the discharge and the HeII (468nm) spectral line was observed in first order.

Results

Fig 1 shows the evolution of the CV toroidal plasma velocity through a discharge with a large density rise (hard puff), together with the position of the spectral feature on the OMA and its HWHM both in pixels (1pix=25μm).

Fig 1 Evolution of the (a) toroidal rotation and (b) ion temperature through a discharge
To determine the absolute velocity, toroidally the plasma was viewed in and against the plasma current direction and for poloidal measurements the plasma centre was viewed to define the zero velocity reference. The spectrometer was also regularly calibrated against a spectral lamp and a correction applied based on the spectrometer temperature, measured by an array of thermistors. The measured scatter from observing a spectral lamp during a plasma discharge was ~0.5km/s. The scatter of ~1.5km/s in the CV data shown in Fig 1 is mostly due to larger Doppler broadening of the plasma spectral line. The crosses show the fitted position of each 2ms integration and the solid line is a weighted smooth across three points which has an estimated statistical error of ~0.5km/s. The deduced ion temperature has been corrected for the instrumental function. Comparing this with a peak value measured by neutral particle analysis of ~330eV the emission region is ~2a/3.

Before we approached the effects of AWH we studied rotation with normal and hard gas puffs in which higher final densities can be achieved without plasma disruption. Fig 2 shows the evolution of the toroidal rotation of HeII and the plasma density.

Fig 2. The evolution of the toroidal rotation velocity, in the direction of plasma current, together with the plasma density through a discharge. The rotation velocity ceases to follow the plasma density at the end of the discharge when the plasma current starts to fall.

Fig 3. Traces of toroidal rotation for CIII and CV, for discharges with a lower density ramp (β and δ) and over double the same ramp rate (α and γ), where the velocity locus saturates and tends to zero. This effect is more pronounced at the plasma edge (CIII). The ion temperatures show that the discharges are otherwise similar.
The toroidal rotation closely follows the plasma density. In discharges where the density reaches a plateau, the toroidal velocity also stabilises. This dependence is less direct when a hard gas puff is used.

Fig 3 plots the measured toroidal velocity against average plasma electron density for CV and CIII for discharges with different density ramps. For the plasma edge (CIII), two discharges with different ramp rates are shown. For the higher ramp rate, the toroidal velocity ceases to rise and returns to zero. For CV, the same effect is present, although we note, as in all the velocity measurements, that the CV velocity reacts less and more slowly to gas puffing than at the edge. The temperatures are similar in both cases and it is clearly the ramp rate, and not the value of the density, that is the cause.

Fig 4 shows the poloidal velocity measurements with AWH and a hard gas puff. Positive velocities correspond to the electron diamagnetic drift direction. To demonstrate the shot to shot differences, two discharges with AWH are shown together with a hard gas puff, where the range of other observations is indicated by the shaded region. There is no clear difference which distinguishes the AWH traces, and the width of the shaded region is dominated by changes in the density ramp, and the cooling effect of the stronger gas puffs, rather than AWH effects.

![Fig 4](image-url)
Discussion

The radial electric field may be deduced from the force balance equation on an ion:

\[
E_r = \frac{\nabla \cdot p_1}{n_i Z_i e} = v_\phi B_\theta - v_\theta B_\phi
\]

The contribution from the gradient is estimated to be less than 10% in all our observations, and much less in the plasma edge. Operation time restrictions precluded a measurement of toroidal velocity with high power AWH, but since its contribution to the deduced electric field is ~10%, the error introduced is small. With AWH, no significant increase in the poloidal velocity, and thus radial electric field, was measured that could explain the large associated density rise. Good agreement between the values measured from HeII and CIII (-1kV/m at low density rising to -6kV/m) was observed. The value deduced from CV was the same at low density and ~1.5x lower at high density. This is considerably lower than that observed on DIII-D [2] which is a larger machine.

Conclusion

With optimal observation geometry we have measured both the toroidal and poloidal rotation velocities in the edge and in the bulk of the TCA plasma. Regular calibration and correction for variations in the spectrometer temperature permitted a measurement with an error of ~0.5km/s which is an order of magnitude smaller than the range of measured velocities.

In general, changes in the velocities are observed to be stronger and faster in the plasma edge than in the plasma bulk. With increasing density, the toroidal velocity is observed to change sign and follow the plasma density, while the poloidal velocity increases. These two effects lead to an increase in the absolute value of the radial electric field. With very strong gas puffing, the toroidal velocity is observed to again reverse and tend to zero, an effect which is stronger as the gradient of the density ramp is increased. Comparison between gas puffing and high power AWH does not show a significant difference in the radial electric field that could be responsible for the large associated density rise, which still remains unexplained.

Acknowledgments

We wish to thank the whole TCA team for the technical support. This work was partially supported by Le Fonds National Suisse de la Recherche Scientifique.

References

ENHANCED TOROIDAL ROTATION IN HOT-ION MODE
WITH NEARLY-BALANCED NEUTRAL BEAM INJECTION IN JT-60

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1. Introduction

So far, a number of plasma rotation studies for tokamaks have been carried out with strong toroidal momentum input from tangential neutral beam injection (NBI) to investigate the relation between plasma rotation and viscosity processes [1]. Otherwise, it is also important to study the toroidal rotation characteristics with much less torque input as compared to the heating power to see how the plasma rotation is coupled with the heat/particle transport. In JT-60, nearly-balanced perpendicular NBI equipped with eight co- and six counter-injection units in the direction of ±15° allows us to study the rotation characteristics under low beam torque input. In this paper, the toroidal rotation characteristics are investigated for hot-ion mode plasmas with enhanced confinement obtained at high q regime in JT-60 [2], and compared with L-mode plasmas produced by co-only and nearly-balanced injections. All these plasmas were operated with a lower single-null divertor configuration; a major radius of R=2.9 cm, a minor radius of a=0.65 cm, an ellipticity of \( \kappa = 1.3 \). Neutral hydrogen beams up to ~21 MW with a beam energy of 65 keV were injected for 3-5 s. Profiles of the ion temperature \( (T_i) \) and toroidal angular velocity \( (\Omega_\phi) \) were measured by eight-channel charge-exchange-recombination-spectrometers (CXRS) with a time resolution of 50 ms.

2. Ion Temperature and Angular Velocity Profiles

Figure 1 shows \( T_i, T_e, n_e \) and \( \Omega_\phi = V_\phi / R \) (toroidal angular velocity) values as a function of flux co-ordinate for a typical hot-ion mode plasma: \( I_p = 0.57 \text{ MA}, B_t = 4.5 \text{ T}, q_{cy} = 5a^2 B_t/(R_{pl}(\text{MA}))(1+\kappa^2)/2 = 8.0, \beta_p = 2.4, P_{NB} = 21 \text{ MW} \) (\( \text{H}^0 \rightarrow \text{H}^+ \)), the power-balanced ratio of \( \Gamma = (P_{co}+P_{cr})/(P_{co}+P_{cr}) = 0.19 \) (nearly balanced injection), and \( n_e(0) = 3.0 \) (<— denotes the volume average). In association with the enhanced toroidal rotation near the axis, the angular velocity profile is found to be remarkably peaked just inside \( r/a \approx 0.3 \) or \( r/a \approx 0.4 \) with the peaked profiles of \( T_i \) and \( n_e \).

In Fig.2, a strong correlation between the ion temperature and the angular velocity in the hot-ion mode is manifested by comparison of \( T_i \) and \( \Omega_\phi \) as a function of normalized flux co-ordinate with the L-mode: the same shot as in Fig.1 for nearly-balanced injection in hot-ion mode; \( P_{NB} = 21 \text{ MW}, q_{cy} = 4.2 \) and \( B_t = 3.0 \text{ T} \) for nearly-balanced injection in L-mode; \( P_{NB} = 12 \text{ MW}, q_{cy} = 4.1 \) and \( B_t = 4.0 \text{ T} \) for co-only injection in L-mode. For the hot-ion mode and the L-mode with co-injection, the observed \( \Omega_\phi \) profile was compared with the absorbed torque density profile \( T_\phi(r) \) calculated by an orbit-following Monte Carlo code. It is found that the \( \Omega_\phi \) profile for the hot-ion mode is more peaked than for the torque density profile (see Fig.3), while the \( \Omega_\phi \) profile for the L-mode is broader than for the torque density profile. These observations imply the presence of some effect except for a direct response of the torque input on the rotation profile through viscosity process.

In the hot-ion mode, the central rotation velocity \( V_\phi(0) \) is correlated with the central ion temperature \( T_i(0) \) as previously reported [2]. The \( V_\phi(0) \) values also represent an offset linear correlation with the degree of ion temperature gradient \( (\Delta T_i/\Delta r) \) in the region where the \( \Omega_\phi \)
profile is peaked, as shown in Fig. 4 for the conditions of $I_p=0.48 \text{ MA}, 0.57 \text{ MA}, B_t=4.5 \text{ T}, a=0.65 \text{ m}, R_p=2.9 \text{ m} \text{ PNB}=19-21 \text{ MW} \text{ (nearly-balanced injection)}$ and $\eta_{\text{cy}}=8-10$. This suggests a correlation between ion thermal and angular momentum confinement with its improvement. In terms of causality, increase in the toroidal rotation velocity was observed to precede increase in the ion temperature in some hot-ion mode discharges.

3. Radial Electric Field and Angular Momentum Confinement Time

From the radial force balance equation coupled with neoclassical theory, the radial electric field is expressed as $E_r=r B_\phi V_\phi(q R)+(\partial T_i/\partial r)(\eta_1+1-K(v_*)i)/(eZ_i)$, where $eZ_i$ is the charge, $E_r$ is the radial electric field, $\eta_1=d(\log T_i)/d(\log n_i)$ and $K(v_*)$ is the neoclassical coefficients depending on the ion collisionality parameter $v_*$ [3]. The potential difference to the plasma edge, $\Delta \Phi(r)=\Phi(r)-\Phi(a)$, can be calculated from the $E_r$. As the calculated $E_r$ and $\Delta \Phi(r)$ profiles for the hot-ion mode of Fig. 1 are shown in Fig. 5, the large negative electric fields are generated in the core plasma in association with the deeply negative potential. From this evaluation, $-e\Delta \Phi(0)/T_i(0)=0.7-1.2$ for hot-ion mode plasmas, so that it can be said that the negative potential at center is formed to be of the order of the $T_i(0)$ value. Thus, the large negative electric fields are found to exist in the rotational equilibrium, while the plasma substantially rotates in the co-direction near the axis.

The toroidal angular momentum confinement time for the plasma region inside $r'=r$ is approximately constructed from $\tau_\phi(r)=<m_i q R^2 \Omega_\phi(r')>/<T_\phi(r')>$, where $<->$ denotes the volume integral from $r'=0$ to $r'=r$. Figure 6 shows the $\tau_\phi$ profiles as a function of normalized flux co-ordinate for the hot-ion mode (balanced injection), L-mode (balanced injection) and L-mode (co-injection) plasmas discussed in Fig. 2. With respect to the profile shapes, the $\tau_\phi$ value for the hot-ion mode is greatly enhanced in the center, but decreases with the minor radius, though the slight radial increase in $\tau_\phi(r)$ is seen for L-mode plasmas. The global toroidal angular momentum confinement times $\tau_\phi(a)$ are compared with the theoretical values from the gyro-viscosity theory [4], $\tau_\phi^{\text{gyro}}=2R_0 Z_{\text{eff}} B_0 /<T_\phi>$, as shown below:

<table>
<thead>
<tr>
<th>Mode</th>
<th>$\tau_\phi(a)$ [ms]</th>
<th>$\tau_\phi^{\text{gyro}}$ [ms]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Hot-ion mode (balanced)</td>
<td>20</td>
<td>76</td>
</tr>
<tr>
<td>L-mode (balanced)</td>
<td>10</td>
<td>68</td>
</tr>
<tr>
<td>L-mode (co)</td>
<td>75</td>
<td>100</td>
</tr>
</tbody>
</table>

In the balanced injection cases, the experimental values are found to be much smaller than the theoretical values.

4. Comparison with a Modelling Calculation

Recently, a toroidal rotation study with perpendicular NBI was reported in the JIPPTHIU tokamak [5]. They observed negative electric fields for both co- and counter-injections and proposed a model of the ambipolar electric field induced by ion losses to explain the observations. The expression of the ambipolar radial electric field was formulated as $E_r^{\text{am}}=\alpha_i (\partial T_i/\partial r)/e-(2C_v/a^2)r T_i/e$, assuming $D_e=D_i$ (ion and electron anomalous transport coefficients), where $\alpha_i$ is a neoclassical numerical coefficient around unity and $C_v$ is the peaking parameter. To see whether this model can predict the hot-ion mode profiles in JT-60, the calculated $E_r^{\text{am}}$ profiles for $C_v=0$ and 1.0 in the hot-ion mode are compared with the $E_r$ obtained in Fig. 5 as shown in Fig. 7(a). Furthermore, to see whether the observed rotation profile can be properly reconstructed from the $E_r^{\text{am}}$, the toroidal rotation velocity profile balanced with the $E_r^{\text{am}}$ is calculated from $V_\phi^{\text{am}}=-q R/(eB_0)\{\xi_1 (\partial T_i/\partial r)/(r+2C_v T_i/a^2)\}$, where $\xi_1=(\eta_1+1-K(v_*)i)/Z_i-\alpha_i$. Figure 7(b) shows the comparison of the angular velocity profile between the observation and the modelling calculation for $C_v=0$ and 1.0. These results suggest qualitative agreement in terms of negative electric field and peaked rotation profiles observed in the hot-ion mode. However, there is a discrepancy in the central region that the model
might predict the decrease in the $V_\phi$ with the $T_i$ in contrast to the experimental results discussed in the previous section. Thus, while it is possible for this model to interpret an important part of the mechanisms in the observed enhanced toroidal rotation, a comprehensive model including angular momentum transport would be required in JT-60.

5. Discussion and Conclusions

The enhanced toroidal rotation is observed in the JT-60 hot-ion mode plasma. The observed features are characterized by the approximate linear correlation between $V_\phi(0)$ and $\Delta T_i/\Delta r$ or $T_i(0)$ and the deeply negative potential at center of the order of the central ion temperature. The formation of the enhanced toroidal rotation localized near the center is considered to arise from two effects: 1) increase of co-toroidal rotation due to improvement of angular momentum confinement, 2) increase of counter-toroidal rotation due to ambipolar electric fields induced by ion losses. These different effects appear to affect the rotational equilibrium in the hot-ion mode plasma. As the ion temperature and the density profiles are peaked, the observed toroidal rotation profile can not directly reflect a momentum diffusion as predicted in the theory. Thus, for weak torque input experiments, the toroidal rotation characteristics should be discussed taking into account the effects of the radial electric fields induced by different ion/electron transport processes.

References


Fig.1 Profiles in a typical hot-ion mode plasma

Fig.2 Comparison of $T_i$ and $\Omega_\phi$ profiles among hot-ion mode and L-mode plasmas
Fig. 3 Absorbed torque density and angular velocity profiles in the hot-ion mode

Fig. 4 Correlation between central rotation velocity and the ion temperature gradient in hot-ion modes

Fig. 5 Radial electric field and potential difference in the hot-ion mode

Fig. 6 Comparison of $\tau_\phi$ profiles among hot-ion mode and L-mode plasmas

Fig. 7(a) Comparison of the $E_r$ profile between experiment (solid line) and modelling calculations (dotted lines)

Fig. 7(b) Comparison of the $\Omega_\phi$ profile between experiment (solid line) and modelling calculations (dotted lines)
The initial experiments on the JT-60 Upgrade tokamak

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Abstract
The physics design review of the upgraded JT-60 tokamak, which can accommodate a large D-shaped plasma with \( I_p = 6 \, \text{MA} \) and the volume of 100 \( \text{m}^3 \), is described in this paper. The aspect ratio ranges from 3.4 to 3.8 for single-null diverted plasmas with the ellipticity of 1.5 - 1.8. According to the ITER-89 power-law scaling, which predicts the fusion gain proportional to \((R/a)^{1.4}\), substantial contribution of the upgraded JT-60 is anticipated for the confinement study among other divertor tokamaks such as JET and DIII-D (\( R/a \sim 2.5 \)). The upgrading modification was just completed, and the initial experiment is starting in March.

Introduction
Confinement improvement and steady-state operation are the major subjects in tokamak research. Experiments on JT-60, which was designed for hydrogen operation with the maximum plasma current of 2.7 MA, focused on these studies. Major results obtained were: (1) fusion product \( n_e \cdot \tau_E \cdot \tau_{i_e} \) of \( 1.3 \times 10^{20} \, \text{m}^{-3} \cdot \text{s} \cdot \text{keV} \), (2) driven current of 2 MA by lower hybrid current drive (LHCD), and (3) current drive efficiency \( \eta_{cd} = \frac{n_e \cdot I_P \cdot \eta_R}{PLH} \) of \( 3.4 \times 10^{19} \, \text{m}^{-2} \cdot \text{A} / \text{W} \). However, improvement of plasma performance was required for further extension of the fusion research.

Since the recent experiments indicate that both energy confinement and beta improve by increasing the plasma current, the maximum plasma current of 6 MA for divertor configuration and 6.5 MA for limiter configuration were aimed at, which are twice those of the original JT-60. These current limits are given by the condition of \( q_{eff} > 2 \). Single null open divertor configurations, together with deuterium operations, were adopted, since these conditions seem to be the most suitable for high quality H-modes. Torus input power of ion cyclotron resonance heating is also increased for the central heating of beam heated plasmas. These modifications allow substantial improvement of the plasma performances, which enable us to investigate characteristics of plasma close to the break-even condition.

The beta limit does not restrict the operation regime, since the strong toroidal magnetic field of JT-60 Upgrade allows the maximum possible pressure \( nT_{\text{max}} \) of \( 7 - 9 \times 10^{20} \, \text{keV} \cdot \text{m}^{-3} \) for the TROYON parameter \( g = 2.5 \), which is considered to be a typical limit for favorable plasma confinement. This maximum pressure will be the highest among large tokamaks.
In the modification process, the original poloidal field coils, its support structures and the vacuum vessel were superseded by new ones to allow single-null open divertor plasmas with plasma current of up to 6 MA. The existing toroidal field coils and their support fixtures are used after reinforcement. The existing high power heating system and the power supply system are used after minor modification. Neutron shields for deuterium operation are being prepared [1] to start D_{beam} \rightarrow D_{2} operation in July 1991. The upgrading modification was just completed, and the initial experiment is starting in late March.

**Operation regime of the JT-60 Upgrade tokamak**

Poloidal field coils are designed to produce 3 types of divertor configurations as illustrated in Fig. 1 [2]. In the standard mode (case A in Fig. 1), the divertor plasma with elongation of 1.5 and an aspect ratio of 3.4 can be produced. In the elongated mode (case B), the divertor plasma with elongation of 1.76 and an aspect ratio of 3.8 can be produced. Case C permits continuous control of the elongation and the triangularity. However, this option requires an additional power supply. The objective of these options is to study the dependence of diverted plasma confinement on the aspect ratio A, major radius R_{p}, minor radius a_{p} and elongation \kappa. Goldston scaling claims that

$$\tau_{E} \propto \frac{R_{p}}{a_{p}}^{0.75} \frac{a_{p}}{R_{p}}^{0.37} \frac{1}{0.5^{0.5}}$$

(1)

On the other hand, ITER-89 power-law predicts

$$\tau_{E} \propto \frac{R_{p}}{a_{p}}^{1.2} \frac{a_{p}}{R_{p}}^{0.3} \frac{1}{0.5^{0.5}}$$

(2)

JT-60 Upgrade will provide the database of a large tokamak with a high aspect ratio, since the aspect ratio of DIII-D and JET are as small as 2.7 and 2.4, respectively. Figure 2 shows the predicted fusion product based on the recently re-established ITER-89 power-law confinement scaling. Since the scaling predicts the fusion gain proportional to \( (R/a)^{-1.4} \), relatively high aspect regime of JT-60 Upgrade (\( R/a = 3.4 - 3.8 \)) will substantially contribute to the confinement database for the next step device.

**Fig. 1** Typical equilibrium configurations of the JT-60 Upgrade plasma for I_{p} = 6 MA, q_{eff} = 2.2: (A) Standard mode with \( \kappa = 1.52 \) and \( R/a = 3.44 \), (B) Elongated mode with \( \kappa = 1.76 \) and \( R/a = 3.83 \), and (C) Continuous mode with \( \kappa = 1.37 \) and \( R/a = 3.09 \).
To obtain stable discharges with high elongation, it is important to stabilize the vertical instability. This instability of elongated plasmas can be analyzed with an average n-index, which is defined by

\[
\langle n \rangle = \frac{\int j_\phi \frac{\partial B_z}{\partial R} dS}{\int j_\phi dS}
\]

(3)

where \( R_{axis} \) is the major radius of plasma magnetic axis, \( B_z \) is the externally applied vertical magnetic field and \( j_\phi \) is the toroidal current density. The average n-index of the elongated JT-60 Upgrade reaches up to -1.5. The passive index \( n_s \) of coils and a vacuum vessel must be larger than this average n-index for stable operations of plasma. In JT-60 Upgrade, the vacuum vessel is not effective in increasing the value of \( n_s \) due to the short time constant of dipole current component of the vessel. Therefore, the horizontal field coil were distributed around the vacuum vessel and conductors of the vertical field coil are designed to serve as passive stabilizers in the elongated mode to yield the \( n_s \) value of 1.7. The vertical position controllability is also improved by changing the horizontal field coil power supply from the original 12 phase convertor to a 24 phase convertor and also by accelerating the plasma feedback system.

Disruption control, together with particle and impurity control, is crucial to obtain steady-state discharges. Four sets of sector coils named DCW were provided both at the inboard side and at the outboard side of the vacuum vessel. Although the toroidal length of two coils at the outboard side is short, it is used to produce the non-axisymmetric magnetic field with the dominant toroidal mode number of 2. This field can produce \( m = 3 / n = 2 \) small magnetic islands at \( q = 1.5 \) surface as shown in Fig. 3. Because of the ergodization of \( m = 2 / n = 1 \) magnetic islands by the small-scale disruption, the temperature profile will be changed and studies of disruption control will be possible as demonstrated in JIPPT-T-II U. On the other hand, successful particle control during H-mode was shown in JFT-2M by using
externally applied resonant magnetic field of high-m modes [3]. Although the poloidal mode number produced by DCW is low, an attempt will be made to control particle and impurity concentration during H-mode and to control edge plasma phenomena, such as ELMs.

**Research prospect of the JT-60 Upgrade experiment**

The primary objective of JT-60 Upgrade in early phase of its experiment is the confinement improvement under favorable impurity control. The other primary objective from 1994 will be the non-inductive current drive by ECH assisted NBCD with 500 keV negative NB in combination of bootstrap current and LHCD. The ability to generate $\alpha$-particles due to $D + ^3He \rightarrow ^4He$ (3.6 MeV) + p (14.7 MeV) reaction with 500 keV D-beam at which the fusion cross section peaks will provide a good opportunity to study $\alpha$-particle behavior prior to the D-T operation. A real fusion reaction of 1 - 1.5 MW range is expected. This will provide more confident confinement properties of $\alpha$-particles compared with the present experiment producing fusion reaction of ~100 kW using ICRF. The energetic single-particle confinement concerns the effect of non-axisymmetry of the magnetic field of the reactor grade plasmas. The stochastic ripple diffusion effect may cause rapid energetic ion loss. The edge of the JT-60 Upgrade plasma will have ~2% ripple. Detailed orbit-following Monte-Carlo calculations of this process shows ~30% of 120 keV fast ion loss for perpendicular injection in high density discharge, compared with less than 10% loss for tangential injection. The experiment will provide an important data base to evaluate the stochastic ripple diffusion process. The other important issue of burning-plasma physics concerns the collective stability of the $\alpha$-particle population and its effect on confinement. The negative NB injection experiment may be able to simulate the relevant physics before attempting a burning-plasma experiment. The expected parameter range will be $V \sim 1.4$ VAlfven and $\beta_0 \sim 3\%$ at $B_T = 3$ T for 500 keV, 10 MW H-beam, which seems to be well in the TAE unstable regime.

![Magnetic islands with m = 3 / n = 2 produced by the DCW coil.](image)

**References**


COMPARISON OF DIMENSIONALLY SIMILAR DISCHARGES WITH SIMILAR HEAT DEPOSITION PROFILES*

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Dimensionally similar scaling is a technique to scale existing tokamak discharges to larger size, \(a\), or larger magnetic field, \(B\), by fixing dimensionless parameters such as the plasma \(\beta\) and collisionality as well as safety factor \(q\), elongation, and aspect ratio while varying the normalized gyroradius \(\rho_* = \rho/a\). Diffusion mechanisms can be classified by their dependence on \(\rho_*\). The confinement time scales as \(\tau \propto \Omega^{-1} \rho_*^{-3} \propto B_{eq}^{5/2}\) for gyroBohm-like mechanisms or \(\tau \propto \Omega^{-1} \rho_*^{-2} \propto B_{eq}^{1/2} a_{eq}^{5/3}\) for Bohm-like mechanisms. Establishing that either mechanism dominates transport allows one to predict performance for next generation devices with reasonable confidence without knowledge of the complex and likely coupled transport mechanisms governing confinement. Statistical studies of global confinement time \(\tau^0\) have consistently shown that the standard \(L\)-mode scaling is close to or worse than Bohm-like whereas the majority of theories suggest a gyroBohm-like local diffusion process.

During the past year experiments to resolve this difference were performed on DIII-D^3 and subsequently on TFTR^4. In each case fixed size and shape \(L\)-mode discharges were compared at low and high \(B\) field while adjusting the density \(n_e\) and power \(P\) to maintain the similarity conditions \((I \propto B, n_e \propto B^{4/3}, T \propto B^{2/3})\). In all cases, the global confinement time was nearly Bohm-like or worse \((\tau \propto B^0)\). However, standard transport analysis with the simple diffusion equation \(-n_e \partial T/\partial \tau = P_{tr}/S\) found a gyroBohm-like effective diffusivity \(\chi_{eff} \propto B^{-1}\) in the DIII-D case and in a high density TFTR scan. In these cases, the apparent discrepancy between the \(\tau\) scaling and \(\chi_{eff}\) scaling was attributed to the markedly poorer penetration of the neutral beam heating profile at the higher \(B\) (and density). This was thought to be a direct consequence of the simple heat diffusion equation. In apparent contradiction, a low density TFTR scan with much less change in the beam penetration appeared to be more consistent with a Bohm-like scaling \(\chi \propto B^{-1/3}\).

Based on the expected result of a simple diffusion model that \(\tau \propto a^2/\chi_{eff}\), we speculated that \(B\)-scaled dimensionally similar discharges with similar heating profiles should have global confinement time \(\tau \propto B\) if gyroBohm-like diffusion prevails. This speculation has been tested recently in a set of DIII-D experiments designed to control the heating profile by vertically shifting the plasma off-axis away from the centerline of the neutral beam heating. This allowed shifted low field, low density discharges to have heating profiles similar to the high field, high density discharges.

The controlled heating profile experiments were performed in \(L\)-mode discharges limited on graphite tiles on the inside wall of the vacuum vessel with deuterium neutral beams injected into deuterium plasmas. The plasma elongation \(\kappa = 1.5\) and safety factor at 95% of the edge poloidal flux \(q_{95} = 3.4\) were held fixed. A vertically asymmetric discharge shifted up 37 cm at 1 T, 0.7 MA, \(3.6 \times 10^{19}\) m\(^{-3}\) with a total power of 3.9 MW was compared with a vertically symmetric discharge at 2 T, 1.4 MA, and \(9 \times 10^{19}\) m\(^{-3}\). Equilibrium flux plots of the two discharges are shown in Fig. 1. The vertical half width of the neutral beam at the 1/e power point is 26 cm at \(B_0\). A total power of 15.4 MW was required to achieve the dimensionally similar plasma parameters required. As in previous experiments the total power required to achieve the similarity condition was closer to that predicted by Bohm scaling \(P \propto B^{3/2}\) than by gyroBohm scaling \(P \propto B\). The global thermal confinement time again remained about constant at 55 ms for 1 T and 61 ms for 2 T.

* This work was sponsored by the U.S. Department of Energy under Contract No. DE-AC03-89ER51114.
The dimensionless parameters $\beta$ and collisionality were held reasonably constant. The temperature and density profiles, normalized to the appropriate scaling with $B$ to achieve constant dimensionless parameters, are shown in Fig. 2. The electron and ion temperatures were equilibrated in these discharges so a single temperature profile was determined by combining the electron and ion measurements. Ion temperature measurements are made along viewing chords through the midplane of the vessel. As a means of obtaining ion temperature measurements closer to the center of the discharge, the shifted plasma was quickly recentered in the vessel on a time scale ($\leq 20\text{ms}$) short compared to a fast ion thermalization time ($\sim 50\text{ms}$). The $Z_{\text{eff}}$ profile was similar in both $1\text{T}$ and $2\text{T}$ discharges varying from 1.5 on axis to near 3 at the plasma edge. The fraction of radiated power $\sim 45\%$ was similar for both discharges.

Fig. 1. Equilibrium flux plots of centered and shifted discharges.

Fig. 2. (a) Temperature and (b) density profiles normalized to the appropriate $B$ scaling for dimensional similarity.
The normalized profile of radially integrated transport power $P_{tr}$ within a normalized radius is reasonably well matched as shown in Fig. 3(a). In contrast to the previous DIII-D experiment with differing heating profiles, the effective thermal diffusivity was not consistent with a gyroBohm model $\chi \propto B^{-1}$. The single fluid effective diffusivity near the half radius [Fig. 3(b)] remains unchanged and has a nearly Bohm-like scaling. This appears to remove the remaining experimental discrepancy with the low density TFTR scan. For comparison, Fig. 3(c) and 3(d) show $P_{tr}$ and effective diffusivity profiles for the previous DIII-D experiment. The values of $\chi_{eff}$ differ by about a factor 2 over most of the profile where $P_{tr}$ differs by a corresponding factor of about 1/2. This is clearly not observed when the heat deposition is more similar. Thus it now appears that the effective diffusivity varies with the heat deposition profile, i.e., simple diffusion does not describe the transport.

To further study the impact of heating profiles on transport and global confinement, off-axis and central heating was performed during the same discharge. The parameters were the same as those for the shifted 1T discharge. After establishing the shifted discharge, the plasma was then recentered and allowed to reach equilibrium again. The density profile remained unchanged and the core temperature increased only slightly. The large variation of the heat deposition profile had little impact on the global confinement time indicating a simple diffusion model is not adequate. This result is entirely consistent with results obtained earlier on ASDEX. The experiments differ only in that the new DIII-D experiment has a much larger variation in heat deposition comparing a hollow power density with no deposition on axis in the shifted case with a peaked power density in the centered case. Further there is no change in the working gas. Again observe in Fig. 3(e) and 3(f) that while the confinement time is largely unchanged by variations in the heating profile, the effective diffusivity can be changed by a factor of 2 near the half radius and even more near the center. Furthermore there is evidence in recent DIII-D off-axis ECH heating experiments that a centrally peaked temperature profile can be maintained with little or no net power flow. This could indicate either the presence of a “heat pinch” or strong marginality to a critical temperature gradient.

In conclusion our initial dimensionally similar discharge B-scaling experiments on DIII-D showed a Bohm-like or worse scaling of global confinement time ($\tau \propto B^0$) but gyroBohm-like scaling of the effective local diffusivity ($\chi_{eff} \propto B^{-1}$). We speculated that the apparent paradox was due to the fact that the heating profiles were not held constant and that a gyroBohm-like confinement time ($\tau \propto B$) could be obtained with similar heating profiles. Our recent experiments refute this speculation and do not support the conclusion that the effective diffusivity is consistent with a gyroBohm-like scaling. Both sets of experiments indicate that while global confinement time is insensitive to changes in the heating profile, the effective diffusivity is very sensitive to such changes. These experiments suggest a breakdown in the notion of a simple local diffusion process and pose an important challenge to theoretical models which exhibit gyroBohm-like scaling.

Fig. 3. Integrated power and effective diffusivity profiles. (a) and (b): comparison of dimensionally similar (DS) discharges with similar heating profiles, (c) and (d): DS discharges with dissimilar heating profiles, and (e) and (f): same discharge, 1 T, with dissimilar heating profiles.
THE SIMULATION OF ENERGY AND PARTICLE TRANSPORT, 
HEAT AND DENSITY PULSE PROPAGATION AND H-MODE 
CONFINEMENT IN JET AND A REACTOR

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Introduction

The simulation of tokamak plasmas is made difficult by a strong and fundamental coupling between energy and particle transport. In particular, it is practically impossible to study energy transport without addressing particle transport. It is necessary to describe correctly not only convective energy losses but also the evolution of impurities and the effect of $Z_{\text{eff}}$. An overall, coherent and quantitative transport model is therefore needed, with the interactions between all elements included. The aim of this paper is to compare such a model with JET results and apply this model to a reactor.

The underlying phenomenon responsible for anomalous transport

Experimental observations support a model for anomalous transport based on a single phenomenon and MHD limits. One such single phenomenon is turbulence in the magnetic field topology which occurs above a certain threshold [1]. In addition, the resilience of the electron temperature to additional heating suggests that electrons are primarily responsible for confinement degradation. The critical electron temperature gradient model exhibits these features and may be formulated as follows: above a threshold in the electron temperature gradient, the anomalous electron conductive heat flux, $Q_e$, is defined [2]:

$$Q_e = -n_e \chi_e \nabla T_e = -n_e \chi_{an,e} \left( \nabla T_e - \left( \nabla T_e \right)_{cr} \right) H(\nabla q)$$

$$Q_i = -n_i \chi_i \nabla T_i$$

$$\chi_i = 2 \chi_e \frac{Z_i}{\sqrt{1+Z_{\text{eff}}}} \sqrt{\frac{T_e}{T_i}}$$

Semi-empirical expressions for $\chi_{an,e}$ and $(\nabla T_e)_{cr}$ have been obtained on the basis of laws of similarity and a theoretical model for the turbulence in the magnetic topology [3,4].

Without changing the form of the above expressions, or the numerical coefficients used, the global plasma properties, temperature profiles and their temporal evolution have been well-simulated for different plasma conditions in JET (L- and H-modes, ohmic, NBI and ICRF heating, pellet injection, etc.).

Furthermore, particular experimental details can be tested against specific features of the model. For example, the temperature perturbation that follows the collapse of a sawtooth propagates with the diffusivity, $\chi_{an,e}$, which largely independent of the mode of operation (L-, H- or ohmic) and is larger, in general, than that obtained from power
balance considerations, $\chi_e$, because of the presence of the critical electron temperature gradient, $(\nabla T_e)_{cr}$. Furthermore, since $\chi_{an,e}$ and $(\nabla T_e)_{cr}$ have different dependences on plasma parameters, the equilibrium $\chi_e$ will exhibit a different behaviour (and hence a different global scaling law) for ohmic discharges (where $\nabla T_e$ is comparable with $(\nabla T_e)_{cr}$) as opposed to additionally heated discharges (where $\nabla T_e$ becomes larger than $(\nabla T_e)_{cr}$).

In an independent analysis of heat pulse propagation [5], it was also conjectured that a critical temperature gradient would lead to the significant difference observed between the equilibrium $\chi_e$ and $\chi_{e,HP}$ deduced from the equation for the perturbed temperature.

Modelling heat and particle pulse propagation

Fig. 1 shows the comparison with experiment of the simultaneous simulation of temperature and density perturbations following the collapse of a sawtooth in JET Pulse No: 19617, with 7.5MW of ICRH. The decay of the signal amplitudes with increasing radius and the time traces agree well with experiment for both temperature and density perturbations. The inward propagating density pulse is successfully modelled with a particle diffusion coefficient, $D_{an,\alpha}$ taken to be proportional to $\chi_{an,\alpha}$ where $\alpha$ represents the different species (electrons or ions). The constant of proportionality needed (\approx 0.7) is also consistent with the evolution of density following pellet injection. Note that in the simulations the transient depression (dip) observed in the evolution of the density perturbation (indicative of coupling between energy and particle transport [6]) is well represented.

Steady state density profiles and density transport

A more complete model for anomalous particle transport, including also the behaviour of impurities, must address more than just the diffusive contribution emphasised by the study of density perturbations. Consistent with the above model for energy transport in a chaotic magnetic topology, radial particle transport should result, in part, from the radial projection of classical particle fluxes which flow parallel to the (chaotic) magnetic field. This will give rise to enhanced radial transport. In particular, impurities, which are subject to parallel thermal forces, will tend to migrate towards the regions of highest temperature and modify the electron density profile. In fact, this force acts like an inward "pinch" and offers an explanation of the peaked density profiles which are observed when $D_{an}$ is large and impurities are present.

This anomalous "pinch" can be checked by a study of density profiles obtained under various plasma conditions. As shown in Fig. 2, very similar density profiles are obtained under the quite different conditions obtained in an ohmic plasma with a $Z_{eff}$ of 1.4 and a plasma with 3.5MW of additional heating and a $Z_{eff}$ of 1.7. The model reproduces well the temperature profiles in the two cases and shows that $\chi_e$ is larger by a factor of about two for the additionally heated plasma ($\chi_e \approx 0.8$ m$^2$s$^{-1}$ in the ohmic plasma, and $\chi_e \approx 1.5$ m$^2$s$^{-1}$ with additional heating; both at $r/a = 0.85$). $D_{an}$ is therefore a factor of about two higher for the plasma with additional heating (and this is confirmed by the $D_{\alpha}$ emission) but since the density profile changes little, the inward "pinch" must increase also by a factor of about two. This is in line with the model of particle transport which predicts the correct density profiles in both cases (see Fig. 2).

Application to the H-Mode

The above model for energy and particle transport is now used to study confinement in the H-mode. No attempt is made to achieve a spontaneous transition from the L- to H-mode. Rather, the features of H-mode confinement are studied by triggering the transition artificially, by suppressing the anomalous transport in the high shear region near the plasma edge, and then observing the consequent effect on plasma parameters.
Fig. 3 shows the effect of reducing transport to its neoclassical value over the first 2 cm inside the separatrix. With no other adjustment, confinement improves in the core of the plasma. Furthermore, the establishment of the edge transport barrier leads to greater impurity retention and, as a result, increased electron density and $Z_{\text{eff}}$, as observed.

**Implications for a reactor**

The implications of this model of energy and particle transport and H-mode confinement for a reactor is shown in Fig. 4. Initially, the improved energy confinement is beneficial, leading to more alpha-particle heating. Subsequently, greater helium retention due to better particle confinement offsets this benefit and eventually ignition is quenched.

**Conclusions**

In addition to the satisfactory simulation of temperature profiles and their evolution in JET, it is now possible to simulate and explain detailed experimental features in both energy and particle transport. The most pronounced example is the simultaneous simulation of temperature and density pulse propagation and the coupling between density and temperature equations. This general behaviour is a direct consequence of including a critical electron temperature gradient in the transport equations. Another logical consequence of the assumed magnetic topology is an anomalous particle "pinch" on impurities which can lead to peaked density profiles. Furthermore, all experimental features of H-mode confinement can be simulated by the model provided the anomalous transport is suppressed in the high shear region near the plasma edge. However, application of the model to a reactor then shows that H-mode confinement does not expand the ignition domain because ashes and impurities are better retained.

**References**


I-180

Fig. 1: Heat and density pulse propagation following a sawtooth collapse (Pulse No: 19617)

Fig. 2: Comparison between experimental and simulated density profiles

Fig. 3: Radial profiles of $\chi_e$ and $\chi_i$ in L and H-modes

Fig. 4: Long term deficiencies due to helium poisoning in the H-mode
LOCAL TRANSPORT ANALYSIS IN L AND H REGIMES

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Introduction

In this paper we analyse the transport properties of L-mode and high density H-mode JET discharges with the same current and very close total auxiliary power input. The analysis extends previous work [1,2] by considering limiter L-mode discharges that do not undergo an L-H transition and by including particle transport. We also report on results for the evolution of equilibrium, current density, temperatures and particle density following pellet injection.

Our study is essentially of predictive nature, based on the predictive equilibrium-transport code JETTO and assuming as a reference the model of Rebut et al. [3]. Comparisons with results of the interpretive code FALCON are also shown.

The main characteristics of the pulses considered in this paper are shown in Table 1. Units in the table are s, MA, T, MW, \(10^{19}m^{-3}\), KeV and MJ.

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<th>(P_{NBI})</th>
<th>(&lt;n_e&gt;)</th>
<th>(T_{eo})</th>
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Table 1

Results of L- and H-mode transport simulations

The times shown correspond to the final times of our simulations which start in the ohmic phase and give time dependent results consistent with the experimental ones. We limit our study to the region which is not directly affected by sawtooth activity (normalised radius \(\rho \geq 0.25\) in these cases). The simulations considered here are not fully predictive because the measured effective charge \(Z_{efp}\) is used as described in [4] to derive the impurity concentration. This is done to avoid our results to be influenced by the uncertainty still existing in impurity transport models. The price to be paid is the uncertainty introduced by errors in the experimental profiles of \(n_e\) and \(Z_{eff}\). The model of Rebut et al. [3] (RLW model in the following) is our reference model. In our code it has been modified in the external region \(\rho \geq 0.75\) as described in [4]. The code updates the equilibrium configuration according to the evolution of plasma profiles and to the magnetic fluxes measured at the surface of the vessel. A consistent geometrical configuration is important for detailed comparisons with experimental results in x-point configurations and in pellet cases with strongly shifted magnetic axis.

Figures 1-3 show the values of the thermal and particle (hydrogenic ions) conductivities \(\chi\) and \(D\), respectively, at the final time of our simulations. The range of values of
\( \chi_{\text{eff}} = -q / 2n_e \nabla T_e \), where \( q \) is the total heat flux, obtained by the interpretive code FALCON from experimental rather than “predicted” profiles is also presented. An anomalous inward particle pinch was required only for the L-mode RF case [4]. In the case of the H-mode strongly inverted \( n_e \) profiles appear, with large uncertainties in \( \chi_{\text{RLW}} \) and \( \chi_{\text{eff}} \). Numerical instabilities may appear when the factor \((\nabla T_e / T_e + 2 \nabla n_e / n_e)\) in \( \chi_e \) becomes as small as observed (see Fig. 4). To avoid this we use a constant value of \( \chi_{\text{RLW}} \) consistent with FALCON results and allowing to reproduce the measured plasma profiles. \( \chi_e \) is derived a posteriori and compared with the “experimental” one (Fig. 3). The very low values of \( D \) found in the H-mode are consistent with the values of the impurity diffusion coefficient in the same shot [5]. Note also that \( \chi_{\text{neo}} \), predicted by the neoclassical theory may become comparable to the “experimental” \( \chi_i \) in the external region, due to the large local density and \( q \).

The difference in the values of transport coefficients in the ICRH and NBI L-modes, which have similar global confinement properties as shown in Table 1, may look surprising but it is entirely accounted for by the difference in the power deposition profiles. These are centrally peaked in the ICRH case but peaked externally in the NBI high density case. External peaking of the power deposition profile occurs also in the H-mode case, but now transport does not increase in the external region.

Results of the simulation of the pellet case

Shot 17749 is a typical representative of the so-called PEP (pellet enhanced performance) L-modes and it is reasonably well diagnosed. However, one problem has complicated past analysis of this and similar discharges. This is the observed discrepancy, of the order of 0.1m, (\( \geq 25\% \) of the radius of the enhanced confinement region) between the magnetic axis found by equilibrium identification codes and the observed maximum of \( T_e \) and \( T_i \) profiles, measured by diagnostics not related to the equilibrium configuration, such as Lidar and charge exchange recombination spectroscopy. We find that the time evolution of the equilibrium configuration, computed consistently with the evolution of plasma profiles, does show the required outward shift of the magnetic axis, and generates at the same time a poloidal field \( B_p \) at the “pick up coils” in agreement with the measured one (Figs. 5 and 6). The shift of the magnetic axis results from the relatively high central pressure and low central current density. The latter is mainly an effect of the bootstrap current. A slightly hollow current density profile just after pellet injection becomes very hollow 1.2s later, immediately before an MHD event reduces the PEP to a normal discharge. A strongly inverted q-profile may imply stability of the PEP against ballooning modes and suggests that the MHD event terminating the PEP phase is related to an instability corresponding to low \( m, n \) modes (e.g. 2,2 or 3,2). A MHD stability analysis based on the equilibria and q-profiles found here is required to elucidate these points.

From the point of view of transport modelling, the evolution of the q-profile is consistent with the one proposed in [3] to explain the local confinement properties of PEPs. Our simulations have been carried out by reducing transport in the region where \( Vq \) is positive. However, we find that while \( \chi_i \) is neoclassical, within the uncertainties related to the evaluation of the ICRH power deposition profiles, \( \chi_e \) is close to \( \chi_i \), rather than being neoclassical. These results are in agreement with previous studies (see e.g. [6,7]). Energy transport in the region of non-inverted q-profile is consistent with the RLW model and similar to standard L-mode discharges. Preliminary results of particle transport analysis show that \( D \) is also reduced (\( \leq 0.1m^2/s \)) in the region of inverted q-profile.

Conclusions

Our study confirms that in the region \( \rho \leq 0.7 \) the local transport coefficients of JET plasmas depend directly on plasma profiles and, indirectly, on power deposition profiles. They do not characterise the plasma regime as L or H according to their value. The H regime is rather characterised by the absence of increase in local transport coefficients, in the external region \( \rho \geq 0.7 \). The q-profile is likely to play an important role in determining the local transport coefficients, as implied by the RLW model.
Electron, ion thermal conductivities and hydrogenic particle conductivity, $\chi_{e}^{RLW}$, $\chi_{i}^{RLW}$, and $D^{RLW}$ in m$^2$/s vs. $p$ for the high density L-mode case of Table 1 ($t=47.5s$); the bars represent the range of the “experimental” $\chi^{eff}$.

Same as Fig.1 for the intermediate density L-mode case of Table 1 ($t=50.0s$).
Thermal and hydrogenic particle conductivities, \( \chi_e = \chi_e^{\text{eff}}, \chi_i, \chi_i^{\text{neo}} \) and \( D \), for the high density H-mode case of Table 1; for comparison \( \chi_e^{\text{RLW}} \) derived from the computed \( T_e \) profile is plotted as well (t = 51.5s).

Measured and computed electron density and temperature profiles for the case of Fig. 3 (t = 51.23s when Lidar \( T_e \) is available).

Measured and computed electron temperature profiles (t = 44.2s) vs. major radius R for the pellet case of Table 1 taken just before the end of the PEP phase; the q-profiles shown are taken at the beginning and at the end of the PEP.

Measured (dots) and computed poloidal magnetic field \( B_p \) along the pickup-coil line for the case of Fig. 5 (t = 44.2s); \( \chi^2 \) is the non-normalised square deviation.
Simulated ash transport experiments in jet using helium neutral beams and charge exchange spectroscopy


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Introduction. The helium pumping requirements of a tokamak fusion reactor will depend on the transport of the thermalised He\(^{2+}\) fusion products (ash) from the core of the plasma where they are formed, to the edge from where they are removed. In order to achieve a steady-state ignited plasma such that the \(\alpha\)-particle production rate is balanced by the He removal rate, it may be shown that the condition \(\tau_H(\text{He})/\tau_F \lesssim 10\) must be satisfied (in the usual notation), where \(\varepsilon\) is the helium pumped fraction (eg. / 1). \(\tau_H(\text{He})\) is governed by helium particle transport and also by the helium particle source profile, including recycling. Thus from the reactor physics viewpoint the relationship between particle and heat transport is of fundamental importance; the behaviour of the ratio of particle to heat diffusivity \(D/\chi\) also has significant implications for theories of tokamak transport.

The successful operation of one of the two neutral beam injection (NBI) heating systems on JET in \(^4\text{He}\) and \(^3\text{He}\) at 120kV has facilitated the simulation of the production of helium ash in the plasma core, by virtue of the excellent penetration of the He\(^\alpha\) beams (Fig. 1). The use of NBI as a helium particle source is intrinsically more satisfactory than gas modulation experiments, since the deposition profile of the He\(^\alpha\) beams is precisely known and is spatially distinct from the source due to helium recycling. Once the source profile is known, the radial flux of He\(^{2+}\) ions can be determined from the temporal and spatial evolution of the measured He\(^{2+}\) density profile \(n_{\text{He}}(r,t)\).

Measurement of He\(^{2+}\) density by charge exchange spectroscopy on JET. JET has initiated a comprehensive experimental and theoretical assessment of atomic data required for a quantitative analysis of CX spectra. Particularly in the energy range below 50 keV/amu, different theoretical approaches and new input from experimental sources have led to a significant revision (up to factors of 3) of excitation cross-sections for the HeI transition at 4686Å (cf./ 2). The JET atomic database - including multistep processes for the neutral beam stopping in an impurity dominated plasma - has been completed in 1990.

Two fiducial cases for the measurement of absolute He\(^{2+}\) densities were tested successfully in the 1990 campaign. In the first instance (He\(^{2+}\) being a minority) the total number of He\(^{2+}\) measured in a radial profile was compared to the total particle number
deposited by He⁺ NBI. In the second case, using a pure He discharge the total number of electrons contributed by He was compared to electron density profiles measured by LIDAR Thomson scattering. The results indicate that absolute measurements accurate to within 30% of the expected He²⁺ densities are attainable.

In addition to CX excitation processes, competitive processes such as beam halo and plume effects /3/ need to be taken into account. The plume effect may - in principle - lead to a significantly enhanced signal and therefore to a distortion of deduced radial profiles. For the JET viewing geometry the plume effect is estimated to be less than 20% in the standard magnetic field configuration, as used in the present experiments, but may reach enhancement factors up to ~5 above the true value of n_{He}(r) when B_φ is reversed.

Results
1. The simplest type of experiment in the present work was conducted as follows. A short (~0.5s) He⁺ NBI pulse (~3-6MW) was applied to a low to moderate density D plasma (<n_e>~1-3x10¹⁹m⁻³), sufficient to produce an average minority concentration n_{He}/n_e~10%. The evolution of n_{He}(0) was measured via a central CX viewing chord whose sensitivity had been absolutely calibrated. Fig.2 shows time traces of parameters of interest in this type of experiment, for a sawtooothing L-mode discharge. The main features of the data are the fast decline of n_{He}(0) after the He⁺ NBI pulse, followed by approximately steady-state conditions (the graphite top X-point tiles with which these plasmas predominantly interacted did not pump helium to any appreciable extent). The central decline is consistent with a redistribution of He²⁺ due to the combined effects of sawteeth and transport. Fig.3 shows a fit to n_{He}(0,t) calculated assuming only that the He²⁺ flux be given by the usual model i.e. \( \Gamma_{He} = -D_{eff} \nabla n_{He} + n_{He} v \). The fit shown is for D_{eff}(r) = 0.4(1 + 2r/a²) m²s⁻¹ and v = 0. The radial form of the diffusion coefficient follows that adopted in previous modelling of electron density transport in JET /4/. The fit to n_{He}(0) is not particularly sensitive to realistic values of v (assigned a radial form also given by /4/, which necessarily vanishes on-axis). The model calculation is performed in 1-D cylindrical geometry for a He²⁺ source profile corresponding to the He⁺ NBI deposition profile; 100% of the edge efflux of helium is redeposited according to an exponential-like radial profile of scale-length ~10cm. The temporal behaviour of n_{He}(0) is well reproduced, and also the absolute value. On the reasonable assumption that n_{He}(r) relaxes to an approximately flat shape in the steady-state phase when there is no central source (as is necessarily predicted by the calculation in Fig.3) the absolute agreement gives confidence in the measurement of n_{He}(0). Fig.3 also shows the edge Hel emission which is a measure of helium recycling, and the model calculation of the edge efflux normalised to it; the time dependence is well reproduced. The helium recycling increases from a low level to a steady-state value, reflecting the change in n_{He}(a) resulting from the combined effects of NBI fuelling (which although centrally peaked, still deposits He²⁺ near the edge, Fig.1) and the relaxation of the n_{He}(r) profile. The basic result of this experiment is to demonstrate a favourable characteristic time for transporting He²⁺ from the plasma core to the edge. This is evident both from the time constant of the decline of n_{He}(0) and also from the value of
based on the results of a heat-flux analysis of a similar L-mode discharge, \(\chi_{m/D\text{in}} \sim 2 - 3\). It is expected that sawteeth, present in the discharges of Figs. 2 and 4 (and indeed in all the experiments in the present work) play a significant role in the redistribution of the injected helium; this will be further investigated in future experiments in sawtooth-free plasmas.

2. The same experiment as (1) was also performed in the H-mode, for which the observed behaviour was more varied. In general, the decline of \(n_{\text{He}}(0)\) was less marked than in the L-mode case. In some ELM-free cases there was a tendency for \(n_{\text{He}}(0)\) even to increase after the central helium NBI source was stopped (Fig. 4). The H-mode in Fig. 4 is also interesting because of the low heating power after the He\(^9\) NBI was switched off, and the steady or even slightly falling \(n_e\). It may be noted in Fig. 4 that the HeII edge emission exhibits a reduction in level during the H-mode phase, similar to the usual \(D\alpha\) signature (also shown). This gives a direct indication of increased helium particle confinement. For the H-mode, the edge transport barrier is important in determining the global particle confinement time for He\(^{2+}\), in addition to the bulk transport properties. In this context, ELMs may provide a mechanism for enhancing the helium loss rate; in the present experiments the helium density behaviour of H-modes with ELMs was more like that of the L-mode in character.

3. Minority helium density profiles have been obtained in D plasmas fuelled by He\(^9\) NBI similar to the discharges of Figs. 2 and 4. In general it was difficult to relate the measured \(\Gamma_{\text{He}}\) to \(n_{\text{He}}\) and \(\nabla n_{\text{He}}\) (or any other driving term) due to the effect of sawteeth occurring between successive measurements of \(n_{\text{He}}(r)\). By comparing the time-averaged \(n_{\text{He}}(r)\) profiles before, during, and after the He\(^9\) NBI pulse, the changes \(\Delta n_{\text{He}}(r)\) and \(\Delta(\nabla n_{\text{He}})(r)\) may be determined. These quantities can be related to the (known) changes \(\Delta \Gamma_{\text{He}}(r)\) resulting from the application/removal of the He\(^9\) beams. The results show \(\Delta(\nabla n_{\text{He}})\) is negative in response to positive \(\Delta \Gamma_{\text{He}}\), implying a conventional diffusion coefficient, of order \(1 \text{m}^2\text{s}^{-1}\). More experiments, particularly in sawtooth-free discharges, are required in order to deduce accurate transport coefficients and to identify additional driving forces.

Conclusions. Effective removal of He\(^{2+}\) from the core of L-mode discharges has been demonstrated, using centrally deposited He\(^9\) NBI particle source and local CX measurements of \(n_{\text{He}}\). Preliminary indications of different behaviour in H and L mode plasmas have been observed, and the roles of both sawteeth and ELMs are identified as important issues for helium removal.

References

/2/ M G von Hellermann et al. JET-P(90)59
Fig. 1 NBI particle deposition profile for 3MW 120kV $^3$He$^0$ injected into a D plasma of volume average electron density $<n_e>$ = 2x$10^{19}$m$^{-3}$.

Fig. 2 Time traces for JET discharge #22977 (3.5MA, 2.5T, double-null X-point) showing L-mode behaviour of $n_{He}(0)$ and He recycling following a 120kV $^4$He$^0$ NBI pulse.

Fig. 3 Measured (broken line) and fitted (solid line) evolution of $n_{He}(0)$ and edge HeI recycling signal, for JET pulse #22977.

Fig. 4 Time traces for JET discharge #23198 (3.5MA, 2.5T, double-null X-point) showing H-mode behaviour of $n_{He}(0)$ and He recycling during and after application of 120kV $^3$He$^0$ NBI.
LOCAL CONFINEMENT IN NEUTRAL BEAM HEATED JET DISCHARGES

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Abstract

A database containing some 150 predominantly neutral beam heated discharges has been analysed. A scaling of the central ion temperature $T_i(0)$ with $P_{nbi}/E/N$ has been found ($P_{nbi}$ is the beam power, $E$: the (diamagnetic) energy confinement time and $N$: the number of particles). The central confinement of hot ion H-mode plasmas is found to be much better than that of hot ion L-mode plasmas. Density profile peaking increases with $P_{nbi}/N$, reflecting the effect of beam fuelling.

1. Introduction

Up to 21 MW of deuterium Neutral Beam Injection (NBI) heating has been applied to JET plasmas, in discharges limited either materially (Carbon or Beryllium surfaces) or by a magnetic separatrix (single or double null X-point). The deuterium beams have been injected at energies of 80 or 140 keV. For typical JET plasmas 60-80% of the beam power is absorbed by the ions. This leads to ion temperatures higher than electron temperatures and at low density it has been possible to generate hot ion plasmas with $T_i(0)$ $\geq$ 20 keV while at the same time $T_e(0) \leq$ 10 keV.

Ion temperature profiles have been measured with the charge exchange recombination spectroscopy (CXRS) diagnostic at JET [1]. A typical exposure time is 100 ms which is too long to observe the ion temperature sawteeth. Measurements yield some sawteeth averaged ion temperature in the centre and for the local transport analysis we shall distinguish between sawtoothing and sawtooth-free discharges.

Some 150 pulses have been analysed, using a local power balance analysis, which used PENCIL beam power deposition profiles. Time derivatives are included, but not convection. Convection is, however, estimated ignoring changes in density profile shape and $Z_{eff}$ because the data are insufficient for all these pulses to calculate the convection with confidence. The results of this estimate have not been included in the calculation of $\chi_i$.

This paper extends earlier work on a limited number of pulses [2,3].

2. Ion Temperature Scaling

Figure 1 gives the central ion temperature versus beam power per plasma ion. Clearly, the H-modes reach higher temperatures with less power per ion. They are actually close to temperature saturation: If we assume that 30% of the beam power is lost to the electrons and the remaining 70% is thermalized without any loss, than the predominantly 80 kV beams lead to a temperature of $2/3 \times 0.70 \times 0.78 \times 80 = 29$ keV, where 22 keV has actually been achieved with 80 kV beams only. The factor 0.78 stems from the fractional composition of the beam at the $E_b$, $E_b/2$ and $E_b/3$ values ($E_b$: Ion source acceleration voltage in kV). Figure 1, with its huge differences in $T_i$ for the same $P/N$ suggests a scaling with confinement: Indeed, Fig. 2, which gives $T_i$ versus $P/E/N$ ($c_E$ is the energy confinement time) gives a better relation. The PEP (Pellet Enhanced Plasma) featuring in these graphs is an enhanced central confinement mode as a result of pellet injection [4]. This mode features a very peaked central density ($> 8 \times 10^{19}$ m$^{-3}$) and has been achieved in both predominantly NBI and RF heated discharges [5].
3. Local Transport

The fact that $T_I(0)$ scales linearly with global energy confinement time shows that not only edge transport is improved during the H-mode, but also confinement well inside the plasma. A pedestal in $T_I$ is not sufficient to explain the difference in $T_I(0)$ between L and H-mode. To prove this, we give in Fig. 3 the local power balance for all these plasmas:

$$\left( V T_I(0.3a) \right) \text{ versus } \frac{1}{n(0.3a)} \int_0^{0.3a} \left( \frac{3}{2} n_k T_I \right) \frac{dV}{dt}$$

where $P_{\text{ions}}$ is the beam power density to the ions and $P_{\text{te}}$ the ion-electron equipartition. The slope of a data point with the origin is $1/(A \chi_I)$ where $A$ is the flux surface area ($A \approx 42 \text{ m}^2$ for $r = 0.3a$).

A slope of 10 in the units of the graph corresponds to $\chi_I = 1.5 \text{ m}^2/\text{s}$. To avoid too much uncertainty in the analysis, only plasmas with $P_{\text{nbi}}>6\text{MW}$, $n_I(0.3a)>10^{19} \text{ m}^{-3}$ and $<n_e><4 \times 10^{19} \text{ m}^{-3}$ have been plotted. Sawtooth free plasmas are plotted in black.

Most L-mode data all have a very similar slope, hence a similar $\chi_I$, independent of the amount of power per particle. $\chi_I \approx 1.6 \text{ m}^2/\text{s}$. Some inner wall plasmas, the so-called "Jolly good shots" [8], have a significantly better central confinement and often reach H-mode values.

The H-mode data fall roughly into two groups. The first group has a low power per plasma ion, a low $T_I$ and $\chi_I$ is better (i.e. lower) than for the L-mode. The second group are the good hot ion H-modes: high $T_I$, high power per plasma ion which is, however, offset by a steep rise in $n$ and often $T_I$ as well. They are all in the upper left corner of Fig.3. $\chi_I$ is extremely low for these plasmas similar to values obtained in reference shots with the TRANSP code [6]. Convection (i.e., particles flowing outwards, taking their thermal energy with them) is dominant in the centre, which means that conductive transport is even lower and could reach neoclassical values. Most of the beam power goes to increasing stored energy in the central plasma or is carried away with the particles. These plasmas are terminated by carbon or Beryllium catastrophe [7] or by loss of H-mode (often triggered by impurity influxes). Ion thermal transport apparently improves with increasing $T_I$ for H-modes.

Figure 4 gives radial profiles of $\chi_I$ for hot ion plasmas in L- and H-mode, sampled from two separate experimental campaigns. It is clearly seen that H-mode transport is not only reduced in the outer plasma (the "pedestal" idea of H-mode confinement) but also in the centre of the plasma. These central values are overestimated because convection is not subtracted from the power balance. FALCON benchmark runs on a hot ion L-mode and a hot ion H-mode suggest that convection is twice as large as conduction for both cases at $r = 0.3a$ and reduces strongly at larger values of $r$. That would bring H-mode confinement in the centre down to $0.1 \text{ m}^2/\text{s}$. Neoclassical conductivity is around $0.03 \text{ m}^2/\text{s}$.

A similar graph for confinement at $r = 0.3a$, can be made for the electron power balance. Because there is less power flow to the electrons, the uncertainty in transport coefficients is large, of the same magnitude as $\chi_e$ itself. Typical values are $0.5 \text{ m}^2/\text{s}$ for hot ion H-modes and $1.1 \text{ m}^2/\text{s}$ for hot ion L-modes. Data are given in Fig.5. To obtain the best accuracy possible, only electron temperature profiles measured with the LIDAR Thompson scattering diagnostic have been plotted in the graph. Further, only points with $n_e(0.3a)>10^{19} \text{ m}^{-3}$, $<n_e><4 \times 10^{19} \text{ m}^{-3}$ and $P_{\text{nbi}}>8\text{MW}$ have been included.

4. Density Profile Peaking

The beams provide a significant particle source in the centre of the plasma, especially at low plasma density. 10MW of 80kV deuterium beams provide $8 \times 10^{20}$ atoms per second in the whole plasma, compared with a typically $10^{21}$ ions target plasma for a hot ion shot. As the deposition profile is peaked, the centre of the plasma is rapidly filled with hot ions resulting in high $T_I$ and a degree of density profile peaking. Figure 6 gives the profile peaking (defined as $n_e(0)/<n_e>$; $<n_e>$ is the volume averaged electron density) as a function of $P_{\text{nbi}}/N$. Density data are from LIDAR. The increase of density peaking with $P/N$ is obvious, both for L-modes and for H-modes.
A dependence of $\chi_i$ on $n_e(0)/\langle n_e \rangle$ could not be established. The L-modes all have similar $\chi_i$ for a wide range of $P/N$ and $T_i$, whereas $n_e(0)/\langle n_e \rangle$ changes strongly over the $P/N$ range. Also for H-modes, which have a more complex behaviour of $\chi_i$, no trend with density peaking can be established. Density peaking is purely a beam fuelling feature, and in itself useful, because it enhances the beam power deposition in the centre of the plasma.

Conclusions

We have demonstrated a crude scaling relation for the central ion temperature as $T_i(0) \sim P_{\text{bene}}/N$. We have shown that the central ion confinement does not depend on power per particle. Hot ion H-mode plasmas can enter a regime where $\chi_i$ is extremely low and all input power either flows away as convection or is stored as plasma. Finally, density peaking in hot ion plasmas is associated with beam fuelling.

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Fig. 3: Central ion energy confinement for L- and H-modes

Fig. 4: Effective ion heat diffusivities for hot ion L- and H-modes

Fig. 5: Central electron energy confinement for L- and H-modes

Fig. 6: Density profile peaking related to beam power per plasma ion
INTERPRETATION OF HEAT AND DENSITY PULSE PROPAGATION IN TOKAMAKS

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INTRODUCTION

This paper addresses two key issues in current research on sawtooth induced heat and density pulse measurements in Tokamaks and their interpretation. First, heat and density pulses in JET and TEXT show different qualitative behaviour implying substantially different transport coefficients [1]. Second, a new description of the heat pulse has been used to describe measurements on TFTR, and it has been claimed that these measurements cannot be simulated with the widely used diffusive model [2]. In this paper, we show that consistency between all these measurements can be obtained assuming a diffusive propagation for the heat and density pulses and using linearised coupled transport equations.

ANALYTICAL TREATMENT: LINEARISING THE TRANSPORT EQUATIONS

By linearising the transport equations around a steady state, the contributions of the diffusion terms become dominant, since the scale lengths of the perturbations are small compared to the gradient lengths of the steady state profiles. As a result the relaxation of a localised perturbation of the electron temperature and density profiles can, to good approximation, be described by a set of two coupled diffusion equations:

$$\frac{\partial z}{\partial t} = A \psi^2 z$$

with: 
$$z = \left[ \begin{array}{c} n_1/n_0 \\ T_1/T_0 \end{array} \right]$$

(1)

Only dependences of the particle diffusion coefficient (D) and the thermal diffusivity (x) on Vn and VT appear in the linearised matrix; A11 = Dinc, A12 = (\partial D / \partial T)(T_0/n_0)\n, A21 = (2Dinc/3) + (2\partial x / 3\partial n)(n_0/T_0)VT, and A22 = 2/3(xinc + A12). Using the incremental diffusivities in the diagonal terms:

$$D_{inc} = D + \frac{\partial D}{\partial n} V n_0$$

$$\chi_{inc} = \chi + \frac{\partial \chi}{\partial T} V T_0$$

(2)

The spatial distributions of the initial perturbations of the electron density and temperature, induced by the sawtooth are to good approximation identical. Hence, the ratio \(\alpha = (n_1/n_0)(T_1/T_0)\) is independent of the minor radius, and the initial perturbations can be represented by a vector in the plane of n_1/n_0, T_1/T_0 (see Fig.1). The eigenvectors of the matrix A are also shown in Fig 1a. In general, a sawtooth excites a combination of the two eigenmodes.

All measurements of heat and density pulses in JET have been analysed using the model described above [3]. The coupling between the heat pulse and density pulse in JET is observed as an initial dip on the main density pulse which evolves with the slow eigenvalue of the coupled equations (Fig 1a).

INTERPRETATION OF DENSITY PULSES IN TEXT

The heat and density pulses in TEXT propagate at the same speed and have identical shape at every radial position. This behaviour is maintained in a wide range of plasma conditions [1]. Hence
\[ n_1/n_0(t) = \alpha \left( T_1/T_0 \right)(t) \] (3)

where \( \alpha \) is a constant. This implies that the initial condition for the heat and density pulses is an eigenvector of the coupled equations. In this case, the heat and density pulses are governed by a single diffusion coefficient (eigenmode) and propagate at the same speed through the confinement zone. Using this relation in the coupled equations (1), we find:

\[ A_{22} - A_{11} = \alpha A_{21} - A_{12}/\alpha \] (4)

where the value of \( \alpha \) at TEXT is normally quoted as 0.3 [1]. It is not possible to find a unique solution for the matrix elements from this equation.

The measurements in TEXT have been simulated numerically, using linearised coupled equations, to assess the possible values for the matrix elements. The off–diagonal term \( A_{12} \) is determined by a least squares fit of simulated pulses to the data, for fixed \( A_{22} \) (\( 2 \text{ m}^2/\text{s} \)) and for the ratio \( A_{22}/A_{11} \) ranging from 1 to 10. The results thus obtained are in good agreement with the analytic expression in eq. 4 (Fig. 2).

The experimentally observed identical behaviour of the heat and density pulses over a wide range of discharge can be explained from the results of these simulations when \( A_{22} > A_{11} \). Consequently, the density pulse is governed by the fast eigenmode of the coupled equations. Any analysis which does not take the coupling into account will lead to incorrect conclusions (e.g.: \( D_{\text{inc}} \approx \chi_{\text{eff}} \)). Also, the observed discrepancy in TEXT between density pulses launched by the sawtooth and density pulses induced by other perturbative techniques [4] can be explained, since the sawtooth driven density pulses are governed by the large eigenvalue (\( \approx 2 \text{ m}^2/\text{s} \)), whereas other type of density perturbations are determined by \( D_{\text{inc}} \approx 0.8 \text{ m}^2/\text{s} \).

By expressing the particle and energy fluxes in terms of \( V_{\text{p}} \) and \( V_{\text{T}} \) the transport matrix \( M \) can be derived from \( \Lambda \). This transport matrix \( M \) is very similar for TEXT and JET.

**Fig. 1:** The initial perturbation of the density and temperature represented in the \( n_1/n_0 - T_1/T_0 \) plane. In JET (Fig a) both fast and slow eigenmodes are launched, whereas in TEXT (Fig b) the sawtooth launches the fast eigenmode.

**INTERPRETATION OF HEAT PULSES IN TFTR**

The heat pulse can also be described using the full transport equations and temperature profiles. It is generally found that a simulation using the diffusion coefficient (\( \chi_{\text{eff}} \)) derived from power balance analysis leads to a slower propagation of the heat pulse than experimentally observed. Therefore a temporary enhancement of the thermal diffusivity needs to be introduced to explain the measurements.

One type of temporary enhancement of \( \chi \) has been used by TFTR and is based on the assumption that the sawtooth crash induces a transient increase in the level of turbulence in the
plasma [2]. For TFTR pulse 30904 $\chi$ is enhanced in the centre of the plasma, a few ms after the sawtooth crash:

$$\chi = \chi_{\text{eff}}(1 + 150 \exp(-t/0.001) \exp(-9.6r^2/a^2)),$$

with: $\chi_{\text{eff}} = (0.3 + 3.8r^2/a^2) \exp(-1.1r^2/a^2)$ \quad t:0: sawtooth crash \quad (5)

In Fig. 3 the numerical simulations of the heat pulse, reproduced at JET using the above equations are shown. There are clear discrepancies between the measured amplitude of the heat pulse at $r = 0.65$ m. This discrepancy was not obvious in [2], since the simulations were normalised to the measurements [5]. In addition, the enhanced $\chi$ decreases the temperature in the centre significantly after the sawtooth collapse, yielding an incorrect initial condition for the relaxation of the temperature.

There is a further discrepancy with the measurements 200 $\mu$s after the sawtooth collapse and the simulations (Fig. 4). It is not possible to modify the model for $\chi$ in order to describe the heat pulse, even when the enhancement of $\chi$ is limited to a region around the mixing radius.

Generally, the time dependence required to describe the heat pulse should be different at different radii outside the mixing radius, such that the enhancement is slower further away from the mixing radius. This type of enhancement is similar to the change in $\chi_{\text{eff}}$ implied by the results using linearised equations: the thermal diffusion coefficient is a function of the local temperature gradient [6].
A good fit to the TFTR data is obtained with the diffusive model using linearised coupled transport equations [3]. The measured $T_e$ evolution just outside the mixing radius can be used as a boundary condition, or the heat pulse can be reproduced starting from an initial localised perturbation in the centre of the plasma, with a mixing radius equal to the measured radial extend of the initial perturbation of the $T_e$ profile (Fig. 5).

**Fig. 5:** Fit to the TFTR heat pulse using coupled linearised transport equations. We find that $\chi_{\text{inc}}^\text{eq} = \frac{1}{1 + 14(r/a)}$ for $r > r_{\text{mix}}$ ($r_{\text{mix}} = 0.34 \text{ m}$). Coupling to the slow density pulse (dashed curve), no coupling (dotted curve).

**CONCLUSIONS**

Even though the heat and density pulses in JET and TEXT appear to be substantially different, the measurements can be simulated consistently assuming a diffusive propagation of the heat and density pulses, using linearised coupled transport equations. The transport coefficients obtained are very similar for JET and TEXT. Similarly, the measurements of the heat pulse in TFTR can be described using the same model. This model has an implicit time dependence for the $\chi_{\text{eff}}$ on the time-scale of the heat pulse. It is not necessary to invoke an explicit time dependence for $\chi_{\text{eff}}$.

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COMPARISON OF THE IMPURITY AND ELECTRON PARTICLE TRANSPORT IN JET DISCHARGES

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The transport of impurities in tokamak discharges is known to be dominated by a large diffusivity, which is between one and two orders of magnitude above the neoclassical values. Analysis of steady state impurity profiles yields a uniform diffusion coefficient $D_{\text{imp}} \geq 1 \text{ m}^2/\text{sec}$. However, steady state profiles are only sensitive to the ratio of $D_{\text{imp}}$ to the convective velocity ($v_{\text{imp}}$) in the absence of sources. Temporal and spatial resolution of transient events are needed to separate $D_{\text{imp}}$ and $v_{\text{imp}}$ in the plasma core. Cases have been reported [1] where a strong reduction of $D_{\text{imp}}$ is required in the central region to account for the experimental data. In JET these cases extend to a wide variety of plasma conditions. This paper sets out to draw together the disparate observations of reduced transport in the plasma core into a general picture. A consequent prediction was that the electron density profile should be peaked during the early current rise, which has been found to be the case. This and some results of modeling with reduced core diffusivity are described in the final section.

Reduced impurity transport in the core of JET discharges

This reduction of the impurity diffusion in the central region of the discharge is detected unambiguously in JET when impurities are seen to accumulate in the centre of the discharge. This occurrence is correlated with electron density profiles well peaked on axis, with the absence of sawtooth activity and with central values of the safety factor $q(0) > 1$. These conditions are often met after pellet injection [2]. In such cases, where high deuterium density gradients are attained in the central region, large neoclassical inward velocities arise for impurities leading to their progressive accumulation in the plasma centre over a time scale of 1 to 2 sec. The high level of peaking reached by the impurity density profiles implies that $D_{\text{imp}}$ be lower than values around 0.1 m$^2$/sec within $r = 0.4\cdot a$.

On a longer time scale and in connection with less pronounced electron peaking, impurity accumulation has also been detected in the early phases of the discharge before the sawtooth activity starts [3].

No impurity accumulation is observed during monster sawteeth. However, this cannot be interpreted as evidence of high central $D_{\text{imp}}$. In fact, as in these cases the electron density profiles are nearly flat in the centre, there is no substantial neoclassical pinch on impurities: the intrinsic impurity density profiles, therefore, are not expected to peak markedly during the sawtooth period.

Positive evidence of a reduced diffusivity zone in ICRF heated pulses during monster sawteeth, and in sawtoothing Ohmic ones, is obtained in impurity injection experiments [4]. In such cases a fast propagation of the injected impurities from the plasma periphery to about $r = 0.4\cdot a$ is observed. From this radial position the propagation of impurities proceeds at a much slower pace, showing a sharp transition of $D_{\text{imp}}$ down to levels 20 to 30 times lower than outside over a distance not larger than 0.15 m. A sharp transition over such short distances is also required in the cases described before.
Sawtooth free targets for additional heating are produced occasionally in JET. During the heating, the same pattern in the profile of the diffusion coefficient has been identified by the analysis of an impurity injection experiment in one of these pulses.

An important transition to low levels for $D_{\text{imp}}$ in the plasma centre is also deduced from the analysis of impurity injection experiments in H-mode pulses [5]. Another notable case where this feature is clearly identified is the impurity depletion phase following the back-transition to L-mode after long lasting H-modes [6].

To evaluate the evidence summarized above, it is important to consider that the central zone of reduced transport and the sharp transition for $D_{\text{imp}}$ are only likely to appear directly in the impurity density profiles a) during transient events (e.g. imp. injection) or b) when important convection effects in the centre lead to strong shaping (accumulation) of those profiles. In most cases the electron density profile is not peaked. Therefore the distribution of impurities is expected to be at most very mildly peaked in the centre. Furthermore the sawtooth activity counteracts the development of any structure in that region because of its well known spreading effect on particles.

Two conclusions can be drawn about the core reduction of $D_{\text{imp}}$: i) the marked transition of the anomalous effects on impurity diffusion is clearly demonstrated in a variety of confinement modes, operation scenarios and plasma parameters including values of $q(0)$, $q'$ in the central region, $n_0$ and $T_e$ and their gradients; ii) no cases have been identified yet without the reduction of $D_{\text{imp}}$ in the core. It should be added that the central values of $D_{\text{imp}}$ are generally close to neoclassical predictions [2,4].

Reduced electron transport in the core of JET discharges

Evidence of a marked reduction of the electron particle transport in the centre after pellet injection has already been reported from JET [7,8]. In those cases a particularly low thermal diffusivity in that region has also been observed. The modifications induced by the pellet on the plasma conditions have been interpreted as a necessary condition for the achievement of reduced energy transport in that region. If this same assumption were made also for electron particle transport, a marked contrast would emerge with the impurity transport that we have shown to be affected by such an important reduction of the anomalous transport in a large variety of cases.

In fact it is difficult to prove or disprove the existence of this zone of reduced diffusivity. The considerations described above for the impurities are equally true for the electrons. However the difficulties are greater because the only neoclassical convective effect expected is the Ware pinch whose strength (measured in terms of the parameter $v \cdot r/D$) relative to the neoclassical diffusivity of electrons is much smaller than the corresponding parameter for impurities: therefore much smaller levels of peaking in the centre are expected for the electron density, if any. Furthermore a correct analysis in terms of diffusion demands a high space and time resolution. In particular the profiles obtained from Abel inverted interferometric data, due to the limited space resolution of this instrument, often involve large distortions on the density gradients, which induce large systematic errors in the deduced transport parameters.

In particular central protuberances as those in the LIDAR profiles of fig 1 are strongly attenuated in general in the profiles from the interferometer. It is worth noting the similarity between the two peaks observed during the early phase of the discharge and after the pellet
injection respectively. While the latter (t=47.5 sec) results from the external fuelling of the central region by a 4 mm pellet injected at 45.5 sec in an ohmic pulse, the former (t=41.5 sec) must have developed spontaneously and persists until a first small pellet is injected 3 sec later. This feature is not just an occasional occurrence: a high central peaking comparable to that of 41.5 sec is observed in all LIDAR profiles available in the JET database for the early phase of discharges with modest edge fuelling. If the Ware pinch is causing such a peak, the particle diffusion coefficient D_p must be suffering a strong reduction between r/a = .5 and .3 and its value in the centre must be about 0.03 m^2/sec, the same value that can be deduced for the post-pellet phase.

Results from a simulation of a long time interval during this pulse are indicated in the figs 1 to 3. The code used integrates the particle and current transport equations simultaneously. Neoclassical resistivity is used and the Ware pinch is calculated selfconsistently. The shape and values of the Ware velocity profiles are not strongly sensitive to the particular shape chosen for the initial current distribution. Furthermore the pinch itself has only a minor influence on the time evolution of the post-pellet profiles. The absolute independence on time of the D_p profile required in the simulation is probably only a coincidence and would almost certainly be removed if the precision and the repetition rate of the LIDAR profiles were higher than presently available. More important it seems to us the finding that a sharp transition in the value of D_p is found before the injection of a pellet and that this feature is very general in the early phase of JET discharges. A comparison is shown in fig 3 between the D_p profile and and the profile for D_imp required for the description of the accumulation of the intrinsic impurities (Ni and C) in the same pulse after the pellet injection. The lower levels of the two diffusion coefficients D_imp and D_p in the centre are close to their neoclassical values for Ni and for the electrons respectively.

According to our previous discussion on the possibility of unambiguously identifying a zone of reduced diffusivity in the centre of the discharge, other candidate pulses for our search are those with monster sawteeth and low density pulses with deep neutral beam fuelling. Due to the high central temperature in monsters, to the consequent low parallel E field and low Ware pinch, only a minimal density peaking is expected which is consistent with the experimental data. For the low density beam fuelled pulses, due to the frequent sawteeth and to the low acquisition frequency of LIDAR profiles, the experimental data analyzed so far are not conclusive on the existence of a reduced diffusivity zone in the centre. It must also be noted that we have no data from impurity injection experiments in such cases for a detailed local study of the impurity transport.

Conclusions
We have shown that a sharp reduction of the anomalous diffusivity of impurity ions in the plasma core is a very general feature in JET discharges: indeed no cases have been identified yet where the absence of this feature can be ascertained. Furthermore we have demonstrated that the post-pellet phases are not the only cases where the same feature is present in the profile of the electron diffusion coefficient: indeed this is generally present in the early phases of the JET discharges and preexists to the pellet injections. The analysis of other plasma conditions from the presently available JET data base is made difficult by unfavourable experimental conditions. Specifically devised experiments and diagnostic arrangements seem to be required.
References

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**Fig. 1** Experimental (solid lines) and simulated (dashed lines) electron density profiles. JET pulse 13572

**Fig. 2** Time evol. of experimental (from LIDAR scattering and Abel inverted interferometric data) and simulated electron density at the centre. JET pulse 13572

**Fig. 3** Radial profile of the the electron and impurity diffusion coefficients. Neoclassical levels at the centre are also indicated for comparison.
PARTICLE AND ENERGY TRANSPORT PROPERTIES DEDUCED FROM THE PLASMA DYNAMIC RESPONSE

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To face up the difficulties encountered in understanding the tokamak transport phenomena, the study of the plasma dynamic response appears as a promising approach. Very often the use of this dynamic response is limited to the evaluation of the characteristic coefficients of a given transport model, by matching the simulated and the actual temporal evolution of the plasma parameters. In this paper we will present a more objective and more deductive methodology. It will be shown that the identification of the transfer function of the system in an appropriate form allows to elicit fundamental properties of the transport processes and to validate or discard potential transport models.

Formalism and analysis method

The basic step is the identification of the transfer function of the system. Details about the employed method are reported in [1]. We recall here that the continuity equation for a transported quantity $y$ takes the general form:

$$ \frac{\partial y(r,t)}{\partial t} = x(r,t) + Ly(r,t) $$

(1)

where $x$ is the corresponding source term and $L$ the transport operator describing the associated flux. $L$ is assumed to be a linear time independent integro-differential operator acting on the spatial variable. This general form covers the usual convective and diffusive terms and we will show that its linearity and stationarity are not severe constraints. If the source term is separable ($x(r,t) = x(r)u(t)$), the transfer function of the system is given in the Laplace variable by:

$$ H(r,s) = \frac{y(r,s)}{u(s)} = \sum_{n=1}^{\infty} \frac{y_d(r)}{s - p_n} $$

(2)

The poles and the residues $p_n$ and $y_n$ coincide with the eigenvalues and eigenfunctions of the transport operator with an appropriate normalisation condition:

$$ p_n y_d(r) = L y_d(r) , \quad x(r) = \sum_{n=1}^{\infty} y_d(r) $$

(3)

It is important to note that (a) the poles do not depend upon the radius, so that the formalism can be applied directly to line integrated or non calibrated measurements, that (b) the eigenmode profiles do not depend on the source term distribution, so that the analysis extend also where the source does not vanish and that (c) any waveform for the excitation signal $u(t)$ can be used, so that a wide variety of perturbations can be considered. Using the system identification technique

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described in [1], the transfer function can be experimentally deduced in the prescribed form (2). This yields to a diagonalised representation of the transport processes, without invoking any particular transport mechanisms. In tokamak physics context, this constitutes a blind and therefore less biased approach to the problem.

Experimental results

4 to 20 pellets (600 m/s, 1-4 Hz) are injected in ohmic Tore Supra plasmas ($R = 2.38 \text{ m}$, $a = 0.75 \text{ m}$, $B = 3.8 \text{ T}$, $I_p = 1 - 1.5 \text{ MA}$, $n_{e0} = 3 \times 10^{19} \text{ m}^{-3}$, $D_p$). The electron density $n$ and temperature $T$ are measured with a 5-channel interferometer and a 6-channel Fabry-Pérot interferometer (fig. 1). The ablation duration (0.5 ms) is shorter than the sampling period (1-2 ms) allowing to consider the excitation $u(t)$ as a Dirac function. Column displacement is negligible.

The transfer function identification is limited by the sampling rate and the noise level to 3 poles. It has been experimentally checked that (a) the plasma response can be reproduced with poles independent of the radius, according to eq. (2) and that (b) the system is linear by showing that identical poles can simulate both small and large amplitude perturbations (particle content increase from 5 to 100%).

Transport modelling

We will first establish two simple properties of the transport operator that will elicit a coupling between particle and heat transport. In such a case the eigenvalue equation is:

← fig. 1 : Raw density and temperature dynamic response and their simulation using a 3 pole transfer function with poles common to both parameters.

↓ fig. 2 : Density and temperature time constants for different discharges obtained with independent poles (+) and poles common to both parameters (o).
The first peculiarity of this system is that its poles are shared by both n and T; this is experimentally verified by comparing the poles obtained with independent transfer functions (fig. 2) or by trying to simulate the dynamic response with identical poles for both parameters (fig. 1). In addition if the four operators of eq. (4) are homotetic, then the eigenmodes group in pairs, in the sense that \((n_1, n_2)\) and \((T_1, T_2)\) display the same first order radial eigenmode profile, while \((n_3, n_4)\) and \((T_3, T_4)\) display the next order one. Inspection of the raw experimental profiles (fig. 3) demonstrates this second property for the three first eigenmodes. These observations are strong evidence for a coupling between particle and heat transport, whatever the underlying mechanisms are.

\[
p_n \begin{pmatrix} n_n \\ T_n \end{pmatrix} = \begin{pmatrix} L_m & L_mT \\ L_{TT} & L_n \end{pmatrix} \begin{pmatrix} n_n \\ T_n \end{pmatrix} \tag{4} \]

Under these circumstances, the following form for the particle and heat fluxes will be evaluated:

\[
\begin{pmatrix} \Gamma_q \\ \nabla \theta_q \end{pmatrix} = \begin{pmatrix} D, D_+ \\ \kappa, \kappa_+ \end{pmatrix} \begin{pmatrix} n_n \\ T_n \end{pmatrix} + \begin{pmatrix} 0 \\ 0 \end{pmatrix} \begin{pmatrix} V_n \\ U \end{pmatrix} + \begin{pmatrix} 0 \\ \gamma \end{pmatrix} \begin{pmatrix} 0 \\ T \end{pmatrix} \tag{5} \]

The diffusive term coefficients \(D, D_+, \kappa, \kappa_+\) and the abnormal pinch velocities \(V, U\) depend not only on the radius, but also on the local parameters \(n, V_n, T, U_T\), since these parameters vary during transients. These coefficients are appropriate to treat complex situations such as turbulence induced transport or critical gradient model, which recently gained a large credit. This general model is then linearised for small perturbation around a steady state (indexed 0), leading to the following eigenvalue equations:

\[
p_n \begin{pmatrix} 1 & 0 \\ \frac{2}{T_n} & 2 n_0 \end{pmatrix} \begin{pmatrix} n_n \\ T_n \end{pmatrix} = \begin{pmatrix} 0 \\ 0 \end{pmatrix} \begin{pmatrix} n_n \\ T_n \end{pmatrix} + \nabla \begin{pmatrix} \nu_m \nu_T \\ \nu_{TT} \end{pmatrix} \begin{pmatrix} n_n \\ T_n \end{pmatrix} + \nabla \begin{pmatrix} D_m D_{TT} \\ D_{TT} \end{pmatrix} \nabla \begin{pmatrix} n_n \\ T_n \end{pmatrix} \tag{6} \]

The elements of the \(\nu\) and \(D\) matrices are easy to deduce combinations of the transport coefficients and their functional derivatives [2]. The damping due to the variation of the ohmic power is introduced by \(\frac{2}{T_{TT}}\). A modification in the recycling conditions is limited to the neutral
penetration depth and does not influence the analysis in the observed region.

Eq. (6) is repeated for each identified pole and solved to evaluate a selected set of transport coefficients. In addition, a given model can be validated if eq. (6) is reasonably satisfied for all these poles, in which case this model can explain the plasma dynamic response in detail. This validation criterion is illustrated for the purely diffusive model ($\gamma = 0, \psi = 0$) which is approximately the situation retained to study the propagation of the sawtooth induced perturbations in JET [3]. The diffusion coefficients (fig. 4) compare qualitatively well with those of JET, featuring also a large ratio between $D_{nT}$ and $D_{nn}$ a negative $D_{nT}$ and a positive $D_{nn}$. However the difference in these coefficients evaluated with fast or slow poles (fig. 4) questions the validity of this model, or at least its linearity.

Applying this kind of argument on all the possible combinations of the transport coefficients and their functional derivatives, we established that the models that can most satisfactorily explain simultaneously the slow and fast dynamic response involve either $D$ and $D_n$ or a particle diffusion coefficient $D$ depending on $n$ and $VT$, both with $\gamma = 5/2$. The latter model leads to the following linearised transport coefficients (fig. 5):

$$
D_{nn} = D; D_{nT} = \nabla n_0 \frac{\partial D}{\partial VT}; \nu_{nn} = \nabla n_0 \frac{\partial D}{\partial n}; D_{TT} = k + VT_0 \frac{\partial k}{\partial VT} + \gamma T_0 \nu_{TT}
$$

(7)

← fig. 5 : Linearised transport coefficients of eq. (7).

**Discussion**

This choice is further supported by the following points:
(a) the coefficients $D_{nn}$ and $\nu_{nn}$ should be very small in the flat density profile region ($r/a < 0.4$), a property satisfied without having been imposed (fig. 5); (b) the positive sign of $\nu_{nn}$ indicates a decreasing particle diffusion coefficient when $n$ is increased, a feature that has been elicited by a lot of scaling laws; (c) similarly the sign of $D_{nT}$ implies that the same coefficient increases for steeper temperature profiles, a dependence suggested by theoretical models involving turbulence thresholds; typical values for $\partial (\ln D)/\partial (\ln VT)$ and $\partial (\ln D)/\partial (\ln n)$ lies between 1 and 2; (d) these dependencies end in a non linear transport model, corroborated by steady state observations such as confinement degradation or profile resiliency. Note that this linearised analysis does not allow a separate evaluation of the thermal conductivity and its temperature gradient dependence (eq. (7)). In our case a vanishing combination has been found appropriate.

In summary we elaborated a formal link between the experimentally identified transfer function of the plasma and the underlying transport mechanisms that allows to deduce properties of these mechanisms and to test transport models.

**Acknowledgements**

The experiments analysed here rely on the full dedication of the entire Tore Supra group.

**References**

Particle and heat transport analysis is an important task in order to make meaningful comparisons between experiment and theory. In this paper a method to obtain transport coefficients from experimental stationary profiles and fluxes is presented. This method can give values of particle and heat transport coefficients and its coupling. We assume that the linearized kinetic equation is a good description of fusion plasmas, even when we include non-neoclassical effects. From this linearized equation and assuming that the effect of anomalous processes over Ware pinch, and related bootstrap current, is not important we can obtain the following general approximate expressions for the total particle and heat anomalous fluxes:

$$\Gamma_{AN} = \Gamma - \Gamma_{NC}$$

$$Q_{AN} = Q - Q_{NC}$$

$$\frac{\Gamma_{AN}}{n} = -D_{AN} \left( \frac{n'}{n} + \alpha \frac{T'}{T} \right) + V_{AN}$$

$$\frac{Q_{AN}}{nT} = -\left( \alpha + \frac{3}{2} \right) \left[ D_{AN} \left( \frac{n'}{n} + (\alpha + 1) \frac{T'}{T} \right) + V_{AN} \right]$$

With these expressions we have for each radial point two equations and three parameters. To obtain a solution one more assumption is necessary. Several assumptions are applied:

a) Minimization of anomalous coefficients: $\Sigma(D^2 + V^2 + \alpha) = \text{minimum}$

b) Constant coefficients: $D_0, \alpha = \text{constant}$, $V + \alpha DT/T = V_{0p}/a$

c) Fluctuations, from quasilinear theory in the plateau regime[1]: $\alpha = -0.5$

d) Non ambipolar transport: $V = -D_{0p}/T$ with $E_r$ experimental radial electric field

As an example this method is applied to the experimental profiles[2] and fluxes in the TJ-I tokamak ($R = 30\,\text{cm}$, $a = 10\,\text{cm}$, $I = 50\,\text{kA}$, $B = 1.5\,\text{T}$, ohmically heated) shown in figures 1 and 2.

1. Electronic Density and Temperature in TJ-I

<table>
<thead>
<tr>
<th>n (m$^{-3}$)</th>
<th>T (eV)</th>
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<tr>
<td>$4 \times 10^{19}$</td>
<td>350</td>
</tr>
<tr>
<td>$3 \times 10^{19}$</td>
<td>250</td>
</tr>
<tr>
<td>$2 \times 10^{19}$</td>
<td>150</td>
</tr>
<tr>
<td>$1 \times 10^{19}$</td>
<td>50</td>
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2. Particle and Heat Fluxes in TJ-I

<table>
<thead>
<tr>
<th>Q (Wm$^{-2}$)</th>
<th>$\Gamma$ (m$^{-2}$s$^{-1}$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$3 \times 10^{21}$</td>
<td>$1 \times 10^{21}$</td>
</tr>
</tbody>
</table>
The results are strongly dependent on the assumption made, as can be seen in figures 3 to 5. It is concluded that in order to make comparisons experiment-theory, it is necessary to use experimental coefficients obtained with an assumption coherent with the theory that is being compared. So, it is not adequate to compare coefficients given by fluctuation theories with experimental values obtained with the assumption of constant coefficients.

Fluctuations

From the results of ref[1] for transport induced by fluctuations and the values of $D$ and $V$ obtained in this work with the assumption c, we can estimate the poloidal phase velocity and the intensity of the electrostatic fluctuations needed to explain the transport:

$$\frac{V}{D} = -v_0 \frac{eB}{T} ; \quad D = \sqrt{\frac{\pi}{4} \rho^2 v_i} \left( \sum k^2 \left( \frac{e\Phi}{T} \right)^2 \right)$$
Phase velocity is shown in figure 6 together with its value corrected by plasma rotation with the experimental $E_r[3]$ and the electron diamagnetic velocity. The agreement between these two later velocities is very good, and we can think in drift waves as a candidate to explain the transport. Values for $k_e$ are estimated from $v_\theta$ and frequency values given by reflectometry experiments, while $k_\parallel$ is taken as $1/\rho_R$. The parameter $k_0 \rho_3$ is around 0.5 in the bulk plasma and lower (0.01) in the border. The resulting fluctuation levels are shown in figure 7 together with the predicted by the Kadomtsev limit for drift waves taking $k_r=k_0$. As can be seen the agreement is good except in the region close to the plasma border, were the quasilinear theory it is not applicable. In order to make a drastic conclusion as 'transport can be explained by drift waves in the plasma bulk' we need a more precise determination of experimental fluxes and profiles together with complete measurements of $k$, $\omega$ and fluctuation levels inside the plasma. This analysis with estimated values is used only as an example of the capabilities of the method.

Ion loss cones

Ions with kinetic energy larger than a certain critical value, can be trapped in a banana orbit which hit the limiter (a toroidal belt in TJ-I) and become lost. This critical energy depends on the distance to the limiter, on the detapping collision frequency and on the plasma rotation. The strength of this process on ion particle transport near the last closed magnetic surface can be calculated as:

$$\frac{\partial N}{\partial t}_{ILC} = -\frac{4}{\sqrt{\pi}} \eta v_{90} (1 - \exp(-u)) \frac{\exp(-u^2)}{u^2} ; \quad \Gamma_{ILC} = \frac{1}{r_0} \int_{r_0}^{L} \left( \frac{\partial N}{\partial t} \right)_{ILC} \, dr$$

being $u = v_\text{critical}/v_t$. Expressions for calculating $u$ can be obtained from ref[4] plus the condition in collisionality, and easily extended to include the effect of radial electric field. In this work it is shown that the ion loss cones can take account of a significant part of particle transport although the plasma ions are not in the banana regime. The main parameter is the edge ion temperature that will be high enough to make important the ion loss cones.

Assuming than ion temperature is given by Artsimovich scaling law in the center and follows a smooth transition to $T_i = 2T_e$ in the border, the obtained ion loss cones flux is shown in figure 8, calculated with the experimental $E_r$ and with $E_r = 0$, together with the experimental particle flux. This ionic temperature is compatible with experimental estimations based on charge-exchange spectroscopy and Hα wings. The experimental radial electric field $E_r$ was obtained from the measured rotation by doppler shifts of impurity lines. Due to the experimental difficulties, the values close to the border could be affected of large errors.
As we can see the ion loss cones can only explain feasibility one portion of the flux. We can assume that the other portion is due to the intrinsically ambipolar mechanism and that the electron transport is made of these ambipolar mechanism plus another one non ambipolar. These second and unknown mechanism is needed to balance the ion loss cones flux and to determine the radial electric field. The analysis of transport coefficients can be made separately for the ambipolar and no ambipolar fluxes, but we need a further assumption: the contribution to the energy flux of the non ambipolar transport mechanism.

References
DETERMINATION OF THE ELECTRON HEAT DIFFUSIVITY FROM TEMPERATURE PERTURBATIONS IN FT AND FTU TOKAMAKS

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EXPERIMENTAL RESULTS. The evolution of electron temperature perturbations following the sawtooth collapse has been measured at different radial positions by Electron Cyclotron Emission and used to determine the electron heat diffusivity $\chi_{e,HP}$ in FT and FTU tokamaks. The $\chi_{e,HP}$ values obtained on FT, using the time-to-peak (tt) method, have been compared to the ones ($\chi_{e,FB}$) resulting from steady state power balance. Power balance analysis performed under the assumption of neoclassical ion transport yields a $\chi_e$ profile which shape is a strong increasing function of radius, while its product with density slowly increases; for any fixed radius in the transport dominated region, $\chi_{e, FB}$ is an inverse function of the line average density $\bar{n}$ in the "linear"confinement regime, i.e. for $\bar{n} < 10^{20} \text{ m}^{-3}$ /3/., while in the "saturated" regime it remains roughly constant at $\chi_e = 0.5 \text{ m}^2 \text{ sec}^{-1}$. The value to be compared to $\chi_{e,HP}$ is the volume average over the diffusion region $\chi_{e, FB}$; it increases strongly for decreasing edge safety factor, since the diffusion region is pushed outwards by the increasing sawtooth mixing radius $r_m$.

<table>
<thead>
<tr>
<th>$n \times 10^{20} \text{ m}^{-3}$</th>
<th>$q_a$</th>
<th>$\chi_{e,HP}$</th>
<th>$\chi_{e,FB}$</th>
<th>$t_{ex}(\text{ms})$</th>
<th>$t_{OH}(\text{ms})$</th>
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The work presented in ref. 2 has been revisited, since a systematic error in the measurement of electron temperature profiles has been removed. A good agreement has been found between $\chi_{e,HP}$ and $\chi_{e,FB}$, the ratio ranging between 0.8 and 1.6 (Tab.1 and Fig.1) for all discharge.
types excepting one characterized by high edge safety factor ($q_a = 4.5$) and low density ($\bar{n} = 6.4 \cdot 10^{19}$), for which $\chi_e^{HF}/\chi_e^{PB} = 2.6$.

Data from FTU have been collected at high $q_a$ and low $\bar{n}$; power balance calculations give $\chi_e^{PB} up to four times smaller than $\chi_e^{HF}$, but they are still in a preliminary stage. $\chi_e^{HP}$ shows an inverse dependence on $\bar{n}$ and is well aligned with data from FT at $q_a > 3$ (Fig. 2).

**Figure 1.** +) Raw experimental values of $\chi_e^{HF}/\chi_e^{PB}$ from the ttp method. △ Values corrected with formula (4) when appropriate.

**Figure 2.** $\chi_e^{HF}$ vs. line averaged density from FT and FTU.

**LIMITS OF THE TTP METHOD FOR FT AND FTU.** During the relaxation following a sawtooth crash, electron temperature at $r > r_m$ reaches a maximum in a finite time $t_p(r)$; it is well known [1] that $\chi_e^{HF}$ can be inferred from the slope of $t_p$ vs. $r^2$:

$$\chi_e^{HF} = \Delta r^2 / 9 \Delta t_p$$

(1)

The main conditions underlying expression (1) are: 1) the sawtooth crash only gives prompt heat deposition for $r < r_m$; 2) there is no electron-ion heat exchange; 3) the perturbation to the local ohmic power input is negligible. 4) $\chi_e$ is independent from temperature and temperature gradient.

The violation of condition 1) would give a strong overestimation of $\chi_e^{HF}$ as observed on TFTR [4]; we analyzed some FTU shots at high time resolution and we did not find any evidence of such a behavior. Conditions 2) and 3) have been examined together in the past, assuming that both give rise to a dissipative term in the perturbed heat transport equation [5]; it turned out that eq. (1) can overvalue $\chi_e^{HF}$. Finally if condition 4) is violated the effect of terms associated to the $\chi_e$ perturbations should be accounted for; this will be discussed later.

In order to check expr. (1) when conditions 2) to 4) are not satisfied, we solved numerically the coupled equations for electron ($\tilde{T}_e$) and ion ($\tilde{T}_i$) temperature perturbations:
in which $t_{ex} = m_i \tau_{ei}/3m_e$, $t_{OH} = nT_e/\eta_j^2$ and $A' = \partial A/\partial r$. We have considered a deuterium plasma with $Z_{eff} = 1$. Considering first the case with constant coefficients, it can be shown that if $\tau = t_{ex}/t_p \ll 1$, electrons and ions behave as a single fluid with a heat diffusivity $\chi_e^{HP} = (\chi_{e+} + \chi_i)/2$, so that if $\chi_i \ll \chi_e$, then $\chi_e^{HP} < \chi_e$.

In order to assess the importance of ion temperature perturbations, we solved (2) and (3) keeping only the first two terms in the rhs. Initial conditions for $T_e$ followed the Kadomtsev sawtooth model, while two different i.c. for ions were considered: a) $\tilde{T}_i(r,0) - T_e^0(r,0)$; b) $\tilde{T}_i(r,0) = 0$. Expression (1) was then applied to the solution $\tilde{T}_e(r,t)$; it turned out that for $\tau > 1$, expr (1) is accurate to within 3%, while for $\tau < 0.2$ the asymptotic condition holds; for intermediate $\tau$ the result depends on the i.c. on ions (Fig. 3): in particular assuming cond. b) for $\tau >> 1$, the results are qualitatively similar to the ones from /5/, while for smaller $\tau$ the deviation is on the opposite side. A simple correction formula valid for $\tau > 1$ is:

$$\chi_e = \chi_e^{HP} \left[ \frac{2}{8} t_{ex} \chi_e^{HP} + \left[ 1 + \left( \frac{2}{8} t_{ex} \chi_e^{HP} \right)^2 \right]^{1/2} \right]^{-1}$$

This equation can be easily generalized in order to include the ohmic power perturbation replacing $t_{ex}^{-1}$ by $td^{-1} = t_{ex}^{-1} + t_{OH}^{-1}$.

**Figure 3.** $\chi_e^{HP}/\chi_e^{PB} : \Delta$ for i.c. a); +) for i.c. b); *) from JETTO code.

**Figure 4** Effect of an explicit temperature dependence in $\chi_e$ with exponent $\alpha_1$. 

$$\frac{3}{2} n \frac{\partial T_e}{\partial t} = \frac{1}{r} \left( n \chi_e \frac{\partial^2 T_e}{\partial r^2} - \frac{3}{2} \left( \frac{1}{t_{CH}} - \frac{T_e - T_i}{T_e t_{ex}} \right) \frac{\partial T_e}{\partial T_e} - \frac{1}{r} \left( n \frac{\partial^2 T_e}{\partial T_e^2} + n \frac{\partial \chi_e}{\partial T_e} \frac{\partial T_e}{\partial T_e} \right) \frac{\partial T_e}{\partial T_e} \right)$$

$$\frac{3}{2} n \frac{\partial T_i}{\partial t} = \frac{1}{r} \left( n \chi_i \frac{\partial^2 T_i}{\partial r^2} - \frac{3}{2} \frac{T_e - T_i}{T_e t_{ex}} - \frac{3}{2} \frac{T_e - T_i}{T_e t_{ex}} \right)$$

(2)
The effect of variable coefficients (with $\chi_e$ following the Neo- Alcator law), of the third term on the rhs of (2) and (3), and of realistic initial conditions was investigated by using the JETTO code /6/; the results are shown in fig. 3.

The minimum $\tau$ value reached in FT was 0.7, so we conclude from Fig. 2 that the accuracy of the ttp method was to within 30%.

EFFECT OF TEMPERATURE AND TEMPERATURE GRADIENT. The effect of a dependence of $\chi_e$ on $T_e$ and $T_e'$ has been studied for a low density case in which the effect of the ohmic and exchange power is negligible. For the sake of simplicity we have assumed a model ($\chi_e = c_0 \xi^4 T_e^2 (T_e')^2 a^2$ with $q$ being the safety factor) which reproduces the most important features of microinstabilities based models. It can be shown that due to the temperature gradient dependence, $\chi_e^{HP}$ overestimates $\chi_e$ by a factor $(1 + \alpha_2)$, while the $T_e$ dependence produces the convective term in Eq. (2).

Assuming a narrow propagation region, such a term yields to a correction similar to Eq. (4) with $t_\mu$ replaced by $t_\mu$ with:

$$t_\mu^{-1} = \alpha_1 \left[ \frac{P_{OH} - P_{ei}}{n T_e} + \chi_e \left( \frac{T_e'}{T_e} \right)^2 \right]$$

(5)

$$t_\mu^{-1} = \alpha_1 \left[ \frac{P_{OH}}{n T_e} + \chi_e + \chi_d \left( \frac{T_e'}{T_e} \right)^2 \right]$$

(6)

where Eq. (5) is suited for low density ($T_e >> T_i$) and Eq. (6) for high density ($T_e = T_i$). The effect of varying the exponent $\alpha_1$ is shown in Fig. 4. Note that, upon increasing $\alpha_2$, the equilibrium temperature profile becomes broader and the correction associated to $t_\mu$ becomes smaller.

An analysis of the heat pulse propagation has been also made for a more general microinstability based model /7/. In such a model the electron heat flux is the sum of an outward convective term and of a diffusive term. It can be shown that the temperature gradient dependence is associated to the ratio between the convective and the diffusiv terms and corresponds to a moderately negative exponent $\alpha_2$, except for the case of flat density profile and high $q$ (peaked $T_e$ profiles), for which $|\alpha_2| < 1$.

Therefore while $T_e$ and $T_e'$ dependences tend to compensate each other in most of the cases (yielding a ratio $\chi_e^{HP}/\chi_e^{FB}$ close to one), at high $q$ and flat density discharges $\chi_e^{HP}/\chi_e^{FB}$ turns out to be a factor 2-2.5 for $\alpha_1 = 1 - 1.5$

CONCLUSIONS The electron heat diffusivity as measured from heat-pulse propagation agrees with the one from steady-state power balance in almost all of the FT operation space. The agreement is lost at low density and high $q$, in both FT and FTU; these results are consistent with a microinstability based transport model.

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A STUDY OF THE ION SPECIES DEPENDENCE OF $\chi_e$
BY HEAT PULSE PROPAGATION

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1. INTRODUCTION

An investigation of the isotope dependence of $\chi_e$ on ASDEX revealed that the values of $\chi_e$ in hydrogen and deuterium were the same within the limits of experimental accuracy [1]. This study in hydrogen, deuterium and helium has been continued on TEXTOR. The 11 channel ECE diagnostic measures the temperature perturbations generated by sawtooth crashes in an ohmically heated plasma. Averaging over the one second flat top phase improves the signal to noise ratio to the extent that differences in the radial profile of $\chi_e$ are able to be inferred. Even though the values of $\chi_e$ found in each of the three gases are greater than the values calculated from power balance [2,3], the basic relationship between the energy confinement time and the value of $\chi_e$ deduced by heat pulse propagation can still be explored.

2. EXPERIMENTAL RESULTS

Shown in Fig. 1 is the density dependence of the energy confinement time in hydrogen and deuterium in the TEXTOR tokamak. The saturation of the energy confinement time with increasing density in deuterium and much better confinement in deuterium plasmas has been of interest to experimentalists [4,5] as theory predicts a deterioration in confinement with an increase in isotope mass. Using heat pulse propagation analysis to investigate the role of electron heat transport thus complements this activity.

Numerical modelling of heat pulse propagation with the forced boundary value method [6] uses the temperature perturbations generated by a sawtooth crash measured by the innermost channel of the ECE diagnostic outside the mixing radius as the driving term. Typically 2 channels on the inside and three channels on the outside of the plasma major radius are available for analysis. The radial positions of the ECE channels on the outside of the plasma major radius were typically located at $r = 0.47a$, $r = 0.58a$ and $r = 0.72a$. The radial profile of $\chi_e$ is the free parameter to be found with the time evolution of the temperature perturbations being the constraint to the best fit procedure. Shown in Fig. 2 are two radial profiles of $\chi_e$ that are representative of the class of profiles considered. For the flat profile of $\chi_e$, it is found that the predictions for the outermost channel at $r = 0.72a$, in particular the decay of the temperature perturbation, is only marginally effected the magnitude of $\chi_e$ at $r = a$. In deuterium, hydrogen and helium the relative temperature perturbation, $\delta T/T$ for a density of $\bar{n}_e = 1.8 - 2.0 \times 10^{19} \text{m}^{-3}$ at $r = 0.5a$ was 0.05, 0.03 and 0.04 respectively.
In Fig. 3 the sensitivity of the analysis is demonstrated. Using the innermost channel as the forced boundary, then the calculated value of the temperature perturbation (dashed line) and the measured temperature perturbation (solid line) can be compared at the outer radial positions for various radial profiles of $\chi_r$. For this deuterium discharge at low density ($\bar{n}_e = 1.4 \times 10^{19} \text{m}^{-3}$), it is clear that a flatter profile of $\chi_r$ is necessary to obtain a good fit. In contrast it is clearly shown in Fig. 4 that the opposite is the case for a discharge in hydrogen. Here, a broad radial profile of $\chi_r$ that markedly rises with minor radius is necessary to obtain a good fit. Consistent observations for the inner and outer set of channels outside the mixing radius are found.

At low densities where the energy confinement time increases almost linearly with density, it is found that the flat profiles of $\chi_r$ yield the best fit to the measured temperature perturbations. In deuterium, the value of $\chi_r$ at $r = 0.5a$ decreases from $4.0 \text{m}^2/\text{s}$ to $3.0 \text{m}^2/\text{s}$ as the line averaged density is increased from $\bar{n}_e = 1.4 \times 10^{19} \text{m}^{-3}$ to $\bar{n}_e = 1.8 \times 10^{19} \text{m}^{-3}$ and the energy confinement time increases from $57 \text{ms}$ to $76 \text{ms}$. A similar trend in helium at low densities is also found, with a decrease in $\chi_r$ at $r = 0.5a$ as the density is increased and a flat profile of $\chi_r$ giving the best fit to the measured temperature perturbations. As the density is increased in deuterium to $\bar{n}_e = 2.7 \times 10^{19} \text{m}^{-3}$, increasingly broader profiles of $\chi_r$ are needed to fit the measured temperature perturbations, with a value of $\chi_r = 3.0 \text{m}^2/\text{s}$ at $r = 0.5a$.

Over the density range analysed so far ($\bar{n}_e \geq 2.0 \times 10^{19} \text{m}^{-3}$), broad profiles of $\chi_r$, with values of $3.0 \text{m}^2/\text{s}$ at $r = 0.5a$, were needed to fit the temperature perturbations in hydrogen. Further experiments at lower densities in hydrogen are planned.

It is worth noting that the relative temperature perturbations in deuterium are larger than those in hydrogen and therefore the perturbation in the temperature gradient is larger. Averaging a millisecond before and after the crash in discharges at $2.0 \times 10^{19} \text{m}^{-3}$ it is found that the temperature gradient rises by 6% and 4% respectively with an absolute value before the crash of $26 \text{eV/cm}$ and $17 \text{eV/cm}$ respectively. The stronger perturbation of the temperature gradient in deuterium discharges could be the reason that the lower values of $\chi_r$ for deuterium discharges obtained from global power balance are not found.

Despite the stronger increase in temperature gradient in deuterium it is observed that a broad profile in $\chi_r$ and therefore larger values of $\chi_r$ at the plasma boundary are deduced for the hydrogen discharges. It is accepted that transient transport analysis yields values above that calculated by power balance. In this case however the situation is that the perturbation with lower amplitude yields higher values of $\chi_r$ at the plasma boundary. Such a simple comparison in this case is complicated by the higher frequency of the sawtooth oscillation in the hydrogen discharge. Recent results from ECRH modulation experiments suggest that $\chi_r$ deduced from transient transport analysis may be a function of both the frequency and amplitude of the perturbation [7].

When comparing deuterium discharges with low density and discharges where the energy confinement time saturates again the relative temperature perturbation decreases and in addition the sawtooth frequency decreases. Despite the tendency towards milder plasma perturbations, the transition from flat $\chi_r$ profiles to broad $\chi_r$ profiles is observed. Such deductions from transient transport analysis must be necessarily cautious given that comparisons of absolute amplitude are complicated by the assumptions used which are expected to give $\chi_r$ values larger than that calculated by power balance. However the indications are that a comparison of discharges under conditions where the absolute amplitude would be expected to be less effected by the perturbation still leads to the deduction of an increase in the absolute value of $\chi_r$ at the plasma boundary.

In the X mode the optical depth is approximately proportional to the product of the electron
density and temperature. For a given ion species this results in the small variation of optical depth as the density range is scanned. In any case, variations of optical depth would lead to errors in the assigned temperature and consequently to false estimates in the amplitude of the temperature perturbation rather than the observed change in propagation speed of the perturbations which are then interpreted as being due to broader $\chi_e$ profiles.

3. CONCLUSION

Experimental determination of the radial profile of $\chi_e$ by heat propagation analysis would suggest that the increased values of $\chi_e$ at the plasma edge may be the governing factor in the saturation of the energy confinement time with increasing density seen in deuterium and would also suggest that a deterioration in electron confinement at the plasma boundary is the cause for the difference in energy confinement time between hydrogen and deuterium discharges.

References

Fig. 3 A comparison of predicted and measured time evolution of the temperature perturbation for flat (left) and broad (right) $\chi_e$ profiles in a deuterium discharge at $n_e = 1.4 \times 10^{19} \text{ m}^{-3}$, with $\chi_e (r=0.4a) = 4.0 \pm 0.4 \text{ m}^2/\text{s}$.

Fig. 4 A comparison of predicted and measured time evolution of the temperature perturbation for broad (left) and flat (right) $\chi_e$ profiles in a hydrogen discharge at $n_e = 2.0 \times 10^{19} \text{ m}^{-3}$, with $\chi_e (r=0.4a) = 3.0 \pm 0.3 \text{ m}^2/\text{s}$.
COMPARISON OF ANOMALOUS MOMENTUM TRANSPORT WITH PARTICLE AND ENERGY TRANSPORT ON ASDEX


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1 Introduction

The experimentally observed anomalous fluxes of energy, angular momentum and particles in a tokamak still lack of a quantitative theoretical description. A possible explanation lies in the imperfect description of nonlinear transport phenomena by quasi-linear theories. While the calculation of absolute values of local transport coefficients fails, the comparison of predicted and measured parameter dependences and relations among different transport coefficients may help to identify the nature of the underlying physical mechanisms. In this paper, we investigate the momentum diffusivity and especially its relation to the diffusion coefficient of impurity ions.

2 Experimental results

2.1 Momentum transport

The radially averaged momentum diffusivity \( \chi_\phi \) \((8 < r < 34 \text{cm})\) for neutral-beam-heated ASDEX L-mode discharges can be written as

\[
\chi_\phi = 1.7 \cdot P_{\text{tot}}^{0.65} \cdot \bar{n}_e^{-0.73} \cdot A_{\text{eff}}^{-0.8} \cdot I_p^{-0.65} \quad (\text{m}^2/\text{s}, \text{MW}, 10^{19} \text{m}^{-3}, \text{amu}, \text{MA}),
\]

\text{rmse}=12\%.

where \( A_{\text{eff}} \) is the mean plasma mass per electron. The search for a simple expression for the local momentum diffusivity (the spatial resolution of the velocity measurement, 5 points/minor radius, is quite low) led to the expression (regression analysis of \( \chi_\phi(r = a/4, a/2, 3/4a) \))

\[
\chi_\phi = 5 \cdot P_{\text{tot}}^{0.63} \cdot (\bar{n}_e \cdot A_{\text{eff}})^{-0.84} \cdot (q/B_t)^{0.8} \quad (\text{m}^2/\text{s}, \text{MW}, 10^{19} \text{m}^{-3}, \text{amu}, T),
\]

\text{rmse}=18\%.
The radial as well as the current dependence are quite well reproduced by the term \( \frac{r}{B_i} \) which corresponds to the poloidal magnetic field via \( B_p = \frac{r}{B_i} B_i \).

2.2 Comparison to particle transport

Radially averaged diffusion coefficients for various impurity species were determined by harmonic analysis of gas oscillation experiments\(^2\). Fig. 1 shows a comparison of \( D(S^{14+}) \) with \( \chi_\phi \) from Eq. 1 as a function of heating power. Good coincidence is found for the absolute value as well as the power dependence. The density dependence of the diffusion coefficient of various impurity species in ohmic discharges is compared to the prediction of Eq. 1 in Fig. 2. Again, good coincidence is found, independent from charge and mass of the impurity species. Momentum and ion diffusion seem to underly the same physical mechanism. The electron diffusion coefficient, obtained from gas oscillation measurements\(^3\) is about a factor of 2 smaller than the momentum diffusivity, but exhibits similar parameter dependences.

While the neoclassical diffusion can be neglected in the comparison of Figs. 1 and 2 (\( D_{neo} < \chi_\phi /10 \)), it is interesting to check the relevancy of the neoclassical inward drift velocities of impurity ions. Under the assumption that \( D_{imp}(r) \) is equal to the momentum diffusivity \( \chi_\phi(r) \), we have calculated \( v_{imp} \) of \( O^{8+} \) for steady state conditions of three L-mode discharges with different values of \( I_p \) and \( B_t \) via

\[
v_{imp} = D_{imp} \cdot \frac{\partial n_{imp}}{\partial r} \cdot \frac{1}{n_{imp}} .
\]  

(3)

Profiles of \( \chi_\phi(r) \) obtained from measured rotation velocities are shown in Fig. 3. Density profiles of fully stripped oxygen are taken from CXR spectroscopy and bremsstrahlung measurements (see Fig. 4). The drift velocities obtained are well above, but of the same magnitude as the neoclassical values (Fig. 5).

2.3 Comparison to energy transport

Coincidences of momentum and energy transport are found in a comparison of global confinement times\(^1\), which are always equal within a factor of 2, and in local transport studies using the TRANSP code. While the effective thermal and momentum diffusivities are about equal in the transport region, the latter takes lower values in the central part of the plasma. A possible explanation may be the lower neoclassical value of \( \chi_\phi \) in comparison with \( \chi_t \) and the smaller influx of sawtooth activity on momentum transport. However, owing to the large experimental uncertainties, more data have to be analyzed to obtain systematic results.
3 Conclusion
The comparison of measured diffusion coefficients of impurity ions with an empirical scaling law for the radially averaged momentum (L-mode) diffusivity showed agreement in magnitude and experimental parameter dependences. Both transport processes can be attributed to an anomalous transport mechanism: The neoclassical momentum diffusivity is generally (with the exception of some controversy concerning the gyroviscous contribution) believed to be small, and the measured diffusion coefficients of impurity ions are more than an order of magnitude higher than the neoclassical values. Radial drift velocities of O\(^{8+}\) ions were calculated under the assumption that even the local values of \(\chi_\delta\) and \(D_{\text{imp}}\) are equal. The drift velocities obtained are higher, but of the same order of magnitude as the inward drift velocities calculated from neoclassical theory. From the comparison, it can be concluded that neoclassical effects may influence the impurity transport even in L-mode plasmas where the diffusion is completely dominated by anomalous effects.

References
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Figure 1  Power dependence of the diffusion coefficient of S\(^{14+}\) ions from gas oscillation experiments\(^2\) in comparison to the prediction of the momentum diffusivity scaling Eq. 1 (solid line).

Figure 2  Density dependence of \(D_{\text{imp}}\) for various impurity species in ohmic discharges in comparison to the prediction of Eq. 1 (solid line).
Figure 3  Measured momentum diffusivities (assumed to be equal to $D_{\text{imp}}$) for three discharges with different experimental parameters.

\[ q_a = 0.484 \frac{B(T)}{I(\text{MA})}, \ P_{\text{tot}} \approx 1.2 \text{MW}, \ \bar{n}_e = 2.7 \times 10^{19} \text{m}^{-3}. \]

Figure 4  Profiles of $O^{8+}$ ions from CXR spectroscopy (fat dots) and bremsstrahlung measurements (solid lines) using the C/O ratio from CXRS. ($\Delta Z_{\text{eff}}(O) \approx 2 \cdot \Delta Z_{\text{eff}}(C)$). The $O^{8+}$ profiles derived from the bremsstrahlung are used for the calculations due to their better spatial resolution.

Figure 5  Radial drift velocities of $O^{8+}$ calculated for the conditions of Figs. 3 and 4 using Eq. 3 in comparison to the neoclassical values. The 'measured' values are higher than the neoclassical inward drift, especially in the low-$q_a$ case, but of the same order of magnitude.
STATISTICAL ANALYSES OF LOCAL TRANSPORT COEFFICIENTS IN OHMIC ASDEX DISCHARGES


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1. Introduction
Tokamak energy transport is still an unsolved problem. Many theoretical models have been developed, which try to explain the anomalous high energy-transport coefficients. Up to now these models have been applied to global plasma parameters. A comparison of transport coefficients with global confinement time is only conclusive if the transport is dominated by one process across the plasma diameter. This, however, is not the case in most of the Ohmic confinement regimes, where at least three different transport mechanisms play an important role. Sawtooth activity leads to an increase in energy transport in the plasma centre. In the intermediate region turbulent transport is expected. Candidates here are drift waves [1] and resistive fluid turbulences [2]. At the edge, ballooning modes or rippling modes [3] could dominate the transport. For the intermediate region, one can deduce theoretical scaling laws for $\tau_E$ from turbulent theories. Predicted scalings [4] reproduce the experimentally found density dependence of $\tau_E$ in the linear Ohmic confinement regime (LOC) and the saturated regime (SOC), but they do not show the correct dependence on the isotope mass [5]. The relevance of these transport theories can only be tested in comparing them to experimental local transport coefficients. To this purpose we have performed transport calculations on more than a hundred Ohmic ASDEX discharges. By Principal Component Analysis we determine the dimensionless components which dominate the transport coefficients and we compare the results to the predictions of various theories.

2. Experimental Data Base
For all discharges measured radial profiles for both electrons and ions are available. The density and temperature of the electrons are measured with a Thomson scattering laser system. The ion density profile is calculated from the bremsstrahlung $Z_{eff}$ profile and the quasineutrality equation under the assumption that the main impurity components are carbon and oxygen. This assumption is well fulfilled for Ohmic discharges with line averaged densities above $1.5 - 2 \cdot 10^{13} \text{cm}^{-3}$. A neutral particle analyzer with active beam technique is used to determine the local ion temperature from the measured charge exchange fluxes.

The current profile, which gives the local Ohmic power input onto the electrons, is calculated from neoclassical resistivity including the neoclassical bootstrap current. In the energy balance, profiles for radiative losses of the electrons, ionization and charge-exchange losses, treated by a Monte Carlo code, are included. All data have been taken in stationary plateaus and averaged over a few hundred ms. Local and global results of experiments and transport calculations are stored in a data bank and used for statistical analyses.

3. Scaling of Global Confinement Time $\tau_E$
Using this data bank we obtained scaling laws for the global confinement time $\tau_E$. The linear confinement regime, with $\bar{n} < 3 \cdot 10^{13} \text{cm}^{-3}$, and the saturated regime were
fitted separately. In these regimes \( \tau_E \) scales as

\[
\tau_E^{\text{LOC}} \sim \frac{n}{0.81} \cdot I_p^{-0.65} \cdot B_t^{-0.22} \cdot A_i^{0.48} \cdot \left( \frac{Z_{\text{eff}}(0)}{I_p} \right)^{0.54} \]

\[
\tau_E^{\text{SOC}} \sim \frac{n}{0.10} \cdot I_p^{-0.10} \cdot B_t^{-0.25} \cdot A_i^{0.55} \cdot Z_{\text{eff}}^{-0.10}
\]

where \( n \) is the line averaged density, \( B_t \) the toroidal field, \( I_p \) the plasma current, \( A_i \) the isotope mass and \( Z_{\text{eff}}(0) \) the central value of the effective charge. Since in the LOC-regime \( Z_{\text{eff}} \) varies strongly with \( n \) and \( I_p \) we used the principal component \( \left( \frac{Z_{\text{eff}}(0)}{I_p} \right)^{0.54} \), which roughly corresponds to the impurity content. In the SOC \( Z_{\text{eff}}(0) \) does not depend on other parameters and is used without changes. In both regimes \( \tau_E \) depends on the isotope mass with similar exponents opposite to predicted values from drift wave theories.

4. Linear Ohmic Confinement LOC

At low densities, electron and ion energy transport can be separated. In Fig.1. we compare measured \( T_i \) to predictions of neoclassical theory, using multipliers for the ion heat conductivity, \( \chi_i = \chi_i^F \cdot \chi_i^{\text{neo}}, \chi_i^F = 1 \) to 2. In the linear regime the ratio between experimental and predicted values is close to unity for \( \chi_i^F = 1 \). Higher than neoclassical transport can be excluded.

As for the electron heat transport a large number of interfering theories exists, statistical analyses of the local \( \chi_e \) of the intermediate region have been performed to get the dominating parameters. At least five different radii have been used for each discharge. Following the idea that the heat diffusivities can be described by a leading term with the correct dimension and a function of dimensionless local quantities, models of the form

\[
\chi_e(r) \sim D_B(r) \cdot \nu(r)^* \cdot \beta_p(r)^{a_1} \cdot \rho(r)^{a_2} \cdot q(r)^{a_3} \cdot \epsilon(r)^{a_4} \cdot A_i^{a_5} \cdot Z_{\text{eff}}(r)^{a_6}
\]

are used, where \( D_B \) is the Bohm diffusion coefficient, \( \nu^* \) the electron collisionality, \( \beta_p \) the poloidal beta, \( \rho^* = \frac{\rho}{R_w} \) the normalized gyroradius at the sound velocity, \( q \) the safety factor and \( \epsilon \) the inverse aspect ratio. In the intermediate region an analysis of the principal components leads to the result, that only four parameters are needed to fit the database. The other parameters can be cancelled or their exponents can be fixed at a certain value, except the electron collisionality and the isotope mass, which are the dominating parameters. Omitting the collisionality or the isotope mass gives unfavourable fits. Fixing an exponent of another variable only leads to a rearrangement of the rest of the exponents. Following the 'gyro-reduced' formalism with the leading term \( D_B \cdot \rho^* \) [6] the exponent of \( \rho^* \) is chosen to unity. The best statistical fit for the intermediate region is, see Fig.2.,

\[
\chi_e \sim D_B \cdot \rho^* \cdot \nu^* \cdot 0.75 \pm 0.06 \cdot \beta_p \cdot 0.83 \pm 0.07 \cdot A_i \cdot 0.86 \pm 0.07 \cdot Z_{\text{eff}} \cdot 0.92 \pm 0.10 \quad R = 0.98
\]

As the poloidal beta is strongly correlated in the LOC regime to \( \epsilon, A_i \) and \( Z_{\text{eff}} \), an alternative fit with similar quality using \( \epsilon \) instead of \( \beta_p \) is

\[
\chi_e \sim D_B \cdot \rho^* \cdot \nu^* \cdot 0.90 \pm 0.06 \cdot \epsilon \cdot 1.62 \pm 0.10 \cdot A_i \cdot 1.19 \pm 0.08 \cdot Z_{\text{eff}} \cdot 0.54 \pm 0.11 \quad R = 0.97
\]
Including more than four varying parameters does not improve these fits furthermore. Comparison of these scalings to theoretical relations shows, that only two candidates are compatible with the experimental results (Tab. 1). Resistive fluid turbulences with 

\[ \chi_e \sim D_B \cdot \rho^* \cdot \nu^* \cdot \beta_p^{-1} \]

have similar dependences on \( \nu^* \) and \( \beta_p \). The theory coming next to the second formula is the collisional drift wave theory with 

\[ \chi_e \sim D_B \cdot \rho^* \cdot \nu^* \]

and without any \( \beta_p \)-dependence. The isotope dependence is wrong for these models, as it is for all used turbulent theories. Due to the coupling of the parameters in this regime a single leading transport mechanism cannot be deduced from the fits above.

5. Saturated Ohmic Confinement SOC

For higher densities (SOC), electrons and ions are strongly coupled and \( T_i \) is close to \( T_e \) over the whole cross-section. The error bars are too large to separate the two channels. Although a comparison of experimental data and predicted ion temperatures in Fig. 1, shows good agreement with neoclassical transport. However, higher ion energy transport is possible, owing, for example, to the presence of the \( \eta_i \) mode. In this domain we analyze \( \chi_{e+i} = \chi_e + \chi_i \) with the motivation that it reflects the characteristics of the dominant loss channel. The same method as above gives the following least squares fit of the added heat conductivities in the intermediate region

\[ \chi_{e+i} \sim D_B \cdot \rho^* \cdot \nu^* \cdot 0.71 \pm 0.06 \cdot \beta_p^{-0.83 \pm 0.03} \cdot A_i^{-0.67 \pm 0.05} \cdot Z_{eff}^{-0.72 \pm 0.10} , \quad R = 0.97 . \]

An analysis of \( \chi_e \) under the assumption of neoclassical ion transport leads to a very similar relation. As it can be seen by principal component analysis, the poloidal beta is one of the dominant parameters. Replacing \( \nu^* \), \( \beta_p \) or \( A_i \) only creates unfavourable fits. Comparing this fit to our result for the LOC shows that they are only slightly different with small deviations for \( A_i \) and \( Z_{eff} \). Here resistive fluid turbulences agree, except for the isotope dependence.

6. Summary

The results of these analyses show that the ion heat transport can be described by neoclassical theory in the LOC regime. The electron heat conduction shows a negative isotope effect, giving the positive exponent in the \( \tau_E \)-scaling. Statistical analyses show that in both regimes resistive fluid turbulences with an inverse \( \beta_p \)-dependence could play an important role in the intermediate confinement region. Trapped electron modes, predicting zero or negative exponents for \( \nu^* \), seem to be more unlikely in both regimes. In the LOC also collisional drift waves may contribute to the electron heat transport. Comparing the results of the LOC and SOC regimes gives an indication, that the electron transport mechanism does not change from low to high densities, but currently it is not yet clear which mechanism causes the saturation in the energy confinement time.

References
Fig. 1: Ratio of experimental $T_i$ values to predicted temperatures from neoclassical theory at $r = 0$ cm and $r = 20$ cm; calculated for $\chi^{F}_i = 1$ (circles) and 2 (crosses).

Fig. 2: Fit of experimental electron heat conductivity for both regimes. Fit parameters are $\nu^*$, $\beta_p$, $A_i$ and $Z_{\text{eff}}$ as described in Chap.4 and 5.

Tab. 1: Scaling of the electron heat conductivity with the dimensionless parameters $\rho^*$ ('gyro-reduced' formalism), $\nu^*$ and $\beta_p$, according to various transport theories and to statistical analyses.

<table>
<thead>
<tr>
<th>Theories and statistical results</th>
<th>$X_e \sim$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Collisionless drift wave</td>
<td>$D_B \rho^* \nu^{+0.0} \beta_p^{-0.0}$</td>
</tr>
<tr>
<td>Collisional drift wave</td>
<td>$D_B \rho^* \nu^{+1.0} \beta_p^{-0.0}$</td>
</tr>
<tr>
<td>Collisionless trapped electron mode</td>
<td>$D_B \rho^* \nu^{+0.0} \beta_p^{-0.0}$</td>
</tr>
<tr>
<td>Dissipative trapped electron mode</td>
<td>$D_B \rho^* \nu^{+1.0} \beta_p^{-1.0}$</td>
</tr>
<tr>
<td>Resistive fluid turbulence</td>
<td>$D_B \rho^* \nu^{+1.0} \beta_p^{+1.0}$</td>
</tr>
<tr>
<td>Resistive ballooning mode</td>
<td>$D_B \rho^* \nu^{+0.75} \beta_p^{-0.83}$</td>
</tr>
<tr>
<td>Least squares fit for the LOC regime</td>
<td>$D_B \rho^* \nu^{+0.9} \beta_p^{-0.0}$</td>
</tr>
<tr>
<td>Second fit for the LOC regime</td>
<td>$D_B \rho^* \nu^{+0.71} \beta_p^{-0.83}$</td>
</tr>
<tr>
<td>Least squares fit for the SOC regime</td>
<td>$D_B \rho^* \nu^{+0.71} \beta_p^{-0.83}$</td>
</tr>
</tbody>
</table>
Reassessment of the interpretation of sawtooth induced pulse propagation

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Introduction

Since it was realised in 1977 [1] that the heat pulse induced by the sawtooth collapse can be employed to measure the electron thermal diffusivity (χ), the interpretation of the results has been the subject of discussion. This discussion was kindled by the discrepancy between the heat pulse deduced χ_{HP} and the value obtained from power balance analysis (χ_{PB}). Many different ideas have been brought forward to explain this discrepancy. These involve: i) the technique of deriving χ_{HP} from the data; ii) the basic difference between a dynamic (heat pulse) experiment, which yields the incremental diffusivity (χ_{inc}) and a static (power balance) evaluation, which measures the effective diffusivity (χ_{eff}); and iii) the possible effect of off-diagonal elements in the transport matrix.

It has also been suggested that after a sawtooth collapse residual turbulence enhances transport for some time. A model based on this idea has been put forward by Fredrickson et al [2] who describe the heat pulse in a high power TFTR discharge with a χ that has an explicit time dependence, found by trial and error. They conclude that this analysis proves that sawtooth induced heat pulse propagation does not give information on the normal transport processes.

In this paper we first analyse the difference between χ_{inc} and χ_{eff}, and we show that through the VT-dependence χ_{eff} has an implicit time-dependence: during the heat pulse χ_{eff} increases. We make a brief investigation into the various methods that are used to derive the diffusivities from pulse propagation, and we evaluate the effect of off-diagonal terms in the transport matrix can have on the evaluation of the incremental diffusion coefficients. Finally, we show that the hypothesis that residual turbulence due to the sawtooth collapse governs the propagation of the heat pulse can be rejected on the basis of experimental data.

Perturbative vs Steady state analysis

A power balance analysis evaluates fluxes and gradients, e.g. the effective electron thermal diffusivity is determined as the ratio of the electron heat flux (q_{e}) and the product of the electron density (n_{e}) and temperature gradient (VT_{e}): χ_{eff} = q_{e}/(n_{e}VT_{e}). q_{e} represents the total heat flux carried by the electrons, in general it is not possible to distinguish convective and diffusive contributions to the flux. Perturbative experiments, on the other hand, yield an evaluation of the incremental diffusivity χ_{inc} = ∂q_{e}/∂(n_{e}VT_{e}) [3]. This basic difference is illustrated in Fig.1.

Fig.1 If the heat flux is not proportional to nVT, χ_{inc} ≠ χ_{eff} and χ_{eff} varies during the heat pulse.
If $q_e$ is a nonlinear function of $nVT$, which is generally the case, it can be linearised near the equilibrium value by separating a diffusive and a convective flux:

$$\dot{q}_e = n\chi_o VT + q_{\text{conv}}$$

yielding

$$\chi^{\text{eff}} = \chi_o + \frac{q_{\text{conv}}}{nVT},$$

$$\chi^{\text{inc}} = \chi_o.$$

It is clear that the perturbative experiment evaluates only the diffusive flux, and that the effective diffusivity is a function of VT.

Local power balance analysis in a series of discharges with different power levels can yield $q_e$ as a function of $nVT$. Studies in JET show that $\chi^{\text{inc}}$ exceeds $\chi^{\text{eff}}$ by a factor of up to 5 in ohmic plasmas to less than 2 for high power [3,5].

The dependence shown in Fig. 1 implies that $\chi^{\text{eff}}$ increases during the heat pulse, in first order proportional to the change of the local VT. Fig.2 shows that the evolution of VT is similar in shape to the heat pulse and travels slightly ahead of it. The variation of VT can be much larger than the variation of T itself, due to the short gradient length of the perturbation.

![Graph showing variation of VT and temperature change](image)

Many authors have attempted to find $\chi^{\text{HP}}=\chi^{\text{PB}}$ for heat pulse propagation. Yet, the first expectation should be that $\chi^{\text{HP}}$ is an estimate of $\chi^{\text{inc}}$, whereas $\chi^{\text{PB}}$ is an evaluation of $\chi^{\text{eff}}$. As shown above, these cannot be equated, and the experimental result $\chi^{\text{inc}} > \chi^{\text{eff}}$ is consistent for all known methods to evaluate local transport. A different question is how accurate the experimental determination of $\chi^{\text{inc}}$ using heat pulse propagation is. Below we discuss sources of systematic errors in the evaluation of $\chi^{\text{inc}}$.

**Analysis Methods**

The analysis of heat and density pulses assumes that the fluxes can be linearized in the gradients near the equilibrium values. Originally the sawtooth heat pulse was treated as an initial value problem, and it is questionable if the linearization is correct in the mixing region. In principle it is more correct to treat the heat pulse as a boundary value problem, taking measured temperature evolutions near the mixing radius and the limiter as the boundaries [6]. However, it is generally found that with the initial value solution a very good fit to the $T_e(t)$ just outside the mixing radius can be obtained, so that both methods outside of this radius become practically identical. The reason is that the recovery of the temperature in the central part of the profile is a process that is slow compared to the heat pulse: in the central region the change of VT is but small and linearisation is justified. However, the value found for $\chi^{\text{inc}}$ in the centre is for the situation with very low VT. At JET, this value is found to be low, typically $\chi^{\text{inc}} < 0.5 \text{ m}^2/\text{s}$. This result is corroborated by modulated ICRH experiments [7].
Various different methods have been developed to analyse sawtooth heat pulses. These include the time-to-peak method [1] which uses the phase velocity of the heat pulse to determine $\chi^{HP}$, the extended time-to-peak [4] which uses both the phase velocity and the radial decay. The latter method includes the possible effect of electron-ion energy exchange and gives a local evaluation of $\chi^{HP}$. This method has been checked against full transport simulations [8]. More detailed analysis can be performed by full numerical modelling. For this purpose a comprehensive code (ACCEPT) has been developed at JET, which includes the effects of coupling between heat and density pulses [9]. Finally, by Fourier analysis of the heat pulse an evaluation of $\chi^{HP}$ can be obtained as a function of the harmonic frequency in the signal [10,11].

It has been shown for JET-data that the extended t-t-p method, the numerical simulation with the ACCEPT-code and analysis with the Fourier method all give consistent results [11]. In summary, the techniques of deriving $\chi^{HP}$ have been cross-checked where possible and found in agreement. The systematic errors in these analyses are estimated $<10\%$, the experimental error is typically 10 to 20\%. Not discussed here are systematic errors due to calibration errors in diagnostic systems, such as errors in the identification of the position of the measuring volume.

**Effect of Coupling of heat and particle transport on the determination of Incremental diffusivities.**

The coupled heat and density pulses are described by two coupled diffusion equations: $\partial u = A V^2 u$ where $u = (n/n_0, T/T_0)$ and $A$ is a 2x2 matrix [4,12]. If $A$ is diagonal, $n$ and $T$ are eigenvectors and $D^{inc}$ and $\chi^{inc}$ can be determined directly from the measured heat and density pulses: $\chi^{HP}$ is an estimator for $\chi^{inc}$, and $D^{DP}$ is an estimator for $D^{inc}$. If $A$ is not diagonal, there are two effects. First, the shape of the heat pulse is deformed by the influence of the density pulse, and vice versa. This may introduce an error in the determination of the transport coefficients, which can be different for the different methods of evaluation. Second, the heat and density pulses are now governed by the eigenvalues of $A$, and these do not coincide with $\chi^{inc}$ and $D^{inc}$ any more. Both effects have been evaluated in a study of coupled transport in JET[12], which showed that $\chi^{HP}$ gives a systematic underestimate of $\chi^{inc}$ by 10 - 20\%, whereas $D^{DP}$ underestimates $D^{inc}$ by $\approx 30\%$. Coupling has been invoked to explain the discrepancy between $\gamma^{PB}$ and $\gamma^{HP}$ (implicitly equating $\chi^{eff}$ to $\chi^{inc}$), but as shown above, the effect of coupling on the value of $\gamma^{HP}$ is small. The same conclusion was reached with Axdes data [13]. Under some circumstances, the effect of coupling may be very important for the density pulse: in TEXT, both the heat pulse and the density pulse appear to be governed by the same eigenvalue of $A$, so that the density pulse gives an evaluation of $\chi^{inc}$ rather than $D^{inc}$ [14,15].

Summarizing, with simultaneous measurements of heat and density pulses the matrix $A$ can be resolved, and $\chi^{inc}$ and $D^{inc}$ can be derived without systematic error. If only the heat pulse is measured, coupling introduces only a small systematic error in the estimation of $\chi^{inc}$. The derivation of $D^{inc}$ from the density pulse, however, can be impossible due to coupling.

**Temporal enhancement of $\chi$ due to residual turbulence after the sawtooth crash.**

A general question not addressed so far is whether transport just after a sawtooth collapse is representative for the plasma or not. Observations such as bursts of density fluctuations and rapid loss of runaway electrons during the crash indicate that the transport may be temporarily enhanced. The time resolution of such measurements has been insufficient to determine if these phenomena coincide with the sawtooth crash, or may remain for a significant time. A careful analysis of sawtooth heat and density pulses can answer this question.
First, the values of $\chi^\text{inc}$ and $D^\text{inc}$ found from pulse propagation agree with those obtained from other perturbative experiments, such as modulated ICRH and pellet injection [16]. This rules out the possibility that a specific enhancement of the turbulence level associated with the sawtooth collapse dominates the heat and density pulses.

The same conclusion is reached by analysis of models that describe the evolution of the full $T_e$-profile. Since in general $\chi^\text{HP} > \chi^\text{PB}$, solving the transport equations with $\chi = \chi^\text{PB}$ leads to a slower propagation of the heat pulse than is observed experimentally. Therefore a temporary enhancement of $\chi$ should be introduced. A generic form of $\chi$ is given by:

$$\chi = \chi^\text{PB} \{ 1 + f(t) \exp(-t/\tau(t)) \}$$

An attempt to fit the TFTR heat pulse with this form is published in [2]. The best result was obtained with $\tau = 1\text{ms}$ and $f(t) = 150 \exp(-9.6t^2/a^2)$, but this still gives a very poor fit to the data [15]. We have varied $f(t)$ and $\tau(t)$ in a trial-and-error procedure. From our investigation we concluded that to fit the data it is necessary to have a relatively small enhancement, which travels with the heat pulse. This is precisely the implicit time dependence of $\chi^\text{eff}$ in the diffusive modelling. This is no surprise, because we already knew that the diffusive modelling gave a very good fit to the data, and implied this time-dependent $\chi^\text{eff}$. Thus by making a best fit using a generic time-dependent form for $\chi$, we recovered the behaviour that is implicit in the diffusive modelling.

**Discussion**

In principle one could maintain that the sawtooth collapse induces an enhanced diffusivity which happens to be such that the resulting heat pulse has a diffusive character. In this case, the sawtooth heat pulse would not give valuable information on transport in the plasma in general. However, this interpretation is highly unlikely, because a) other perturbative methods give the same value for the incremental transport coefficients, and b) the values found for $\chi^\text{inc}$ are corroborated by local power balance analysis.

In conclusion, we find that the sawtooth heat pulse is adequately described by linearising the heat flux near the equilibrium. No other temporary enhancement than that implicit in the diffusive model need be invoked to explain the measurements. Ballistic contributions as proposed in [2] are not supported by the data. Comparison of different perturbative experiments, and comparison with local power balance analysis, shows that heat pulse analysis does measure the incremental diffusivity characteristic for the plasma.

**Acknowledgement**

This work has profited from stimulating discussions with our colleagues at JET and FOM.

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ROLE OF BOUNDARY PLASMAS ON THE ENERGY AND PARTICLE TRANSPORT IN JT-60

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Introduction

Since the energy deposition onto the divertor in the fusion reactor is huge, we must understand the behaviors of boundary plasmas both in steady state condition and during the MHD instabilities. The scrape-off layer plays an important role on the energy transport to the divertor and the particle fuelling to the main plasma. In this paper, we discuss the roles of the scrape-off layer on the energy and particle transport in the JT-60 L-mode discharges.

Energy transport due to MHD instabilities

The heat flux measurement to the divertor has shown that the energy transport from the core plasma to the divertor plate due to MHD instabilities, such as locked modes and disruptive instabilities, is large enough to be seriously considered when a fusion is designed. In beam-heated lower-X point divertor discharges, resistive instabilities often show quasi-stationary behaviors which are referred as locked modes. It is found that this MHD instability degrades the core energy confinement significantly due to the enhanced energy loss channel to the plasma periphery when the amplitude is as large as $\delta B_\theta/B_\theta \sim 1\%$. (Fig.1) The energy transport measurement during the quasi-steady mode shows that 40% of input power is lost by the helical structure of the MHD instability which sweeps the divertor plates. The fluctuation of the heat flux density is 4 times larger than the heat flux at the steady state at the divertor plates. In the case of the termination of the discharge by the major disruption, more than half of thermal energy stored in the core plasma can be released.
around separatrix in a few milliseconds due to a resistive MHD instability. More than 300 MW/m² is released to the divertor at the energy quench and the surface temperature of divertor plates rises from 300°C to more than 1200 °C in 1 millisecond.

![Figure 1](image.png)

**Energy and particle behavior onto the divertor**

We found in beam-heated discharges that the ratio of the power deposited at the divertor plates to the input power is up to 45% in the electron density $n_e < 3.5 \times 10^{19}$ m⁻³ regime and it is reduced to less than 30% in high electron density regime where the remote radiative cooling is effective at the divertor. The width of scrape off layer deduced from the FWHM of heat flux is typically about 1.5 cm in the discharges with $P_{NBI} \sim 12$ MW and $n_e \sim 3 \times 10^{19}$ m⁻³. In LHCD plasmas at the steady states, 40% of injected power are deposited at the divertor. The width of the heat flux decreases with auf power up to 4 MW. The width of scrape off layer is estimated to be about 1.5 cm with $P_{RF} \sim 4$ MW and $n_e \sim 3 \times 10^{19}$ m⁻³.

It is found that the particle and energy transport are effectively controlled by changing X-point height in beam heated discharges. As large as 20% reduction of the temperature rise of divertor plates is obtained by vertical swing by 5 cm at 2 Hz in the discharge with 20 MW of beam power for three seconds because the effective heat flux density is reduced. It is also found that more than 10% of the core plasma density can be controlled
in beam heated discharges by changing X point height. Since longer distance from the plate causes the larger shielding efficiency of neutral particles at high X point operation, lower plasma density is obtained than lower X point operation. (Fig. 2)

![Graph](image)

*Fig. 2*
The time evolution of the plasma parameters during X-point swing discharge.

In beam-heated discharges, the asymmetry in particle recycling between inner divertor and outer divertor is enhanced and it is reversed when the ion grad-B drift direction is reversed. We also found that the heat flux to the divertor shows the reversed asymmetry which also depends on the direction of ion grad-B drift. These asymmetries in particle recycling and heat flux indicate the enhanced particle fluxes due to collisional transport driven by the temperature gradient along the magnetic field in the scrape-
off layer play an important role on the particle recycling[1].

The IDC regimes (Improved Divertor Confinement) were found in JT-60 [2] and were obtained in the discharges with high density and high neutral beam power when ion grad-B drift is toward the X-point. The growths of the IDC phenomena are associated with the reduction in heat flux and the increase in the radiation loss due to the carbon accumulation only at the inner divertor. The particle pinch due to the collisional particle transport is expected to be enhanced by the asymmetry in heat flux which is caused by the radiation loss at the inner divertor. For the experimentally observed asymmetry in heat flux density at the divertor $\delta Q = 300$W/cm$^2$, the growth rate of the plasma density is estimated to be an order of seconds. The value is consistent with the experiment. Thus the collisional transport seems to be a dominant process to improve particle transport in the IDC phenomena.

Conclusions

The behavior of the boundary plasmas on the energy and particle transport is studied in JT-60 L-mode discharges. The MHD instabilities degrade the core energy confinement significantly due to the enhanced energy loss channel to the plasma periphery. The particle and energy transport are effectively controlled by X-point height in the beam heated discharges. The enhanced particle fluxes driven by the temperature gradient along the magnetic field in the scrape-off layer play an important role on the particle recycling and the IDC behaviors.

References.

LOCAL TRANSPORT ANALYSIS OF L-MODE PLASMAS IN JT-60

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Introduction

Global confinement characteristics of L-mode are well summarized as a power scaling, like ITER89[1]. In order to extrapolate reliability of the scaling to future plasma regimes, however, one needs to understand physics for local confinement in tokamaks, capable of describing global confinement scaling. The present work is intended to contribute to this important topic by analyzing profile data more systematically and by attempting to find mechanisms which are essential for the L-mode transport. In JT-60, beam heating experiments use hydrogen beams and nearly perpendicular, balanced NB injection ((P_{CO}-P_{CTR})/(P_{CO}+P_{CTR})=0.2). These features enable a direct study of transport properties of the L-mode plasma. Indeed, data for hydrogen plasmas should be useful in considering the isotope effect presented in the global confinement scaling. Perpendicular injection results in centrally peaked heat deposition profiles up to medium density so that we can study the density dependence of thermal transport under almost constant heat effectiveness. Furthermore, because of balanced injection, the local transport properties can be inferred without taking into account the plasma rotation effect.

Experiments

In order to study the thermal transport properties of L-mode plasmas, systematic experiments were performed in a wide range of plasma parameters. For divertor discharges, major radius, R_0 = 2.9 m, minor radius, a = 0.64 m, elongation, \( \kappa = 1.4 \), the plasma current, \( I_p = 1.0 \sim 1.8 \text{MA} \), the line averaged electron density, \( \bar{n}_e = 1.2 \sim 5.0 \times 10^{19} \text{m}^{-3} \), total input power including Joule and neutral beam heating power, \( P_{abs} = 1.3 \sim 16.7 \text{MW} \), and for limiter discharges, \( R_0 = 3.0 \text{m} \), \( a = 0.88 \text{m} \), \( \kappa = 0.96 \), \( I_p = 1.0 \sim 2.7 \text{MA} \), \( \bar{n}_e = 1.2 \sim 6.5 \times 10^{19} \text{m}^{-3} \), and \( P_{abs} = 3.0 \sim 17.4 \text{MW} \). The toroidal field, \( B_t \) was 4.5 T in all experiments. The plasma configuration was slightly shifted 0.1 \sim 0.2 \text{m} inside, compared with the normal discharge configuration [2], for better accessibility of diagnostics to the central region of the plasma. The characteristic feature of the consistent data set for limiter discharges is that the electron densities are increasing with heating power due to the neutral beam particle source. On the other hand, for divertor discharges, the power scan could be performed with almost constant electron densities \( (3.0 \times 10^{19} \leq \bar{n}_e \leq 4.0 \times 10^{19} \text{m}^{-3}) \).

Transport Analysis

First, the density dependence of the thermal component of stored energy, \( W_{th} \) is considered. Figure 1-(a) shows the diamagnetic measured stored energy, \( W_d \) as a function of
the density for 1.0 MA, divertor discharges with fixed beam power of 2.1 to 2.6 MW. As the L-mode scaling indicates, $W_d$ is independent of the density. However, $W_{th}$ which is estimated by subtracting the beam component from the diamagnetic one, apparently increases with a rise of the density, as shown in Fig. 1-(a). So there should be some improvement in the plasma confinement as the density increases. The thermal diffusivities at half the plasma radius indicate the reduction of the ion thermal diffusivity, $\chi_i$, with increasing density (Fig. 1-(b)). On the other hand, the electron thermal diffusivity $\chi_e$ does not seem to be reduced. Even with much higher heating power of 11.1 to 13.7 MW in 1.5 MA divertor discharges, we can also see $W_{th}$ increase with the density [3]. $\chi_i$ still seems to be a decreasing function of the density, although the separation of the transport loss channel becomes slightly difficult because of the high density. Since the JT-60 neutral beams are nearly perpendicularly injected, the deposition profiles are centrally peaked and remain essentially unchanged, during the density scan. Therefore, the reduction in $\chi_i$ suggests that the density dependence of $W_{th}$ is mainly due to the improvement in the ion thermal transport with increasing density. This density dependence of the ion thermal transport, however, confuses the correct parameter dependence of the thermal diffusivity. Accordingly, in the local transport study presented, we have paid careful attention to the density.

Next, the power dependence of the confinement, which is the most basic parameter dependence shown in the L-mode scaling, is studied. Figure 2-(a) presents the kinetic stored energy, $W_k$ including the beam component, $W_{th}$ against $W_d$ for 1.0 MA, divertor discharges in the density region of $\tilde{n}_e = 3.0 - 4.4 \times 10^{19}$ m$^{-3}$, only two shots have a density higher than $4.0 \times 10^{19}$ m$^{-3}$, accompanied by higher heating power of around 15 MW, and the other shots are in the relatively narrow density region of $3.0 - 4.0 \times 10^{19}$ m$^{-3}$. The reasonable agreement shown in Fig. 2-(a) gives us confidence that we have an accurate estimate of the thermal stored energy, as well as the evaluated beam component. In order to separately determine global confinement characteristics of ion, electron, and fast ion, the kinetic stored energy is decomposed into those three components, and they are shown as a function of the total heating power in Fig. 2-(b). From this figure, it is found that the electron and beam components increase with a rise of heating power. The ion stored energy, however, slightly increases, and the increment of ion energy to the total power is only one-third that of the electron energy. The slight increase in the ion stored energy cannot be ascribed to the proton dilution, because values of $Z_{eff}$ are still less than two for powers less than 15 MW.
Local thermal transport analysis provides a better description of confinement properties of plasmas with substantially different heating powers. Figure 3 shows ion and electron thermal diffusivities at half the plasma radius as a function of the total heating power, for those shots referred to in Fig. 2 with $P_{\text{abs}} < 15 \text{ MW}$. We can see that $\chi_i$ significantly increases in proportion to the heating power, while $\chi_e$ gradually increases with the power. In this regard, $\chi_i$ is about three to four times $\chi_e$ at around 10 MW. Therefore, the enhancement of $\chi_i$ with increasing heating power represses the increment in $W_i$, resulting in the degradation in confinement of L-mode plasmas.

**Comparison with transport driven by $\eta_\parallel$ - mode**

The ion temperature gradient instability, $\eta_\parallel$-mode, has been proposed as one of the highly probable candidates which are responsible for such ion anomalous transport[4]. A comparison of measurements with theoretically predicted ion temperature profiles has been done for NB-heated discharges using analysis and simulation codes. We employed a model proposed by Dominguez and Waltz [5]. Figure 4 shows comparisons of calculated ion temperature profile with measured ones, for heating power scan in 1.0 MA divertor discharges. The predicted ion temperature profiles are in good agreement with observations. In general, reasonable agreements between predicted and observed ion temperature profiles are achieved for various discharges with higher $I_p$ [6]. Accordingly, the contribution of strong electrostatic turbulence produced by the driven modes to the ion heat diffusivity is very successful in explaining the observed response of ion temperature profiles to heating power changes.

In limiter discharges with low $I_p$ of 1.0 MA, however, the predicted ion temperature profiles by the $\eta_\parallel$ - mode turbulence model becomes higher than observations. Ion transport enhanced by a factor of four is needed to explain measured profiles. This discrepancy might be caused by the possibility that $\chi_i$ for limiter discharges essentially has strong $I_p$ dependence, because the transport model used does not explicitly have
the $I_p$ dependence. At present, this problem is still open.

**Conclusion**

It is now generally acknowledged that heat flux through the ion channel may be important in determining confinement. The deterioration in the energy confinement time with increasing auxiliary heating power is mainly due to the degradation in ion energy transport. It is found that the ion thermal diffusivity has clear dependence on the density. This cannot be excluded even at medium density. The ion anomalous transport is partly explained by the ion temperature gradient mode. On the other hand, parameter dependence of $\chi_e$ is essentially indeterminate, because in hydrogen plasmas, the ion anomalous transport is so dominant in the power balance that the electron transport is masked by the ion transport characteristic.

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TRANSPORT SIMULATIONS USING THEORY-BASED MODELS

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A combination of theoretically derived transport models,[1–5] called the Multi-Mode Model, is used in the 1-1/2-D BALDUR transport code[6,7] to simulate tokamak discharges. The present version of Multi-Mode Model consists of effective thermal diffusivities resulting from a combination of trapped electron and ion temperature gradient ($\eta_i$) modes, which dominate in the core of the plasma, together with resistive ballooning modes, which dominate in the periphery. It is found that L-mode and H-mode plasmas can be simulated with the same core transport model, with changes only at the edge.

Several different Multi-Mode Models have been successfully calibrated against a wide range of experimental data. In these models, the trapped electron mode contribution is based on the model used by Dominguez and Waltz[8] with modifications suggested by Rewoldt, Tang, et al.[9,10] The ion temperature gradient driven turbulence ($\eta_i$-mode) contribution is modeled by either the Dominguez and Waltz form[8] or the Hamaguchi and Horton theory.[11] Resistive ballooning and resistive interchange modes are modeled by the Carreras and Diamond theories.[12,13] In each version of the Multi-Mode Model, the coefficients and thresholds in these contributions were adjusted to match the experimental data and were then held fixed for a consistent set of simulations.

Figures 1 and 2 show the profiles and effective thermal diffusivities for a TFTR L-mode simulation. For this simulation, the parameters were $R = 2.58$ m, $a = 0.92$ m, $I_p = 1.8$ MA, $B = 3.76$ tesla, $q_{\text{edge}} = 4.3$, $P_{\text{aux}} = 16.8$ MW from 2.8 seconds on, $n_e = 4.26 \times 10^{19}$ m$^{-3}$ at 3.65 seconds, and $Z_{\text{eff}} = 1.8$ with carbon impurity. Figure 1 shows the simulated temperature and density profiles as a function of major radius compared with experimental data from TFTR. (For this simulation, the theory-based models were used to predict the temperature profiles while an empirical particle transport model was used to produce a density profile to match experimental data.) Figure 2 shows the electron and ion effective thermal diffusivities as a function of flux surface half-width and the contributions due to trapped electron modes, $\eta_i$ mode, and resistive ballooning modes. A detailed description of the transport model for this simulation is given in reference [5].

It is an essential feature of this simulation that the sum of the three contributions produces an effective thermal diffusivity that increases monotonically from the center.

Work supported by the U. S. DoE Contract No. DE-AC02-76-CHO-3073 and DE-FG02-88-ER53269.
of the plasma to the edge. Removing any contribution would produce a local thermal barrier, which is not observed in the experimental data for L-mode plasmas.

It is found that the scaling of confinement with heating power results from the $T_e^{3/2}$ dependence, which is characteristic of collisionless trapped electron modes and $\eta_i$ modes. The role of resistive ballooning modes and the value of $\eta_i$ relative to the threshold value does not change noticeably in the outer half of the plasma as the heating power changes.

The scaling of confinement with current in this model results mostly from resistive ballooning modes, which intrude further into the plasma as the magnetic q-value increases at the edge. As the plasma current is reduced, the effective thermal diffusivity due to resistive ballooning modes is increased while the diffusivity due to trapped electron and $\eta_i$ modes is reduced because the central temperatures are lower. The shear dependence in the Hamaguchi-Horton model for $\eta_i$-mode transport[11] plays very little role in the L-mode shots simulated.

Fig. 1. Electron temperature, ion temperature, and electron density as a function of major radius from BADLUR simulation TL02Y1 compared with experimentally measured data from Thomson scattering and CHERS for TFTR L-mode shot 41325 (SNAP try 5).

Fig. 2. Contributions to the electron and ion effective thermal diffusivities as a function of flux surface half-width. Contributions are shown from trapped electron modes (TEM), $\eta_i$ mode, and resistive ballooning modes (RB). The total (TOT) also includes the neoclassical thermal diffusivity.
Figure 3 shows the temperature and density profiles as a function of major radius for the simulation of the JET H-mode #15894. In this case, both the temperature and density profiles were simulated using a combination of theory-based transport models. The boundary conditions at the top of the pedestal were determined according to the time-dependent model for H-modes presented in reference [15]. The magnitude (but not the shape) of the temperature profile is somewhat sensitive to the value of the boundary temperatures. It is found that H-mode, L-mode and OH plasmas can be simulated using the same core transport model, with changes only at the edge.

The objective of this research has been to find a combination of theory-based transport models that can be used to predict temperature and density profiles for a wide range of experimental data. It is clear that better theory-based models are needed, especially to improve the prediction of transient effects and supershot profiles. Also, a more accurate model is needed for the boundary values of temperature and density at the top of the pedestal for both H-mode and L-mode plasmas. As new models become available, they will be incorporated into our work. However, the present class of theory-based models has been found to work remarkably well in predictive simulations of essentially steady state Ohmic, L-mode, and H-mode plasmas.

Fig. 3. Temperature and density profiles as a function of major radius for the simulation of JET H-mode #15894 compared with experimental data from ref [14].


NON-DIFFUSIVE HEAT TRANSPORT DURING ELECTRON CYCLOTRON HEATING ON THE DIII-D TOKAMAK*

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Of central importance to magnetic confinement fusion is the understanding of cross-field heat transport, which is usually modeled as a diffusive process down a temperature gradient with a small additional convective term due to particle transport. This paper reports results from off-axis electron cyclotron heating (ECH) experiments which cannot be adequately described in this framework. In particular, net heat appears to be flowing up the temperature gradient in the electron channel.

Electron cyclotron heating experiments at 60 GHz have been carried out in the DIII-D tokamak with launched power levels up to 1.4 MW. The ECH launch system, located on the inside wall at \( z = +13 \) cm, launches the extraordinary X-mode in a Gaussian pattern with a 12° half width. Eight antennas direct their power at 15° and two antennas direct their power at ±30° with respect to the major radius. The orientation is such to drive current aiding the Ohmic current for normal operation.

For the off-axis ECH experiments described here, the waves launched from the inside wall must first pass through the fundamental cyclotron resonance before they reach the plasma center. A 3-D ray tracing code calculates that single pass absorption near the fundamental cyclotron resonance is strong (60%–90%), while absorption near the plasma center is negligible. The ECH power has a very narrow spatial deposition profile (full-width at half-maximum 0.05–0.10 \( a \)); the location of the ECH power deposition can be changed by changing the toroidal field of the plasma. If the electron heat transport is purely diffusive, then off-axis ECH should result in an electron temperature profile with a sharp change in gradient at the heating location, as observed for off-axis ICH on JET. In DIII-D, however, the electron temperature profile is found to remain strongly peaked, even in the extreme case of 69% of the total input power deposited outside the the normalized radius \( p = 0.7 \).

The electron temperature profile and ECH deposition profile for an L-mode plasma where 63% of the ECH power is deposited between \( p = 0.5–0.6 \) on the first pass are shown in Fig. 1. Also shown is the simulated electron temperature profile assuming that the electron heat transport is purely diffusive (details of this simulation will be discussed later). Strongly peaked electron temperature profiles for off-axis ECH deposition are observed for fundamental heating in both L-mode and H-mode plasmas, and for low-field O-mode as well as high-field X-mode launch. The peaked temperature profile of Fig. 1 cannot be explained by "profile consistency" since the shape of the electron temperature profile varies with the ECH deposition location. In general, the electron temperature profile becomes more peaked the more on-axis the ECH deposition is.

The electron heat balance has been analyzed using the power-balance transport code ONETWO, which determines the primary ion diffusion coefficient, the ion conductivity, and the electron conductivity by solving the ion particle balance, ion heat balance, and electron heat balance equations, given the particle densities and the ion and electron temperatures. The heat balance for electrons can be written as

\[
\langle \nabla \cdot q_e \rangle = \langle Q_{\text{ion}} + Q_{\text{aux}} - Q_{ei} - Q_{\text{rad}} \rangle ,
\]

* This work was sponsored by the U.S. Department of Energy under Contract Nos. DE-AC03-89ER51114 and W-7405-ENG-48.
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where \( \langle \rangle \) denotes an average over a magnetic flux surface, and

\[
q_e = -n_e x_e^{PB} \nabla T_e + q_{flow}
\]

(2)

is the total electron heat flux. The term \( q_{flow} \) is assumed by ONETWO to be the convective heat flux, a negligible loss term for the core region in DIII-D. The total electron and ion heat fluxes determined by ONETWO are shown in Fig. 2 for the same case of off-axis ECH as in Fig. 1. The ECH power not absorbed on the first pass is assumed to be distributed over the portion of the plasma where there is an ECH resonance. The indicated error bars are due to uncertainties in the electron and ion temperature, the electron and ion density, and the radiated power. The total electron heat flux is negative within error bars from \( \rho = 0.34 \) to \( \rho = 0.48 \). The electron and ion heat diffusivities determined from the electron and ion heat fluxes are shown in Fig. 3. In order to support the observed electron temperature profile, \( x_e^{PB} \) drops dramatically inside of \( \rho = 0.55 \) (the resonance location) and becomes negative. The ion channel does not display this unusual behavior.

The plasma current-density profile is modeled by ONETWO using the magnetic diffusion equation and is consistent with the measured electron temperature profile. Since a sawtooth model is not included in the analysis, the actual current-density profile is probably broader than the modeled profile, and the total electron heat flux therefore more negative near the plasma center. Although the power balance analysis is performed assuming steady-state conditions, a simulated time evolution of the response of the current-density profile to electron heating displays no transient effects that would significantly alter the steady-state results. No change in the plasma internal inductance is measured during ECH, and the ECH deposition is well outside the \( q = 1 \) surface. Even if the Ohmic power profile is actually more peaked than believed, which could explain the measured central \( T_e \), a sharp change in the electron temperature gradient at the ECH heating location should still be observed (\( P_{\text{Ohm}} = 0.25 \text{ MW} \) compared to \( P_{\text{ECH}} = 1.10 \text{ MW} \)).

Because diffusion opposite the temperature gradient is unphysical, a different method of determining the conductive loss from the relaxation of the electron temperature profile\(^5\) is used. When the ECH is switched on, the electron temperature profile evolves to a new stationary state. The electron heat balance equation during the temperature rise can be expanded assuming \( T_e(\rho, t) = T_0(\rho) + T_1(\rho, t) \), where \( T_0(\rho) \) is the equilibrium temperature profile during ECH. Assuming that the first order heat flux is purely diffusive, and neglecting perturbations to the Ohmic source and electron-ion coupling terms, results in the first-order electron heat-balance equation

\[
\frac{3}{2} \frac{\partial}{\partial t} (n_e T_1) = \nabla \cdot (n_e x_e^{TD} \nabla T_1)
\]

(3)

Since \( n_e \) is independent of time for ECH L-mode and since \( T_1(\rho, t) \) is measured, Eq. (3) can be solved for \( x_e^{TD}(\rho) \), assuming \( x_e^{TD} \) is independent of \( T_e \). This "time dependent" \( x_e \) is shown in Fig. 3 for comparison with \( x_e \) determined from power balance analysis. Note that there is no unusual behavior in \( x_e^{TD} \) near the ECH deposition location.

Assuming that \( x_e^{TD} \) governs the equilibrium conductive loss allows the electron temperature profile to be simulated for off-axis ECH deposition, as shown in Fig. 1. The measured central \( T_e \) is twice that expected if the electron heat transport were purely diffusive. In addition, using \( x_e^{TD} \) in place of \( x_e^{PB} \) in Eq. (2) allows Eq. (1) to be solved for the "auxiliary power" profile (i.e., the apparent input power other than Ohmic) required to support the measured electron temperature profile against known losses. The total auxiliary input power determined in this way is nearly identical in magnitude to the actual ECH power, but the apparent deposition location is shifted significantly inwards from the ECH resonance (Fig. 4).

It is important to remember that the deposition profiles from ray tracing are well supported experimentally from local heating rates determined from ECE and soft x-rays, as well as from
Fig. 1. Measured electron temperature profile and calculated ECH deposition profile for a L-mode plasma with \( I_p = 0.5 \text{ MA}, B_T = 1.7 \text{ T}, n_e = 2.1 \times 10^{13} \text{ cm}^{-3} \), and \( P_{\text{ECH}} = 1.1 \text{ MW} \). A simulated electron temperature profile for a purely diffusive model is also shown.

Experimental data is predicted to absorb most of the power, and then be damped near the plasma center by some anomalous process (there is no cyclotron resonance inside of \( \rho = 0.5 \) for this plasma). This explanation is clearly unsatisfactory.

Fig. 2. Total electron and ion heat flux for the same case of off-axis ECH as in Fig. 1.

phase lag measurements during pulse modulated ECH. To get such an inward shift of the deposition profile by direct absorption, the waves would first have to propagate unattenuated through a layer which

Fig. 3. Electron and ion heat diffusivities determined from power balance analysis \((\chi_{e,i}^{\text{PB}})\) and a time dependent method \((\chi_{e,i}^{\text{TD}})\) for the discharge shown in Fig. 1.

Fig. 4. Power deposition profiles for the deduced “auxiliary” heating, first pass ECH, and Ohmic heating sources for the discharge shown in Fig. 1.
Using $Q_{ECH}$ from ray tracing along with $\chi^{TD}$ in Eq. (1) allows the non-diffusive term $q_{\text{flow}}$ in Eq. (2) to be evaluated. The non-diffusive power flowing through a flux surface with surface area $S$ is then determined from $P_{\text{flow}} = q_{\text{flow}}S$. The non-diffusive power profile is shown in Fig. 5 for a three point scan of the ECH resonance location. Negative power means that the flow is towards the plasma center, thus $P_{\text{flow}}$ has the characteristic of a "heat pinch". The non-diffusive power is seen to increase monotonically with $\rho$ up to the ECH resonance location, outside of which it is zero. Because the ion heat balance appears to be unaffected, and because the density profile is unchanged during the ECH pulse, it is unlikely that this inward power flow can be explained by an $E \times B$ particle pinch.

The ECH results presented here indicate a general lack of dependence of the global energy confinement time on the auxiliary power deposition profile. This result may ease technology requirements for ECH systems because off-axis heating on the low field side of a tokamak allows the use of lower frequency (and therefore less expensive) gyrotrons. It may also be beneficial to heating scenarios in which monochromatic ECH is applied during the toroidal magnetic field ramp up since it allows peaked temperature profiles to be achieved at an earlier stage of the discharge.

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![Fig. 5. The non-diffusive power flowing through a flux surface for a three point scan of the ECH resonance location. The peak ECH absorption is at $\rho = 0.32, 0.47, 0.55$ for $B_T = 2.0, 1.8, 1.7$ T respectively.](image-url)
A MODEL TO EVALUATE COUPLING OF ION BERNSTEIN WAVES TO TOKAMAK PLASMAS

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Introduction. Higher order Bernstein waves have wavelengths comparable or shorter than the thermal ion gyroradius. To describe their propagation in inhomogeneous plasmas therefore one should solve an integro-differential wave equation[1]. High order differential wave equations valid in the immediate vicinity of a cyclotron harmonic have also been proposed[2]. Use of these equations to evaluate coupling of high order ion Bernstein waves with external antennas is difficult, however, since it is not clear which boundary conditions should be imposed at the plasma edge.

Model wave equations. To circumvent this difficulty, we have used the following set of model wave equations ($x$, $y$ and $z$ are the radial, poloidal and toroidal directions, respectively; since curvature and shear are neglected, the last two are ignorable and allow a Fourier decomposition; lengths are measured in units of $c/\omega$):

$$- \frac{d}{dx} \left( \frac{dE_x}{dx} - i\delta \frac{dE_y}{dx} \right) + n_y^2 (\sigma E_x - i\delta E_y)$$
$$+ in_y \left[ \frac{d}{dx} \left( i\delta E_x + \sigma E_y \right) - \left( i\delta \frac{dE_x}{dx} + \sigma \frac{dE_y}{dx} \right) \right] + in_x \frac{dE_x}{dx} - in_y n_x \xi E_x$$
$$+ (n_y^2 + n_z^2 - S) E_x + in_y \frac{dE_y}{dx} + iDE_y + 2in_y \tau \left( \frac{dE_y}{dx} - in_y E_x \right) = 0$$

$$- \frac{d}{dx} \left( i\delta \frac{dE_x}{dx} + \sigma \frac{dE_y}{dx} \right) + n_y^2 (i\delta E_x + \sigma E_y)$$
$$- in_y \left[ \frac{d}{dx} \left( \sigma E_x - i\delta E_y \right) - \left( \sigma \frac{dE_x}{dx} - i\delta \frac{dE_y}{dx} \right) \right] - ny n_x E_x + n_z \frac{d}{dx} (\xi E_z)$$
$$+ in_y \frac{dE_x}{dx} - iDE_x - \frac{d^2E_y}{dx^2} + (n_z^2 - S) E_y - 2 \frac{d}{dx} \left[ \tau \left( \frac{dE_y}{dx} - in_y E_x \right) \right] = 0$$

$$in_x \frac{dE_x}{dx} - n_x \xi \left( \frac{dE_y}{dx} - in_y E_x \right) - ny n_x E_y - \frac{d^2E_x}{dx^2} + (n_y^2 - P) E_x = 0$$

where $S$, $D$, $P$ are the elements of the local dielectric tensor in the zero Larmor radius.
approximation, while the hot plasma terms are defined as follows:

\[
\sigma = \sigma_e + \sum_{i} \beta_i \sum_{n=\pm \infty} \frac{n^2}{\lambda_i^2} \Gamma_n^{(\sigma)}(\lambda_i)(-x_{0i} Z(x_{n_i}))
\]

\[
\delta = -\sum_{i} \beta_i \sum_{n=\pm \infty} \frac{n}{\lambda_i} \Gamma_n^{(\delta)}(\lambda_i)(-x_{0i} Z(x_{n_i}))
\]

with \(x_{n_i} = (\omega - n\Omega_e) / k_{\parallel} v_{thi},\) and

\[
\Gamma_n^{(\sigma)} = I_n(\lambda_i) e^{-\lambda_i} - \frac{\lambda_i}{2} \delta_{|n|,1} \\
\Gamma_n^{(\delta)} = [I'_n(\lambda_i) - I_n(\lambda_i)] e^{-\lambda_i} - \frac{1}{2} \delta_{|n|,1}
\]

Also, \(\beta_e = \omega_{pe}^2 v_{thi}^2 / 2\Omega_e^2 c^2,\) and \(\lambda_e = k_{\perp}^2 v_{thi}^2 / \Omega_e^2,\) with \(k_{\perp}^2\) evaluated by solving the local dispersion relation for the Bernstein wave. The electron contributions are limited to first order in the electron Larmor radius:

\[
\sigma_e \approx \frac{3}{4} \beta_e \\
\tau \approx \tau_e \approx -\beta_e x_{e0} Z(x_{e0}) \\
\xi \approx \xi_e \approx -\beta_e \frac{\Omega_{ce}}{\omega} x_{e0}^2 Z'(x_{e0})
\]

All other electron terms and the ion contributions to \(P, \tau\) and \(\xi\) have been neglected.

Except for these minor simplifications, the system (1) reproduces the exact dispersion relation in the homogeneous limit. Although not derived rigorously from the Vlasov–Maxwell equations, it should be a good approximation provided that the Lower Hybrid resonance is not in the domain of integration. At the same time it has the same structure as the familiar Finite Larmor radius wave equations. Thus boundary and radiation conditions are easily imposed. In particular, at the plasma–vacuum boundary continuity of the tangential field components and vanishing of the kinetic part of the power flux provide the correct number of boundary conditions, and ensure conservation of energy.

**Applications.** As an example, we have implemented the model wave equations in a modified version of the FELICE code to evaluate coupling with a metallic loop antenna oriented parallel to the toroidal magnetic field. One can distinguish two coupling regimes, depending on whether the density at the plasma edge is well below or above the density corresponding to the Lower Hybrid resonance \(n_{LH} \approx 5.3 \times 10^{15} (\omega^2 / \Omega_e^2 - 1) B^2 / A_i\) (MKS units). With metallic antennas the latter situation is probably not realistic, but it might be achievable with waveguide launching. A code to deal with this case is being developed.

In our examples the magnetic field at the plasma boundary is 6.45 Tesla, and the frequency is 450 MHz, intermediate between the fourth and fifth cyclotron harmonic of \(H^+\). The antenna is located 0.5 cm away from the edge itself, behind an ideal Faraday screen, and 7.5 cm from the wall. It consists of a T–loop feeded at the center and shorted at both ends, so that the current distribution is antisymmetric, as required to optimise coupling to slow waves.
In the first case a scrape-off thickness of 5 cm has been assumed, with \( n = 10^{19} \text{ m}^{-3} \) and \( T = 100 \text{ eV} \) at the separatrix, and e–folding length of 2 cm. The outward radiation conditions have been imposed 1 cm outside the LH resonance, which in turn is located 2 cm outside the separatrix. While the model is not able to predict what happens at the resonance, reflection from this layer is not expected; due to the very short wavelengths the validity of the outward radiation conditions should not be questionable.

Under these conditions, cold slow waves with \( k_\parallel c/\omega \ll 1 \) are launched with good efficiency. Fig. 1 shows \( E_x \) for a partial wave with \( n_\parallel = 2 \): wavelength and amplitude are just as expected from the local dispersion relation. Fig. 2 shows the radiated \( n_\parallel \) spectrum: it closely recalls the one from a two-waveguide Grill with phase \( \Delta \phi = 180^\circ \). The Fourier spectrum of the current is appreciably broader, implying that waves with large \( n_\parallel \) are poorly matched. Interestingly, the total radiated field (superposition of 61 toroidal \( \times 11 \) poloidal modes) has little to do with the partial field of Fig. 1: instead, one recognised resonance cones, although not very sharp ones because low frequency and low \( n_\parallel \) imply that electromagnetic effects are important. The loading resistance reaches several Ohms, and increases rapidly with decreasing edge density (Fig. 3): again a similar behaviour would be expected for a waveguide launcher in this regime. Note that these results could have been obtained with a cold plasma model, but not with the FLR model, which predicts complete reflection from the LH layer due to a spurious very short wavelength root at low densities which does not exist in the full hot plasma dispersion relation.

If on the other hand the edge density exceeds the LH resonance, the efficiency of coupling is strongly reduced, because the ratio \( E_x/E_z \) of the BW is much larger than that of vacuum waves. As a consequence a large amplitude evanescent cold slow wave is first excited, which rapidly dies out leaving a Bernstein wave with much smaller amplitude. This is clear from Fig. 4, which shows \( E_x \) assuming a density \( n_e = 10^{19} \text{ m}^{-3} \) at the plasma boundary. Since the perpendicular wavelength of BWs is practically independent from \( n_\parallel \), in this case the total radiated field shows all the features predicted by the dispersion relation. The radiated spectre is similar to the low density case, but the resistance is only 0.015 Ohm: any non–ideal absorption effect not included in our model would completely wash out this tiny radiation loading.

We conclude that coupling of BWs is only possible at edge densities below the LH resonance. Whether absorption at the resonance layer can be avoided cannot be answered with the present model.

Fig. 1 – Low density coupling: real and imaginary part of $E_z$ for $n_\parallel = 2$.

Fig. 2 – Low density coupling: $n_\parallel$ spectrum from a $T$- antenna.

Fig. 3 – Low density coupling: resistance vs density at the separatrix (the edge density is lower by a factor $e^{-5/2} \approx 0.0808$).

Fig. 4 – High density coupling: real and imaginary part of $E_z$. 
FLUCTUATION MEASUREMENTS BY LANGMUIR PROBES DURING LHCD ON ASDEX TOKAMAK

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ABSTRACT The level of edge electrostatic fluctuations decreases and the global particle/energy confinement improves during lower hybrid current drive (LHCD) regimes on ASDEX, when the total power remains below the initial OH power level. For higher powers, the fluctuations increase noticeably, whereas the global confinement is returning to its OH value. The observed increase of fluctuations is poloidally asymmetric and is caused by local power deposition in front of the grill antenna.

INTRODUCTION Generally, deterioration of the global confinement is observed during LHCD at sufficiently high LH powers. However, at low densities, an improvement of the global energy \(1/1\) and particle \(3/3\) confinement has been observed for LH powers lower than the initial OH value. A significant reduction of electrostatic fluctuations at the plasma edge was observed in \(4/4\) under similar conditions.

Here, the edge electrostatic fluctuations and the global particle/energy confinement are investigated simultaneously during the LH power scan.

EXPERIMENTAL ARRANGEMENT Experiment was performed in the LHCD regime at \(f = 2.45 \text{ GHz}\). The LH antenna (phasing \(\Delta \phi = 900\)) consists of two independently powered parts, denoted as the upper and lower grills.

The electrostatic fluctuations are monitored by a Langmuir probe, located near the equatorial plane, 3 cm outside the separatrix (to protect the tips from damage by suprathermal electrons) and 210 toroidally away from the LH antenna. The flux tube connected with the probe crosses the plasma in front of the lower grill (Fig. 1). In connection with an analog correlator, the probe (with three poloidally separated tips) allows us to monitor the rms-values of density \(n_e\) and poloidal electric field \(E_p\) fluctuations as well as the radial turbulent flux \(\Gamma\), induced by the cross-field drift \(E_p \times B_{\text{tor}}/4\).

The fluctuation data could be qualitatively compared with the global particle confinement, assuming the fluctuation-induced
flux as the dominant loss channel \( \tau_p \sim a_n e / \Gamma \).

**EXPERIMENT** The fluctuations have been studied in regimes with \( B_{\text{tor}} = 2.8 \, \text{T}, \, I_p = 420 \, \text{kA} \) \( (q(a) = 3.2) \). The density \( n_e = 1.3 \times 10^{13} \, \text{cm}^{-3} \) is far below the ASDEX LHCDD density limit \( n_e(2.4 \, \text{GHz}) = 4.5 \times 10^{13} \, \text{cm}^{-3} \).

The experiment was aimed to determine the variation of the electrostatic fluctuations with LH power. To avoid scattering of experimental data, the LH power scan was performed during a single shot. The LH wave was launched in eight power steps, covering a broad range of powers with respect to the initial OH input, starting from \( P_{\text{LH}} < P_{\text{OH}} \) (Fig. 2a). Duration of each step was sufficiently longer than the global particle/energy confinement times.

During the LH power scan, the loop voltage, and consequently the residual OH power \( (P_{\text{OH}}) \), decreases. The total power \( P_{\text{tot}} = P_{\text{LH}} + P_{\text{OH}} \) decreases at first below \( P_{\text{OH}} \), reaching its minimum at \( P_{\text{LH}} \approx 74 \, \text{kW} \) (Fig. 2b). At the same time, the central electron temperature decreases only slightly. The line average \( T_{\text{eff}} = 1.5 \) is nearly constant during the major part of the power scan.

**Electrostatic fluctuations** (Fig. 2f - g):

For the first four power steps, the fluctuations are reduced noticeably. The reduction is maximal at \( P_{\text{LH}} = 74 \, \text{kW} \), where the relative level of density fluctuations \( \tilde{n}/n \approx 0.22 \) represents about 75% of its OH level. The \( E_p \) fluctuations are reduced by 10% compared with \( E_{\text{OH}} \approx 10 \, \text{V/cm} \).

The fluctuation levels increase substantially for the last four power steps, when \( P_{\text{LH}} \geq 100 \, \text{kW} \). The increase is more pronounced for \( E_p \)-fluctuations, which rise up to \( E_p \approx 50 \, \text{V/cm} \). The relative level of density fluctuations reached \( \tilde{n}/n \approx 0.8 \) the last power step.

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**Fig. 2** Evolution of the shot 
\# 31601 with the LH power scan.
A similar increase is observed in the fluctuation-induced flux \( \Gamma \).

However, it should be emphasized that the strong jump of the fluctuation levels occurs together with the switching of the LH power from the upper to the lower grill. In order to clarify the importance of the relative position of the probe and of the active grill, the sequence of the waveguides power-supply was reversed. It has been found that for \( P_{PLH} > 100 \text{ kW} \), the fluctuation level is considerably higher, when the lower grill, directly connected to the probe through a magnetic flux tube, is activated. This indicates a strong poloidal asymmetry in the enhancement of the fluctuations.

The local modification of the edge plasma during LHCD can be observed directly from photos (Fig. 3), showing a bright zone just in the front of an active grill. This effect is attributed to local power deposition inside the LH wave cones due to nonlinear effects /5/. In addition, some enhanced light emission is seen from bands extending in toroidal direction. The poloidal pitch corresponds to a helix with the rotational transform \( q(a) \). The probe lays inside the bright band connected with the lower grill. Therefore, the parameters measured with the probe during the activation of this grill should be strongly influenced by local nonlinear processes.

**Global confinement** (see Fig. 2c-e)

The feedback controlled density increases slightly during the first four power steps. The total flux \( \Phi \text{ GAS} \) of neutral atoms is reduced by a factor of two at \( P_{PLF} = 74 \text{ kW} \). Simultaneously, the \( \Omega_{\alpha} \)-spectral line intensities from the main chamber and divertor regions decrease during the first part of the LH-power scan as well. Light and heavy impurity radiation decreases as seen from CIII and soft X-ray radiation. Moreover, \( \text{Zeff} \) drops near the plasma edge. Therefore, an improvement of the **global particle confinement** is concluded for \( P_{PLH} \leq 74 \text{ kW} \).

The signals, characterizing the global particle confinement, start to return to their initial OH values during the second half of the power scan, indicating qualitatively the same behaviour of \( \Omega_{\alpha} \).

The **global energy confinement time** \( \tau_E \) (Fig. 4) is determined as a ratio of the total energy stored in the plasma (derived from magnetic measurements /1/) and the total power. An increase of the total energy \( W_{TOT} = W_{BULK} + W_{TAIL} \) appears due to the creation of suprathermal electrons by the LH wave. The laser light scattering data show that the energy stored in the bulk electrons does not change appreciably.
The resulting global energy confinement time starts to increase at low LH power levels, reaching its maximum at the fourth power step. Then, it decreases to its initial value. For still higher power, $\tau_E$ is already deteriorated with respect to the $\tau_{OH} /1/$. 

CONCLUSIONS Simultaneous measurements of edge fluctuations and global confinement during the LHCD show:

i) For the low LH powers ($P_{LH} < P_{OH}$), the level of electrostatic fluctuations is considerably lower than in the target plasma. Simultaneously, the global particle/energy confinement improves. The observed increase of $\tau_E$ can be attributed to superior confinement of the suprathermal electrons. The energy confinement of the bulk electrons seems to remain close to the OH level. The maximal reduction of the fluctuations and the best confinement is reached at the minimum of the total power, when the product $nE_p$ (proportional to the fluctuation-induced flux $\Gamma$) drops by a factor of 1.5. It is consistent with the estimated increase of $\tau_p$ (up to factor of 2) and $\tau_E$ (by a factor of 1.6). It suggests a link between the level of fluctuations and the global confinement also for LH current drive regimes.

ii) When the LH power increases above $P_{LH} = 74$ kW, the global confinement times decrease gradually to their initial OH values. At the same time, a turbulent region, localized in a magnetic flux tube crossing the mouth of the active grill, appears at the plasma edge. It seems to be caused by local power deposition in front of the grill antenna.

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EXPERIMENTAL RESULTS OF RUNAWAY ELECTRONS AND MAGNETIC FLUCTUATIONS ON HL-1 TOKAMAK

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INTRODUCTION

The runaway electron has been studied. It shows that the safe factor $q_A$ passes through $10.0, 9.0, ..., 5.0$ values at plasma margin. It is also measured that the perturbation of the runaway electron is correlated sawtooth oscillation of soft X-ray and sawtooth oscillation of surperthermal electron radiation.

In the second part of this paper we'll discuss that correlation between the mironov perturbation and density fluctuation shows that the macroscopic MHD instabilities is responsible for the particles and energy loss and a new magnetic fluctuation at the margin of the plasma which was found in the experiment.

BEHAVIOUR OF THE RUNAWAY ELECTRONS

The HL-1 tokamak is a medium size tokamak ($R=1.02M, a=0.2M$) with thick copper shell ($d=0.05M$), which has good confinement on plasma and runaway electrons which are usually with energy $E_x=1-2$MeV.

The hard X-ray from limiter bursts during the plasma current rising phase when there are very severe perturbations of magnetic configuration at the margin of plasma column and the safe factor $q_A$ is near a number of rational values. It shows a close correlation between runaway electrons and plasma MHD instabilities. When the $q_A$ passes through the different rational values, such as $10.0, 9.0, 8.0, 7.3, 6.7, 5.7$ and $5.0$ on HL-1 tokamak. The runaway electron loss indicated by the hard X-ray burst increases very rapidly (Fig.1, No:9384). This indicates that the local disruption could change the magnetic energy into the kinetic energy of some electrons, this leads to the hard X-ray burst, and the confinement time[1] of the runaway electrons depends a bit on $q_A$ and the diameter of the limiter.

In long pulse discharge ($\geq 900$ms) the energy of the hard X-ray increases gradually, but the intensity of it is almost a constant at the begining of the hard X-ray rising stage as shown in (Fig.2(a), No:9216). It could be explained by non-collisions model between particles[2]. They are still accelerated when the runaways are no longer produced. On the contrast the intensity keeps still almost
constant while the energy decreases obviously at the tail part of it (Fig.2(b), No: 9365). Maybe the reason is that hard X-ray intensity doesn't change, but there are some decelerating mechanism or the lost part are compensated by the continuously produced runaways which have lower energy. Thus they reach the dynamic balance. Especially, the newly produced runaways are more than the lost ones.

The strong correlation between sawtooth oscillation of soft X-ray and sawtooth oscillation of the hard X-ray shows that the internal disruption could bring about the runaway electrons and the delay time is about 180us (Fig.3, No: 7967).

We also determined the correlation between oscillation of the hard X-ray sawtooth and radiation oscillation of the superthermal electrons (8.2-10GHz) (Fig.4, No: 8217)[3]. The mean period is about 1.2ms, generally the correlation between them is good.

**MAGNETIC FLUCTUATIONS**

The magnetic coils, which was fixed in the plasma SOL region, could use to study the edge magnetic fluctuations. It has been shown[4] that the macroscope MHD instabilities and the edge magnetic fluctuations associated with the density fluctuations in previous work. The high level correlation between density fluctuation and mirnov perturbation indicated that the vast of particles and energy was lost by the macroscope MHD instabilities. The density fluctuations was great in this case. Compare with this, the density fluctuations was restrained when the discharges was without mirnov oscillation.

We have found another kind of macroscope MHD perturbation which is different from the mirnov's in the experiment also. This kind of MHD perturbation, and, the relations between density fluctuations has been studied in detail. The typical character frequency of the mirnov perturbation is about 5kHz in HL-1 tokamak. Its amplitude was about \( \frac{B_0}{B_0} = 1\% \). Compare with this, the character of frequency of this kind of MHD perturbation is about 20kHz, the mode numbers of poloidal is in the region of 8<\(m<40\), and the amplitude is about 1/10 of mirnov's (as shown in Fig.5). From the Fig.5, we found that its amplitude is unchanging at one shot of discharge. The spectrum of this kind of magnetic fluctuation and density fluctuation show that its cross-correlation value is about R=0.7 (given in Fig.6) and the density's amplitude is greater than without the mirnov's and smaller than the density fluctuation within mirnov perturbations. The Fig.5(d) shown that \( \frac{\Delta n}{\omega} << 0.1 \), it has the characters of typical macroscope MHD perturbation.

From the Fig.7, we found that, this kind of MHD perturbation's frequency is dependence on \( 1/\bar{n}_0 \), and \( q_1 \), where the \( \bar{n}_0 \) is the line-average of electron density and \( q_1 \) is the safe factor in the edge. It conform to the electron diamagnetic drift frequency:
then, we notice that $f_{ed} = 1/n_e$ and $1/r_f$. Of course, the $r_f = 1/q_a$ in commonly. But, we could not observe the relations between frequency of this kind of MHD perturbation and toroidal field $B_T$. In factor, the existence of drift-tearing mode make the magnetic field oscillated which at the local singular surface in the diamagnetic drift frequency. The amplitude of this oscillation do not grow up when the input and output energy is balanced.

The phenomena of this experiment indicated this kind of macroscope MHD fluctuation came from the plasma column. In factor, the developed magnetic island which in the plasma could emit the drastic field ($\partial B/\partial t = -\nabla \Phi$). This field have the electrons compelling motion, and lead to the density fluctuations. Usually, the density fluctuation's frequency spectrum was flatten rapidly by wave-wave interaction and Landau damping. We also found a frequency peak in the density fluctuation spectrum which was compelling motioned is in the mirnov perturbation frequency region, when the magnetic island was too large. However this kind of MHD instability would grow to the mirnov perturbations if the plasma supply enough energy to the magnetic island.

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Fig. 3 Sawtooth oscillations on hard X-ray and central soft X-ray. The sawtooth period is about 2 ms.

Fig. 4 Hard X-ray sawtooth compared with superthermal electron radiation.

Fig. 5 The magnetic fluctuations (a), (b) and (c) are the magnetic fluctuations vs. time (d) and (e) are the frequency spectrum of magnetic fluctuations.

Fig. 6 The spectrum of fluctuations (a) and (b) is the power spectrum of magnetic and density fluctuation respectively. (c) is the cross-correlation of them.

Fig. 7 The magnetic fluctuation's frequency vs. $F_n$ and $\rho_n$.
ABSTRACT: Two-dimensional intensity distribution mappings of pellet ablation cloud trajectories reveal irregular shapes which are interpreted as the first direct visual observation of pre-existing magnetic island structures in the hot core of tokamak plasmas. The apparent erratic character of these features is compared to a theoretical model of magnetic turbulence.

EXPERIMENTAL METHOD
The experimental observations are based on the following facts: when hydrogen (deuterium) pellets are injected into tokamak plasmas, their ablation undergoes strong fluctuations in space and time. Time-integrated photographs of the ablation clouds show a series of radially displaced luminous striations extending on the average along the local magnetic field lines. The emission is essentially due to Balmer line radiation from neutral H (D) atoms. With neon doped pellets, the striations extend over longer distances since neon ions diffusing further also contribute to the radiation. When an ablation cloud is viewed with an oblique angle $\gamma$ relative to the injection plane, the striations appear at radius $r$ with an angle $\alpha'(r)$ relative to the toroidal field direction. By measuring $\alpha'(r)$ at different radii it is possible in most cases to obtain the true magnetic field angle $\alpha$ and, thus the radial profile of the safety factor $q(r)$ [1-3], in good agreement with $q(r)$ obtained by other methods. Pellet emission photographs taken on TFR from 1984 to 1986 show cases where $\alpha$ from one or more striations changes more rapidly along the toroidal direction $z$, but on average each $\alpha$ follows the local $q(r)$ value. Recent pictures of ablation clouds of D2 pellets injected into TORE SUPRA (TS) also exhibit these features. For TFR, the pictures were taken on 'Polaroid' films, while for TS a CCD camera is used. - For calibration purposes a screen with lines in both toroidal and radial direction was introduced into the toroidal chamber and photographed at the same angle as the ablation clouds. For both the TFR and TS ablation clouds, the light intensity has been scanned in radial direction at different toroidal positions and the intensity maxima have been plotted leading to a two-dimensional intensity mapping in the $(r,z)$-plane.

EXPERIMENTAL RESULTS
Figure 1 shows an intensity mapping with 4 striations in TFR discharge Nr. 94351 for a deuterium pellet doped with 1% neon injected during the current plateau. Injection is from left to right where $x$ and $y$ are the photometer scanning co-ordinates. The broken line between the two dots indicates the true toroidal direction; the distance between the dots corresponds to 23 cm in the toroidal chamber in the injection plane. The scanning axis is at an angle of 12 deg. relative to the true radial direction which is perpendicular-
lar to the broken line. \( y = 120 \text{ mm} \) corresponds to \( r = 12.05 \text{ cm} \), \( y = 140 \text{ mm} \) to \( r = 13.57 \text{ cm} \) in TFR. The striation on the right side is relatively straight and its inclination relative to the toroidal direction leads to \( q(r) \approx 2.3 \). The next striation to the left clearly shows radial deviations from a straight line which are outside the experimental uncertainties, but the 'overall direction' still shows an inclination relative to the toroidal direction associated with a resonant surface near \( q = 2 \). The intensity maxima of the two striated features on the left side show radial excursions which resemble magnetic islands as they are known from projections of magnetic surfaces on the \((r,\theta)\)-plane. Within the experimental errors, they are situated near \( q = 2 \). The toroidal length of the two island-like forms is approximately 6 cm with a radial extension of approximately 0.7 cm. The experimental conditions of this discharge are identical to those of discharge Nr. 94349 described in [1b]. No irregular striations where found in this discharge in the same radial region.

Multipellet injection is currently being carried out on TS. Figure 2 shows intensity mappings of 5 striations of the 4th pellet injected into TS discharge Nr. 2687. (There are more striations, but only the most interesting have been plotted). Also shown are the toroidal and radial reference lines. Injection is from left to right. The two striations on the right side are straight lines. Since observation is nearly perpendicular to the injection plane the angle \( \alpha' \) is too small to see the pitch angle. The 3rd striation shows radial excursions and a splitting of its lower half. The two striations on the left yield a picture of an apparent optical island-like structure. From the measured values of \( B_t, I_p, R_0, N_e(r,t) \) and \( T_e(r,t) \) one calculates \( q(r,t) \). It is found that the island-like feature appears at \( q \approx 2 \) within the experimental errors. - It should be mentioned that the picture of the first pellet in this series, which has a different penetration depth than the 4th one, also shows an island-like structure near \( q \approx 2 \).

Further results of our experiments are: the striations seen on pictures of ablation clouds can optically form an x-point due to optical crossing of striations. It seems, that these apparently irregular structures occur more often at low than at high electron density. The exploited pictures yield irregular structures in the region of \( q \approx 1.5 \) and 2. The features seem to be erratic.

INTERPRETATION - A possible theoretical model
It is very likely that these structures are due to small magnetic islands. The crossings would not be true x-points, but projections of radiating fluid tubes rotating around a magnetic island, see Fig.3

It is well known that high-\( m \) number tearing modes and micro-tearing modes are stable under normal tokamak conditions. Other possible instabilities have been considered in the literature, with negative results. The case of the thermal filamentation instability [4-5] has been reexamined, with promising results: in [6], we consider a system of non-linearly interacting chains of islands created by local cooling of island cores and, thus negative current perturbations. The island cores consist of nested KAM torii and are imbedded in a stochastic sea (the Chirikov parameter being slightly above unity). The stochastic sea exhibits anomalous transport, while only collisional transport acts inside the island cores. Our analysis can be sketched briefly as follows (see Fig.4).
1. We relate the half width $\delta r$ of the island core to the half width $\delta_o$ of the virtual separatrix (i.e. the separatrix which would exist in the absence of adjacent perturbations).

2. We establish the relation between $\delta_o$ and a current perturbation $\delta J$ which can be related to a negative temperature perturbation $\delta T(p)$ inside the island using the Spitzer resistivity.

3. We make an energy balance of the island core (i.e. for $p < \delta r$) taking into account a) thermal conduction with a collisional $Xe$ value inside the island remnant (consisting of nested KAM torii); b.) Joule heating which is weaker than outside due to the negative $\delta J$; c.) Impurity radiation of power $P_{\text{rad}}$ with one (several) impurity species: $P_{\text{rad}} = \int N_e f_{\text{rad}}(T_e) dV$.

We find that if $f_{\text{rad}}(T_e)$ is maximum for a temperature $T_o < T_{\text{out}}$, the local temperature of the stochastic sea, a stable solution of the core non linear energy equilibrium state may exist, with $\delta r$ such that the Chirikov parameter is slightly above unity, ensuring the coherence.

If we use values of the parameters (density, temperature, loop voltage, etc.) typical of present day tokamaks, we find island widths of order $\sim 1 \text{ cm}$, in agreement with observations. The erratic character of the observations is in qualitative agreement with dynamical models of non-linearly interacting systems of islands [7].

The pitch of the radiating fluid tubes, computed from the observed inclinations, is too short to be explainable by magnetic field line rotations around the islands (it would need a huge $\delta J$). Although a full understanding is not reached, we believe that a poloidal fluid velocity around the island cores is present. The diamagnetic drift is the more likely candidate, but momentum balance has to be verified.

ACKNOWLEDGMENTS: It is a pleasure to thank M.S.M. BENKADDA and T. EVANS for contributions and stimulating discussions.

Fig. 1. Mapping of intensity maxima in the \((r,z)\)-plane of 4 striations for TFR discharge Nr. 94351. \(B_t=4.5\ T, I_p=183\ kA, R_0=0.98\ m\)

Fig. 2. Mapping of intensity maxima in the \((r,z)\)-plane for 5 striations of the 4th pellet injected into TS discharge Nr. 2687. \(B_t=3.8\ T, I_p=0.994\ MA, R_0=2.40\ m. z\) is in toroidal direction

Fig. 3. Schematic representation of observation conditions which can explain the formation of island-like structures in the focal plane

Fig. 4. Interacting island chains showing residual KAM cores
HIGH FREQUENCY MAGNETIC MODES AND PARTICLE TRANSPORT IN ROTATING AND LOCKED TOKOLOSHE PLASMAS


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Introduction

We report studies of edge $\dot{B}_\theta$ activity and electrostatic fluctuation induced particle transport in Tokoloshe tokamak plasmas during controlled mode locking with external coils. The measurements were made during strong Mirnov activity ($\dot{B}_\theta / B_\theta > 1\%$) with two tearing modes present in the plasma ($m/n = 2/1$ and $3/1$) and $q_\psi(a) \simeq 3.3$. The $\dot{B}_\theta$ measurements were made with a set of twelve coils at different $\theta$, which could be positioned at one of two toroidal locations, another set of eight coils with closer spacing ($\Delta \theta = 10.4^\circ$), and a single $\dot{B}_\theta$ coil which could be moved radially. The particle transport was estimated from $\dot{n}$ and $\phi$, measured with a four tip Langmuir probe. Data was normally recorded with a 1 MHz sampling rate and for transport and $\dot{B}_\theta$ coherency analyses at least eight blocks of data of length 128 $\mu$s were used.

Results

Typical traces from four $\dot{B}_\theta$ coils during external coil induced mode locking are shown in Fig. 1, with the inset showing a distinctive frequency, $f_H \simeq 60 \text{ kHz}$, well above the Mirnov frequency ($f_M < 8 \text{ kHz}$). The amplitude of this high frequency component is a maximum at $\theta \simeq \pi$, unlike the Mirnov amplitude which peaks at $\theta \simeq 0$. Measurements at two toroidal positions suggest these $m = 1$ shifts are not helical (i.e. $n = 0$).

Typical $\dot{B}_\theta$ power spectra from rotating and locked island phases for $q_\psi(a) \simeq 3.3$ are shown in Fig. 2(a) and (b) respectively. The former spectrum is completely dominated by the low frequency Mirnov activity which is broadened by the external coil induced slowing down of the rotation frequency. The locked phase shows the absence of Mirnov activity and a broadband spectrum, typical of many tokamaks, peaking at $\simeq 50 \text{ kHz}$ and falling off at higher frequency as $\simeq f^{-n}$ with $n \simeq 3$. Superimposed on this spectrum is a peak at $\simeq 70 \text{ kHz}$ (c.f. Fig. 1) which is clearly seen to be a mode from the high coherency between two adjacent $\dot{B}_\theta$ coil signals ($\Delta \theta \simeq 25^\circ$). The measured cross phase yields a poloidal mode number $m \simeq 5.5$ (see also below) and the coherence time is very long ($> \tau_E$, where $\tau_E \simeq 1.5 \text{ ms}$). The high frequency mode is much less significant in the rotating phase, as seen from the reference coherency calculated from random test data for the same analysis intervals. Nevertheless a second harmonic is also apparent on many shots (e.g. Fig. 2).

The Mirnov activity and the HF mode clearly both have an $m = 3$ structure strongly distorted by the low aspect ratio toroidicity (Fig. 3). The measured separation
of island 'O' points symmetric about \( \theta = \pi \) is in good agreement with a toroidal calculation, \( \Delta \theta_0 = \frac{2\pi}{m} - 2\epsilon \sin(2\pi/m) \). For \( m = 3 \) this gives an effective \( m \) near \( \theta = \pi \) of \( \frac{2\pi}{\Delta \theta_0} \approx 5 \), consistent with the correlation analysis result. The measured radial decay of \( \bar{B}_\theta \) for the HF mode is also consistent with \( m \approx 3 \). Unlike the Mirnov oscillation, the HF mode is clearly non-propagating (Fig. 3), \( \bar{B}_{\theta H} \sim \cos(\Omega_H t) \cos(3\theta + \delta) \), during both locked and slowly rotating Mirnov activity.

The electrostatic fluctuation induced transport, measured 5 mm into the plasma from the limiter, showed no correlation with the HF mode (or the Mirnov frequency) but a significant increase in particle flux following locking (Fig. 2 and Fig. 4). There is also a large change in plasma potential but a less significant change in \( n_e \).

Discussion

The dominance of the HF mode during the locked phase suggests it could arise from 'bouncing' of a tearing mode about the locking angle, i.e. \( \bar{B}_\theta = \bar{B}_{\theta_0} \sin(m\theta + n\phi + \delta) \) where \( \delta = \delta_0 + \Omega_H t \) for a rotating mode, \( \delta = \delta_2 \) for a locked mode and \( \delta = \delta_2 + \Delta \phi \sin \Omega_H t \) for a locked mode bouncing toroidally about \( \delta_2 \) with amplitude \( \Delta \phi \) and frequency \( \Omega_H \).

For \( \Delta \phi \ll 2\pi \) and \( \Omega_0 = 0 \), this would give \( \bar{B}_\theta \approx \Delta \phi \Omega_H \bar{B}_{\theta_0} \cos(\Omega_H t) \cos(m\theta + n\phi + \delta_2) - \Delta \phi^2 \Omega_H \bar{B}_{\theta_0}/2 \sin(2\Omega_H t) \sin(m\theta + n\phi + \delta_2) \). Thus, for the fundamental frequency, \( \bar{B}_\theta \) is modulated in time by \( \cos(\Omega_H t) \) and in \( \theta \) by the edge tearing mode helicity, as observed. The ratio of the HF mode amplitude to that of the Mirnov mode before locking is then \( \bar{B}_{\theta H}/\bar{B}_{\theta M} = \Omega_H \Delta \phi/\Omega_0 \). From the observed ratio (Fig. 2) we obtain \( \Delta \phi \approx 1^\circ \). The amplitude of the first harmonic is \( \Delta \phi/2 \) times smaller than the fundamental, which is also roughly consistent with observations.

The \( \vec{j} \times \vec{B} \) restoring force for a small displacement in \( \phi \) leads to oscillation with frequency \( \Omega_H \propto \sqrt{(I_2 \bar{B}_r)/\alpha M} \) about the locking angle, where \( I_2 \) is the locking current and \( \alpha \) is the proportion of the total mass, \( M \), of the plasma involved in the motion. For \( \alpha = 1 \) we obtain \( \Omega_H \sim \Omega_0 \). This implies that only a small fraction of the plasma around the edge 3/1 island may be involved in the 'bouncing' (\( \alpha \lesssim 0.01 \)). Bouncing in the poloidal direction would also be consistent with the observations. The fact that the HF mode can be seen before complete locking may indicate that the edge 3/1 island is then locked, while the rotating field is due to the 2/1 mode. The latter has a large 3/1 field component at the wall, due to the low aspect ratio of 2.17, and it is possible that this component is not reconnecting at the 3/1 surface due to the differential toroidal (or poloidal) rotation between the 2/1 and 3/1 rational surfaces.

A striking feature of the locking with external coils, is the (usually) rapid onset of a minor disruption, which appears to be triggered by loss of wall stabilisation of the edge 3/1 mode, followed by a long delay (often \( \gg \tau_e \)) before a major disruption which is triggered by rapid growth of the 2/1 mode. This delay is clearly far too long for the disruption to be the result of a direct MHD effect like loss of wall stabilisation. The transport measurements so far indicate it is precipitated by the marked change in the plasma edge parameters as a result of the locking.
Fig. 1 $\dot{B}_θ$ traces, from four coils at different $θ$, showing mode deceleration induced by 1.5 kA in an $\ell = 2$ external helical coil and (enlarged) the high frequency (HF) mode.

Fig. 3 Measured poloidal structures of Mirnov and HF modes at different times within the period ($T \approx 16\,\mu s$) of the HF mode.

Fig. 4 Particle flux, 5 mm into the plasma, as a function of time for rotating ($Ω = Ω_0$) and $\ell = 2$ coil induced locked phases ($Ω \approx 0$) of Mirnov activity. A $\dot{B}_θ(m = 2)$ Fourier coil trace for one of the shots is shown on the same timescale for reference.
Fig. 2 $\dot{B}_\theta$ and particle transport results, for the time intervals shown in the top traces, during (a) rotating and (b) locked phases ($\ell = 2$ external coil, 1.5 kA, switched on at 10.0 ms). $\dot{B}_\theta$ power from the coil at $\theta = 140^\circ$ and the coherency is between coils at $\theta = 140^\circ$ and $\theta = 165^\circ$. Random data results are dashed.
DENSITY FLUCTUATIONS AT THE SAWTOOTH CRASH IN TFTR

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Introduction

Broadband transient density fluctuations at the sawtooth crash, propagating in the electron diamagnetic drift direction, have been observed in the interior of ICRH and low power NBI heated deuterium plasmas in TFTR. Density fluctuations during the sawtooth crash were measured using an X-mode scattering system operating at 60GHz [1]. Within the spatial and temporal resolution of the scattering system (\(\Delta z \approx 30\text{cm for } k_\theta \approx 7\text{cm}^{-1}\)), transient (100\(\mu\text{s}) enhanced density fluctuations are modulated with the period of the sawtooth precursor suggesting fluctuations are enhanced in an extended region around the x-point of the m=1 island.

In a number of theoretical studies, it has been suggested that enhanced turbulence may account for the rapid crash times and anomalous heat pulse propagation rate observed on many tokamaks [2,3]. In a previous study using O\(_2\) laser scattering, enhanced density fluctuations were found to be correlated with the period of sawtooth oscillations (but not with the much faster period of the m=1 precursor), and it was argued that the heat pulse propagation may be enhanced by an increased level of drift wave turbulence [4]. In a recent study in TFTR, it was suggested that \(\chi_e\) need only increase transiently in the core of the plasma to account for the rapid evolution of the heat pulse following the crash [5]. The present observation of transiently enhanced fluctuations in phase with the m=1 precursor is qualitatively consistent with such a model.

Experimental Overview

Fig. 1 illustrates the typical scattering geometry used in this experiment. The relevant plasma parameters are \(B_T=3\text{T}, I_p=1.2\text{MA}, n_e(0)=3.5\times10^{13}\text{cm}^{-3}, (0)=4.0\text{keV}, R_0=260\text{cm} (\text{with } 15\text{cm Shafranov shift}), a=95\text{cm} \text{ and there is } 6\text{MW of ICRH with no NBI. These plasmas were chosen for their long sawtooth crash times (}\sim400\mu\text{s}) \text{ which allowed for a complete poloidal rotation of the m=1 island.}
mode through the crash. Central density and electron temperatures collapse by approximately 3% and 20% respectively through the crash. The q=1 radius is located at \( r/a \approx 0.2 \) (close to the inversion radius) as inferred from the midplane grating polychromator (GPC) measurements of the radial extent of rotating helical hot spot before the sawtooth crash.

In these plasmas, the period of the \( m=1 \) mode is essentially unchanged through the crash. The GPC is displaced \( \sim 56 \) deg. toroidally from the scattering volume. For the data presented in this paper, the scattering volume is displaced approximately 50 deg. in the poloidal direction from the outer midplane in such a way that a field line near the scattering volume maps into the outer midplane at the GPC toroidal position. The approximate poloidal location of the x-point relative to the scattering volume is inferred from GPC and horizontal X-ray camera measurements of the location and helical pitch of the hot spot at the inversion radius.

**Results**

Fig. 2 shows the correlation between high frequency bursts of fluctuations observed on the scattering system and peak electron temperature at the outer midplane \( R=293 \) cm near the inversion radius as measured by the GPC. Note that no enhancement of fluctuations is observed at the time of the central temperature collapse at \( R=276 \) cm or when the temperature peaked at the inner midplane. Furthermore, two distinct bursts of density fluctuations are observed during the crash corresponding to the x-point appearing twice in the region around the scattering volume. For core fluctuations, \( \dot{n}/n \sim 0.005 \) [1], while the peak enhanced fluctuation level during the crash can be \( 5 \times \) this level averaged over the entire scattering volume.

Comparing power spectra for two very different wavenumbers both during and just before the crash (c.f. Fig. 3) indicates a strong enhancement fluctuations with approximately a linearly dispersive poloidal propagation relative to the plasma motion as may be observed from the bandwidths of spectra for two oppositely directed wavenumbers. The propagation is in the electron diamagnetic drift direction with an inferred phase velocity of \( V_\theta = 2.5 \times 10^5 \text{ cm.s}^{-1} \). The scattered power is enhanced \( \sim 10 \times \) at the peak of fluctuations in narrow time window of \( \sim 10 \mu s \). Fluctuations before the crash propagate in the electron diamagnetic drift direction but with a phase velocity consistent with bulk toroidal rotation. Fig. 4 shows the correlation between the location of the scattering volume and the inferred location of the x-point during enhanced fluctuations at the sawtooth crash for a range of shots with similar characteristics. The correspon-
ence is very good given the various uncertainties associated with the spatial resolution of forward scattering diagnostics operating in the poloidal plane.

Discussion

Whether these transient fluctuations have any role to play in the anomalously short crash times or in the enhanced heat pulse propagation times observed on TFTR is still an open issue. The transient $m=1/n=1$ distribution of enhanced fluctuations at the inversion or mixing radius is certainly consistent with time dependent $\chi_e$ models of heat transport where $\chi_e$ may rise dramatically as long as the helical $m=1$ perturbation persists. In studies on TFTR for a wide range of discharges, up to a 40$x$ enhancement in local $\chi_e$ was needed to model the rapid heat deposition outside the mixing radius [5].

A recent theory of the sawtooth crash, the anomalous current diffusion is mediated by a current convective instability at the x-point [3]. However, the fluctuations observed in our experiment are in the drift wave range ($k_{parallel} \rho_i \sim 1$) which are most likely enhanced by a strong reduction in $L_n$ near the x-point in the crash time, assuming some mixing length condition holds. Enhancement of the drift wave turbulence at the inversion radius is qualitatively consistent with models of core localized transiently enhanced thermal diffusion. However, further information is required on the precise spatial distribution of fluctuations before reasonably accurate estimates of the local variation in $\chi_e$ can be made. Also, future work will need to focus on how the level of transient fluctuations scale with the crash time over a wide range of plasma conditions.

References

[1] Work supported by DOE contract No. DE-AC02-CH03073.


Fig. 1. 60 GHz X-mode scattering geometry for the analysis of core fluctuations in TFTR. The scattering volume gives Δz~30 cm for K₀ ~ 7 cm⁻¹.

Fig. 2. Temperature variations measured at R=293 cm at the midplane at φ=56 deg. correspond to temperature variations near the scattering volume at φ=0 for θ~60 deg. The location of the X-point closely corresponds to the location of the peak temperature near the inversion radius through the crash. Fluctuations peak in the region around the hot spot through the crash.

Fig. 3. Scattered power (200μs integration) increases by ~10x through the crash measured at two different wavenumbers. Enhanced fluctuations propagate in V_de direction relative to plasma.

Fig. 4. Within experimental uncertainties, the inferred location of the X-point through the crash corresponds to the location of the scattering volume during transient bursts of density fluctuations.
LONG WAVELENGTH DENSITY TURBULENCE MEASUREMENTS IN TFTR BEAM-HEATED DISCHARGES

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Introduction

Despite extensive studies over the past decade, the basic mechanism which causes the anomalous cross-field transport observed in all tokamaks has not been conclusively identified. There is considerable expectation that the cause lies in low-amplitude fluctuations of the equilibrium plasma, possibly due to drift-wave turbulence [1,2]. Previous measurements [3] in TFTR from microwave scattering have indicated that the density fluctuation spectrum is a monotonically decreasing function of $k_L$ for all $k_L \geq 2$ cm$^{-1}$, and the peak power in the turbulence spectrum occurs in the long-wavelength range of $k_L \leq 2$ cm$^{-1}$, which has been inaccessible with standard fluctuation measurements. In an effort to further characterize the density turbulence in TFTR and its relation to local plasma transport and confinement, a novel Beam Emission Spectroscopy (BES) [4] density fluctuation diagnostic system has been employed to study long-wavelength density fluctuations and turbulence in TFTR discharges with high-power auxiliary heating. The purpose of this experiment is to measure and characterize plasma density turbulence in the range $k_L \rho < 1$, and thus complement the microwave scattering measurements.

The Beam Emission Spectroscopy diagnostic measures the fluctuations in light emitted by an injected neutral beam as it penetrates into the hot plasma core region. For TFTR plasma parameters and a 60 keV/amu hydrogenic neutral beam, the beam neutrals are collisionally excited by proton and impurity ion impact; hence the observed fluctuations in the light intensity reflect fluctuations in the local plasma ion density in the volume defined by the intersection of the neutral beam and the optical sightline used to observe the beam. With careful choice of beam-plasma viewing geometry, resolutions of 1-2 cm in both the radial and poloidal directions are possible. A 15-channel detector system has been installed on TFTR [5], and it allows simultaneous observations with spatial separations as low as 1.5 cm in the radial and poloidal directions. Each channel integrates all wavenumbers in the fluctuation spectrum from $k_L = 0$ to $\sim 1.5$ cm$^{-1}$ ($k_L \rho \leq 0.5$). Cross-correlation measurements between channels allows determination of the coherence lengths of the observed fluctuations and their apparent direction of propagation in the plasma. The 15 channels can be redirected between shots to allow full coverage over the plasma minor radius.

The BES system became operational near the end of the 1990 experimental campaign on
TFTR, and the emphasis of the first studies has been on the characterization of the observed density turbulence. We concentrate here on the results of a detailed radial scan of the fluctuation spectra obtained for a standard L-mode plasma in TFTR, and on variations of core and edge turbulence as a function of injected beam power.

**Experimental Results**

The results discussed here involve a standard low-β L-mode TFTR discharge, with $I_p=1.0$ MA, $B_t=4.8$ T, $P_{inj}=9.6$ MW (nearly balanced), $R=2.45$ m, $a=0.80$ m, and $q(a)=7.5$. The density fluctuation measurements were made in a 256-msec integration time, during which all discharge parameters were constant. No MHD or sawtooth activity was present.

Power spectra of the local density fluctuations are shown in Fig. 1 for several radii. The power spectrum decreases monotonically from $f=5$ kHz until it merges with the photon statistical baseline in the 50-100 kHz range. As the position is decreased from the limiter radius, the spectrum initially increases in amplitude and broadens in frequency. In a narrow region near $r/a=0.9$, a broadband structure centered around 40 kHz appears. In the core region, $r/a<0.9$, the detectable fluctuation spectrum narrows to the 5-50 kHz range.

Figure 2 shows the total integrated density fluctuation amplitude as a function of radius in the plasma cross section. At the plasma edge region, $0.9<r/a<1.1$, fluctuation levels reach up to $\dot{n}/n=10\%$. As $r/a$ decreases further to $0.5<r/a<0.9$, $\dot{n}/n$ drops abruptly to levels of 1% or less, and then either remains constant or drops slowly as the central core region is approached. Perhaps coincidentally, the electron density scale length shows a sharp drop in the region of the sharp rise in $\dot{n}/n$ at the plasma edge region.

The radial correlation length of this turbulence is less than 1 cm at $r/a=1.0$, and increases as $r/a$ decreases, reaching values of 5.6 cm at $r/a=0.5$. Poloidal correlation lengths have been measured at $r/a=0.95$ and 0.70. At both positions, their values are comparable to the radial values at those positions (i.e., 0.7 and 2.3 cm, respectively), implying roughly isotropic turbulent structures.

Cross-phase ($\Delta \phi$) measurements between channels separated in the vertical (i.e., poloidal) direction indicate propagation in the ion drift direction for all radii in the plasma frame of reference. However, the measured phase velocity in the poloidal direction, $v_{ph}=\omega\Delta \phi/\Delta z$, is found to be equal to the projection of the toroidal rotation velocity on the vertical direction, $v_\psi = v_\psi r/q R$, to within the 10% uncertainty in the measurement of the rotation speeds. Thus, to lowest order, the observed spectrum is generated by nearly fixed structures in the plasma frame traversing the fixed observation point in the lab frame, and our observations are consistent with no propagation in the plasma frame.

Finally, a scan of $P_{inj}$ allows the study of any correlations between this low frequency density turbulence and the global plasma confinement time. Since $\tau_E \sim P_{inj}^{-1/2}$ in these L-mode discharges, a scan in $P_{inj}$ effectively provides a $\tau_E$ scan. Figure 3 shows plots of the total density fluctuation amplitudes versus $\tau_E$ in the core region ($r/a=0.7$) and the edge plasma region ($r/a=0.95$) for $P_{inj}=2 - 20$ MW. No effort was made to keep a constant
plasma density in this scan, and \( n_e \) increases monotonically as \( P_{\text{inj}} \) increased. In general, the global confinement time decreases as the core fluctuation level increases, while there is no evident relation between \( \tau_E \) and the edge turbulence level.

**Discussion**

Overall, the long wavelength spectrum in the hot plasma core of these L-mode discharges is dominated by relatively large-amplitude structures whose spatial extent is large and increasing as one goes into the hot plasma core region. The amplitude of these density perturbations in the plasma core region are correlated to the plasma confinement time. These structures appear to nearly purely growing in the plasma frame, although the present data set does not allow a determination of their average lifetime. This picture of large, slow disturbances in the plasma leading to enhanced local transport contrasts with the conventional model of small-scale drift-wave turbulence at relatively high frequencies in the plasma frame.

The relatively simple picture of long wavelength turbulence presented above is complicated in the plasma edge region by the appearance of a mode propagating in the electron-drift direction with large poloidal and short radial correlation lengths. More complex turbulence spectra can also arise in other operating regimes in TFTR. While low frequency turbulence similar to that described above is evident in the enhanced confinement hot-ion (i.e., supershot) operation regime in TFTR, higher frequency semi-coherent modes also arise in the core region of supershot plasmas, and these modes may be coincident with high \( \beta_{\text{pol}} \) operation. In a similar vein, the low-field, low-current discharges used to study excitation of the Toroidal Alfvén Eigenmode in TFTR [6] show the presence of both the nearly stationary modes at low frequency in the plasma frame and moderate frequency (20 - 80 kHz) turbulence propagating in the electron drift direction deep in the plasma core region.

Future studies in TFTR will consist of detailed parameter scans in different confinement regimes to quantify the relation of these long-wavelength fluctuations to local plasma transport, and to provide comparison to relevant theoretical models of plasma turbulence. Comparisons to other fluctuation diagnostics on TFTR, such as the microwave scattering system and the newly-installed microwave reflectometer, are also planned over a wide range of TFTR operating conditions.

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#Supported by U.S. DOE Grant No. DE-FG02-ER53296 and Contract No. DE-AC02-76-CHO-3073

Fig. 1. Density fluctuation power spectra at different radii in TFTR L-mode. The horizontal line indicates the zero level due to photon statistics.

Fig. 2. Radial profile of rms density fluctuation amplitude in TFTR L-mode discharge. The limiter is at R = 325 cm.

Fig. 3. Local density fluctuation amplitude as function of total confinement time for locations in plasma core and at plasma edge.
The 119 µm laser scattering device on ASDEX was used to investigate the direction of propagation and temporal development of density fluctuations. A description of the scattering system and a summary of previous results is given in [1].

Results from the ohmic phase:

Frequency spectra show an asymmetry which becomes more pronounced with increasing electron line density. By matching the scattering system to the magnetic field pitch on the torus inside and outside, respectively, the dominant part of the scattered power could be localised on the low field side. The resulting propagation is in the electron-diamagnetic direction and changes sign when the toroidal field is reversed, which is in accordance with the assumption of an electron-drift wave. At high densities (SOC-regime) a mean propagation velocity of 1500 m/s is observed in the dominant k-range comparable to the mean el.diamagnetic drift velocity in the gradient region. Making use of the decrease in phase velocity predicted for drift waves at sufficiently high wavenumbers we estimate an upper limit for the poloidal rotation of 500 m/s. Over the whole density range no separate η₁-feature was observed.

Results with NI-heating:

The analysis of the broadband microturbulence is complicated by
the presence of coherent MHD activity which has to be identified by independent diagnostics. In general the frequencies of these coherent magnetic oscillations are seen with laser scattering and in many cases also with microwave reflectometry [2].

a) The L-phase:
Disregarding the contributions from coherent MHD-activity we observe that the wavenumber spectrum shifts to lower k-values compared to the ohmic phase. The total scattered power increases and with co-injection the frequency spectra are shifted to the apparent ion-diamagnetic direction. If modelocking occurs, toroidal rotation is stopped and the frequency shift is drastically reduced.

b) The L-H transition:
In the frequency- and wavenumber range characteristics of the L-phase the scattered power is reduced in both the horizontal channel (equatorial plane) and the outer vertical channel (separatrix region). Fig.1 shows the L-H-transition as indicated by the Dα radiation in the upper divertor chamber. In this shot transport into the divertor is repeatedly reduced before the transition finally occurs. The scattered power in the dominant wavenumber and frequency range follows closely the Dα signal within 100-200 μs. In Fig.2 measurements in the horizontal channel (poloidal orientation of the scattering wavevector) show an increase in the frequency shift towards the apparent ion-diamagnetic direction. Fig.3 shows the development of the high frequency component which begins ≈ 300 μs after the build-up of the transport barrier when the level of Dα radiation in the divertor has reached the H-mode level.

At present there is no theory for drift wave turbulence in steep gradient regions where the gradient length is comparable to, or smaller than the wavelength of the fluctuations.

c) The H-phase:
The fluctuation level can strongly increase again in the course
of the quiescent phase to levels considerably above those of the L-mode without significant changes in confinement. This increase is due to coherent MHD-modes with frequencies up to 30 kHz, and/or narrowband incoherent fluctuations at frequencies around 50 kHz observed close to the separatrix.

Poloidal, perpendicular and toroidal velocity components for impurity ions have been determined spectroscopically in the edge region of ASDEX [3]. The consequences of the radial E-fields thus determined for the propagation of drift waves are not clear at the moment.

ELMs are characterised by a precursor oscillation and the broadband turbulence during the ELM [4]. The appearance of precursor oscillations at 100–200 kHz causes a drastic reduction of the high frequency short wavelength component around 1 MHz, even if no ELM evolves as seen in Fig.2. During the ELM the fluctuation level of frequencies up to 300 kHz increases as seen in Fig.1.

Conclusions:

In ohmic discharges the density fluctuations propagate predominantly in the electron-diamagnetic direction and change direction with Ni co-injection. A strong drop in total scattered power together with a further increase in the frequency shift is observed after the build-up of the transport barrier. Similar observations have been reported on other tokamaks (e.g. [5]). Due to the finite spatial resolution of the scattering system the variation of the fluctuations with local parameters cannot be sufficiently resolved to confirm their nature.

References:
Fig. 1: top: $\text{D}_\alpha$ monitor in the divertor.  
bottom: Scattered power in the vertical channel close to separatrix (15 - 300 kHz). Time window = 8 ms.

Fig. 2: Scattered power as a function of frequency showing the existence of a high-frequency feature propagating in the ion-diamagnetic direction. Frequency range 0 - 2.5 MHz.

Fig. 3: top: $\text{D}_\alpha$ monitor in the divertor.  
bottom: Scattered power in the horizontal channel (960-1460 kHz). Time window = 4.5 ms.
ANALYSIS OF COUPLED TEMPERATURE AND DENSITY PERTURBATIONS USING FOURIER METHODS

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1. INTRODUCTION

In the study of transport in magnetically confined thermonuclear plasmas, perturbative methods have become an important tool. In contrast to power balance analysis, which only measures effective diffusivities, perturbative experiments yield incremental diffusivities, and allow the determination of off-diagonal elements in the transport matrix. Here we consider three methods for the analysis of sawtooth induced heat and density pulse propagation. The first derives the diffusivity from the time-to-peak of the pulse, making also use of the radial decay of the amplitude [1]. The second solves the linearized transport equations numerically and makes a full fit to the measured traces [2], treating the perturbation either as a forced boundary problem or as an initial value problem. The third is based on Fourier analysis of the signals [3,4] and provides the diffusivity as a function of frequency from the phase velocity and amplitude decay of each harmonic component. The three analysis methods are sensitive to different aspects of the data: the time-to-peak method uses only the part of the signal around the maximum, the full curve fitting gives more weight to the tail of the pulse, whereas Fourier analysis allows to separate contributions of different parts of the pulse into different frequency ranges. In Sec.2 of this paper we illustrate the type of information that Fourier analysis can provide and we compare the three methods for JET data. In recent years, it has become clear that diffusive cross-coupling between particle and energy transport may be important [5]. In Sec.3 methods are developed to reconstruct the transport matrix through Fourier analysis of temperature and density perturbations in cases of significant coupling.

2. FOURIER ANALYSIS IN THE LIMIT OF NO COUPLING

Previous work [4] showed that Fourier analysis of JET sawtooth induced heat pulses gives reliable results in a frequency range limited by distortions of the pulse tail at low frequencies (i.e. below 20 Hz) and by lack of cross-coherency at high frequencies (i.e. above 50-120 Hz, depending on $\chi$). The high $\chi/D$ values ($\chi/D \approx 10$) and low $D$ values ($D \approx 0.1-0.4 \text{ m}^2/\text{s}$) observed in JET [2] imply that a slow component in the temperature perturbation due to coupling would be barely visible above the low frequency limit. Temperature perturbations can then be Fourier analyzed under the assumption of negligible coupling. In this limit, the linearized transport equation for a single harmonic component of the temperature perturbation, $T_\omega$, takes the form

$$n_0 \left( \frac{3}{2} i \omega + \frac{1}{1+\tau} \right) T_\omega - V (n_0 \chi VT_\omega + n_0 UT_\omega) = 0 \quad (1)$$

Here $\chi$ is the incremental diffusivity, $n_0$ the equilibrium density, $1/\tau$ an effective damping rate and $U$ the incremental heat pinch velocity, which are all functions of the plasma minor radius. Approximate expressions for the derivatives of the phase ($\varphi$) and amplitude ($A$) of $T_\omega$ can be derived in cylindrical geometry [6] when the effects of damping, heat pinch velocity and spatial gradients are small:

$$\varphi' \approx \frac{k}{\sqrt{2}} \left( 1 - \frac{A}{2} \right) \quad (2)$$

$$-\left[ \frac{A'}{A} + \frac{1}{2} \left( \frac{1}{1+\tau} - \frac{1}{r_n} \right) + \frac{1}{2\chi} (U + \chi) \right] \approx \frac{k}{\sqrt{2}} \left( 1 + \frac{A}{2} \right) \quad (3)$$
where \( r_n = -n_0 / n_0 \), \( k = (3\omega / 2\chi) \)^{1/2} and
\[
\Delta = 2 \frac{2}{3\omega r} + \frac{1}{k^2} \left[ \left( \frac{1}{2} \left( \frac{1}{r} - \frac{1}{r_n} \right) + \frac{U + \chi'}{2\chi} \right)^2 - \frac{U}{\chi} \left( \frac{1}{r} - \frac{1}{r_n} \right) - \frac{U'}{\chi} \right]
\]
is a generalized damping term which scales as \( 1/r \). We can then understand that, while the effect of damping disappears at high frequencies, other terms may still affect \( A' / A \), inducing a difference between \( A' / A \) and \( \phi' \). A more refined treatment of all effects in Eq. 1 [7] yields the following expression for \( \chi 
\]
\[
\chi = \frac{3}{4} \omega \left( \phi' \left[ \frac{A'}{A} + \frac{1}{2} \left( \frac{1}{r} - \frac{1}{r_n} \right) + \frac{U}{2\chi} \right] \right)^{-1}
\]
valid over a wide radial range. This expression for \( \chi \), with \( U = \chi' = 0 \), has been used for the analysis of JET heat pulse data. An example is shown in Fig. 1, where the \( \chi \) values obtained with the extended time to peak method [1] and with the ACCEPT code [2] are also shown. The results of all methods agree within their uncertainties. The absence of frequency dependence of the \( \chi \) values from the Fourier method is consistent with the assumption of negligible coupling effects in the observed frequency range. This type of result is what we generally find for JET data; a different result obtained from the preliminary application of this method to old JET data [4] has not been confirmed by a more extensive analysis [6] of more recent and better diagnosed heat pulses. The influence of the \( (U + \chi') \) term can be determined from the difference between \( \phi' \) and \( A' / A \). As shown in Fig. 2, these are equal within error bars (\( \sigma_{\phi'} = \sigma_{A' / A} = \pm 0.7 \) m/s, corresponding to \( \pm 10\% \) around 50 Hz), which implies \( U + \chi' < 6.5 \) m/s. This is an interesting result that could be improved when data of better quality become available. Also, in this particular case, the effect of damping is seen to be negligible above 20 Hz.

3. FOURIER ANALYSIS IN THE PRESENCE OF COUPLING

The analysis of JET density perturbations or of heat pulses in machines with smaller \( \chi / D \) values may require an extension of the Fourier method to include the effect of coupling. Below we show how the elements of the 2x2 transport matrix can be determined. We consider the coupled transport equations for heat and particle transport, taking only the highest spatial derivatives [5]:
\[
\partial_t u = \mathcal{A} \nabla^2 u
\]
where \( u = (\nu, \tau) = (\tilde{n} / n_0, \tilde{T} / T_0) \) and \( \mathcal{A} \) is the diffusive transport matrix. The solution of Eq. 5 is a superposition of slow and fast purely diffusive components
\[
u = \sigma \hat{\sigma} + \phi \hat{\phi}
\]
where \( \hat{\sigma} = (\epsilon_{\sigma}, 1) \) and \( \hat{\phi} = (\epsilon_{\phi}, 1) \) are the eigenvectors of matrix \( \mathcal{A} \) relative to the eigenvalues \( \lambda_{\sigma} \) and \( \lambda_{\phi} \). From the knowledge of \( \epsilon_{\sigma}, \epsilon_{\phi}, \lambda_{\sigma} \) and \( \lambda_{\phi} \), matrix \( \mathcal{A} \) can be reconstructed[5].

Two methods for the determination of \( \mathcal{A} \) based on Fourier analysis are considered. First, we observe that, if \( \mathcal{A} \) is diagonal, \( \nu \) and \( \tau \) are eigenvectors. Fourier analysis of any perturbation will in this case show frequency independent \( \chi \) and \( D \), which coincide with the eigenvalues of \( \mathcal{A} \). If \( \mathcal{A} \) is not diagonal, the diffusivities \( \lambda \) found by Fourier analysis of the temperature and density perturbations will be frequency dependent (see Fig. 3). The coordinate transformation required to diagonalize \( \mathcal{A} \) can be determined by applying the Fourier method to a weighted superposition of \( \nu \) and \( \tau \) measurements. There are two superposition ratios for which the resultant diffusivity \( \lambda \) is frequency independent. These ratios yield the eigenvectors of \( \mathcal{A} \), the corresponding diffusivity are the two eigenvalues. Thus \( \mathcal{A} \) is determined. This reconstruction method has been tested with simulated heat and density pulses, using the transport matrix \( A_{11} = 1, A_{22} = 4, A_{12} = -0.2, A_{21} = 4 \) and a ratio of initial perturbations \( \alpha = \nu_0 / \tau_0 = 0.3 \). Fig. 4 shows that the eigenvectors are well determined by the minimum of the standard deviation of
the frequency dependent diffusivity: \( \sigma^2 = \langle \lambda^2 \rangle - \langle \lambda \rangle^2 \) where \( \langle \rangle \) denotes averaging over frequency. The test matrix was accurately reconstructed from the simulated data (see Table I).

An alternative method is based on the simultaneous best-fit of the frequency behaviour of the \( \chi(\omega) \) and \( D(\omega) \) obtained from Fourier analysis of temperature and density perturbations. The harmonic components of density and temperature are modelled by the sum of the two slab-model exponential solutions for the slow and fast components

\[
\nu_\omega = \nu_{\omega_0} \left[ \delta_{\omega_0 \nu} \exp(-z_{\sigma}(x-x_0)) + \exp(-z_{\phi}(x-x_0)) \right]
\]

\[
\tau_\omega = \tau_{\omega_0} \left[ \delta_{\omega_0 \tau} \exp(-z_{\sigma}(x-x_0)) + \exp(-z_{\phi}(x-x_0)) \right]
\]

where

\[
z_{\sigma,\phi} = \sqrt{\frac{\omega}{2\lambda_{\sigma,\phi}}}(1 + i)
\]

and \( \delta_{\omega_0 \nu} = \nu_{\omega_0} / \sigma_{\omega_0} \), \( \delta_{\omega_0 \tau} = \tau_{\omega_0} / \tau_{\omega_0} \) are the (a priori complex) ratios of the slow and fast components of density and temperature at frequency \( \omega/2\pi \) and at \( x = x_0 \) where the boundary condition is assigned. \( \delta_{\omega_0 \nu} \) and \( \delta_{\omega_0 \tau} \) are linked to the analogous ratios \( \delta_{\nu} \) and \( \delta_{\tau} \) defined for the eigenmode analysis in the time domain [5] by

\[
\delta_{\nu} = \delta_{\omega_0 \nu}(\sigma_\sigma / \phi_\phi)^{1-\gamma} \quad \delta_{\tau} = \delta_{\omega_0 \tau}(\sigma_\sigma / \phi_\phi)^{1-\gamma}
\]

where the exponent \( \gamma \) depends weakly on the spatial shape of the initial temperature and density perturbations and is found to be \( \gamma \approx 0.72 \) for typical perturbations. The parameters \( \sigma_\sigma, \phi_\phi, \sigma_{\omega_0 \nu} \) and \( \tau_{\omega_0 \tau} \) (and hence \( \delta_{\nu} \) and \( \delta_{\tau} \)) are identified by the best-fit together with the mixing radius position \( x_0 \). From these and from the independent knowledge of \( \alpha \), matrix \( A \) can be reconstructed, according to [5]. The results of the application of the method to the same simulated case are reported in Fig.5 and Table I. The first method has the advantage of being independent of the type of the density and temperature perturbations, the second relies on the assumption that they are equal but allows, in principle, the determination of \( \sigma_\sigma \) and \( \phi_\phi \) from either temperature or density measurements. Clearly, both methods give the correct results when applied to simulated data. The sensitivity of both methods to experimental uncertainties in the data is presently under investigation.

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Acknowledgement

The authors express their gratitude to JET, in particular the electron temperature group, for the use of heat pulse data, and to A.C.C.Sips and G.Kramer for help with the ACCEPT code. This work was performed under the Euratom-FOM and the Euratom-ENEA-CNR association agreements, with financial support from NWO, CNR and Euratom.
Fig.1 $\chi$ values determined by the Fourier method for JET shot 15011. The lines represent the $\chi$ values obtained by the time-to-peak method and the ACCEPT code.

Fig.2 $\phi'$ and $A'/A$ values obtained by means of the Fourier analysis for JET shot 15011.

Fig.3 Frequency dependence of $\lambda$ for different ratios of the slow and fast components. $\lambda_\phi$, $\lambda_\sigma$ are the same of Tab.I.

Fig.4 Standard deviation of $\lambda(\omega)$ from its average value as a function of the density and temperature superposition ratio, for the simulated case of Tab.I.

Fig.5 $\chi(\omega)$, $D(\omega)$ values obtained by the Fourier analysis of the simulated case of table I. The continuous lines represent the result of the best fit obtained by means of the second method.

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<th>Method 2</th>
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Table I
TWO PHASE REDUCTION OF MICROTRUBULENCE AT THE TRANSITION INTO H-MODE MEASURED ON DIII-D*


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At the L-H transition in the DIII-D tokamak, a two phase reduction of microturbulence has been observed. A first, rapid (∼100 μs) suppression phase of microturbulence in a narrow region of the plasma edge is followed by a second, bulk turbulence reduction phase in the plasma interior, with a relative turbulence reduction (n'/n) of >50% below L-mode level. This second, slow reduction phase (within 10's of ms) correlates with the results of recent transport studies on DIII-D, suggesting that transport coefficients decrease in a comparable timescale in the same region in the plasma. We have established a quantitative relationship between frequency shifts in far infrared scattering measurements and $E \times B$ plasma velocities measured with charge exchange recombination spectroscopy. The measured profiles correspond to a positive radial electric field, in the interior of the plasma, which predicts a frequency shift in the ion diamagnetic drift direction of the same magnitude as that actually observed in the scattering data. This analysis, together with results from reflectometer measurements, radially locate the density fluctuations and reflect their time evolution.

DIAGNOSTIC INSTRUMENTS

Density fluctuations are measured using FIR collective Thomson scattering and reflectometry. The far infrared scattering system$^1$ is based on a 245 GHz (λ = 1.22 mm) twin frequency laser and uses a heterodyne detection technique which enables us to determine the propagation direction of fluctuations. Three receiver channels are available to study poloidally propagating fluctuations along the entire midplane of DIII-D in a wave number range of 2 ≤ $k_θ$ ≤ 16 cm$^{-1}$. This system, with its good wave number resolution but moderate spatial resolution at low wave numbers, is complemented by reflectometer systems,$^2$ which have a larger wave number sensitivity, but very good spatial resolution. The homodyne reflectometer system utilizes seven discrete O-mode channels spanning from 15 to 75 GHz and a frequency tunable X-mode channel (50–75 GHz) with corresponding critical densities of 0.28–7.0 × 10$^{19}$ m$^{-3}$.

Ion temperature and toroidal and poloidal rotation profiles are measured using a multichord, high spatial resolution charge exchange recombination$^3$ (CER) system. From these measurements, the radial electric field $E_r$ is inferred using the lowest order force-balance equation for a single-ion species

$$\left(1/Z_1 n_I \right) \partial p_I / \partial r = e E_r - e/c \left( v_\phi B_θ - v_θ B_\phi \right),$$

where $Z_1$, $n_I$, and $p_I$ stands for the charge state, density and pressure, respectively. Frequency spectra from poloidally propagating density fluctuations, recorded by scattering diagnostics in the laboratory frame of reference, are known to be Doppler shifted$^4$ due to $E_r$ induced plasma flows. Calculating $E_r \times B$ plasma velocities enables us to determine these frequency shifts. The turbulence mode frequency in the plasma frame $\omega_\rho$ and its propagation direction and phase velocity can now be determined, subtracting the Doppler shifted part $\omega_D$ from the measured spectra $\omega$.

---

* Work supported by the U.S. Department of Energy under Contract Nos. DE-FG03-86ER53225 and DE-AC03-89ER51114.
† General Atomics, San Diego, CA.
‡ Japan Atomic Energy Research Institute.
\[
\omega_o = \omega - \omega_D \quad \text{with} \quad \omega_D = -clE_r/B \quad .
\]

However, in the measurements shown, any reference to ion or electron propagation refers to a propagation in the ion or electron diamagnetic drift direction measured in the “lab frame”.

**FREQUENCYhifts Related to \(E_r \times B\) Flows**

General characteristics of density fluctuations measured in Ohmic DIII-D plasmas\(^5\) are comparable to results from other machines and are in qualitative agreement with density gradient driven drift wave turbulence. Figure 1 shows scattering spectra in Ohmic discharges as well as before and after the L-H transition. Neutral beam injection increases the fluctuation level at low wave numbers. Due to increasing plasma rotation in the L-mode phase, fluctuations are spread over a broad frequency range and the center of the bulk of the fluctuations is located on the ion diamagnetic drift side. After the L-H transition, a rapid reduction of low frequency fluctuations is observed and the center of residual fluctuations is rapidly shifted to the ion side. Spectra measured at \(k_\theta = 2.5 \text{ cm}^{-1}\) show a smaller frequency shift (relative to \(5 \text{ cm}^{-1}\)) as would be expected from Eq. (2). These fluctuation spectra as well as reflectometer and rotation data were measured during a series of single-null divertor discharges with \(I_p = 1.6 \text{ MA}, B_1 = 2.1 \text{ T}\) and 8 MW of D\(^\circ\) neutral beam power co-injected (same direction as \(I_p\)) into a D\(^+\) plasma.

A negative (inward pointing) radial electric field was inferred in the plasma edge, between normalized \(\rho (\rho_n)\) of 0.85 and the separatrix. Field values of 0 to \(-100 \text{ V/cm}\) in L-mode [Fig. 2(a)] are changing to more negative values of 0 to \(-400 \text{ V/cm}\) in H-mode. However, in the interior of the plasma, \(E_r\) was found to be positive (outward pointing) in these co-injection plasmas. In the interior, an increase of \(E_r\) after the transition is mainly due to an increase of the toroidal rotation and to a lesser extent to the flattening of the pressure profile in H-mode. Positive \(E_r\) together with the total magnetic field induces \(E_r \times B\) plasma velocities and hence Doppler shifts on fluctuation measurements in the ion diamagnetic drift direction. Negative electric fields, as observed in the plasma edge, produce shifts in the electron diamagnetic drift direction. Figure 2(b) shows profiles of the calculated frequency shifts to be expected on scattering spectra for \(k_\theta = 5 \text{ cm}^{-1}\).

Mixing length estimates and experimental evidence locate the bulk of the fluctuations outside \(\rho_n = 0.6\), where according to our calculations, large, radially dependent \(E_r \times B\) velocity shifts are to be expected. Frequency shifts observed in measured scattering spectra, are known to be composed of \(E_r \times B\) velocity shifts and possible shifts due to phase velocity changes of the turbulence. However, changes in phase velocity are expected to be small compared to \(E_r \times B\) velocities and are, therefore, neglected. Hence, assuming a fixed mode frequency for Ohmic, L-mode, and H-mode turbulence, a comparison of measured and calculated frequency shifts allows us to radially locate the bulk of the density fluctuations.

In L-mode, 7 ms before the transition, measurements at \(k_\theta = 5 \text{ cm}^{-1}\) show the center of the bulk of the fluctuations to be shifted by about 200 kHz to the ion side. Calculations predict that amount of shift around \(\rho_n\) of 0.8. Four milliseconds after the H-mode transition, measured spectra show the center of the remaining fluctuations shifted by about 600 kHz to the ion side which according to the calculations, locates them to \(\rho_n = 0.78\). Error bars on the inferred electric field of ±70 V/cm result in error bars of about ±250 kHz on calculated shifts at \(k_\theta = 5 \text{ cm}^{-1}\). Accordingly, by comparing the measured and calculated frequency shift and including the above mentioned error calculations, we locate the bulk of the fluctuations to be within \(\rho_n\) of 0.7 and 0.9.
TWO PHASE REDUCTION OF TURBULENCE

An initial fast suppression phase of microturbulence at the plasma edge is best observed by reflectometer measurements\(^6\) (Fig. 3). At 1862.7 ms, a fast drop within \(\approx 0.1 \mu s\) together with the first decline of \(D_\alpha\) emission is shown by all channels \(\leq 40 \text{ GHz}\). The 50 GHz channel is delayed by more than 1 ms before its reflecting layer is shifted into the suppression zone by the steepening of the density profile. Reflectometer measurements indicate the suppression zone to be between \(\rho_n = 0.9\) and the separatrix. The far infrared scattering channels which integrate over a larger plasma region appear to see a smaller drop of the fluctuation level during this first suppression phase. However, these observations are in agreement with the result of our rotation analysis which indicate that the bulk of the fluctuations is inside \(\rho_n = 0.9\).

A second phase of microturbulence reduction, subsequent to the fast build-up of the transport barrier, reduces the bulk of the turbulence in a much slower time scale (Fig. 4). Within tens of milliseconds \(\bar{n}/n\) drops more than 50% from its pre-transition level. A significant increase of fluctuations to about I-mode level is observed during the ELM, and in a transition-like time frame the fluctuations resume the H-mode level. A predominant increase of turbulence at low wave numbers in I-mode, and a stronger suppression of high wave number fluctuations in H-mode suggest a narrowing of the wave number spectra with neutral beam injection.

This second slow reduction phase correlates with recent transport measurements made on DIII–D.\(^7\) An improvement in local transport, predominantly in the electron channel, has been found in the same plasma region, with an evolution over a similar time scale. The increasing \(E_r \times B\) velocity shear in the plasma interior after the L-H transition, correlates with the measured broadening and frequency shifts of the turbulence spectra, and may be responsible for the observed reduction of microturbulence.

SUMMARY

A two phase reduction of microturbulence at the transition into H-mode has been measured on DIII–D. These turbulence reductions have been related to the evolution of radial electric field induced plasma flows which may cause the suppression. A fast first suppression phase of fluctuations is observed with the formation of the transport barrier in a narrow region at the plasma edge. In a second slow phase, the bulk of the turbulence, localized in a region of \(\rho_n\) of 0.7 to 0.9, is measured to be markedly reduced within 10's of milliseconds after the transition. This slow reduction phase of microturbulence correlates with improved local transport, measured on a similar time scale and in the same location in the plasma.

REFERENCES

Fig. 1  FIR scattering frequency spectra measured at edge of DIII-D outer midplane.

Fig. 2  Radial electric field a) and corresponding frequency shifts b) in the outer midplane before and after the transition, inferred from pressure profiles and plasma rotation measurements.

Fig. 3  Frequency integrated power (40 kHz-1MHz) from Reflectometer and FIR scattering around the L-H transition.

Fig. 4  Frequency integrated normalized fluctuation level \( \tilde{n}/n \) in Ohmic, L-mode and H-mode up to the first giant ELM.
ELM PRECURSORS ON DIII-D*

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On DIII-D, density and magnetic fluctuation precursors are routinely observed
prior to the occurrence of a variety of edge localized modes (ELMs). Identification of these instabilities may provide a way to identify the limiting plasma mechanism(s) responsible for ELM activity, and thus significantly improve our understanding of the H-mode. Such an understanding is highly desirable as H-mode operation is one of the approaches to breakeven on current machines (JET), while the design of the next generation of ignited devices assumes H-mode like confinement (BPX and ITER).¹ ELMs² are currently an unavoidable feature of H-mode operation, and they play a large role in determining the quality of the plasma discharge; ELMs are marked by sudden, repetitive spikes in recycling light, during which the H-mode transport barrier is reduced and particles and energy are expelled from the edge plasma region into the scrape-off layer. Consequently, ELMs degrade the plasma confinement and are thought to be the limiting mechanism on the improved performance observed during H-mode.¹ At the same time, however, ELMs can also be beneficial in controlling impurity accumulation; by controlling ELM frequency, a 10 s quasi steady state discharge has been achieved on DIII-D.³ Thus, an improved understanding of the ELM trigger mechanism may lead to advances in ELM control and hence improved plasma operating parameters.

The precursors observed on DIII-D are of different types depending on the
variety of ELM: Before giant ELMs, where the edge plasma is close to the ideal ballooning limit, both an increase in the general broadband density turbulence, and the appearance of quasi-coherent modes, have been observed. These features appear on a time scale of up to 30-50 ms before the start of the ELM, and the increase in density fluctuations occurs in the high density gradient edge plasma region. Before "Type III" ELMs, where the edge plasma is well below the ideal ballooning limit, coherent modes of 50-120 kHz center frequency are observed, on time scales of up to 10-20 ms before the ELM. The radial extent of these modes appears to grow as a function of time until the edge plasma is perturbed, triggering the ELM.

INSTRUMENTATION

As the results presented below depend largely on the spatially localized measurements provided by the reflectometer system, the fluctuation diagnostics employed on DIII-D are briefly reviewed. The reflectometer system normally used for fluctuation studies on DIII-D is an O-mode, homodyne system utilizing 7 discrete channels spanning 15-75 GHz, with corresponding critical densities of $2.8 \times 10^{18}$ to $7 \times 10^{19} \text{ m}^{-3}$.(⁴)

* This work was sponsored by General Atomics Subcontract SC120536 under U.S. Department of Energy Contract No. DE-AC03-89ER51114, and by U.S. DOE Grant No. DE-FG03-86ER53225

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The GUNN sources in this system are only narrowly tunable in frequency, so the critical densities are essentially fixed. The spatial resolution of reflectometer measurements has been both theoretically predicted, and experimentally determined, to be localized to the vicinity of the critical layer. This spatial localization can also be directly observed in the reflectometer data on DIII-D, where reflectometer channels with critical layers lying within ~1 cm of each other generally exhibit very low cross-coherence. Thus, by using the multichannel reflectometer system, a detailed picture can be formed of the radial variation of density fluctuations in the DIII-D edge plasma region during H-mode operation. The other relevant fluctuation diagnostics on DIII-D are an FIR scattering system, and both poloidal and toroidal arrays of Mirnov probes. The multichannel, heterodyne FIR system operates at 1.2 mm and provides wavenumber and spatially resolved measurements over a poloidal wavenumber range of 2.5–16 cm⁻¹. At \( k = 10 \text{ cm}^{-1} \), the radial spatial resolution is ±20 cm.

**TYPES OF ELM**

A description of ELM activity is complicated by the large variety in size, frequency, and type of ELM encountered. On DIII-D, at least three different types of ELM have been identified:

1. **Type I, "Giant ELMs."** These occur when the edge plasma reaches the ideal ballooning limit. Their repetition frequency increases with power and target density, and decreases with increasing current. They are marked by large \( \text{D}_\alpha \) spikes (above L-mode levels).

2. **Type II, "Grassy ELMs."** These are irregular, high frequency, low amplitude ELMs, which appear when the plasma is in the connection region between the first and second stable ballooning regimes.

3. **Type III.** These occur in plasmas well below the ideal ballooning limit and are of medium amplitude. As the input power is increased, their repetition frequency decreases and they disappear.

All the available data on DIII-D suggest that giant ELMs are a transient return to L-mode; the fluctuation levels, edge rotation, density and temperature profiles all return to L-mode like conditions.

**ELM PRECURSORS**

In contrast to the L-H transition, edge fluctuation spectra are generally observed to change prior to (rather than coincident with) the first modification in \( \text{D}_\alpha \) emission at an ELM. For giant ELMs, the general broadband turbulence level on the edge reflectometer channels is observed to slowly increase before the ELM, with time scales of up to 30–50 ms before the start of the ELM itself. This may be due to the fact that the quantities governing the reduction in edge turbulence during the H-mode, \( E_r \) and \( v_\theta \), saturate shortly after the L-H transition, whereas possible fluctuation driving terms, such as density and temperature gradients, continue to increase. It should be noted that this increase cannot be explained by the (theoretical) dependence of the reflectometer signal on the density gradient, as this would predict a reduction in the measured signal level during this period. In addition, quasi-coherent precursors are on occasion observed before giant ELMs, sometimes with a "bursting" character.
By contrast, quasi-coherent fluctuations are routinely observed before Type III ELMs. These precursors have center frequencies of 50–120 kHz and are observed up to 10–20 milliseconds before the ELM. An example of the coherent features observed in reflectometer spectra when these precursors are present is given in Fig. 1. Again, these precursors are only observed on edge reflectometer channels, indicating that they are localized to the high gradient edge region. These coherent reflectometer signals are strongly correlated with Mirnov probe signals, from which toroidal mode numbers of 6–12 have been deduced, and are also observed by the FIR scattering system. In addition to being localized to the plasma edge, the signals also appear to have a ballooning character, i.e. the magnetic signals are strongly weighted to the outboard midplane. While the precursors referred to above clearly occur on a resistive time scale, the changes at the initiation of the ELM itself occur on a more ideal time scale (typically \( \leq 300 \mu s \)). Another observation about both forms of ELM precursor is that the maximum time before an ELM at which a precursor is observed is proportional to the time between ELMs. This suggests that when ELMs are less frequent, the relevant plasma instability is simply being approached more slowly; e.g. giant ELMs occur more frequently as input power increases, the pressure profile builds more rapidly, causing the instability to grow more quickly.

However, the most important feature of the quasi-coherent precursors is that their radial extent appears to grow until the edge plasma is perturbed, triggering an ELM. This is illustrated in Fig. 2, where the precursor can be seen to initiate on the 1.3 and \( 2.0 \times 10^{19} \text{ m}^{-3} \) channels, and then spread sequentially to lower density channels, until the perturbation appears to reach the separatrix, triggering the ELM. Note that the highest density channel displayed \( (3.1 \times 10^{19} \text{ m}^{-3}) \), does not observe either the precursor or the ELM itself, thus demonstrating the edge localized nature of both phenomena. Attempts to identify these modes are currently underway, and will be greatly facilitated by the recent commissioning of a new multipulse Thomson scattering system on DIII-D. This system can provide profiles every 25 ms, which combined with other profile diagnostics (CER etc.), and reproducible plasmas, can provide a very detailed time history of \( T_i, T_e, n_e \) and pressure profile evolution in the critical period leading up to the ELM initiation. Detailed data on profile evolution should be available by the time of the EPS conference.

REFERENCES

Fig. 1. Power spectrum of 32 GHz reflectometer channel (critical density of $1.3 \times 10^{19} \text{ m}^{-3}$), showing typical quasi-coherent Type III ELM precursor features.

Fig. 2. Spatial growth of a quasi-coherent precursor to a Type III ELM. Illustrated is the frequency integrated power in five reflectometer channels, plus a $D_\alpha$ monitor. The precursor can be seen to initiate on the 1.3 and $2.0 \times 10^{19} \text{ m}^{-3}$ channels, and then spread sequentially to lower density channels, until the perturbation appears to reach the separatrix, triggering the ELM.
STUDY OF EDGE ELECTRIC FIELD AND EDGE MICROTURBULENCE AT THE L–H TRANSITION IN DIII–D*


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Experimental studies of the L–H Transition in the DIII–D tokamak clearly show that the edge poloidal rotation velocity and the shear in the edge electric field increase dramatically at the L–H transition, and that the edge density fluctuations are simultaneously suppressed in a narrow region just within the separatrix. The greatest changes in the poloidal rotation and electric field are within 3 cm of the separatrix and overlap with the region of suppressed microturbulence. Furthermore, this zone of sheared electric field and suppressed fluctuations also marks the location of the transport barrier, as indicated by the increased gradients of $T_i$, $n_e$, $T_e$ and the carbon density after the transition. The shear in the edge electric field is maintained later into the H–mode with the sign of the electric field remaining negative as a result of the pressure gradient contribution to the total electric field becoming more significant than the rotational term, which becomes more positive with time into the H–mode.

Theoretical studies of the L–H transition have focussed on the role of the radial electric field in suppressing microturbulence at the plasma edge which then leads to improved edge confinement. The theory of Shaing and Crume predicts that a negative radial electric field or a positive value of $dE_r/dr$ could suppress the microturbulence. Biglari, Diamond and Terry indicate that an increase in the magnitude of the shear of the poloidal rotation could suppress microturbulence, although they do not specifically point to a bifurcation condition relevant to the L–H transition. Previous experimental data from DIII–D and JFT–2M indicate increased poloidal rotation, $\psi_p$, and a more negative radial electric field at the transition, and reflectrometry measurements on DIII–D indicate suppression of density fluctuations in a narrow region at the plasma edge at the transition. These observations of the importance of the radial electric field at the transition are substantiated by experiments on CCT and TEXTOR where transitions to H–mode plasmas have been produced by inducing a radial electric current by intersecting a biased electrode into the plasma.

The measurements of the ion temperature and poloidal and toroidal rotation in the DIII–D plasma are carried out by active Charge Exchange Recombination (CER) Spectroscopy of Doppler shifted and broadened line emission from the C VI ion resulting from single charge exchange interactions with atoms from the injected neutral beams used for plasma heating purposes. The CER system comprises 32 viewing chords of which fifteen intersect the plasma edge and span the separatrix. Eight viewing chords intersect the edge in the poloidal plane in an inter-weaving arrangement. Seven chords lie in the toroidal midplane with a chord to chord separation of 0.3 cm for $T_i$ measurements and 0.6 cm for poloidal and toroidal rotation measurements. This is an improvement in spatial resolution over previous experimental results presented in Ref. 7 where the chord to chord separation was 1.5 cm for the rotation measurements and 0.8 cm for the ion temperature measurements. The nominal spatial resolution is 0.5 cm. The minimum integration time is 1 ms, limited by the signal strength of the C VI emission lines for the given plasma and neutral beam parameters studied.

The electric field is determined from the CER rotation data using the lowest order force balance equation for a single species,

$$E_r = \frac{1}{n_i Z r e} \left( \frac{\partial P_i}{\partial r} \right) - \left( \vec{u} \times \vec{B} \right)_r,$$

* This work was sponsored by the U.S. Department of Energy under Contract No. DE–AC03–89ER51114.
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where \( n_f \) is the ion density, \( Z_i \) is the ion charge number, \( e \) is the electron charge, \( P_f \) is the ion pressure, \( v_f \) is the ion fluid flow velocity and \( B \) is the total magnetic field. Whereas the pressure gradients and poloidal rotation velocities of different ion species may be different,\(^{16}\) the value of \( E_r \) determined from the above equation applies to all the plasma species.

RESULTS

The evolution of the profiles of the radial electric field through the L-H transition is shown in Fig. 1. The distinctive feature is that the shear in the radial electric field at the plasma edge increases substantially at the L-H transition. The time steps are 3 ms and reflect the integration time of the measurements. The times of the profiles are shown with respect to the time of the L-H transition. The shaded region marks the 0.5 cm uncertainty in the magnetic determination of location of the separatrix at the midplane. The radial electric field is zero or slightly negative in the L-mode, but becomes substantially negative within a shear layer of 3 cm inside the separatrix at the L-H transition. Simultaneously with the increased shear of the radial electric field at the L-H transition, the gradients of the ion temperature, \( T_i \), and the carbon density are observed to increase as shown in Figs. 2 and 3. These profiles are characteristic of neutral beam heated hot-ion H-mode plasmas,\(^{16}\) with relatively low edge electron densities giving rise to high edge \( T_i \) values. Within 2 ms of the L-H transition, the edge \( T_i \) gradient has increased significantly from the L-mode value and continues to increase such that the \( T_i \) values are nearly doubled 23 ms after the transition. Measurements of \( T_i \) at 100 ms after the L-H transition show \( T_i \) gradients of 1 keV/cm just within the separation. The gradients of the edge carbon density also exhibit substantial increases after the transition as shown by Fig. 3, which incorporates a relative chord to chord intensity calibration of the detection system. The spatial steepening of the profiles occurs over the region which corresponds to the shear in the radial electric field and marks the location of a transport barrier formed at the transition.

Figures 2 and 3 also indicate that the edge carbon pressure gradients become very large with time into the H-mode. The importance of this behavior is shown in Figs. 4 and 5, which show the edge radial electric field profiles and just the value of \( E_r \) determined from only the \( v \times B \) term, respectively, as they develop in time later into the H-mode. With the evaluation of \( E_r \) from just the edge rotation (Fig. 5), the edge radial electric field increases and actually becomes positive about 65 ms after the transition mainly as a result of the increasing toroidal velocity at the edge. However, with the inclusion of the pressure gradient term of \( E_r \), the edge radial electric field remains negative just within the separatrix and the shear in the electric field is maintained through the H-mode (Fig. 4). The spatial extent of the shear layer in the radial electric field and the transport barrier formed at the edge are then obviously interdependent later in the H-mode and this result emphasizes the importance of including the pressure gradient term in the determination of \( E_r \). At major radius \( R \leq 2.26 \text{ m} \), the radial electric field becomes positive as a result of the increased toroidal velocity.
Fig. 2. The ion temperature profiles as a function of the major radius, $R$.

Fig. 3. The carbon density profiles as a function of the major radius, $R$.

Fig. 4. The edge radial electric field as a function of the major radius, $R$, for times later into the quiescent H-mode plasma.

Fig. 5. The edge radial electric field determined from only the rotation term of the force balance equation for the same times with respect to the L-H transition as shown in Fig. 4.
rotation velocity towards the center of the plasma where a frequency shift due to strong positive electric field is consistent with FIR scattering data.\textsuperscript{17}

The simultaneous occurrence of electric field shear, fluctuation suppression, and transport reduction is in agreement with theories that predict the stabilization of edge turbulence by sheared poloidal flow.\textsuperscript{3,4} However, the bifurcation theory of Shaing and Crume predicts a scaling of the width of the stabilization layer with poloidal gyro-radius, i.e., with the plasma current, but the width of the shear layer determined from reflectometry data was basically the same in plasmas with 0.8 MA and 1.6 MA current. Therefore, the mechanism for the production of the edge radial electric field through ion orbit losses at the plasma edge needs to be improved.

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FRAMING CAMERA STUDIES OF THE EDGE IN ROTATING AND LOCKED TOKOLOSHOE PLASMAS.


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Introduction

We report on the use of a framing camera to study the edge plasma of Tokoloshe tokamak during both high \( \frac{\dot{B}_\theta}{B_\theta} \gtrsim 0.5\% \) and low \( \frac{\dot{B}_\theta}{B_\theta} \lesssim 0.02\% \) Mirnov activity. The purpose of this study is to determine the position and structure of tearing modes locked at the plasma edge by an external resonant \( \ell = 2 \) or \( \ell = 3 \) winding, as well as possible coil induced islands and to characterise macroscopic bright filaments in the plasma edge. The camera is used to look tangentially at the fully poloidal limiter over the entire minor cross-section of the plasma. The framing rate was \( \lesssim 1000 \) frames per second with shutter speeds between 20 and 200\( \mu \)s.

Observations

Framing camera pictures of a reference shot show very clearly the contraction and inward motion of the hot central region together with the expansion of the \( H_\alpha \) annulus during a thermal collapse which ultimately leads to a major disruption [1]. Also seen in the \( H_\alpha \) halo are poloidally distributed, stationary \( (f_\theta \lesssim 10 \text{ Hz}) \) localised, bright filaments in the plasma edge. These filaments are seen most clearly in the outer half of the plasma edge, appearing soon after current profile peaking. They persist to the disruption, often becoming more intense closer to the disruption. To enable a quantitative analysis, cine frames were digitised and computer enhanced [2]. These were analysed for the harmonic content of the poloidal intensity distribution profile at the limiter radius (Fig 1). A broad spectrum of modes centred about \( m = 24 \) is seen. This spectrum of modes is due partly to the low aspect ratio of Tokoloshe \( (R/a = 2.17) \), and the asymmetrical interaction of the plasma with the limiter. The filaments appear to lie along the field lines with a helicity consistent with \( q_\phi(a) \sim 3.5 \). A preliminary correlation analysis of the intensity distribution for consecutive frames puts a lower limit on the coherence time of individual filaments of \( \simeq 1 \) ms.

The \( \ell = 3 \) coil is energised during high Mirnov activity by a pulse with peak current of 400\( \text{A} \) and duration \( \gtrsim 10 \text{ ms} \), which is sufficient to lock the tearing mode for \( \simeq 2 - 7 \text{ ms} \) before inducing an early major disruption. Cine pictures show localised stationary bright regions with an \((m = 2)\)-like mode structure (Fig 2) which persist to the disruption and often become more intense just prior to the disruption. The disruption occurs within one frame interval \( (\lesssim 1 \text{ ms}) \). These bright regions show a change of position and intensity when the helical coil current is reversed. For the \( B_\phi, I_p, I_\ell = (-++,+) \) configuration [3] the preferred locking orientation results in poloidally localised bright regions on the inside, whilst the \((-+-)\) configuration gives more diffuse regions of greater poloidal extent on the outside. Fieldline tracing calculations show that these distortions are consistent
with a low aspect ratio machine. Similar locking positions and intensity distributions are obtained if the \( \ell = 2 \) coil is energised during high Mirnov activity by a current pulse of 1800A. The locked structure can also show limited poloidal rotation with an effective poloidal velocity \( v_p \approx 120 \text{ ms}^{-1} \) (cf. \( v_p^{\text{Mirnov}} \approx 1.2 \times 10^4 \text{ms}^{-1} \)).

When a short duration (\( \leq 3 \text{ ms} \)) current pulse is used, a locking orientation which is independent of the sense of the helical coil current direction is observed. A probable explanation is that the plasma locks to the coil field, then unlocks, and re-rotates to lock to a stray field.

To study coil induced islands during low Mirnov activity, the \( \ell = 2 \) coil was energised by a current pulse of 1400A. This coil current is just below that normally required to precipitate a major disruption [3]. However, cine frames showed no evidence of changes in the poloidal \( H_\alpha \) intensity distribution compared to a reference shot.

**Discussions and Conclusions**

The filamentary structure seen on Tokoloshe for ohmic, limiter discharges is similar to that reported by ASDEX [4] for ohmic, divertor discharges. However, the analysis so far suggests that the filaments are caused by a high m-number coherent mode(s), possibly due to a stationary island chain at the plasma edge. This could be caused by either a stray field, or a resistive instability, rather than by random density fluctuations [5].

The locking position of an island is determined by the phases of the island and external coil field. However, on Tokoloshe during high Mirnov activity both \( m = 2 \) and \( m = 3 \) islands are present. Also, the \( \ell = 2 \) and \( \ell = 3 \) coils have \( m = 1, 2, 3 \) field components, thus the resulting locking position is determined by a combination of these effects and both coils could give a similar locking position.

The lack of evidence of coil induced islands during low Mirnov activity is probably due to the fact that at the time that the coil is energised, the thermal collapse is well advanced and the \( H_\alpha \) annulus is already very bright and broad so that any additional modulation of the intensity is masked by the dominant bright halo.

**References**


Fig. 1(a) Typical frame 28 during low Mirnov activity showing the filamentary structure, most clearly visible at the top of the image. The dark regions correspond to high $H_\alpha$ intensity on the film (inside refers to the major axis).

Fig. 1(b) Poloidal intensity distribution at the limiter for frame 28 showing that the filaments are present for all poloidal angles, but most noticeable in the second quadrant.

Fig. 1(c) Harmonic analysis of the intensity distribution for frame # 28.
Fig. 2  Computer enhanced images of frame 15 -- + (a) and frame 11 ---- (b) showing the difference in the locking position as well as the poloidal extent of the locked structures.

Fig. 2(c) Poloidal intensity distribution at the limiter for frame 15 -- + configuration showing two very distinct, localised peaks on the inside.

Fig. 2(d) Poloidal intensity distribution at the limiter for frame 11 ---- configuration showing two diffuse peaks of greater poloidal extent and lower intensity.
Introduction
Tokamak, Stellerator and RFP confinement systems have similar topology but differ markedly in their equilibrium configurations. Experiments to date show that the particles and energy losses in these systems [1,2,3] are higher than those predicted from (neo)classical theories. The anomalously high particle and energy fluxes are often attributed to turbulence-induced transport. Although experimental results indicated that fluctuation-induced fluxes are significant in the edge plasma of these systems, it is not known whether the driving mechanisms or the origins of the turbulence are the same. Here, we compare the edge fluctuations and their associated transport by applying similar Langmuir probe diagnostics and analysis techniques to discharges in TEXT Tokamak, ATF Stellerator and ZT40M RFP.

Langmuir Probe Techniques
The Langmuir probe array [4] which contains 4 probe tips provides simultaneous information on density, potential and electron temperature. Two of the probe tips are connected in a double probe configuration and the other two are floating independently. For small level of fluctuations, probe theory gives

\[
\frac{\tilde{n}}{n} = \frac{I_{\text{sat}}}{I_{\text{sat}}} - \frac{1}{2} \frac{\tilde{T}_e}{T_e}
\]

(1)

\[
\epsilon \phi_p = \epsilon \phi_f + \alpha \tilde{I}_e
\]

(2)

\[
\tilde{T}_e = e(\tilde{\phi}_p - \tilde{\phi}_f)/\ln 2
\]

(3)

where \(\alpha = \ln(\sqrt{2m_i/m_e})\) for a perfectly absorbing plane probe and \(\phi_+\) is the floating potential of the positively biased probe tip of a floating double probe system. For equation 3 to hold, it is necessary to have the double-probe bias voltage much larger than the electron temperature to ensure the collection of ion saturation current. These equations assume that all quantities are measured at the same location. In practice there is always some finite separation between probe tips which can introduce phase errors. It is necessary to correct for the phase error unless the probe-tip separation is small compared to the wavelength of the fluctuations. To extend the triple probe method [5] to measure temperature fluctuations, we rewrite eq (3) to first order in \(k_1d_1\) and \(k_2d_2\) assuming \(k_|| \ll k_\perp\) for the probe configuration as shown in Figure 1:

\[
\frac{\tilde{T}_e}{T_e} = \frac{\tilde{T}_e^*}{T_e} - \frac{ik_1d_2}{\ln 2} \frac{\epsilon \tilde{\phi}_f}{T_e} - \frac{ik_1d_1}{2\ln 2} \frac{I_{\text{sat}}}{I_{\text{sat}}}
\]

(4)

where \(\tilde{T}_e^* = e(\tilde{\phi}_+ - \tilde{\phi}_f)/\ln 2\) is the apparent or the uncorrected temperature fluctuations and \(k_\perp\) is the perpendicular wavevector. To illustrate the finite probe separation effect, we use data from the edge of ATF and plot in Figure 2 the \(k_\perp\) dependence of the fractional change in \(\tilde{T}_e^*\) when \(d_2\) is changed by 2 mm (i.e. using \(\tilde{\phi}_f\) from floating probe B instead of A in the triple
probe method, see Fig. 1). Here, we use the value of k⊥ obtained from the cross-correlation of the two floating potentials. The difference which reflects the uncertainty in |T_e*| measurement caused by the finite probe separation increases with |k⊥| and is about 50% at |k⊥| of 200 m⁻¹. The measurement of T_e in ZT40M is not affected by the finite probe separation because of the much smaller k⊥ ~ 10 m⁻¹.

**Edge Parameters**

This comparison is limited to the plasma edge inside the last-closed flux surface (LCFS) away from the scrape-off region. Table 1 summarizes the edge parameters in the three devices at r/ρ_s = 0.95 where ρ_s is to the radius of the LCFS or the shear layer.

<table>
<thead>
<tr>
<th></th>
<th>TEXT I_Φ = 200kA</th>
<th>ATF ECH</th>
<th>ZT40M I_Φ = 120kA</th>
</tr>
</thead>
<tbody>
<tr>
<td>n (10¹⁹ m⁻³)</td>
<td>0.2</td>
<td>0.1</td>
<td>1</td>
</tr>
<tr>
<td>T_e (eV)</td>
<td>30</td>
<td>25</td>
<td>25</td>
</tr>
<tr>
<td>L_n (mm)</td>
<td>30</td>
<td>40</td>
<td>20</td>
</tr>
<tr>
<td>L_Te (mm)</td>
<td>35</td>
<td>40</td>
<td>≥ 80</td>
</tr>
<tr>
<td>v_ph (10³ m/s)</td>
<td>-3</td>
<td>-1.5</td>
<td>-80</td>
</tr>
<tr>
<td>v_exB (10³ m/s)</td>
<td>-2</td>
<td>-0.5</td>
<td>≤ 20</td>
</tr>
<tr>
<td>v_de (10³ m/s)</td>
<td>-1</td>
<td>-1</td>
<td>-20</td>
</tr>
<tr>
<td>k_L (m⁻¹)</td>
<td>250</td>
<td>200</td>
<td>10</td>
</tr>
<tr>
<td>k_L ρ_s</td>
<td>0.1</td>
<td>0.1</td>
<td>0.05</td>
</tr>
<tr>
<td>h⁻¹/n</td>
<td>0.2</td>
<td>0.1</td>
<td>0.5</td>
</tr>
<tr>
<td>f_Te/T_e</td>
<td>~ 0.1</td>
<td>~ 0.2</td>
<td>0.4</td>
</tr>
<tr>
<td>f_φ/T_e</td>
<td>0.3</td>
<td>0.2</td>
<td>1.0</td>
</tr>
<tr>
<td>f_B/T_e</td>
<td>~ 10⁻⁵</td>
<td>~ 10⁻⁶</td>
<td>0.02</td>
</tr>
<tr>
<td>Γ (10²¹ m⁻²s⁻¹)</td>
<td>0.5</td>
<td>0.1</td>
<td>10</td>
</tr>
</tbody>
</table>

Table 1: Edge Plasma Parameters at r/ρ_s = 0.95 in TEXT, ATF and ZT40M.

Negative velocities are in the electron diamagnetic drift direction.

The three types of discharges differ markedly in their equilibrium configurations which translate into differences in the magnetic fields and the edge q_a. There are notable difference in the edge density but not as much in the edge temperature among the three plasmas. The density scale length (L_n) and the temperature scale length (L_Te) are comparable except that L_Te tends to be longer in ZT40M. In both TEXT and ATF, the measured phase velocity of the fluctuations is comparable to the sum of the equilibrium ExB rotation and the electron diamagnetic drift in contrast to that in ZT40M where the phase velocity is many times that of the electron diamagnetic drift and that of the equilibrium ExB rotation. Although k_L is much smaller in RFP, k_L ρ_s = 0.05 - 0.1 in all three cases. They all have significant level of
fluctuations in density, potential and temperature. In ZT40M, the magnetic field fluctuations of the internally resonant m=1 modes cause density fluctuation \( \ln / \rho = (a / L_n) \mathbf{B}_r / \mathbf{B} \) where \( L_n = n^{-1} \rho / \partial r \) and are responsible for a large fraction of the observed density fluctuations. In contrast, the magnetic field fluctuations in TEXT and ATF are too small to cause significant level of density fluctuations through the coupling of flux surface perturbation to the equilibrium gradient.

**Electrostatic Approximation**

The parallel electric field \( \mathbf{E}_|| = -\nabla || \tilde{\phi} - \partial \tilde{\psi} / \partial t \) reduces to \( \mathbf{E}_|| = -\nabla || \tilde{\phi} \) for electrostatic modes for which the induced electric field is negligible. For this approximation to hold, it is required that \( |n_\| || > | \partial \tilde{\psi} / \partial t \) or \( k_\| ^2 > \omega \mu_0 / \eta = (\omega / c_v) (\omega / c_e) ^2 \). This condition is violated in the edge plasma of ZT40M. In TEXT and ATF, the criterion is generally satisfied in the scrape-off-layer (SOL) and is marginally satisfied in the region further in from the LCFS.

Similar conclusion can also be derived from the experimental measurements of fluctuations by comparing the electrostatic part \( |\nabla || \tilde{\phi} | = |k_\| || \tilde{\phi} | = |k_\| x \tilde{\phi} / L_\| \) (where \( x \) is the effective radial mode width) with the electromagnetic part \( | \partial \tilde{\psi} / \partial t | = |w_\| \tilde{B} / k_\| | \). Here, \( \tilde{\phi} \) corresponds to \( \tilde{\phi} \) in eq.2 above and an estimate of it from the measured \( \tilde{\phi} \) requires a proper correction for the \( T_e \) component. With \( k_\| = 10m^{-1}, \omega / k = 5 \times 10^4 m/s, B = 0.13T, |\mathbf{B}_r | / \mathbf{B} = 0.02, | \phi | = 10V \) and assuming \( k_\| = 0.1k \) in ZT40M, we found that the electromagnetic component \( | \partial \tilde{\psi} / \partial t | = 130V/m \) is an order of magnitude larger than the electrostatic component \( |\nabla || \tilde{\phi} | = 10V/m \). For TEXT, the two components can become comparable as shown in the estimate in Table 2 where we have used \( k_\| = 250m^{-1}, L_\| = 1.5m, \omega / k = 3 \times 10^3 m/s, B = 2T \) and the island width for \( x \).

| \( |\mathbf{B}_r | / \mathbf{B} \) (V) | \( |\tilde{\phi} | \) (V) | \( x = w \) (mm) | \( |\nabla || \tilde{\phi} | \) (V/m) | \( | \partial \tilde{\psi} / \partial t | \) (V/m) |
|----------------|----------------|----------------|----------------|----------------|
| data from SOL | 10^{-5} | 15 (1) | 0.5 | 1.2 | 0.06 |
| estimates inside LCFS | 10^{-4} | 2 | 1.6 | 0.5 | 0.6 |

Table 2: Estimates of electrostatic and electromagnetic components of \( \mathbf{E}_|| \) in the scrape-off region and an estimate of the maximum electromagnetic contribution in the edge region inside the last-closed flux surface.

(1) an upper bound

Although the induced parallel electric field may not be negligible in the edge region inside the last-closed flux surface, the extend of its effect on the turbulence has yet to be identified.

**Conclusions**

There are significant levels of density, electron temperature and potential fluctuations in the edge plasma of TEXT Tokamak, ATF Stellerator and ZT40M RFP. Particle transport induced by the electrostatic fluctuations is important in the edge of all three devices. The measured phase velocity is dominated by the equilibrium ExB rotation in TEXT, less so in ATF and not at all in ZT40M. A velocity shear layer is not observed in ZT40M. The edge fluctuations in ZT40M are better described as resistive MHD activities. In TEXT, the criterion for electrostatic approximation is satisfied in the scrape-off region and may become only marginally satisfied in the edge region inside the last-closed flux surface.
Acknowledgements
The authors would like to thank the staffs of ATF and ZT40M experiments. This work is supported by the USDOE under grant DE-FG05-88ER-53295.

References


Fig. 1 Configuration and pin connections of the probe array.

Fig. 2 The fractional change (error) in the triple probe measurement of the electron temperature fluctuations when the separation (d2 in Fig. 1) of the probe tips is changed by 2mm.
MULTI-CHANNEL LANGMUIR-PROBE AND Hα-MEASUREMENTS OF EDGE FLUCTUATIONS ON ASDEX

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Introduction
The anomalous transport observed in tokamaks is caused by turbulent fluctuations, the nature of which is still poorly understood. Diagnostic difficulties are one major reason for this lack of understanding, at least with respect to the bulk plasma. The plasma edge, however, is accessible by several diagnostics permitting localized measurements of different parameters with good spatial and temporal resolution. For this reason one can hope to obtain enough information about edge fluctuations to permit the development of theoretical models. Different ranges of plasma parameters and the lack of closed magnetic surfaces distinguish this plasma zone from the bulk plasma. Edge turbulence might well involve other mechanisms than the turbulence in the bulk. Although transport in the bulk plasma receives more attention transport in the edge plasma and edge physics are very relevant for reactor design. The realistic hope to find a solution and the importance of the problem for the next step in fusion research are reasons for the strong effort in this field on many tokamaks.

Like in many limiter tokamaks Langmuir probes were used in the ASDEX divertor device for measurements of the floating potential and of the ion saturation current. Under certain assumptions the electron density and the plasma potential can be derived from these data. Observation of the Hα-light emitted from the edge in the vicinity of a neutral gas source yields information about the electron density. While probe measurements are more suitable for quantitative evaluations including the calculation of the local particle flux the Hα-method is not calibrated and integrates radially over the edge. It is however applicable in situations where probes fail because of excessive heat load. With 16-channel arrays both methods permit spatial correlations and wavenumber spectra to be determined without any further assumptions.

Experimental findings
A detailed description of previous ASDEX results and references to work done elsewhere are found in /1/ and /2/, a summary is given in /3/. Here we summarize these experimental findings together with new results and our conclusions.

Density and potential fluctuations are observed. Both are sufficiently strong and correlated in such a way that they might well cause all of the gross radial particle diffusion in the edge. A more precise statement is not possible (not only for ASDEX) because of poloidal and toroidal asymmetries and because the assumption of negligible temperature fluctuations necessary for the evaluation is hardly justified as could be shown on ASDEX with a novel method for fluctuation parameter measurement /4/.

In both directions perpendicular to the magnetic field the fluctuations are correlated within about 1 cm. In a direction which agrees with the field direction within the accuracy of measurement a
85% correlation is found between the midplane and the divertor. The connection length is about 10 m. Zero time delay is found for maximum correlation. We cannot exclude a small fraction of fluctuations with a short correlation length along the magnetic field nor can we exclude that the direction of maximum correlation differs from the direction of the field. The most likely explanation and the basis of further considerations however is that the turbulence is essentially 2-dimensional perpendicular to B, that is the structures are flute-like. This leads to the conclusion that drift waves are probably not involved in the edge turbulence. The potential drop necessary to generate the E-field perpendicular to B is most probably localized in or close to the sheath at the divertor plates and any theory must include or should even be based on the physics in front of the target plate.

Frequency spectra are broad with essentially all the power in the range up to a few 10 kHz and a monotonic decay observable up to the 100 kHz-range, also for wavenumbers around zero (see Fig. 1). This broadening reflects the finite lifetime of spatial structures. The temporal correlation length is shorter than the mean period. A mean velocity can be defined in several ways to describe the poloidal propagation of the fluctuations. This velocity is in the order of 100 m/s to 1000 m/s and is directed in the ion drift direction outside the separatrix. Around the separatrix the velocity changes its sign. Fig. 1 shows k-ω-spectra computed from potential measurements with a 16-pin Langmuir probe array at different radial positions in an ohmic discharge. The velocity shear is clearly indicated by the different shear (Doppler shift) of the spectra. While the spectrum 1.5 cm outside the separatrix is narrower there is no change of the spectrum in the separatrix region aside from the different shear. In particular the spatial decorrelation in the velocity shear layer as reported elsewhere is seen to be an artefact of the two point analysis commonly applied.

k-ω-spectra very often show structures apparently indicating different processes. High-frequency fluctuations around 200 kHz with rather short wavelengths and propagation velocities in the order of a few 10^4 m/s in the electron drift direction are often observed. Also low-k MHD-modes are often seen. We consider these effects as independent of the turbulence responsible for transport and try to filter them out when evaluating the data.

There is no indication of an intrinsic time scale or length scale in the spectra or temporal-spatial correlation functions <f(t,x)-f(t+τ,x+δ)>. Frequency integrated k-spectra decay with about 1/k^2 above a certain k-value. The inverse of this edge-k coincides roughly with the width of the scrape-off layer and a decay according to a power law excludes the definition of a scale length.

In double null discharges the fluctuations appear only in the outer boundary layer. In single null discharges the inboard boundary layer is no more separated from the outboard layer and fluctuations appear there also. One might conclude that the inner boundary layer is inherently stable indicating the bad curvature on the outside as an essential component of the driving mechanism. This might however not be the only possible explanation for this asymmetry. If the transport in the bulk plasma is strongly inboard-outboard asymmetric, the inner boundary might simply not be fed with particles or energy in the double-null case and therefore appear quiet.

A study of the dependence of characteristic fluctuation parameters on plasma parameters should allow identification of the physical processes involved. The relative fluctuation level f/n is the simplest fluctuation parameter. From the temporal-spatial correlation function of a fluctuating variable f one can define a poloidal propagation velocity v, a temporal correlation length τ and a spatial correlation length δ. τ is defined as the decay time of the spatial maximum of the correlation function (which shows a nearly exponential decay), that is we measure the decay time in the frame of reference of the moving fluctuations. The k-ω-spectra display the same information as the correlation functions and do therefore not permit to define additional
parameters. Fig. 2 shows the dependence of characteristic parameters derived from H\textsubscript{\alpha}-
measurements in ohmic discharges on the line averaged electron density and on the plasma
current at fixed toroidal field. Reduction of B from 2.17 T to 1.24 T at fixed density and plasma
current results in an increase of δ by a factor of 2.3 and of τ by a factor of 1.3.

**Basics of a proposed model.**

Guided by the experimental observation we propose to consider the following picture as a basis
for a theoretical model.

The scrape off-layer consists of flux-bundles intermixing randomly. Only bundles close to the
separatrix are fed by energy and particles. At some small distance from the separatrix there is no
perpendicular diffusion or heat conduction. Instead the macroscopic diffusion and heat
conduction is caused by the random drift of flux bundles. For the flux bundles at some distance
from the separatrix the models derived for scrape-off layers apply: pressure balance along the
field up to the pre-sheath close to the divertor plates, classical electron heat conduction along the
field, acceleration of the plasma to the local ion sound speed in front of the target plate, energy
losses through the sheath and by atomic physics processes in the pre-sheath. The electrostatic
potential is nearly constant along the magnetic field lines and mainly determined by the electron
temperature at the sheath. In a turbulent state the temperature distribution at the target plate is
random. This causes random electric fields perpendicular to B and a random E×B-drift.

As the simplest assumption we suggest that the plasma has no intrinsic spatial scale length nor
time scale. The scale length of the fluctuations is given by the width of the scrape-off-layer, the
time scale is given by the energy confinement time of the scrape-off layer (energy stored in the
scrape-off layer divided by the power throughput). The turbulent diffusion coefficient is
proportional to the product of the characteristic dimension and a characteristic E×B-drift
velocity, resulting in a Bohm-like dependence on the temperature in front of the plates and on
the magnetic field. The poloidal propagation velocity is caused at least to a major part by the
radial electric field caused by the temperature gradient in the boundary layer.

Applying this Bohm-like diffusion formula and the equations for the physics parallel to the B
field for the scrape-off layer as a whole we get a self-consistent model allowing the prediction
of the width and confinement time of the SOL which we assume to be proportional to the
characteristic parameters δ and τ of the turbulence. Under realistic simplifications the set of
equations can be easily solved. The parameter dependence of δ and τ from the calculation is
indeed in fair agreement with the dependence found experimentally. The model therefore might
contain the most important effects. From the density dependence of the measured propagation
velocity we must conclude, however, that more effects have to be included. The zero-crossing
of v indicates that two independent velocities counteract each other.

**References**

/3/ A. Rudyj et al., 17th EPS Conference on Controlled Fusion and Plasma Heating,
/4/ A. Carlson, L. Giannone et al., this conference.
Fig. 1. Grayscale Plot of wavenumber-frequency-spectra at different distances $d_s$ from the separatrix (potential fluctuations measured with poloidal Langmuir probe array). $d_s$ is positive for positions outside the separatrix.

Fig. 2. Dependence of the correlation length $\delta$, the correlation time $\tau$ and of the poloidal propagation velocity $v$ on the line averaged density (top) and on the plasma current.
ELM STUDIES ON ASDEX

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1. Introduction

The performance of tokamaks operating in the H-mode [1] crucially depends on the ELM (Edge Localized Mode) behaviour of the discharge. On DIII-D [2] and ASDEX [3] it was found that only in ELMy discharges with a carefully optimized ELM repetition rate stationary conditions could be achieved. Nevertheless, the nature of the ELM as an MHD phenomenon is not yet clear. On DIII-D, giant ELMs occur at the ideal ballooning limit while grassy ELMs are found in the connection region of the s − α diagram pointing towards resistive ballooning modes [4]. On the other hand, on PBX-M it was found that ELMs appear in discharges which were analyzed to be ballooning stable and an MHD stability analysis indicated that ELMs could be connected to low-n pressure driven ideal kink modes [5].

The general features of ELMs in ASDEX are summarized in ref. [1]. The MHD phenomena accompanying ELMs on ASDEX are described in ref. [6], [7] and [8]. In this paper we present the results of further studies on the MHD characteristics of ELMs.

2. Phenomenology of ELMs on ASDEX

The occurrence of an ELM leads to a sudden outflux of particles and energy. This outflux manifests itself as a spike on the Dα light in the divertor. The rise time of the Dα signal is typically about 50-100 \(\mu s\), the half-width of the signal ranges between 100 and 400 \(\mu s\). The rapid loss of energy causes an inward motion of the plasma. In case of grassy ELMs the ELM frequency (repetition rate) is comparable to the half-width of the Dα signal.

The pivot point for the change in the electron temperature profile introduced by an ELM is at \(a - r = 4\) cm (the plasma minor radius \(a\) being 40 cm). It is the same for two different values of the edge safety factor \((q_0=3.1\) and 4.1), indicating that the ELM is not connected to a certain (rational) \(q\)-surface.

The energy loss caused by an ELM can be measured in terms of the poloidal beta, \(\beta_p\). A single ELM can cause a loss in \(\beta_p\) of the order of 5%. Such a loss is also typical for a soft minor disruption. During such an event, however, a positive current and negative voltage spike due to the decrease of poloidal flux is observed. ELMs do not show such a behaviour, indicating that the current profile is almost unaltered by the ELMs.

The electron pressure which is measured on ASDEX in intervals of 17 ms reaches the highest values at the edge in quiescent phases. With the onset of ELMs the edge pressure gradients are rapidly reduced. Between ELMs the edge gradients rise but are far from the values obtained in quiescent phases. ELMs can also appear right after the H-transition at still low pressure gradients. In agreement with our preceding studies [1] we therefore cannot attribute even singular ELMs to the ideal ballooning limit. We have, however, no measurement of the edge pressure with a spatial resolution better than 2.5 cm.
If the H-transition is not triggered by a sawtooth, the transition conditions may fluctuate leading to a sequence of short L- and H-phases prior to the final H-transition. This phase is dubbed 'tithering' H-mode. ELMs can diagnostically easily be discriminated from the spiky nature introduced by the 'tithering' H-mode.

In ref. [1] it was stated that ELMs lead to an enhanced toroidal asymmetry of the heat flux to the divertor plates. This statement was due to a misinterpretation of the results. More recent measurements with ELMy H-phases lasting for seconds, rule out large poloidal or toroidal asymmetries [3].

3. MHD Mode Characteristics

In this section we present results on the MHD mode characteristics of ELMs using a fast (up to 3 MHz) data acquisition system. From this diagnostics, the temporal sequence of a single ELM is determined as follows:

Fig. 1 shows $D_{\alpha}$-traces from different locations as well as the soft-X-ray emission from the edge ($r=a$, outside the ELM inversion radius), the electron temperature at $r=a/2$ and a magnetic signal ($B_r$) from midplane outside. A high frequency (110 kHz) magnetic precursor oscillation can be detected followed by a highly turbulent phase. From integrated magnetic signals we find that the growth of the precursor is nearly exponential with a growth time of $\approx 50\mu s$. As only few magnetic probes were available for the fast data acquisition, modenumbers or the propagation direction of the precursor could not be determined. Nevertheless, a rough indication of the $m$ numbers involved can be estimated from the comparison of data from Mirnov coils which integrate the signal over their length of $\approx 15\ cm$ and data from printed circuit coils (integration length 0.3 cm). From this comparison we can state that the precursor, as it is detected with both coils, has an $m$ number below $\approx 20$ while the turbulent phase of the ELM must have higher mode numbers as it is hardly resolved by the Mirnov coils due to the integration effect.

Note that the last cycles of the oscillation are also seen on the soft-X-ray trace. Also, density fluctuation measurements at the plasma edge detect the oscillations. This precursor is not always seen as clearly as shown in the example above. Particularly at low $q$ it is masked by the occurrence of additional MHD activity as described in ref. [6] and [7]. Nevertheless, the precursor was observed prior to all ELMs we examined (about 5 different parameter settings). Moreover, in the quiescent H-phase, sometimes a growing magnetic oscillation with the same frequency is seen. If this oscillation does not reach the amplitude as shown in figure 1, no ELM evolves. In this case, where the precursor grows and then again vanishes, no effect on the confinement properties of the H-mode is detected although the level of density fluctuations at very high frequencies ($\approx 1\ MHz$) is reduced. We therefore conclude that the precursor oscillation does not change the transport barrier whereas the ELM rapidly alters the confinement properties in this region owing to the turbulent phase following the precursor.

The ELM manifests itself on the magnetic fluctuation signal as a broadband turbulent phenomenon with frequencies reaching up to 300 kHz. Fig 2 shows a contour plot of the temporal evolution of the frequency spectrum of the magnetic fluctuations. The precursor is seen as a narrow line at about 180 kHz, the following turbulent phase lasts about 250$\mu$s (roughly the duration of the ELM as deduced from the $D_{\alpha}$ emission).

This phase is identical with the outflux of heat and particles into the divertor. The phase in which the magnetic oscillation changes from the precursor to the turbulent phase is identical with the onset of the drop in $T_e$ at $r=a/2$. The heat pulse arrives at the separatrix about 30$\mu$s later (seen on the outer soft-X-signal). An analysis of several ELMs shows that the transport
barrier is reduced within typically 10-50 μs. The rise of the $D_0$ signal in the midplane and in the divertor is nearly coincident, the lag of the divertor signal, if any, is of the order of 30μs.

The frequency of the precursor oscillation is rising with the plasma current and does not vary between ELMs at different pressure gradients. There is a slight rise of the edge electron pressure with plasma current (at constant $q_a$).

The MHD nature of the ELM is in qualitative agreement with the model of the 'peeling-mode' given in ref. [1] and might therefore be a kink-like motion due to a cluster of resistive ballooning modes. The measured magnetic turbulence suggests that during the ELM the magnetic configuration is altered. On the other hand, measurements show that around the X-points the configuration is not destroyed. It is therefore likely that the turbulent zone is located in the midplane outside where the unfavourable curvature of the magnetic field lines is especially important for pressure driven modes. These modes then lead to the observed increase in transport which we suppose to be mainly convective as it happens on a fast timescale. These transport properties also clearly distinguish an ELM from the L-mode so that in our opinion ELMs are not temporary H-L-H transitions (as e.g. the 'tithering' H-mode). The MHD mode scenario mentioned above has been shown to be very sensitive to wall stabilization effects and might therefore explain the fact that the ELM behaviour of ASDEX H-mode discharges is critically dependent on the plasma position (major radius $R_0$). A radial outward movement of the plasma column from $R_0=1.66$ m by 2 cm changes a steady state ELMy H-mode into a quiescent one.

4. Conclusions

We have shown that the ELM consists of a coherent precursor oscillation followed by a turbulent phase. While the precursor does not change the confinement properties of the H-mode, the turbulent phase leads to a rapid ($\approx 10-50$ μs) increase in transport across the separatrix. The findings are consistent with the picture of a kink-like precursor motion leading to a turbulent state consisting of resistive ballooning modes as described in ref. [1].

References

Figure 1: Temporal sequence of an ELM

Figure 2: Contour plot of magnetic fluctuations during an ELM
Detection of coherent structures in the edge of the TEXT tokamak plasma

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Conventional plasma turbulence theories are based on the assumption of weakly interacting waves and use the random phase approximation. The fluctuation distribution in this case is quasi Gaussian and no long-lived structures, maintained through self binding forces, exist. The difficulty in identifying a turbulence model to describe the edge fluctuation characteristics in tokamaks may be a result of using such oversimplifying assumptions. Coherent structures violate these conditions. They have been found in neutral fluids and have also been studied in a cylindrical laboratory plasma.1 In tokamaks, granular structures have been observed in the edge,2 but no detailed statistical studies have been done so far.

In the edge of the TEXT tokamak, the probability distribution for fluctuations of the floating potential as measured by Langmuir probes, is shown in Fig. 1.a. The distribution of the potential fluctuation is close to Gaussian (solid line in the figure) with a skewness $S = \langle \dot{\phi}^3 \rangle / \langle \dot{\phi}^2 \rangle^{3/2} = -0.2 \pm 0.2$ and kurtosis $K = \langle \dot{\phi}^4 \rangle / \langle \dot{\phi}^2 \rangle^2 = 3.2 \pm 0.3$. With a conditional averaging technique3 using higher order correlation functions, a limited set of experimental data has been analyzed in more detail. This technique can test a condition $C_{\text{cond}}$ which is theoretically predicted to lead to long lived structures (such as a substantially negative potential dip trapping positive ions). For the simplest condition, the occurrence of a potential dip $C_{\text{cond}} = \varphi_{\text{cond}}$ at the reference point $(\bar{r}, \bar{t})$, the conditional potential at position $\bar{r} + \rho$ and time $t + \tau$ is given by

$$\varphi(\bar{r} + \rho, t + \tau)_{\text{cond}} = \langle \varphi(\bar{r} + \rho, t + \tau) | \varphi(\bar{r}, t) = \varphi_{\text{cond}} \rangle.$$  \hspace{1cm} (1)

A Taylor's expansion around $\varphi(\bar{r}, t)$ may be used to approximate $\varphi(\bar{r} + \rho, t + \tau)$:

$$\varphi(\bar{r} + \rho, t + \tau) = \sum_{n=1}^{N} \alpha_n(\rho, \tau) \varphi^n(\bar{r}, t).$$  \hspace{1cm} (2)

This expansion minimizes the mean square error between the estimated value (symbolized by the hat (^)) and the actual (measured) value of the potential $\varphi(\bar{r} + \rho, t + \tau)$ when the unknown coefficients, $\alpha_n(\rho, \tau)$ satisfy a set of $m = 1, N$ equations:

$$\sum_{n=1}^{N} \alpha_n(\rho, \tau) \langle \varphi^{n+n}(\bar{r}, t) \rangle = \langle \varphi^{n}(\bar{r}, t) \varphi(\bar{r} + \rho, t + \tau) \rangle.$$  \hspace{1cm} (3)
Solving the set of equations for the coefficients $\alpha_n(\bar{p}, \tau)$ the conditional average for the condition $\varphi(\bar{r}, t) = \varphi_{\text{cond}}$ is found:

$$
\dot{\varphi}(\bar{p}, \tau)_{\text{cond}} = \sum_{n=1}^{N} \alpha_n(\bar{p}, \tau) \varphi_{\text{cond}}^n.
$$

(4)

Note that the conditional average can be estimated for several conditions, once the coefficients $\alpha_n(\bar{p}, \tau)$ have been computed.

In the analysis presented here, we keep terms to third order, $N = 3$. The measurements are taken during the flat top phase of the discharge, where the fluctuations are stationary, i.e. statistically independent of $t$. We thus take advantage of the long record length with respect to the correlation time and ensemble average over many time segments. To vary the probe separation $\rho_{\text{pol}}$ in poloidal direction we make use of the long correlation length in toroidal direction and rotate the signal probe with respect to the reference probe on the same flux surface. The scan in $\rho_{\text{pol}}$ is obtained over successive and similar discharges.

The structures with substantial negative excursion in the floating potential are slightly longer lived (Fig. 1.b) than the positive potential spikes (Fig. 1.c). The eddy turnover time $\tau_{\text{eddy}}$ for ions in a negative potential dip is comparable to the ambient correlation time $\tau_{\text{corr}}$ of the turbulence and slightly smaller than the lifetime $\tau_{\text{life}}$ of the structures, $\tau_{\text{eddy}} \simeq \tau_{\text{corr}} < \tau_{\text{life}}$. This is indicative of ion trapping.

To optimize the TEXT data analysis technique and to determine its sensitivity we use data generated with a numerical simulation code. The computer code models dissipative drift wave turbulence using the model equations

$$
\frac{\partial}{\partial t} \nabla_i^2 \varphi = \alpha (\varphi - n) + C \nabla_i^2 \varphi , \quad \frac{d}{dt} n = \alpha (\varphi - n) - \kappa \frac{\partial \varphi}{\partial y} .
$$

(5)

The nonlinear terms arise from the $E \times B$ convection as $d/dt = \partial/\partial t + \nu_E \cdot \nabla \varphi$. Here $\alpha = \chi_k^2$ is the adiabaticity parameter depending on the parallel electron conductivity, $C$ is a damping term, and $\kappa$ is an inverse scale length for the unperturbed density profile. The adiabaticity parameter describes the degree of linear coupling between the density and potential equations. In the adiabatic regime ($\alpha >> 1$), the equations for $n$ and $\varphi$ are tightly coupled and the functions are virtually indistinguishable. In the hydrodynamic regime ($\alpha << 1$), the equations behave in a similar manner to two-dimensional hydrodynamical turbulence with the equation for $n$ responding much like a passively advected scalar.

This set of equations allows a control of the coherent structures. For $\alpha << 1$ intermittent structures in density and potential are generated. For $\alpha = 0.01$ the kurtosis of potential fluctuations is $K \approx 4.4$ and the skewness is $S \approx -0.1$, while $\alpha >> 1$ generates fluctuations with $K \approx 3, S \approx 0$ and no visible long lived structures. Fig. 2.a shows the motion of a few structures in the x-y plane over their life time $\tau_{\text{life}} >> \tau_{\text{eddy}}$ for a case with $\alpha = 0.01$. The probability distribution is shown in Fig. 2.b. The conditional potential of the simulation data is computed with the same technique applied to the experimental data. The result is shown in Fig. 2.c. A negative condition was used, but a positive condition gives similar results for the correlation
Figure 1: a) Probability distribution for measured floating potential fluctuations and conditional average for b) $\varphi_{\text{cond}} = 3\sigma$ and c) $\varphi_{\text{cond}} = -3\sigma$.

Figure 2: a) Motion of a few structures b) probability distribution and c) the spatial decorrelation of the conditional potential for three conditions $\varphi_{\text{cond}}/\sigma$ for the numerical simulation.
length and time. This is also reflected in the symmetric probability distribution. We find the size of the structures to be well estimated, but the estimated life time $\tau_{\text{life}}$ is much smaller than the true life time of the structures $\tau_{\text{life}}$. The reason is the random walk of the structures (see Fig. 2.a), which is not accounted for by the simple condition which has been used in the analysis.

To summarize, we find the edge floating potential fluctuations to be close to Gaussian distributed. The conditional sampling technique shows that positive potential spikes are slightly shorter lived than negative spikes. For the negative spikes we find $\tau_{\text{edge}} \simeq \tau_{\text{corr}} \lesssim \tau_{\text{life}}$. Present experimental data thus indicate that longer lived negative potential spikes may be present, but are probably not a dominant feature of the edge plasma of the TEXT tokamak. The analysis of the computer simulation data with known coherent structures reveals that the conditional sampling technique is capable of correctly estimating the size of the structures, but can largely underestimate their lifetime $\tau_{\text{life}}$, because of the random motion of the coherent structures in the x-y plane. More sophisticated conditional sampling is required. Improvements in the data analysis techniques to confirm structure lifetime calculations are in progress.

Work supported by the U.S. Department of Energy.

References

PARALLEL EXPANSION OF THE ABLATION CLOUD DURING PELLET INJECTION IN TORE SUPRA

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The ablated matter propagation along the field lines during pellet ablation is observed with a five chords interferometer toroidally located at $-\pi/3$ of the pellet injector (=2.5m on the discharge axis). The distance between two chords is 16.5cm and their 1/e total width 1.7cm (see Figure 1). The time resolution is 16\mu s and the sensitivity better than $3 \times 10^{17}$m$^{-2}$. The beginning of the fast acquisition is triggered by the pellet itself (nevertheless, due to the uncertainty on the pellet velocity, the time at which the pellet enters the discharge (t0) is only known with an accuracy of $\pm 100\mu$s) and its maximum duration is 20ms. About 500\mu s after the pellet enters the discharge, the experimental signals exhibit a steep increase (in one or two steps, each of a duration of order 100\mu s). Excepted in a few cases for which a strong oscillation at a typical frequency of 0.5kHz was detected during several ms, a new quasi-steady state is reached after $\approx 1$ms.

The importance of the measured perturbation and the details of the sequence described above depends, for each chord, on both the q profile and pellet penetration. The measured quantity is the increase of the line integrated density for times larger than to: $\delta n(t) = n(t) - n_0(t)$. An example is displayed Figure 2 for the shot TS 2696 (chords n° 1, 2, 3 and 5). The main discharge (He) parameters are: $I_p=0.7\text{MA}$; $B_T=2.5\text{T}$; $a=0.75\text{m}$; $T_e(0)=1200\text{eV}$; $n_e(0)=5.75 \times 10^{19}\text{m}^{-3}$. The flight pass of the pellet (D2) is $L_p=0.5\text{m}$ (corresponding to an initial radius $r_p=1.16\text{mm}$) and its velocity is $V_p=607\pm 9\text{m/s}$.

MODEL

These observations are analysed with a self-consistant model whose only input parameters are the magnetic structure of the discharge, the main characteristics of the plasma ($T_e(\rho)$, $n_e(\rho)$, eventually complemented with a poloidal rotation of frequency $\Omega(\rho)$), the pellet velocity $V_p$ and its total penetration $L_p$ (no absolute measurement of the pellet mass is available). The successive steps of the computation are described below:

-1- The pellet radius $r_p$ is adapted to fit the total penetration depth given by the classical NGS model [1] to the experimental value $L_p$. Since, for known pellet radii, the computed flight passes are generally larger than the experimental values, the above procedure can yield an under-estimation of the pellet mass, and thus of the simulated $\delta n$'s.

-2- For 20 radial points along the pellet trajectory, the bulk ionization degree, the expansion dynamics (radius: R and half-length along BT: Z) and the main parameters (temperatures and density) of the plasmoid (i.e. the quantity of pellet material released on an individual magnetic surface) are computed using a MHD model close to that described in [2]. The latter is modified...
to take into account the fact that, on a time scale of less than 1ms, the strong perturbations of the background plasma parameters influence the plasmoid expansion. This can be seen on Figures 3 and 4 where are displayed, for the pellet of fig.2, the time evolution of Z and of the electronic and ionic temperatures in both the plasmoid (Te0 and Ti0) and plasma (Te°° and Ti°°) at the location of maximum ablation. The half-length of the plasmoid Z exhibits several strong oscillations before to steady at an asymptotical value of 260m (fig.3). They are due to the alternative compressions which result from the relaxation of the plasmoid-plasma system inside the finite volume associated with the considered magnetic surface. The global relaxation time (=25ms) is mainly determined by the electron-ion collision time Tei. Indeed, it can be seen on fig.4 that the plasma and plasmoid electronic temperatures are equalized in less than 1.5ms and that the ionic temperatures become comparable in =7ms (after averaging over the few compressional oscillations). Nevertheless, due to the weak collisional coupling between electrons and ions, the thermal equilibrium is only reached after =0.1s, with a 1/e folding time of =25ms. In its present state, two phenomena limit the validity of these computations for long time scales: (1) no dissipation due to the plasma viscosity is considered which likely overestimates the above mentioned oscillations and (2) the global radial transport is neglected. In these conditions, one can consider that this model is no longer relevant for time intervals longer than 2ms after the pellet injection.

For each chord, the increase of the line integrated density is computed in the exact discharge geometry (given by the IDENT C equilibrium code). Due to the fast parallel expansion of the plasmoids, their half-length can reach several hundred of meters in a short time: in the case presented fig.3, Z(1ms)=300m which corresponds to 2Z/2πRo=40 toroidal turns. For the simulation presented in the next section, only the contributions of the first 15 are explicitely taken into account (the contributions of the 25 remaining are averaged on 1 poloidal turn). In a poloidal plane, in addition to their motion due to the rotational transform, the plasmoids are assumed to move with a velocity 2πpΩ(p) immediately after their ionization. In absence of measurements of the plasma poloidal rotation, the profile Ω(p) is the only free parameter of the model.

RESULTS AND DISCUSSION

Two models computed with the parameters corresponding to the data of fig.2 are displayed Figures 5 and 6. In the first simulation (fig.5), Ω(p)=0 everywhere inside the plasma. The main discrepancies between the experiment and this simulation are: (1) the lack of the double peak around (t-to)=600μs on the signal of chord n°1 (θn11) and (2) the poor relative level of the signals of chords n°2 (θn12) and 5 (θn15) for (t-to)>750μs. In the second case (fig.6), the poloidal rotation frequency profile is adapted to best fit the measurements and all the features of the experimental signals can be clearly identified on the simulation. Note the presence of the double peak on θn11 and the nearly equal levels of θn12 and θn15 for (t-to)>750μs. Moreover, the fact that θn12>θn15 for 750μs<(t-to)<900μs and the relative position in time of the different peaks are fairly well reproduced. The residual discrepancies remain: (1) the steepness of the density increase, which is larger for the models than for the data
and (2) the absolute value of the $\delta n_1$'s for large (t-to) which is under-estimated by a factor of $\approx 2$ in the simulations. The first one is easily understood if one's consider that $n_0$ and $n_\infty$ are only volume-averaged values and that the density gradient is therefore infinite on the plasmoid-plasma junction surface. The second results mainly from an under-estimation of the pellet mass and, possibly, from an over-estimation of the plasmoids parallel expansion at large times, due to the fact that the friction forces between two adjacent flux tubes are neglected in the present state of the model.

The poloidal rotation frequency profile used in the second simulation, $\Omega(\rho/a)$, is displayed Figure 7. It corresponds to a rotation in the direction of the electronic diamagnetic speed and exhibits a maximum at $\rho/a=0.55$ which seems well constrained by the data. Such a behaviour of a pellet density perturbation was observed in T10 and led to similar conclusions [3]. The domain of $\rho/a$ investigated is limited by the radius of maximum pellet penetration towards the center of the discharge and by the small amplitude of the density perturbation for $\rho/a>0.85$. The error bars plotted on the figure are estimated on the basis of computational tests. Interpreted as an $E\times B\times r$ drift, this poloidal rotation implies the existence of a radial electric field of order of magnitude 3.5 kV/m around $\rho/a=0.5$. The poloidal rotation frequency associated with the plasma (He) ionic diamagnetic velocity before the pellet injection is also plotted for comparison.

REFERENCES:

Figure 1: Geometry of a poloidal cross-section in the interferometer plane. The injection plane of the pellet is at $+\pi/3$ in the toroidal direction.
Figure 2: Experimental data (t_o=3.1462s). Each curve is labelled with its corresponding chord number.

Figure 3: Time history of Z near the location of maximum ablation.

Figure 4: Time history of the plasma and plasmoid ionic and electronic temperatures at the location of maximum ablation.

Figure 5: Simulation of the data of fig.2 with $\Omega \,(p)=0$.

Figure 6: Simulation of the data of fig.2 with the $\Omega \,(p)$ profile displayed fig.7.

Figure 7: Poloidal rotation frequency profile $\Omega \,(p/a)$ used for the model of fig.6.

For comparison, $\Omega \,_{dia} \,(He)$ is the diamagnetic poloidal frequency of the plasma ions before the pellet injection.
DENSITY CONTROL IN TORE SUPRA WITH
ERGODIC DIVERTOR AND MULTI-PELLET INJECTION


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1. INTRODUCTION

A large part of the experimental programme of TORE SUPRA has been devoted to ohmic plasmas studies in many configurations. These experiments were restricted as far as the particle content is concerned by the "density limit" [1]. This limit appears to be a strong function of the ionic species, the fuelling method and the plasma edge characteristics. Typical examples of this are shown in ergodic divertor and multi-pellet experiments [2]. The experimental procedures such as different type of gas and pellet injection during the ergodic divertor pulse are described in the next section: the fuelling efficiency is reduced if the edge deconfinement produced by the ergodic divertor is accompanied by some pumping by the wall. At the same time, the high-recycling zone behaves like a radiating layer, especially since it becomes cooler with ergodization. The simultaneous use of pellet injection and edge field lines ergodization should yield a convenient way to control both plasma density and radiative losses.

2. EXPERIMENTAL OBSERVATIONS

The general observations concerning the ergodic divertor experiments have been summarized in ref.[2]. In particular, a modification of the energy transport was given evidence in the edge ergodized layer, inducing a strong decrease of the electron temperature there. The edge particle transport will be shown to change to a lesser extent. An important feature of these experiments consists of the new connexion between the plasma and the wall, as imaged in the heat load pattern derived from the wall I.R. thermography [3]. The somewhat intricate particle balance will be described as follows, depending on the fuelling method:

a) Helium fuelling

- In Helium discharges, the recycling coefficient stays always very close to 1. Consequently the averaged density depends essentially of the total amount of gas injected, the fuelling efficiency remaining large (close to 80%). When the ergodic divertor is activated (fig.1a), the volume-averaged density deduced from I.R. interferometry decreases only by 10 to 15%. The measurements (r/a < 0.75) from the Thompson scattering system [4] show a dominant decrease at the edge (35 to 40%) and hence a much smaller effect in the bulk. When Helium is then puffed in at a constant rate, the density increases linearly, its profile remaining unchanged.

The main effect of the ergodic divertor is to decrease the confinement time in the edge $\tau_{edge}$ to a value when ergodized $\tau_{edge}(E.D.)$ as (the bulk confinement staying the same):

$$\tau_{edge}(E.D.)/\tau_{edge} = (N_{edge}/N_{b})(E.D.)/(N_{edge}/N_{b})$$

[1] General Atomics, San Diego, Ca (USA);
[2] Oak Ridge National Laboratory, Oak Ridge, Tn (USA)
where $N_{\text{edge}}/N_b$ is the ratio of the particle numbers in the edge and in the bulk. If the edge channel of the Thompson scattering system is representative of the edge average, $T_{\text{edge}}(E.D.)/T_{\text{edge}} = 0.65$.

The gas puff (up to 7.5 mbar.l/s for 2s) allows an increase of the radiated power by factors up to 2, i.e. from 0.4 to 0.8 MW when the ohmic power is about 1.5 MW. If the gas puff is stronger, a disruption occurs with a sudden and large increase of the radiation loss.

**b) Deuterium fuelling**

In Deuterium discharges, the observations depend on the plasma configuration [5]: if the plasma is leaning on the modular limiter located on the low-field side, the application of the magnetic perturbation causes a density decrease as shown in figure 1b. This depends on the wall status and may reach 50%. An important feature is its reversibility once the magnetic perturbation is off. When leaning on the high-field side graphite wall, the effect is generally the opposite. This is very similar to the observations made when the plasma is moved from one configuration to the other, as encountered in Tore Supra [6] or JET [7]. The ergodic divertor induces a new equilibrium with the wall, which is intermediate between the two configurations. The pumping effect by the wall causes a quasi-immediate decrease at the edge whereas it is delayed in the bulk where the particle confinement is unaffected by the ergodization.

The gas puff is very ineffective in this case. In experiments where up to 10 mbar.l/s were introduced through the valve, the fuelling efficiency, defined as the ratio of the plasma content increase divided by the gas puff, stays close to 1%! For larger gas puff, the marginal fuelling efficiency increases likely due to a better penetration, until a disruption will occur.

The application of the ergodic divertor allows a general decrease of the total radiated power, only the outer edge showing a large increase. This seems related to its decontamination properties discussed at this conference [8]. When the gas is puffed in, the edge radiation is enhanced. This occurs gradually, the size of the edge radiating layer increasing. If the gas is introduced up to the situation when the total radiation loss reaches the ohmic input, a disruption occurs. On the contrary, an early enough stop of the injection may prevent the disruption occurrence even when 80% of the power is radiated at that time. It should be stressed that the edge electron temperature which is affected by the ergodization to a large extent keeps on decreasing during gas injection below 50 eV at $r = 0.94$ a.

**c) Pellet injection experiment**

When the discharge is fuelled with pellet injection by means of a centrifugal injector built by ORNL (600 m/s, up to 10 Hz, $2.5 \times 10^{20}$ D$_0$ per pellet), much larger densities are obtained ($7 \times 10^{19}$ m$^{-3}$) in ohmic plasmas ($I_p = 1.65$ MA ; $B_t = 3.5$ T). The observations are very similar in situations where the ergodic divertor is applied or not. An example is given in fig. 1c where series of pellets is injected in Tore Supra; the ergodic divertor acts from $t = 3.6$ s. The fuelling efficiency is about 1 in that case; the pellets are injected mid-radius and thus beyond the ergodic layer. The radiative loss is generally connected to the density increase but it appears that the wall plays a major role by pumping the additional particles. This can be sustained for about 50 pellets in 2 typical shots, corresponding to $2 \times 10^{22}$ particles. Once this value reached, the discharge may disrupt very suddenly. A means to recover is by achieving 1 or 2 shots in Helium, thus desorbing a part of the implanted gas; the fact this part is very scarce is now under investigation. It should be stressed that the pellet injection produced much more peaked plasma density when the E.D. is on: the core density increases by a factor 6 when the edge one is limited to 3. In fact, this peaking decreases when more and more pellets are injected. The recycling deduced from the H$_\alpha$ measurements behaves like the edge density during the shot as confirmed by the increase of the apparent particle confinement time $T_{P^*}$ deduced from the density decrease rate after the pellet ablation; moreover, a gradual shot-to-shot increase can give evidence of the wall status evolution. The edge electron temperature is rather low as usual when the ergodic divertor is activated but remains larger than 70 eV. At the same time, the central temperature decreases from 2.5 keV to 1.2 keV. The plasma cooling is thus general.
3. DISCUSSION AND CONCLUSIONS

In the various configurations studied in this paper, very different density behaviour were shown. This is true especially for the "density limits". An interesting way to consider it can be deduced from the evolution of the fraction of power radiated in the plasma. This implies of course, a similar level of "natural" contamination in the plasma, which is taken for granted when the conditioning of the wall is considered as good. In that respect, figure 2 shows the ratio of the radiated power deducted from the bolometric measurements to the total ohmic power input as a function of the parameter \( M \cdot q \) where \( M = \frac{n_e R_B}{B_t} \) is the Murakami parameter, \( n_e \) the line-averaged density and \( q \) the edge safety factor. For ohmic plasmas in Deuterium and Helium, this results in curves yielding limits of \( M \cdot q \), respectively 9 and 25. When the ergodic divertor is used, the edge radiated power is further increased when the density is raised, leading to a lower limit for \( M \cdot q \) : 6 in Deuterium and 15-20 in Helium. It should be stressed that the path to this limits is somewhat different as pointed out in the former paragraph. If pellet injection is used, the \( M \cdot q \) limit remains similar in Deuterium. In this case, the ergodic divertor does not affect the results drastically; a far more important factor is the wall status as far as its Deuterium saturation is concerned. This explains why at some time the radiation suddenly increases, the "knee" in the curve depending on the wall preparation and its evolution.

The most interesting feature of these experiments is that they allow to study these limits in cases when either the edge can be drastically modified or the bulk affected without large modification of the edge (at least transiently). The ergodic divertor induces a strong deconfinement in the outer part of the plasma, where the electron temperature is lowered, and the particle transport is somewhat affected. In addition, the connexion to the wall becomes more diffuse, yielding a new equilibrium with the wall. The particle content of the plasma varies accordingly. If one tries to counterbalance this by edge gas puffing, one finds that the screening properties of the ergodic layer plus the wall pumping prevents any efficient density control. Moreover, the high edge recycling plus the already lowered temperature are essential factors for the creation of a strongly radiating layer.

An alternative to raise the density consists of injecting pellets. The limitation comes here from the necessity to control the sum of recycling plus fuelling. Till now, only the graphite wall was used to pump particles, which gives a definite limit of about 50 pellets. Without enough control, the density increase results in the general cooling of the plasma and finally its disruption. This could be alleviated in the future by increasing the power input with auxiliary heating, and by actually pumping out the particles with pump limiters.

Both the density control and limits are strongly dependent on the conditions prevailing at the edge of the plasma. The use of the ergodic divertor combined with bulk and edge fuelling, active and reliable pumping may yield relevant experimental situations, where, in addition, impurities and radiation losses might be controlled.

ACKNOWLEDGEMENTS:

The authors are particularly indebted to the whole TORE SUPRA team.
This research is performed under the USA/France Collaboration Agreement and was sponsored in part by the office of Fusion Energy, US DoE, under contracts DE-AC03-89ER51114 with General Atomics, Inc. and DE-AC05-84OR2400 with Martin Marietta Energy Systems, Inc.

REFERENCES:
Figure 1a: He plasma with He gas puff (plain curve on the lower graph) and ED (dashed)

Figure 1b: D plasma with D gas puff (plain curve on the lower graph) and ED (dashed)

Figure 1c: D plasma with pellet injection and ED (dashed curve on the lower graph)

Figure 2: $P_R / P_\Omega$ versus $M_q$ ($\bar{n}_e / I_p$)

- fit to the limiter data base
- plain line D-plasma, dashed He-plasma, dotted line multipellet injection D-plasma,
gas puff with Ergodic Divertor (full symbols), squares D-plasma, triangles He-plasma,
multipellet injection with the Ergodic Divertor, open circles (D-plasma).
ENERGY CONFINEMENT OF HIGH-DENSITY PELLET-FUELLED H-MODE PLASMAS IN ASDEX

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Introduction:
Previous investigations on the divertor tokamak ASDEX (R = 1.65 m, a = 0.40 m) revealed that repetitive pellet injection (PI) can significantly improve the plasma performance, especially the energy confinement time $\tau_E$. However, the achieved enhancement in $\tau_E$ degrades, as shown earlier, from about 2 in the ohmic regime down to $\sim$1.3 in the high auxiliary power heated L-mode case [1].

H-mode discharges, on the other hand, exhibit even at the highest auxiliary heating powers a confinement gain of up to 2 compared with the corresponding L-mode [2]. The object of this paper is to examine the possibility of combining the favourable confinement properties of both regimes (PI, H-mode) by means of injecting pellets into NBI heated H-mode discharges. Similar investigations have already been performed on JFT-2M [3] and recently on JET [4]. Auxiliary heating by ICRF was also applied on these machines.

Both discharge types are characterized by strong deviations of their electron density profile shape from the "standard" profile. Commonly PI discharges show markedly peaked profiles but gas-puff (GP) fuelled H-mode discharges relatively flat ones. If the intrinsic confinement behaviour of both regimes is linked to the density profile shape, these beneficial properties may be difficult to combine.

Experimental Parameters and Plasma Performance:
During the experiment parameter scans over the range $1.7 \leq B_t/T \leq 2.0$, $320 \leq I_p/kA \leq 420$ in double null magnetic separatrix configuration were conducted with boronized vessel walls. The neutral beam heating power was limited to $P_{NBI}=2.0$ MW. Up to 30 $D_2$-pellets, each contributing $2 \cdot 10^{19} m^{-3}$ to the volume-averaged density, are injected into a single discharge with a time interval of 33 ms. The line-averaged density exceeds $n_e = 1 \cdot 10^{20} m^{-3}$. The resulting peak electron temperatures are below 1 keV.

Typically, H-phases are studied which start during the PI cycle. Owing to the temperature reduction accompanying PI prior to the H-mode transition the pellets penetrate, in contrast to earlier experiments with smaller pellets [5], deeply to about half the plasma radius. The pellets trigger instantaneously high ELM frequencies, i.e. "grassy" ELM's with $\nu_{ELM} \approx 1-2$ kHz, which are similar to those observed on GP H-phases close to the density limit. Although during the H-mode the particle confinement time is normally strongly improved, the high densities and the ELM's lead to a recycling
coefficient smaller than 1. During the density build-up with PI an additional moderate gas-puff (GP) is applied, which is significantly lower than the amount of gas needed to drive a conventional H-mode discharge into the density limit. It is mainly employed to reduce the impurity influx. This motivation is different from the experiments with PI into OH and L-mode discharges where it is found that a sufficient auxiliary GP is the prerequisite to attain successful high densities and improved energy confinement [1]. In contrast to high density GP H-phases, where large sawteeth repetitively eject the impurities from the plasma centre, PI fuelled discharges lose this self cleaning mechanism and the effect of the "grassy" ELM's is not strong enough to prevent accumulation. Nevertheless, the fraction of the total radiated power to the input power $P_{rad}/P_{input}$ is slightly lower for the PI discharges ($\approx 0.4$) than for the conventional fuelled high-density H-modes ($\approx 0.5$). Although PI has a beneficial effect on the energy confinement, low central temperatures of PI discharges result in similar neutron yields as in the corresponding GP case.

Profiles:
The density profiles during the conventional H-mode tend to be very flat with steep edge gradients (fig. 1). The absolute densities closely inside the separatrix reach close to the density limit values of up to $7 \cdot 10^{19} m^{-3}$ despite ELM's. The profile peaking factor $Q_n = n_e(0)/(n_e)_{vol}$ is as low as $\approx 1.1$.

**Fig. 1:** Electron density profiles of PI (solid line) and GP (dashed line) high-density H-mode discharges with "grassy" ELM's.

Deep fuelling by PI produces bell shaped profiles in the central region on top of the edge pedestal as shown in figure 1. The density scale lengths $\lambda_n = -n_e/\nabla n_e$ are reduced in the bulk by about a factor of 5, whereas beyond $r \gtrsim 0.3$ m they are very similar in both fuelling scenarios. $Q_n$ grows with PI to $\approx 1.5$. This is, however, much lower than the maximum peaking attained in the PI L-mode ($Q_n^L \approx 2.5$) [1]. The peaking during the H-mode cannot only be explained by the sawtooth suppression alone but also requires an improvement of the central particle confinement.
In parallel the temperature profile shape is not affected by PI. Only in late phases \( \approx 400 \) ms after the transition into the H-mode strong central radiation can lead to flattened or even hollow \( T_e \) profiles. The central peaking of the \( n_e \)-profile, i.e. the reduction of \( \lambda_n \) alone, does not reduce the ratio of the density to temperature scale length \( \eta_e (\approx \eta_i) \) below 1, which might indicate the suppression of the transport channel associated with the ion temperature gradient driven turbulence. Only in the case of radiation induced \( T_e \) profile flattening \( \eta_e \) falls below 1 in the bulk \( (r \leq a/2) \).

The H-mode particle confinement, further improved by PI and the sawtooth suppression leads despite the ELM’s to undesired strong impurity accumulation. This is clearly demonstrated by the peakedness of the normalized bolometric radiation \( P_{\text{rad}}/n_e \) and \( Z_{\text{eff}} \) profiles shown in figure 2. Spectroscopic data indicate that light impurities (i.e. oxygen) as well as heavy impurities (mainly the target material copper) accumulate.

**Fig. 2:** \( Z_{\text{eff}} \) and normalized radiation profiles \( P_{\text{rad}}/n_e \) of the same discharges as in figure 1 showing clearly the impurity accumulation of PI discharges. Line styles are defined in figure 1.

Confinement:

The resulting H-phases with PI last up to 400 ms (\( \sim 8 \) confinement times) and end, if they do last not too long, with an H- to L-transition otherwise in a radiative collapse. The confinement times are compared with those of high-density GP H-mode discharges which show generally a moderate \( \tau_E \)-improvement compared with the L-mode \( (\tau_E/\tau_E^{\text{Goldston}}[6] \approx 1.4) \). On ASDEX the maximum improvement factors of about 2 are only routinely attained in conventional ELM-free H-modes [2] at significantly lower densities, normally not subject to pellet fuelling experiments. With PI only, \( \tau_E \)-enhancements of 10 % to 15 % could be achieved (fig. 3) despite successful peaking of the central density profile.

The PI H-mode combines some confinement properties of PI L-mode and standard H-mode discharges. On one hand, the \( \tau_E \)-improvement seems to be more favourable at low \( q_a \) as found in the L-mode. On the other hand, the beneficial density dependence
of \( \tau_E \) of the PI L-mode discharges [7] could not be attained, \( \tau_E \) is independent of \( \bar{n}_e \).

Fig. 3: \( \tau_E \) normalized to \( \tau_E^{Goldston} \) of PI (solid symbols) and GP (open symbols) L- (squares) and H- (stars) phases as function of the total heating power.

Conclusion:
It is found that PI into H-mode discharges can produce a regime of improved plasma performance which combines characteristic beneficial properties of the H-mode and the pellet mode. The density and temperature profiles in the edge indicate that the well known transport barrier at the boundary is unaffected by PI. The improvement of \( \tau_E \) with PI results from a further confinement enhancement in the central region. The increased central particle confinement, however, aggravates the problem of impurity accumulation and correlated central radiation losses.

The achieved confinement times with PI do not exceed the values obtained in standard ELM-free H-phases. This lack of further improvement, which is significantly lower than in the PI L-mode, may be attributed to the only moderate peaking of the density profile shape, which might not be enough to suppress the \( \eta_p \)-mode driven transport. This is consistent with the observation that \( \tau_E \) does not correlate with the profile peaking.

Similar low global confinement improvements with PI are also found on JFT-2M [3] and JET [4].

References:
1. INTRODUCTION AND EXPERIMENTAL OBSERVATIONS

Low Z impurity pellets are injected into ohmic TEXT discharges. A persistent m=1 structure, somewhat akin to snakes in JET, is often observed after injection. Observed pellet ablation is compared with the predictions of an improved neutral shielding model. A model for the observed trajectory curvature and striations (modulated ablation) is proposed. This leads to novel diagnostic applications of pellets including the possibility of measuring profiles of the suprathermal density and energy, safety factor q, and suprathermal D⊥ which may be a measure of internal magnetic fluctuations (possibly responsible for anomalous transport).

Cylindrical pellets of Li, B, C, Na, Mg, Al, etc. are injected radially into TEXT (R₀=1 m, a=0.26 m, circular TiC limiter) from the outside midplane using a one stage pneumatic gun. We demonstrate for the first time central injection of LiH (a room temp. solid) which we strongly suggest be considered for reactor fueling using the low T₁ reaction LiT+D [1] since its heat of sublimation, mechanical strength, attainable pellet velocity, and penetration depth are superior to those of frozen tritium. Pellets with an equivalent spherical radius R_p=70-300 μm, mass m_p=3-90 μg, and velocity v_p=40-700 m/s are injected into H, D, or He discharges with prepellet parameters n_e=1-8×10¹⁹ m⁻³, T_e(0)=0.6-2.0 keV, B_T=1.5-2.8 T, q_a=2-8. T_e(0) drops well before the pellet reaches r=0 suggesting a higher than normal X_e (the "cold wave" seen with H₂ pellets). The pellets increase n_e by 0.05-6×10¹⁹ m⁻³ (1-600%); up to 400% without disruption. R_p is determined either by optical microscopy or a 1 μg resolution analytic balance, and verified by noting a density rise commensurate with particle accounting. The pellet trajectory and penetration depth λ is determined from digitized CCD camera images, and v_p is measured with an 8 channel photodiode camera. λ varies from 6 to 52 cm (the plasma diameter) depending on plasma and pellet parameters.

![Graph](image)

Fig. 1. Carbon pellet penetration λ vs. T_e(0) for r_p=120 μ, v_p=240 m/s, n_e=2×10¹⁹ m⁻³.
Experiments in which $R_p$, $v_p$, $\bar{n}_e$, $T_e(0)$, and material (heat of sublimation) are each varied independently yield the expected qualitative behavior in $\lambda$ as in Fig. 1. $T_e$ and $n_e$ profiles are also measured allowing detailed comparison with an ablation model. The pellets have a constant radial velocity but curve toroidally with $0.1<\nu_\phi(r=0)/\nu_r<5$. Reversing $I_p$ and tangential view photos confirm a helical acceleration along field lines in the electron drift direction. Striations in visible continuum and line emission (width $w<10$ mm) are apparent particularly for larger pellets and $\bar{n}_e<3\times10^{19}$ m$^{-3}$. More pronounced dips in ablation light are sometimes observed near $r_\phi=1$ ($w<1.5$ cm) and possibly near $r_\phi=2$.

2. CURVATURE, ABLATION, AND STRIATION PHYSICS

The pellet trajectories can curve due to a heat flux asymmetry $\delta$ resulting from $I_p$. The momentum balance is $P_p=P_e-P_g$ i.e. the momenta of the pellet, electrons stopped by the pellet (negligible), and the ablating gas (rocket effect). The resulting pellet velocity $v_\phi$ is

$$v_\phi = \frac{3}{4} \frac{v_g^2}{r_p} \int_0^r \frac{\delta}{r} \frac{dt}{t}$$

where $v_g$ is the gas velocity. For a Vlasov-Maxwell equilibrium with $L_\|_e=L_T$ we find

$$\delta = 2 \int v^3 f(v) \frac{dv}{\int v^3 f(v) \frac{dv}{dv}} = \frac{45}{4} \sqrt{\frac{\pi}{v_\phi}} \sqrt{\frac{T_e(r)}{T_e(0)}} = 0.05 \text{ to } 0.20$$

where the drift velocity $u_D=J(0)/e\bar{n}_e(0)$. Note $\delta$ which depends on $v^3$ is dominated by suprathermals and is very sensitive to the form of $f(v)$. For a drifted Maxwellian $\delta$ is 3.7 times smaller. A comparison of the predicted final toroidal deflection $\lambda_\perp$ with carbon $\lambda_\perp$ vs. $R_p$ data shows good agreement over the range of $\lambda_\perp=0$ to 10 cm when a constant factor $g$ representing uncertainty in $f(v)$ is adjusted for best fit ($g=1.95$ in this case). $\lambda_\perp$ increases with $R_p$ because $\delta$ is greater near $r=0$. $\lambda_\perp$ for C pellets is greater than for Li, as observed, because the toroidal acceleration $a_\phi=\nu_\phi^2/(m_a r_p)$, with $m_a$ the atomic mass, is higher due to the greater ablating temperature and smaller typical $r_p$.

The ablation data is compared with a neutral shielding [2] simulation code which integrates the heat flux along the observed pellet trajectory using measured $n_e$ and $T_e$ profiles.

$$Q_{\text{code}} = \int_0^\lambda q(\lambda) A_p(\lambda) \, d\lambda \quad Q_{\text{exp}} = \int_0^\lambda q(\lambda) A_p(\lambda_{\text{code}}/\lambda_{\text{exp}}) \, d\lambda \quad R_Q = Q_{\text{code}}/Q_{\text{exp}}$$

where $\lambda_{\text{exp}}$, $\lambda_{\text{code}}$ are the observed, predicted penetration, and $A_p$ is the pellet surface area. The pellets consistently ablate faster than predicted, see Fig. 2 (even using prepellet $T_e$ profiles, not accounting for the cold wave) though $R_Q\to1$ for larger or faster pellets which reach $r=0$ ($\lambda=a$). The anomalous heat flux varies from 10% near $r=0$ to 80% near $r=a$. Consequently we modified the ablation model to account for suprathermal heat flux, variations in magnetic shielding $f_B$ as $r_p\to$ gyroradius, finite connection length $L_c$ and heat flux on rational magnetic surfaces, and nonspherical pellet shape. The effective heat flux $q_{\text{eff}}=(f_B\eta+0.5\zeta)q_\infty$ where $q_\infty$ is the background thermal flux, $\eta$ is the neutral shielding for thermais, $\zeta$ the suprathermal enhancement factor. A Bessel function profile is assumed for $\zeta=n_3 J_0(kp)[E_{3}/T_e(p)]^{3/2}$, where $n_3$ is the suprathermal density at $r=0$ of a monoenergetic population of energy $E_3$. $k=2.2$ is determined using $R_Q$ vs. $r$ from the $\lambda$ vs. $R_p$ data. From the magnitude of $R_Q$ and pellet parameters we can fix $n_3 E_3^{3/2}$ using a Bethe-Bloch model for the electron energy deposition efficiency $e$. For a typical C pellet $e\to0$ above 150 keV. Limited x-ray pulse height data of $n_3$ constrains the relevant $E_3<100$ keV. For 100 keV we find $n_3/\bar{n}_e(0)=2\times10^{-4}$. A lower limit on $E_3$ is provided by noting that low $E_3$ results in pellet surface heating. If the entire suprathermal flux (mainly in the electron drift direction) were at 30 keV, more than 10 times the observed $\nu_\phi$ would result. This consideration yields $E_3>55$ keV. So
\( n_s/n_e(0)=2 \times 10^{-4} \) at \( E_s=55-100\text{ keV} \) with a \( J_0(2.2\rho) \) profile can explain the observed edge ablation enhancement and is consistent with the curvature data. If we produce a runaway discharge with high \( n_e(E>500\text{ keV}) \) but normal \( n_e(E<100\text{ keV}) \) (possible because \( \tau_p \) increases monotonically with \( E \)) no significant effect on ablation or curvature is observed, consistent with the picture that runaways pass right through the pellet and do not contribute to curvature or ablation.

**Fig. 2.** \( R_Q=Q_{\text{code}}/Q_{\text{exp}} \) vs. \( R_p \) for \( \nu_p=350\text{ m/s Li, } B_T=2.2\text{ T, } I_p=220\text{ kA, } \bar{n}_e=3\times10^{19}\text{ m}^{-3} \).

**Fig. 3.** \( \alpha \) vs. \( \lambda, \ E=50\text{ keV, } \nu_p=230\text{ m/s, } R_p=220\mu, B_T=2.8\text{ T, } I_p=320\text{ kA, } \bar{n}_e=8\times10^{19}\text{ m}^{-3} \).

Striations in pellet ablation light can be associated with finite \( L_e \) on rational magnetic surfaces. We extend the model of Pegourie and Dubois [3] by deriving an analytic expression for the heat flux attenuation \( \alpha \) for a realistic \( f(v) \), accounting for \( D_\perp \), and by adding monoenergetic suprathermals. These fast electrons can be depleted in a time shorter than
\( \tau_D = 2 \tau_p / v_p \), but gained by \( D_\perp \) which for high \( E_s \) should be dominated by magnetic fluctuation driven transport. For a Boltzmann energy distribution

\[
\alpha = \frac{1}{\tau_D} \int_0^{\tau_D} \left[ 1 - \frac{8}{3 \sqrt{\pi}} \int_x^{\infty} z^4 \exp(-z^2) \, dz \right] \, dt
\]
where \( x = \sqrt{\frac{m_e}{2} \left( \frac{L_c}{\tau} \right)^2 / T_e} \)

\( L_c = 2 \pi R m \sqrt{1 + (\tau/Rq)^2} \) and \( q = m/n \). \( \alpha \) can be expressed analytically in terms of parabolic cylinder functions. We account for \( D_\perp \) with \( \alpha_D = 1 + (\alpha - 1) \exp[-D_n \tau / \tau_p] \) with \( \tau_p = L_c / v_c \).

The striation brightness modulation is quite sensitive to \( D_n \) (Fig. 3) and may serve as a diagnostic of internal magnetic fluctuation induced transport. Our data is consistent with \( D_n = 0.2 \, m^2/s \) which assuming Rochester-Rosenbluth diffusion yields a mean internal \( b_{rms} / B_0 \leq 1 \times 10^{-4} \). A high precision \( q(r) \) diagnostic may be possible by fitting model \( q \) profiles to ablation rate profiles.

3. PELLET INDUCED \( m=1 \) OSCILLATION

Long duration (\( \tau \leq 160 \, ms >> \tau_E \)), large amplitude (soft x-ray \( \Delta A / A \leq 50\% \)) \( m=1 \) oscillations often occur after injection (Fig. 4). For the TEXT \( m=1 \) events (vipers), sawteeth are invariably suppressed, with the perturbation localized within \( r = q_0 \). Typically \( \Delta n / n e = 3\% \), \( \Delta T_e / T_e = 10\% \) and \( \Delta A / A = 35\% \). Vipers rotate in the electron diamagnetic direction with a frequency \( \approx 7 \, kHz \).

Over 100 vipers \( (=30\% \) had \( \tau > 30 \, ms \) were observed after C, Li, Al, and Mg pellets for a wide range of plasma parameters, \( 2.7 < q_a < 4.7 \), but most reliably \( (50\% \) success) for \( B_T = 2.5 \, T \), \( I_p = 300 \, kA \) with large \( r_p = 200 \, \mu \) carbon pellets. Two distinct growth rates for the soft x-ray \( m=1 \) oscillation are observed; \( \gamma_{\text{slow}} = 0.3 \, ms^{-1} \) and \( \gamma_{\text{fast}} = 0.8 \, ms^{-1} \).

\( \gamma_{\text{slow}} = \gamma_{\text{FKR}} \approx \gamma_{\text{R}}^{3/2} \gamma_{\text{A}}^{2/5} \) the Furth, Killeen, Rosenbluth linear tearing mode rate, while \( \gamma_{\text{fast}} = \gamma_{\text{K}} \approx \gamma_{\text{R}}^{1/2} \gamma_{\text{A}}^{1/2} \) the Kadomtsev nonlinear rate, with \( \gamma_{\text{R}} \) the resistive rate and \( \gamma_{\text{A}} \) the poloidal Alfven rate. It seems likely that vipers are simply a saturated \( m=1 \) resistive kink mode.

![Fig. 4. m=1 mode - Central soft x-ray brightness after a R_p=200 \, \mu \) carbon pellet at 305 ms.](image)

We wish to acknowledge helpful discussions with R.K. Fisher and P.B. Parks

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PELLET ABLATION STUDIES IN THE TCA TOKAMAK

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ABSTRACT Pellets with a variation in size of an order of magnitude and of injection velocity of a factor of 4 were injected into a Tokamak with a range of plasma current and density. The penetration into the plasma has been measured against these parameters. Contrary to most ablation models, only weak scaling was found which suggests that the pellet ablation process is dominated by other factors.

Introduction

The injector installed on TCA (R = 0.615 m, a = 0.18 m) was used to inject frozen hydrogen pellets with a large range of velocities (200 - 800 m/s) and a wide range of size (3 x 10^{18} - 3 x 10^{19} particles), into a wide variety of target plasmas. One of the main goals of this flexible system was to study pellet penetration as a function of the parameters which most frequently appear in scaling laws. Thus, in this paper, we concentrate on the pellet ablation process, rather than the effect on the plasma. Although a wide range of visible diagnostics were used to monitor the injected pellet, the Hydrogen Balmer-alpha (Hα) emission monitored with a diffused fast photomultiplier acquisition is used to define the pellet penetration. The pellet velocity was measured before and during the ablation process and used to transform the ablation time into distance.

Experiment

Several experimental runs were performed in which the pellet size and velocity together with the deuterium target plasma density and plasma current were varied over as wide a range as possible. Most of the theoretical models of pellet ablation can be summed up by a generic scaling law:

\[ \text{Penetration length} = \frac{m_p \xi}{I_p \gamma n_e \delta} \]

where \( m_p \) is the total number of particles contained in the pellet, \( v_p \) is the pellet velocity, \( I_p \) is the plasma current and \( n_e \) is the central density.

In order to extract the dependence of pellet ablation on individual parameters in Fig 1 we selected discharges in which only one of these parameters changes significantly. Apart from the dependence on the plasma current, there is surprisingly little variation in the penetration depth.
The plasma current must be reduced to $\sim 60$ kA ($q(a) \sim 5-6$) for the pellet to penetrate to the plasma centre. At full current ($\sim 130$ kA; $q(a) \sim 3$) most pellets were ablated at or before half the minor radius. Although we have plotted only the greatest penetration achieved by a pellet, the H$\alpha$ line intensity was often highly peaked near $a/2$ indicating a localised region of high ablation. At 60 kA, where we have deep pellet penetration, the H$\alpha$ line intensity shows a peak followed by a low intensity region before the machine axis. Pellets that passed the plasma centre were observed to continue with a far reduced ablation rate, suggesting that this region of the plasma was already cold.

The pellet size is seen to have little effect on the penetration (Fig 1c), and there is only a small tendency to increasing pellet penetration with increasing pellet velocity (Fig 1d). Extrapolating this dependence, a pellet velocity of $\sim 1500$ m/s would be required to reach the plasma centre. However, the validity of this simple extrapolation is cast into doubt by the presence of several nonlinear and highly localised enhanced ablation phenomena which were often observed during pellet ablation and seemed to increase in magnitude with increasing pellet velocity.

We would expect the penetration to decrease with increasing target density, although Fig 1b shows if anything the opposite dependence. Fig 2 highlights this
conclusion, where the normalised Hα line intensity is shown for two pellets of similar mass (~1.5x10^{19} \text{pcles}) and speed (~720 \text{m/s}) injected into two different target plasma densities. If the light intensity is assumed to be proportional to the local ablation rate, injection into lower target density (Fig 2a) shows a monotonic ablation rate, whereas injection into a higher target density (Fig 2b), shows a series of localised features which we term "positive striations". There was little difference between the profiles as measured by a 4 channel FIR interferometer, which leads us to conclude that the density gradient may determine the existence of striations. Fig 3 shows the number of these striations plotted against the central target plasma density where a striation is defined as a maxima which exceeding a certain threshold between two minima. The number of striations increases with the plasma density and the intensity of these striations also increases, reaching 70 % of the maximum intensity at high density. Comparison of shots with the same, and different, plasma currents did not yield a direct dependence of the location to the rational magnetic surfaces.

**Figure 2:** Normalised Hα emission intensity for two similar pellets injected into plasma with different densities.

**Discussion**

The above description becomes less simple as we look at the additional experimental data. We have observed different penetration distances under similar conditions, which cannot be accounted for by errors or uncertainties in the measurements.
Figure 3: The number of striations is seen to increase with central plasma density, a complexity not accounted for in theoretical models of pellet ablation.

Some additional phenomena such as the increased pellet ablation in discharges with a high runaway electron content are qualitatively easy to understand. Other phenomena including curved pellet trajectories and increased penetration in the plasma current descent are not easy to quantify in the above terms. Specifically, the intensity and number of positive striations which accompany the pellet ablation have a strong influence on the penetration. These phenomena can be so strong that a pellet changes trajectory at a very localised radial position, which is often accompanied by plasma current disruption. The analysis of these effects will be greatly aided by the other available diagnostics which give complementary information on the spatial and temporal evolution of the pellet trajectory in the plasma. These additional diagnostics may help to resolve existing discrepancies between a simplistic ablation model and the complex physical processes that actually occur during the pellet penetration.

Conclusion

A large range of conditions for pellet ablation have been studied in TCA and we have investigated the experimental scaling of pellet penetration with plasma current, plasma density, pellet size and pellet velocity. We have found no dominant term in this scaling, although we did observe reduced penetration with increasing plasma current and, rather surprisingly, slightly increased penetration with increasing target plasma density. The pellet size was not found to significantly influence the penetration and there was only a slight increase in penetration with increasing pellet velocity. Other very non-linear phenomena such as positive striations and regions of reduced ablation near the plasma core appear to have a greater influence on the penetration, and these in turn appear to depend on parameters like the plasma density gradient. In general, we conclude that there is no simple formula that can guarantee a given penetration for certain conditions and that seemingly similar discharges can display quite different behaviour.

Acknowledgements

We wish to thank the whole TCA team for its excellent support. This work was partially supported by the Fonds National Suisse de la Recherche Scientifique.
STUDIES ON FAST OSCILLATIONS AND ON PARTICLE TRANSPORT DURING SAWTOOTH CRASHES IN PELLET-INJECTED TEXTOR PLASMAS

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[1] Introduction

A pellet injector of the pneumatic type has been installed on the TEXTOR tokamak in order to investigate details of the behavior of pellet injected plasmas and the overall particle balance of torus plasmas with both particle sink (ALT-II pump limiter) and source (pellet).

A pellet with a size of 1.4 mm both in diameter and in length with velocities in the range of 400~900 m/s has been injected into ohmic D₂ discharges (Iₚ=340 kA, \bar{n}_e=1~4 \times 10^{19} \text{ m}^{-3}). A couple of interesting features has been studied. Fast oscillations of large perturbations in density and smaller ones in temperature have been observed just after the pellet injection. Also, the characteristics of the sawtooth crash during the temperature recovery phase of the pellet injected plasma have been studied by examining the density profiles before and after the crash.


A fast oscillation in the density profiles has been observed by a 9-channel FIR interferometer system and by a soft X-ray mode analysis system just after pellet injection into TEXTOR plasmas. The phenomena seem to be similar to those observed by soft X-ray measurements in Alcator C¹ and JET², and to recent results in JT-60³, although there exist several differences.

A typical result of electron line densities obtained by FIR interferometry is shown in Fig. 1. A fast oscillation of large perturbation is clearly seen on channels No.5 and No.6 just after the pellet injection. The poloidal mode number of this activity is most likely m=1 or at least odd, and the direction of the rotation has been found to be to the electron diamagnetic drift one by analyzing both FIR and soft X-ray signals. The oscillation is localized clearly inside the q=1 surface, which is quite different from the results in Alcator C and JET. In both experiments fast oscillations have been observed after pellet ablation; for Alcator C the oscillation has taken place outside the q=1 surface, and for JET the "snake" has been described to occur at the q=1 surface.

The frequency is typically around 0.7-2 kHz and the oscillation
usually terminates after about ten msec. In some cases, however, the oscillation revives once again after its first termination and a certain interval; this has not been observed on other experiments. The frequency of these phenomena in TEXTOR is between those in Alcator C (~5kHz) and JET/JT-60 (0.03-1kHz). The oscillations start immediately after the injection in Alcator C and TEXTOR, while in JET and JT-60 a certain build up time (~10ms and/or 0.2~0.3 sec) is needed. Also, the duration times differ from each other, and again the TEXTOR case is between Alcator C (~1ms) and JET/JT-60 (0.2~2 sec). The temporal behavior of the equi-temperature profile is shown in Fig.2. Typical values of the relative density perturbation are much larger than those of the temperature perturbation.

The characteristics of the phenomena seem to be qualitatively similar in these machines, however there exist several differences in quantitative features. One explanation of these phenomena might be the persistent trapping of particles ablated from a pellet inside magnetic islands, however more detailed studies will be necessary to identify the mechanism clearly.

[3] Particle Transport during a Sawtooth Crash

In the high density regime of TEXTOR plasmas, sawtooth oscillations generally disappear after pellet injection for about 0.5~1 sec, and improved confinement is obtained in this regime. (τE~200ms)

In the medium or lower density regime (n_e<3×10^{19} m^{-3}) the sawtooth activity persists and the following studies have been carried out. Temporal evolutions of the density profile of pellet injected plasmas have been studied in detail. In a typical sawtooth oscillation, the density usually does not change much in the sawtooth crash phase, as observed in many tokamaks and also as in TEXTOR plasmas before pellet injection, which is shown in Fig.4. Here, the sawtooth crash occurs at t=1528.2 msec (Fig.3), and the density profiles stay almost unchanged before (t=1527/1528 msec) and after (t=1529 msec) the crash. On the other hand, a rather large change in density profile has been observed at the sawtooth crash (between t=1609/1610 msec and t=1611 msec) during the temperature recovery phase of the pellet injected plasma, as shown in Fig.5. That is, an increased particle transport surely exists at the crash (t=1610.4 msec) which is in clear contrast to the usual crash (t=1528.2 msec) before pellet injection. The energy release or transport at the sawtooth crash during the temperature recovery phase of the plasma seems to be comparable or even larger than that before the injection. This means that the sawtooth activity seen in the temperature is not necessarily a measure of the total sawtooth activity in this kind of plasmas.

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Fig. 1(a) Fast oscillation of density perturbation after pellet injection. (A pellet is injected at $t=1542$ msec.)
(b) Time expansion of Fig.1(a).

Fig. 2 Time variation of equi-temperature profile after pellet injection.
Fig. 3 Temporal evolution of central temperature and density of pellet injected plasma. (A pellet is injected at $t=1542$ msec.)

Fig. 4 Behavior of density profile at the sawtooth crash before pellet injection.

Fig. 5 Behavior of density profile at the sawtooth crash after pellet injection.
1. Introduction.

Fast evolution of an electron density perturbation in the time interval between pellet ablation and density toroidal symmetrization is an important issue of pellet fuelling. The studies of the toroidal velocity of density perturbations created by pellets were carried out on TFTR [1,2], JET [3], JT-60 [4] and T-10 [5] using neutron or SXR emission and interferometer data. Both triangle [1] and self-similar [5] profiles of density perturbation with the constant and varying front speed of order of ion-sonic velocity $c_s = \sqrt{T_e/m_i}$ were used for the process description.

The main purpose of this paper was to obtain additional information on density cloud evolution in the case of fast formation of density perturbation by deuterium pellet using high spatial and temporal resolution SXR-technique along with interferometer measurements.

2. Experimental setup.

Deuterium pellets with velocities of 300 - 700 m/s and size $1.35 \times 1.35$ mm were injected at the angle of $30^\circ$ relative to the equatorial plane. Measurements of the electron density were performed using an 8-channel interferometer with 40 $\mu$s time resolution. The vertical chords of the interferometer were located at distances of $\pm 4.2$ cm; $\pm 12.6$ cm; $\pm 21$ cm; $\pm 296$ cm apart from the chamber center. The electron temperature was measured using ECE. The ECE receivers and the interferometer were situated in the pellet injection section. Measurements of SXR emission were carried out using three X-ray cameras about 5 m ($180^\circ$) toroidally displaced from the $D_2$ injector. A vertical camera and two cameras with viewing line angles of $\pm 30^\circ$ relative to equatorial plane were exploited.
3. Experimental results.

It was observed that temporal evolution of the SXR signals during pellet ablation depended on the plasma temperature. Typical behavior of SXR emission signals in ohmic and ECR heating regimes is illustrated by Fig. 1. The OH signal falls with a little delay after the ablation moment. On the contrary substantial bursts of SXR emission correlating with pellet ablation arose surprisingly on strong ECRH (~1 MW) being applied. A strong asymmetry of the SXR emission profile was observed by the vertical camera only (see Fig. 2). The increase of SXR emission was located in a narrow region of ~5 cm near the ~8 cm vertical chord. The delay between maxima of the ablation curve and the SXR signal did not exceed 40 μs.

The fast toroidal expansion of the density exhibited itself on the rf-interferometer as rapid signal increase on internal chords (~42, ~210 cm) with delay times between various chords not greater than 200 μs. Typical time evolutions of electron density perturbation during ECRH regime are shown in Fig. 3.

4. Discussion.

Numerical simulation of the interferometer data with \( T_{\text{cold}} = 10 \) eV and a high value of \( V_0 \sim 3 \times 10^5 \) cm/s (which reasonably described carbon data [5]) is not consistent with the hydrogen pellet experiment (see Fig. 3). The data can be reasonably simulated under an assumption that the electron temperature \( T_{\text{cold}} \) is close to the value of the ambient plasma \( T_e(r) \) and with the poloidal rotation velocity below \( 10^5 \) cm/s (compare dashed and solid curves in Fig. 3).

The same assumptions lead us to understanding of the SXR signals behavior. Tracks of pellet trajectories calculated using \( q(r) \) profile \( (V_0 \text{ assumed to be zero}) \) at the X-ray camera's section are shown in Fig. 4. The main density deposit from the pellet is concentrated near the end of the trajectory - in the region crossed by the ~5 cm chord of the vertical SXR camera. The SXR emission profile observed by the +30° camera should thus be approximately symmetric and rather broad. It is clear that the density perturbations qualitatively correlate with perturbations of SXR signals shown in Fig. 2. Note that the rotation of the picture in Fig. 4 due to plasma poloidal rotation does not change the agreement mentioned above only for velocities below \( \sim 10^5 \) cm/s.

The fact that no bursts of SXR signals occur during pellet ablation in OH plasma and their appearance during ECRH can be explained by the difference between two characteristic times. The first one is the propagation time of the density pulse from \( \text{D}_2 \) injector to the SXR-camera's section

\[
\tau_{dp} \sim \pi R / V_{\text{cold}} \sim \pi R / \sqrt{T_e / m_i}.
\]
The second one is the decay time of the electron thermalization on the magnetic surface $\tau_{ee}$ which can be estimated as inverse e-e collision frequency. In ECRH plasma $\tau_{ee} \sim 100 \, \mu s$ exceeds $\tau_{dp} \sim 20 \, \mu s$ in the region of maximum evaporation rate. It leads to the increase in SXR signal due to the appearance of additional density in the region observed by the cameras while $T_e$ doesn't change significantly. In the case of OH plasma the opposite relation between $\tau_{ee} \sim 20 \, \mu s$ and $\tau_{dp} \sim 100 \, \mu s$ takes place. So no increase in the SXR signal is observed.

In the case of carbon pellet ablation [5] with slow temporal evolution of the interferometer signals both approaches used in this paper and in [5] give results reasonably consistent with experimental data. But now taking into consideration the SXR signals makes us to find the fast propagation approach with ambient $T_e$ more adequate. This conclusion is consistent with the model [1] proposed on the basis of TFTR neutron measurements.

A wave mechanism of density pulse propagation indicated by the TFTR data [2] can also contribute to formation of the SXR bursts and should not be excluded.

5. Conclusions.

1. The model [5] of density expansion with the ion-sonic velocity corresponding to the ambient plasma temperature $T_e(r)$ and low values of poloidal rotation velocity is consistent with rf-interferometer and SXR emission data in the case of deuterium pellet.

2. Bursts of SXR emission during deuterium pellet injection in ECRH plasma along with their absence in OH plasma are observed. This phenomena can be explained by faster propagation of the density pulse than cooling of electrons on magnetic surface in hot plasma.

Acknowledgements.

The authors thank the entire T-10 team for assistance in experiments. Helpful discussions with V.A.Rozhansky are gratefully acknowledged.

References.

Fig. 1. The time evolution of the electron temperature, SXR and Dα emissions after injection: a) OH regime. Shot 49970: \( n_e = 2 \times 10^{13} \text{ cm}^{-3} \), \( T_e(0) = 1.3 \text{ keV} \), \( B_t = 3 \text{ T} \), \( I_p = 200 \text{ kA} \); b) ECRH regime. Shot 51173: \( n_e = 1.1 \times 10^{13} \text{ cm}^{-3} \), \( T_e(0) = 6 \text{ keV} \), \( B_t = 2.8 \text{ T} \), \( I_p = 200 \text{ kA} \).

Fig. 2. Profiles of the SXR emission at the maximum of \( \frac{dN}{dt} \). Shot 51173.

Fig. 3. The time evolution of interferometer signals. Shot 51173: points - experiment; solid curve - \( T_{\text{cold}} = T_e(r) \), \( v_{\text{pol}} = 10^5 \text{ cm/s} \); dashed curve - \( T_{\text{cold}} = 10 \text{ eV} \), \( v_{\text{pol}} = 3 \times 10^5 \text{ cm/s} \).

Fig. 4. Tracks of pellet trajectory calculated using \( q(r) \)-profile at injection section (dashed curve) and at SXR section (solid curve). Shot 51173. Arrows point out the places where \( \frac{dN}{dt} \) reaches maximum.
**PLASMA CLOUD NEAR THE PELLET INJECTED INTO A TOKAMAK**

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**Introduction.** At present neutral shielding models of pellet ablation are well-developed and are widely used for calculating of pellet penetration depth. But there is also experimental evidence that neutral spherical cloud near the pellet is surrounded by the plasma cloud elongated in the magnetic field direction (observations of H-alpha luminosity on ASDEX, TFR, T-10, TEXT for hydrogen pellets and carbon pellets photographs on T-10 and TEXT). Plasma cloud near the pellet leads to the additional shielding of the electron thermal flux from the ambient plasma, so to develop complete shielding model it's necessary to predict form and density of the cold plasma cloud.

The first approach to the problem of plasma shielding was made in /1/ where 1-D numerical modeling was performed. The calculations were interrupted at a time \( t = \frac{c_i}{v_p} \), where \( c_i \) is neutral cloud transverse size, \( v_p \) - pellet velocity. Similar estimate of the optical thickness of the plasma cloud along \( B \) was suggested in /2/: \( \tau \sim \frac{N}{l_i} \), where \( N \) is the ablation rate. According to this estimate pellet penetration depth should be independent from \( v_p \) which is inconsistent with experiment. The idea that \( \tau \) is determined by poloidal rotation velocity \( \Omega_p \), \( \tau \sim \frac{N}{l_i} v_p \) was suggested in /3/. As \( v_p \) is much larger than \( v_p \), the authors thought that plasma shielding should be reduced by poloidal convection.

In fact convection near the pellet is far more complicated. It was shown in /4/ that during ionization mass-loading current across magnetic field arises. The charge separation produced by this current leads to the screening of the ambient electric field so that plasma convection in the poloidal direction is strongly reduced near the pellet. The same
effect leads to the screening of the poloidal electric field produced by pellet motion in radial direction. So plasma cloud velocity across $\hat{B}$ is almost equal to pellet one. At the same time toroidal drift produces additional polarization so that plasma drifts in the major radius direction. Drifts in thermal electric fields lead to the symmetrization of the cloud with respect to the magnetic field.

**Model.** Here results of 2-D modeling of the plasma cloud are presented. Ablation rate $\dot{N}$ and transverse neutral cloud scale were taken as free parameters (chosen from experiment). Neutral cloud was replaced by the boundary conditions ($\hat{Z} \parallel \hat{B}$)

$$\eta|_{Z=0} = \frac{\dot{N}}{8} c_s(0) u_e \exp\left(-|x|/\ell_i - |y|/\ell_i\right), \quad |u_2| |Z=0 = c_s(0) \right) \right)$$

where $c_s(0) = \sqrt{\frac{T_e}{m_i}}$, $n$ -plasma density, $u_2$ - hydrodynamic velocity. So we consider neutrals as plasma source situated at $Z=0$ with the given transverse profile in $X$,$Y$ directions. Electron temperature $T_e(0)$ at the edge of the neutral and plasma clouds was taken as 5eV (the results are not sensitive to this value). Plasma expansion along $Z$ was modelled by set of equations

$$\frac{\partial n}{\partial t} + \nabla \cdot (n u_2) = 0, \quad \rho m_i \left(\frac{\partial u_2}{\partial t} + u_2 \nabla u_2 / \partial Z\right) = -\nabla (n T_e) / \partial Z,$$

$$\frac{3}{2} n \left(\frac{\partial T_e}{\partial Z} + u_2 \nabla T_e / \partial Z\right) + n T_e \frac{\partial u_2}{\partial Z} = -\frac{\partial \delta e_i}{\partial Z}. \quad \right)$$

Heat flux along $\hat{B}$ was taken in the model of non-local heat conductivity /5/. In the collisionless limit this expression is consistent with results of self-consistent electric field shielding model /6/. The boundary conditions at infinity were

$$n(\hat{Z} \rightarrow \infty) = n_0, \quad T_e(\hat{Z} \rightarrow \infty) = T_e(0), \quad u_2(\hat{Z} \rightarrow \infty) = 0 \right) \right)$$

Plasma source was switched on at $t=0$ and switched off at $t=(R/2 \ell_e)^{1/2} \ell_i / c_s(0)$, where $R$ is major radius. The latter value is the time of plasma shift in the major radius direction. For further details see /7/.

**Numerical simulation results.** The calculations were performed for the parameters of $T-10$: $n_0=3 \cdot 10^{13} cm^{-3}$, $T_e=500$eV. For hydrogen pellets $\dot{N} = 6 \cdot 10^{23} s^{-1}$. Parameter $\ell_i$ was taken as 2.5mm and 4mm according to experimental data. At fig. 1
isodensities are presented and at fig.2—longitudinal profiles of density, temperature and hydrodynamic velocity. The cloud is symmetric with respect to the plane $Z=0$ and has axial symmetry. The cloud is cold enough and isodensities in the core are close to the straight lines. The form of the cloud reminds "ace of diamond". At fig.2 the experimental data obtained on TEXT /8/ are compared with numerical results. The calculations for carbon pellet were also performed and the form of the plasma cloud was even more close to the "ace of diamond". Similar "ace of diamond" form of the cloud was observed in experiments on T-10 /9/.

**Analytical model of the plasma cloud.** It's possible to suggest a simple model of the plasma cloud which is based on numerical results. For constant electron temperature of the cloud $T_e = T_e(0)$ a self-similar solution for the plasma expanding into vacuum is well-known. For $t = (R/2L_e)^{1/2} L_e/c_s(0)$ we obtain $Z > 0$

$$n = \frac{N}{8c_s(0) L_e^2} \exp \left( -\frac{Z}{\sqrt{RL_e/2}} - \frac{c}{L_e} - \frac{c_s}{L_e} \right),$$

(4)

$$u_2 = c_s(0).$$

Isodensities corresponding to (4) are straight lines and (4) are in sufficient agreement with core profiles, fig.1,2. So the estimate for the optical thickness of the plasma cloud is

$$S = \sqrt{RL_e/2},$$

(5)

where $S$ depends only on major radius and neutral cloud size.

**Conclusions.** The form of the plasma cloud near the pellet reminds "ace of diamond". 2-D numerical results are consistent with TEXT data for hydrogen pellet and luminosity observations for carbon pellets on T-10.

**References**

Fig. 1 Isodensities for hydrogen pellet, $\ell_i = 2.5\text{mm}$, density in $\text{cm}^{-3}$. Dotted line—iso-densities calculated according to (4).

Fig. 2 Longitudinal density (a), electron temperature (b) and hydrodynamic velocity (c) profiles for hydrogen pellet, $\ell_i = 4\text{mm}$. Dotted lines—self-similar solution (4). Experimental results from TEXT are presented.
1. Introduction.

Information on particle transport is of great interest in connection with fusion reactor physics [1,2]. Taking into account the lack of data on density transport an inverse task technique [3,4] was developed and applied to investigate T-10 plasma. It allowed to obtain information about radial dependencies of diffusion coefficient and pinch velocity in various regimes using both carbon and hydrogen pellets as a perturbation source.

2. Experimental layout.

Carbon pellets with size of 0.3 mm and hydrogen ones of 1.35 mm were injected into the plasma. They formed positive perturbation of the electron density. A 8-channel RF-interferometer with vertical plasma probing was used for electron density measurements. Studies were performed in a regime with $I_p = 200$ kA, $B_t = 3$ T, $a_\perp = 28$ cm and various levels of plasma perturbation. The other parameters of the shots are presented in Table.


The method of particle diffusion coefficient $D$ and pinch velocity $V_{pinch}$ study is based on the assumption that in the region of density perturbation caused by a pellet the electron sources are negligible. In this case the evolution may be described by the continuity equation

$$\frac{\partial n}{\partial t} = - \text{div} \Gamma.$$

Particle flux $\Gamma$ is assumed to consist of the diffusion and convection terms.
\[ \Gamma = -D \cdot v_n + V_{\text{pinch}} \cdot n. \]

It was also supposed that \( D \) and \( V_{\text{pinch}} \) depend on the minor radius only. The values averaged over the evolution time interval (see Table) and describing satisfactorily RFI signals were determined. Such an approach contains definite disadvantages dealing with a possible dependence of transport coefficient on density. However, it develops the earlier used method of plasma simulation with artificially postulated dependencies of transport coefficients on minor radius [5] \((D = \text{const}(r), \ V_{\text{pinch}} \sim r/a \ \text{or} \ V_{\text{pinch}} \sim (r/a)^2 \ \text{etc.})\).

The RFI data were smoothed by convolution with a gaussian kernel, and the inverse transport problem was solved for the symmetrical part of the \( n(r, \vartheta) \) profile. The set of algebraic equations for density fluxes at various time layers

\[ \Gamma_i = -D(r) \cdot v_n_i + V_{\text{pinch}}(r) \cdot n_i, \ i = 1 + k, \ k = 2 \div 10, \]

was solved using the least squares procedure. In order to avoid singularities on the coefficients' profiles the zeros of the numerator and denominator of the solution were driven to coincidence by introducing a regularization factor \( \alpha - 1, \ |\alpha - 1| \sim 0.1. \)

3. Experimental results.

The results of the transport coefficients determination are shown in Fig. 1 for various regimes mentioned in the Table. Both increasing and falling down dependencies for \( D(r) \) were observed. The values of \( D(r) \) vary in a wide range from \( 10^3 \) up to \( 10^4 \) cm\(^2\)/s. Corresponding values of pinch velocity have magnitudes of order \( < 10^3 \) cm/s and do not differ significantly from linear type dependence. Absolute values of transport coefficients lay in the range determined earlier by the simulation approach [5].

The developed technique allowed to observe changes in the profile and magnitude of transport coefficients under the influence of pellet injection. The effect is illustrated by the shot 49960. With a small carbon pellet \((\Delta n_e/n_e \sim 0.3)\) being injected into the ohmically heated plasma with low density a growing profile of \( D(r) \) with high values occurred. However the following injection of a large hydrogen pellet changed the coefficients drastically. It is obvious that the falling \( D \) profile and approximately neoclassical values of the pinch velocity correspond to this part of the discharge.

The results of simulation of the chord signals with two transport coefficients within corresponding time intervals are shown in Fig. 2. An agreement between experimental and simulated
signals is evident. It is interesting to note that a small hydrogen pellet changes the shape of the D(r) profile without substantial variation of the D values in the central zone (compare data for shots 49960 and 49961 in Fig. 1). On testing the plasma with greater electron density ($n_e=3 \times 10^{13}$ cm$^{-3}$, shot 49970) no significant decrease in D(r) in comparison with the case $n_e=10^{13}$ cm$^{-3}$ was observed. So no simple dependence of D or $V_{\text{pinch}}$ on plasma density exists, we think. When comparing the profiles of the electron density shown in Fig. 3 for shots 49970 and 49960 with similar values of $n_e$ and the values of D(r) differing greater than by a factor of three one may come to a conclusion that more subtle effects of the density and/or electron temperature profiles are responsible for the differences in the D, $V_{\text{pinch}}$ observed. To understand this phenomenon a more thorough treatment of the data base is necessary which will be done in the future.


An inverse task technique was developed and applied for analysis of the T-10 density evolution data. Both growing and falling down type profiles of D(r) are observed with the values of diffusion coefficient similar to the ones estimated by the simulation approach.

Substantial decrease in transport coefficients D and $V_{\text{pinch}}$ was observed after pellet injection. This decrease is more significant for bigger pellets.

Acknowledgements.

The authors thank S.M. Egorov, I.V. Miroshnikov, P.V. Reznichenko for assistance in experiments. Helpful discussions with V.A. Rozhansky are also acknowledged.

References.

Table

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<th>time interval</th>
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<td></td>
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<td>627 ± 700</td>
<td>H_2</td>
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</tr>
<tr>
<td>49961</td>
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<td>1.3 ± 1.4</td>
<td>565 ± 615</td>
<td>C</td>
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<tr>
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<td>1.3</td>
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</table>

Fig. 1. Profiles of transport coefficients D and V_{pinch}:
- - - - - - - - - # 49960, H_2;
- - - - - - - - - # 49961, H_2;
- - - - - - - - - # 49960-49961, C;
- - - - - - - - - # 49970, H_2.

Fig. 2. Time evolution of the experimental (solid) and simulated (dashed) data of RFI for two chord distances -12 and -3 cm. Shot 49960.

Fig. 3. Density profiles:
1 - before H_2 pellet, # 49970;
2 - before C pellet, # 49960;
3 - after H_2 pellet, # 49960.
Along with the main task – pellet injection utilization for plasma fueling (plasma density profile shaping) – the experiments with pellet injection are of interest from the viewpoint of studying the energy and particle balance in plasma.

A given study is devoted to the ion energy balance in the T-10 tokamak plasma under injection of a deuterium pellet, 1 mm in size [1]. For this purpose the behaviour of atomic charge-exchange fluxes was studied with a multichannel atomic analyzer [2] with a sufficiently - high time resolution (1 msec) allowing one to obtain the ion temperature behaviour representation in time under injection and to analyze the ion energy balance on the basis of these data. The measurements were realized under Ohmic plasma heating in the modes of operation with I=230 kA, B=3.0 T.

A characteristic feature of these experiments is the fact that a pellet giving an amount of particles comparable with the total number of particles in the plasma column is injected into it. In this case, in principle, one can study the nature of ion energy balance per pulse in a rather wide plasma density range.

Experimental data on the electron density and temperature behaviour at the plasma centre in time, as well as the intensity of Dα line under injection, are given in Fig. 1. As seen in this Figure, the pellet burns up for a time considerably - shorter than 1 msec, and the plasma reaction lasts for a time sufficient for studying the plasma behaviour without direct effect of the adsorption process. The behaviour of atomic charge exchange fluxes with various energies and that of a neutron flux are shown in Fig. 2. The ion temperature behaviour at the plasma centre (obtained
from the atomic energy spectra, taking account of opaque) is shown in Fig. 3. This behaviour represents the reaction of ions upon the pellet injection, one should note that, as seen in Fig.1, the electron temperature and density after the pellet burn-up practically remain to be constant for a long time exceeding 50 msec that can be a result of changes in the nature of transport processes in the plasma under pellet injection effect. In the opposite case, the electron temperature and density should relax for the time of about 5 msec, characteristic for the T-10 plasma, as seen in Fig. 3, the ion temperature relaxes for 50 msec. One should pay attention to the fact that the time behaviour of an ion temperature is determined by an electron plasma component. If the time deuterium temperature drop is determined by the time of changes in the plasma density (about 5 msec, $T_{ii} = 0.3$ msec), the deuterium temperature rise will occur for characteristic times of heat exchange between electrons and ions of deuterium ($T_{ei} = 18$ msec).

For a more precise analysis of the ion temperature behaviour a non-stationary equation of ion heat conduction was solved by numerical methods. The experimental electron temperature and plasma density profiles were initial data. A variably providing the coincidence of the calculated ion temperature with the experimental results is the ion heat conduction anomaly coefficient with respect to the neoclassical one. In more detail this procedure is described in [3]. It is assumed that there is a pulsed source of cool ions, which, by its absolute value, is determined by an experimentally registered change in the plasma density, and its profile is defined by an evaporation curve of the injected pellet. An analysis of the results has shown that the coincidence between the calculated ion temperature behaviour and the experimental one can be achieved at $K_{an} = 2.0$ at the times before and after the pellet injection (see Fig.3). This confirms the fact that a change in the plasma density profile, produced by the pellet injection, does not probably result in a change in the nature of energy confinement within the plasma ion component, in difference from the results obtained at various facilities [4].

In conclusion, the authors express their gratitude to the people involved into the pellet injection group at the Leningrad State Technical University for an opportunity to realize the measurements and to A.A. Bagdasarov and N.L. Vasin for the data on electron temperature and density.
FIGURE CAPTIONS

Fig. 1 The central electron temperature, density and intensity of Dₐ-line behaviour under deuterium pellet injection ($I_p=230$ kA, $B_t=3.0$ T).

Fig. 2 Time dependences of atomic fluxes with various energies and a neutron flux.

Fig. 3 The central ion temperature behaviour under pellet injection.

REFERENCES


Fig. 1

Fig. 2

Fig. 3
INFLUENCE OF VB DRIFT DIRECTION ON H-MODES IN JET


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Abstract

In common with other tokamaks, reversing the toroidal field direction in JET substantially alters the H-mode behaviour in single null discharges [1,2,3]. The power required to achieve an H-mode is approximately doubled and the usually observed asymmetry in power flow to the target tiles is reduced. Nevertheless, the H-mode confinement is very similar with both field directions, as indeed is the L-mode confinement. The target power observations are consistent with suggestions that the X-point induces additional power flows [4,5], but such ideas do not describe the different power thresholds.

Power Required to Achieve an H-mode

Figure 1 shows the input power and stored energy for a discharge with forward toroidal field (VB drift towards the X-point) and two with reversed toroidal field (all at 3MA, 3.4T). The forward field shot achieves an H-mode whilst the first reversed field shot does not, even when the applied power is increased by 40%. By increasing the power further an H-mode is achieved, seen in the third discharge. Figure 2 shows a plot of the power applied to achieve H-mode as a function of toroidal field [6]. When the field is reversed the

Fig. 1: With reversed field more power is needed to achieve H-mode.

Fig. 2: More power is needed with reversed field at all values of toroidal field.
minimum power required is greater by up to a factor of 2 at all values of toroidal field studied. The solid lines represent the threshold power needed to achieve an H-mode. All shots are with carbon target tiles.

**Energy Confinement in H-modes**

Figure 3 shows the plasma energy in two discharges with opposite field directions at 2.2T, 3MA with input power of 11 MW. Both achieve H-modes with confinement times of approximately 1 second. There is no significant difference between the confinement times with forward and reversed toroidal field, as is also shown by the reversed field dataset as a whole (Figure 4).

![Figure 3](image1.png)

**Fig. 3:** In similar discharges H-mode confinement is the same with either field direction.

![Figure 4](image2.png)

**Fig. 4:** Reversed field H-mode confinement is the same as forward field for whole dataset. (Solid line is ITER H-mode scaling).

**Power Deposition at X-point Target Plates**

Figure 5 shows the ratio of power reaching the outer (large major radius) target tiles to power reaching the inner target tiles for the forward and reversed field shots shown in Figure 3. Only the period between application of NBI and the peak stored energy is shown. The powers are measured using an infrared camera looking at one band of tiles (there are 32 such bands around JET). With forward field more power is seen to flow to the outer target, as expected in toroidal geometry [7,8]. When shadowing of the target tiles is taken into account the ratio between powers to the strike zones is approximately 2.5. In contrast, the power deposition in reversed field discharges is much more symmetric. The corrected ratio is approximately 0.7 in H-mode, and 0.9 in L-mode.
Improved Plasma Performance with Reversed Toroidal Field

With the more symmetric power deposition on the target tiles it should be possible to apply high heating powers for a longer time before a carbon bloom and so achieve higher stored energy and neutron yield. In fact this has not been fully exploited [9] but the results previously attained by careful sweeping of the power load across the target tiles, can now be achieved in reversed field with a stationary X-point. The more even power distribution in reversed field reduces the need to sweep the X-point. Figure 6 shows the deuteron reaction rate for the best forward field shot, where the X-point was swept, compared with a reversed field shot, where the X-point was stationary. These experiments were at 4 MA and 3.5 MA respectively, both at 2.8T with 18 MW of NBI. The reversed field discharge reached a neutron rate of $3.7 \times 10^{16}$ s$^{-1}$ and a triple product of $n_B T_I(0) \tau_E = 9 \times 10^{20}$ keV. s. m$^{-3}$.

Comparison with Theory

It has been suggested [4,5,10] that the difference between forward and reversed field X-point plasmas is a result of there being two components to the radial flow of heat near the plasma edge:

i) the flow determined by the anomalous transport (which doesn’t depend on the field direction), and

ii) an additional flow of heat from one side of the X-point to the other, which reverses when the field reverses.
Such a description can explain the observed power flows to the target plates provided the underlying heat flow is larger on the outboard midplane than on the inboard, as expected in toroidal geometry. A ratio of 3 between the power flow to the outer and inner targets in forward field, and a ratio of 1 in reversed field, is reproduced assuming the underlying heat flow to the outboard side is a factor of $\frac{5}{3}$ larger than to the inboard side, with an additional flow from one side of the X-point to the other of only $\frac{1}{8}$ of the total power conducted to the edge. The required additional flow is from inboard to outboard targets in forward field, reversing when the field reverses.

These ideas, however, don't explain the difference between H-mode power thresholds with forward and reversed toroidal field. The threshold in reversed field is twice as high as in forward field, but this is not due to a factor of 2 reduction in confinement time associated with an additional power flow localised near the X-point. Indeed, we have not discovered any significant change in L-mode confinement with reversed field. The critical parameter determining the transition to H-mode appears still to be unidentified in JET.

Conclusions

The power required to achieve an H-mode in JET is increased by a factor $\sim 2$ when the toroidal field is reversed.

Nevertheless, the H-mode and the L-mode confinement are unchanged when the field direction is reversed.

The outer strike point (large major radius) receives a factor $\sim 2.5$ more power than the inner strike point. This ratio is reduced to $\sim 0.8$ when the field is reversed.

The more symmetric power distribution makes high performance plasmas easier to achieve in JET.

The power deposition observations are consistent with the suggestions of an additional power flow in X-point plasmas, carrying heat from one side of the X-point to the other. However such ideas do not describe the difference in threshold power between forward and reversed toroidal field plasmas.

References:

CONTROL OF CARBON BLOOMS AND THE SUBSEQUENT EFFECTS ON THE H TO L MODE TRANSITION IN JET X-POINT PLASMAS

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Introduction
Rapid ingress of carbon from the X-point tiles into the plasma during high power additional heating (‘carbon bloom’) has been a severe operational restriction in the highest performance JET X-point discharges. The X-point tile surfaces in JET have been composed of 32 individual bands of tiles with restricted power capability (a total area of up to 0.6m² being involved). The upgrades to the JET machine currently being installed will consist of continuous, high power capability carbon and beryllium X-point dump plates thus, in principle, alleviating the problem. In the course of trying to solve the problem we have made studies of the dependence of the bloom on discharge parameters, some of which have been already reported [1].

Diagnosis of the carbon bloom
The plasma configuration in this experiment is shown in Fig.1, ie., Single Null X-point discharges with the X-point near the upper carbon fibre composite (CFC) tiles. A visible bremsstrahlung detector established the effective charge (Z_{eff}) of the plasma, in conjunction with ECE electron temperature and Far Infra Red (FIR) electron density measurements. Assumming dominance of carbon impurity (a good approximation in these plasmas [2]), the deuteron content (N_D^{TOT}) of the plasma was available from the total electron content (N_e^{TOT}) via N_D^{TOT} = N_e^{TOT} - (Z_C - Z_{eff})/5.

The ratio of edge carbon light (CIII) to D_α light outside the X-point region was available from the visible spectroscopy line of sight (AA′) in Fig. 1. A section of the band of X-point tiles at one Octant was viewed by a CCD camera able to see both inner and outer X-point strike zones [3] with an Infra-red filter for temperature measurements. Also in the tiles were Langmuir probes for edge density and temperature measurements [4].

For H-mode discharges where the carbon bloom occurs, the Z_{eff} rises rapidly and typically ~ 0.5 seconds later the neutron yield in the discharge shows a sharp decline and the central carbon density (measured from charge exchange recombination spectroscopy) peaks. It has previously been found [5] that the first indication of the carbon bloom is a strong decrease in the total deuteron content, although the plasma electron density can continue to increase for some time afterwards.

An increase in the ratio of carbon III/D_α light outside the X-point region is seen in these bloom discharges. Figure 2 shows the comparison of these aspects for 2 discharges, one of which suffers carbon bloom whilst the other survives until the neutral beam power is turned off.

Effect on the carbon bloom of discharge configuration and density programming
The following parameters were varied to assess their effect on delaying the carbon bloom.

i) The direction of the toroidal field and hence the ion VB drift direction. Experiments were performed with VB drift towards the upper X-point (defined as VB drift positive) and away from the upper X-point (defined as VB negative) (see Fig. 1).

ii) Toroidal field strength (B_T).

iii) Distance X-point to the tiles (ΔX-T in Fig. 1).

iv) Gas puffing before and during the high power heating. v) Plasma isotope (D or H).

For a representative sample of the database, Fig.3(a) shows for a number of discharges the total conducted energy to the tiles up to the time of the bloom, E_{cc} = \int [P_{in} - P_{rad} - W_P]dt.
Here $P_{\text{in}}$ denotes input power, $P_{\text{rad}}$ the radiated power and $\dot{W}_D$ the rate of change of plasma energy. Shots with strong gas puffing before or during the beams are excluded.

No influence of $P_T$ magnitude on $E_{\text{EC}}$ was found. Also it is seen from Fig.3(a) that the influence of the VB drift direction and the distance $\Delta X_T$ is unclear. From toroidal effects the conducted power loading is expected to be higher at the OUTER strike zone [6]. This modified by the observed difference in radiation pattern near the X-point tiles [3]. For VB drift negative it is observed that the peak of the X-point radiation is at the OUTER strike zone, whilst with VB drift positive radiation peaks at the INNER strike zone. The peaking of the radiation at the outer zone in the VB negative case reduces the extra conducted power to this zone and hence the tile temperature. Thus the strike zone temperatures are observed to be lower in the VB drift negative case [3]. This beneficial effect only shows slightly in the bloom behaviour as depicted in Fig.3(a) probably because tile edges not in the field of the CCD camera could be loaded by the various equilibria and get very hot and these edges, and not the strike zones viewed, were often the source of carbon entering the plasma.

A similar line of argument can be employed to explain the lack of systematic dependence on $\Delta X_T$, although an X-point well inside the vessel is slightly favoured by the data.

In summary, all but one of the discharges which survived the H-phase, i.e., where $E_{\text{EC}} > 10$ MJ, either had a dense, low-temperature divertor region set up by strong gas puffing before or during the NBI heating. Some of these shots are shown in Fig.3(b) which plots $\dot{E}_{\text{cond}}$, obtained in a similar manner to $E_{\text{EC}}$ but replacing $t_{\text{bloom}}$ with the H-phase termination time ($t_H$). It has already been observed [1], [5] that strong gas puffing during the high power H-mode phase controls carbon blooms and achieves long Elm-free H-modes (> 5 secs) with steady $Z_{\text{eff}} (\sim 2.5)$. For clarity only two discharges with gas puffing during NBI are shown on Fig.3(b). The point marked ‘D noELMs’ is the best of the ELM-free D plasma H-modes. That marked ‘H/ELMs’ is a discharge with frequent ‘grassy’ ELMs. The grassy ELM discharges could only be achieved by strong hydrogen gas puffing in hydrogen discharges.

**Behaviour of ELMs, hydrogen discharges**

The behaviour of a hydrogen discharges with D$^+$ NBI is shown in Fig.4. On this shot a strong hydrogen gas puff (140 mb/s$^{-1}$) was set up during NBI and frequent ELMs resulted. This stabilised the $N_D/N_e$ ratio; the energy content; the plasma density and the ratio $P_{\text{rad}}/P_{\text{in}}$. The enhancement of energy confinement time ($\tau_E$) over the Goldston L-mode value ($\tau_{\text{EC}}$) was $\tau_E/\tau_{\text{EC}} \sim 1.8$. A similar shot where the gas puffing terminated early, resulted in an early carbon bloom.

Similar ELMs could not be generated with deuterium plasmas and gas puffs for reasons still unclear. The ELMs seem nevertheless efficient at keeping carbon content of the plasma stationary. Although the inner strike zone temperature drops to ~1800°C over a period ~2 secs as the power to the zone is cut off by radiation during the gas-puffing, the outer zone remains at ~2700°C and is hence a strong source of sublimed and radiation-enhanced sputtered carbon [7], [8]. This impurity screening by frequent ‘grassy’ ELMs has previously been seen on other Tokamaks, eg. DIII-D [9].

**Behaviour of high recycling discharges**

Discharges where a dense, low temperature divertor region was set up by strong gas-puffing before the NBI heating, were able to survive without carbon bloom. These discharges tended to have systematically lower $Z_{\text{eff}}$ at the end of the H-mode for $\Delta X_T$ further inside the vessel. This improvement in divertor shielding was only mild however, due to the occurrence of hot edges on tiles at random points in any scan of $\Delta X_T$. Langmuir probe measurements and X-point tile temperature measurements enabled estimates to be made of the sputtered carbon yield from the divertor plates, although the strike zones in the field of view were not average [3]. The shielding factor $S = \phi_c^{\text{tile}}/\phi_c^{\text{plasma}}$ could be estimated assuming the particle confinement time for carbon ions in the plasma $\tau_c$ as being - equal to $\tau_p$ and $\tau_p = 5\times \tau_E$ [10]. $\phi_c^{\text{plasma}}$ was available from $\phi_c^{\text{plasma}} = N_{\text{c Tot}}^\text{TOT} + N_{\text{c Tot}}^\text{TOT}/\tau_c$ and (N$_c = N_{\text{c Tot}}^\text{TOT} - N_{\text{c Tot}}^\text{TOT})/6$. The results for a typical high recycling shot are shown in Table 1, where it can be seen that a high recycling divertor existed (divertor target flux...
\( \phi_{div} \) being >> plasma deuteron flux \( \phi_{plasma} = N_D + N_D/\tau_p - \phi_{NB} \), and that substantial shielding of impurities (\(~\text{factor} 100\)) was also obtained. Given the assumptions regarding target area and particle confinement time used to extract \( S \), the likely value is in the range \( 30<S<350 \).

**Conclusions**

For discharges without strong gas puffing before or during the H-mode, there is only slight systematic extension of the pre-carbon bloom power handling capability of the CFC tiles as the X-point is moved inside the vessel (increased connection length). There is also only slight systematic benefit in delaying the carbon bloom from reversing the 'normal' (towards X) direction of the VB drift in spite of the more favourable tile power loading consequent to this reversal. This lack of strong benefits is probably brought about by power incident on hot tile edges.

All discharges surviving prolonged (\( \geq 3\text{sec} \)) H-phases without carbon bloom have a high-recycling divertor set-up by strong gas puffing before or during the high power heating. For hydrogen plasmas this strong gas puffing creates frequent 'grassy' ELMs which stabilise the carbon content of the plasma in spite of very hot strike zones (up to 2700°C).

For representative high recycling discharges, the divertor shielding factors against carbon ingress into the plasma are estimated to be in the range \( 30<S<350 \).

**References**

[3] R Reichle et al., these proceedings.
[4] L de Kock et al., these proceedings.

<table>
<thead>
<tr>
<th>TABLE 1</th>
<th>Measured and derived parameters for H-mode discharge #22526 (JMA/2.2T/VB-ve) t=2sec. after H-mode start (9MW NB)</th>
</tr>
</thead>
<tbody>
<tr>
<td>( \Delta x ) (m)</td>
<td>Outer Strike Zone</td>
</tr>
<tr>
<td>( L_{\text{flow}} ) (m)</td>
<td>5.41</td>
</tr>
<tr>
<td>( n_{\text{sep}} ) \text{(m}^3\text{)} &amp; \text{[1]}</td>
<td>3.0 \times 10^{19}</td>
</tr>
<tr>
<td>( T_{\text{sep}} ) (eV) &amp; \text{[1]}</td>
<td>17.0</td>
</tr>
<tr>
<td>( \lambda_4 ) (m)</td>
<td>0.16</td>
</tr>
<tr>
<td>( \lambda_{\text{SOV}} ) (m) &amp; \text{[2]}</td>
<td>0.22</td>
</tr>
<tr>
<td>( \lambda_{\text{SOV}} ) (m) &amp; \text{[2]}</td>
<td>0.09</td>
</tr>
<tr>
<td>( \phi_{\text{div}} ) \text{(e}^{-1}\text{)}</td>
<td>~1.0 \times 10^{22}</td>
</tr>
<tr>
<td>( \phi_{\text{plasma}} ) \text{(e}^{-1}\text{)} &amp; \text{[3]}</td>
<td>~1.4 \times 10^{21}</td>
</tr>
<tr>
<td>( T_{\text{flow}} ) (C)</td>
<td>~2150</td>
</tr>
<tr>
<td>( \rho_{\text{sat}} ) (g/cm(^3)) &amp; \text{[1]}</td>
<td>total ~2 \times 10^{22}</td>
</tr>
<tr>
<td>( \rho_{\text{Sat}} ) (g/cm(^3)) &amp; \text{[3]}</td>
<td>total ~1.7 \times 10^{20}</td>
</tr>
<tr>
<td>Divertor Shielding : S</td>
<td>~100</td>
</tr>
</tbody>
</table>

Notes: (1) Edge parameters measured at probe nearest the separatrix (usually within 5cm of separatrix).
(2) SOL values measured poloidally at the tiles.
(3) Assuming particle confinement \( \tau_p \) or \( \tau_e = 5 \tau_p \).
Fig. 2(a) Evolution of total plasma electrons ($N_e^{\text{tot}}$), deuterons ($N_D^{\text{tot}}$) and electrons originating from impurities ($N_e^{\text{imp}}$) for an H-mode discharge suffering carbon bloom. Also the edge CIII and $D_\alpha$ radiation at $R=3.1m$ (outside the divertor) and their ratio $R($CIII:$D_\alpha$). Note increase in $R($CIII:$D_\alpha$) with the bloom.

Fig. 3(a) Comparison of the conducted energy to the tiles, $E_{\text{cond}}$ (see text) producing a carbon bloom against ($\Delta x_T$) and direction of VB drift.

Fig. 3(b) Comparison of total conducted energy to the tiles ($E_{\text{cond}}$) (see text) to the H-mode end for shots surviving without a carbon bloom. All shots have strong gas puffing before or during the H-phase.

Fig. 2(b) As in Fig. 2(a) but for an H-mode phase surviving without bloom. Note the constancy of $R($CIII:$D_\alpha$).

Fig. 4 H-mode produced by $D^+$ NBI into a hydrogen plasma with strong hydrogen puffing, leading to frequent ELMs.
THE EVOLUTION OF $Z_{\text{eff}}$ DURING H-MODE OPERATION IN JET

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1. INTRODUCTION
As part of the 1990 JET experimental campaign, discharges in X-point configuration were studied extensively with the aim of optimising plasma performance during H-mode confinement. Systematic investigations were undertaken to establish the role played by various parameters, e.g. magnetic field strength and direction, separation between X-point and target tiles, auxiliary heating power and plasma current. Of crucial importance to obtaining good performance is the control of the plasma impurity content.

This may be globally monitored by measuring the effective ion charge, $Z_{\text{eff}}$. The evolution of this parameter has been studied during X-point discharges by means of several techniques, yielding chord-averaged values and radial profiles. Three regularly-used techniques are, respectively, based upon absolute continuum measurements at visible wavelengths, charge-exchange recombination spectroscopy and high-resolution X-ray spectroscopy. In this paper, detailed results of $Z_{\text{eff}}$ measurements obtained under a variety of operating conditions are presented and discussed.

2. APPARATUS AND ANALYSIS
The continuum emission at 523.5 nm is measured absolutely using a number of discrete lines-of-sight and a multi-chord telescope array [1]. Each channel measures the line integral of the continuum emissivity, or brightness. Using the technique of Abel inversion the brightnesses recorded by the array are transformed into the radial profile of emissivity, $e(r)$. Since $e(r) \propto Z_{\text{eff}}(r) n_2(r) T_e(r)/T_0^{1/2}$, a knowledge of $n_e(r)$ and $T_e(r)$ permits $Z_{\text{eff}}(r)$ to be determined. The brightnesses determined by the other views are used to derive $Z_{\text{eff}}$.

During NBI, charge-exchange recombination spectroscopy is used to determine the densities of the dominant light impurities, [2], leading to local values of $Z_{\text{eff}}$. The CX spectra are recorded using 3 spectrometers equipped with OMASs. A fan of 12 horizontal sight-lines intersects the path of the octant 8 neutral beam at various radial positions. The intensities of the Be$^{3+}$, C$^{5+}$ and O$^{7+}$ lines recorded are proportional to the neutral beam density and to the density of the appropriate fully-stripped ions, yielding the densities of these impurities.

The high-resolution X-ray crystal spectrometer has been described previously [3]. It is used to measure the plasma central ion temperature. Recently, an extensive absolute calibration of the instrument has been performed using a laboratory X-ray source. This permits the derivation of line-averaged values of $Z_{\text{eff}}$ from absolute measurements of photon fluxes, with time resolution ~ 20 ms. The continuum level is extracted from a line-of-sight integrated synthetic spectrum, which is fitted to the
observed X-ray data. The entire measured spectrum, i.e. continuum and line radiation, is used for the fit, ensuring a good S/N ratio under most conditions [4]. Whilst the calibration is independent of other diagnostics, the deduction of $Z_{\text{eff}}$ from the X-ray spectrum is dependent on radial profiles of $T_e$ and $n_e$ measured by other means, as well as atomic data for X-ray continuum radiation. The motivation for the development of this new technique was the observation during many discharges of an enhanced background at visible wavelengths, from reflected light, complicating the evaluation of data based on absolute measurements of visible continuum.

3. RESULTS AND DISCUSSION

The impurity content of the plasma is dependent on many discharge and machine parameters. In general, machine operation is aimed at maximising the fusion performance under different conditions. Strategies aimed at minimising any increase in $Z_{\text{eff}}$ are important, for two reasons: firstly, to ensure that radiation losses from the plasma centre are kept low and, secondly, to minimise fuel dilution, i.e. to maintain a high ratio of $n_D/n_e$.

The distance separating the X-point and the divertor target tiles $\Delta x$ plays a significant role in determining the plasma impurity content. Figure 1 shows a plot of $Z_{\text{eff}}$ versus $\Delta x$ for constant NBI power. At large separations, $\geq 10\text{cm}$, the values of $Z_{\text{eff}}$ are rather low, $\leq 2$, but as $\Delta x$ is reduced from $\sim 10\text{cm}$ to $\sim 5\text{cm}$ there is a gradual increase in $Z_{\text{eff}}$ to $\geq 3$. Whilst the impurity production rate does not depend sensitively on the separation, at large values of $\Delta x$ the shielding of the plasma to penetration by impurities appears to be improved.

In Figure 2 the dependence of $Z_{\text{eff}}$ on applied NBI and ICRF power is plotted for a number of discharges. At low values of applied power the plasma impurity content at the end of the H mode is not greatly different from that of the target plasmas. However, in the case of high powers there can be significant increases in $Z_{\text{eff}}$. For the same operating conditions, in terms of plasma impurity content there is little difference between the use of either heating method. The application of both methods simultaneously appears to lead to higher values of $Z_{\text{eff}}$ for a given power than either NBI or ICRH alone. To achieve good coupling between antennae and plasma their separation must not exceed $\sim 2\text{cm}$. Consequently, the vessel walls are sufficiently close to the plasma edge that the NBI power deposited there causes enhanced impurity production. With NBI alone the plasma-wall separation can be maintained at a much larger value.

The spread in the values of $Z_{\text{eff}}$ in Figures 1 and 2 may be understood in terms of the rapidity with which the plasma is contaminated at a given power due to evaporation from the C or Be X-point tiles - an extreme example of this is the carbon or beryllium 'bloom' [5]. Strategies which can delay the onset of this phenomenon, or suppress it altogether, lead to smaller increases in $Z_{\text{eff}}$. These include sweeping the X-point vertically, to spread the deposited power on the target tiles, and heavy gas puffing into the X-point region, to establish a regime of high recycling. Also, operating with reversed toroidal field (VB ion-drift away from the target plates) results in a more uniform power deposition between the two strike zones than operation with $B_\phi$ in the forward direction, although the power threshold to achieve H modes in this configuration is higher [6].

Following the L-to-H transition there is a gradual rise in $Z_{\text{eff}}$ throughout the lifetime of the H mode, despite the global reduction in the recycling of the fuel and impurity species [7]. As discussed, the magnitude of the increase depends on a number of operating conditions. The increase
in $Z_{\text{eff}}$ indicates that either the impurity yield at the target plates has increased or that impurities are better confined than deuterons. In Figure 3 the evolution of the $Z_{\text{eff}}(r)$ profile is illustrated. The profile, which is flat or slightly hollow at the transition, becomes increasingly hollow with time as the edge impurity influx steadily rises. The peak of the $Z_{\text{eff}}(r)$ profile is located at approximately two thirds of the minor radius. At the termination of the H-mode $Z_{\text{eff}}(r)$ loses this hollowness and reverts to a peaked or flat profile.

Generally, there is good agreement between the values of $Z_{\text{eff}}$ derived using the 3 techniques. However, during H-mode operation produced by intense heating, under certain operating conditions the collection optics of the visible bremsstrahlung diagnostic can receive significant levels of reflected black-body radiation from hot target tiles (there are no viewing dumps because of mechanical constraints). On many shots this has led to spuriously-high values of $Z_{\text{eff}}$ being derived from these measurements. Nevertheless, before the heating, and within a few seconds of its termination, the measurements are trustworthy. The X-ray measurements are immune from the effects of the black-body radiation and are used to circumvent the shortcomings of the visible bremsstrahlung measurements. Also, when NBI is applied CXRS may be used to derive values of $Z_{\text{eff}}$. Figure 4 shows an example of a discharge in which $Z_{\text{eff}}$ from visible continuum is too high. The other measurements are unaffected. There is reasonable agreement between the $Z_{\text{eff}}$ value derived from CXRS measurements and that obtained using X-rays.

4. CONCLUSIONS
The global plasma impurity behaviour during H-mode confinement in JET has been studied by measuring $Z_{\text{eff}}$ using three independent techniques, under a wide variety of operating conditions. Such measurements have been important in the development of strategies to maximise fusion performance.

Of crucial importance have been techniques to reduce or spread the power loading to the target tiles, to minimise the production of impurities and to avoid or delay the phenomenon of blooming. Gas puffing into the X-point region, vertical sweeping of the X-point across the tiles, operating with large separation between plasma and tiles, and operation with reversed toroidal field have all been beneficial. These have permitted the attainment of plasmas with low levels of contamination at high heating powers, leading to record values of the product nrT [8].

5. REFERENCES
FIG. 1 Variation of $Z_{\text{eff}}$ in H mode with distance between upper X-point and target tiles. $I_p = 3.1$ MA. $P_{\text{NBI}} = 9.5$ MW. Significance of symbols:

- $+: t = 0-0.5$ s, $\Delta: t = 0.5-1.0$ s,
- $*: t = 1.0-1.5$ s, $\varphi: t > 1.5$ s

after $L \rightarrow H$ transition.

FIG. 2 Variation of $Z_{\text{eff}}$ with heating power in H mode: (a) NBI, (b) ICRF. $I_p = 3.1$ MA. Symbol details in Fig. 1.

FIG. 3 $Z_{\text{eff}}(r)$ profiles for pulse #21801. (a) $t = 7.0$ s, OH phase. (b) $t = 10.4$ s, 0.2 s after H mode starts. (c) $t = 12.0$ s, 0.2 s before $H \rightarrow L$ transition. $P_{\text{NBI}} = 10.0$ MW.

FIG. 4 Temporal variation of $Z_{\text{eff}}$ during H mode: (a) CXRS, (b) X-ray, (c) visible continuum. Shot #22323.
Hot-ion and H-mode plasmas in limiter configuration in JET


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The inner wall in JET is a shaped surface of approximately 12 m² protected by reinforced fiber-graphite tiles. Limiter plasmas have been produced in contact with the inner wall with plasma currents up to 5MA.

Additional heating by Neutral Beams injection of these plasmas has shown that in order to avoid overheating of the graphite tiles, with consequent sharp increase of carbon impurities it is essential that the plasma has the same shape as the contour of the inner wall. Fine tuning of the plasma shape feedback allows control within 1cm. The distance between the plasma and the inner wall, as calculated by equilibrium code reconstruction, was less than 1.5 cm along a poloidal inner wall length of 2m. In these optimized discharges it was possible to apply 16.5 MW of neutral beam heating for 2s. without strong influx of carbon.

In this paper the experimental results on the achievement the H-mode and the properties of hot ion L-mode on the inner wall graphite tiles are described, the last section is devoted to the presentation of experimental results showing enhanced confinement in plasmas limited on the Beryllium limiter.

The extension of the JET H-mode to plasmas limited on the inner wall

The H-mode has been achieved in tokamaks with a magnetic separatrix configuration [1,2]. The achievement of the H-mode with plasmas limited on the inner wall was reported in D-III-D [3]. with a magnetic configuration with strong edge shear. The JET inner wall H-mode is similar to that obtained in D-III-D. Improved confinement regimes in plasmas which are limited by a large limiter surface, and do not use a divertor, could be useful for plasma scenarios in tokamaks in which the plasma volume and plasma current is maximized with the aim of achieving transient improved plasma performances.

In the JET experiment the plasma discharges, with current of 2-3MA, were formed on the belt limiter. The magnetic configuration was that of an elongated limiter plasma, b/a = 1.7, with the position of the X-points 10-15cm outside the vessel. Typically the value of field line q at the plasma boundary was 5.0, while the value of the cylindrical q was 2.5, low q discharges were also achieved with value of field line q of 3.0. Following the application of Deuterium 80/140keV neutral beam injection the plasma was moved in contact with the inner wall. By stepping up the injected power an H-mode transition was induced. The H-mode transition had the signatures of the H-modes achieved in the X-point configuration, the duration of the H-mode phase was however at most 0.8s, due to the difficulty of precisely controlling the plasma shape. The time evolution traces of one of those discharges are shown in fig 1, the plasma current was 2MA, and the value of toroidal field was 1.2T. The H-mode phase appears to be ELMs free, starting several hundred ms after the increase of the injected power. The onset of the H-mode is marked by the reduction of recycling and the increase of the global confinement time, with the electron density profiles becoming flat. The ELM's free H-phase persists for a couple of neutral beams slowing time times after the contact with the inner wall had been lost and the plasma is limited by
the graphite tiles in the region at the top of the vessel. The power threshold of the Inner Wall H-modes is higher than that for the H-mode in X-point configuration. The comparative scaling with the applied toroidal field is shown in fig. 2. The trend with toroidal field appears to be consistent with the scaling observed in X-point configuration [4].

The global energy confinement time of the inner wall is similar to that of similar discharges in X-point configuration. Fig. 3 shows the values of global energy confinement versus total loss power for a series of discharges with plasma current of 2MA and toroidal field of 1.2T. The solid line is the value of the global confinement time for a series of JET X-point H-modes at the value of toroidal field of 2.2T. The confinement of the inner-wall H-mode discharges degrades with power in the same way than the X-point H-modes but the confinement time is on average 20% lower. It is possible that the reduction could be due to the lower value of the toroidal field used in the inner wall H-modes, compared with the X-point ones.

_Hot-ion plasmas limited on the inner wall_

Overnight Beryllium evaporation on the inner wall graphite protection plates produced plasmas with good control of plasma density and a significant improvement in plasma purity, therefore allowing low beam target densities $n < 0.7 \times 10^{19} \text{m}^{-3}$. With strong additional neutral beam heating on low density deuterium target plasmas it was possible to achieve an L mode hot ion regime at plasma current of 5MA, typical parameters were: plasma density $2.5 \times 10^{19} \text{m}^{-3}$, central electron temperature 9.0 KeV, central ion temperature 22.0 KeV, ratio between deuterion and electron densities at the plasma center of 0.85, confinement time of 0.6 seconds. In these conditions the best value for the triple product $n_D\tau_T\phi_T$ has been achieved for limiter discharges, of $3 \times 10^{30} \text{m}^{-3}\text{sKeV}$, and at the same time the best value of $Q_{DD}$ of $1.4 \times 10^{-3}$. The time evolution of the plasma reactivity, total input power, central ion temperature are shown in fig 4. The central electron temperature also shown is perturbed by emission of suprathermal electrons. In the hot-ion L-mode the electron density profile and the pressure profiles are quite peaked, in contrast with the hot ion H-mode, therefore the central plasma reactivity is $4 \times 10^{14} \text{neutrons/m}^3$, which is comparable to that of the best H-mode pulses [5].

_Enhanced confinement of plasmas limited on the Beryllium belt limiter_

An increase in confinement in limiter discharges has been observed in JET in a number of deuterium pulses with ICRF heating, at a plasma current of 3/5 MA with a toroidal field of 3/3.5 T. After having reached a quasi-stationary sawtooth-free condition, the plasma suddenly displays a simultaneous improvement in particle and plasma energy, as shown in Figure 5. At the same time, the $H_e$ intensity drops and impurities start to accumulate (increase in $Z_{eff}$ and $P_{rad}$), while the asymmetry in the poloidal distribution of the radiated power is strongly reduced. The transition can occur at different levels of plasma density $(1 - 2.5 \times 10^{19} \text{m}^{-3})$ and ICRF heating power $(6 - 13 MW)$ with $H^+$ or $He^+$ minority and dipole or monopole antenna configuration. In all cases, the cyclotron resonance was close to the magnetic axis. The enhanced confinement phase is in some cases associated with an increase in the $D - D$ reaction rate, and is terminated by the auxiliary power switch-off or by a “monster sawtooth” crash - but it can also persist after the latter. Unlike the usual H-mode, this limiter
regime is characterized by an increased central peaking (rather than a flattening) of the electron density profile, and no effect on the current moments or on the surface loop voltage is observed at the transition. No obvious correlation has been found so far with variations in plasma or operational parameters. Distinctive features of these discharges are the long duration of the heating pulse (the enhanced confinement appears after a time \(\gtrsim 5 \tau_0\)) and the absence of sawteeth. During the L-mode phase that precedes the transition, the plasma current slowly approaches its steady-state distribution, and the internal inductance increases (\(L_i \sim I\) is reached at the transition). The global energy confinement in this improved limiter regime is lower than in H-modes at the same power level. Nevertheless, the energy transport for the thermal plasma is found to be significantly reduced, across the plasma cross-section, with respect to the preceding L-phase. This is shown in Figure 6, where an “effective” thermal conductivity is used to describe local energy transport - ion temperature profile measurements are not available for these discharges. Globally, part of the gain in thermal energy content is offset by a reduction in \(W_{\text{fast}}\) due to the density increase.

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Figure 3 - Global energy confinement time scaling with total loss power for Inner wall H-modes (dots), and X-point discharges (solid line), for a series of 2MA discharges.

Figure 4 - Time evolution of JET inner wall Hot-ion pulse #22202, a) Plasma reactivity, b) Central ion temperature, c) Total plasma input power, d) central electron temperature.

Figure 5 - Time evolution of JET limiter pulse #20123, displaying the transition to an improved confinement regime. The ICRF minority is He³, with antennae in the monopole configuration.

Figure 6 - Local effective thermal conductivity for the discharge in Figure 5, before and after the transition. With respect to the Goldston L-mode scaling law, the thermal energy replacement time varies from $\tau_g$ to $\approx 1.5 \tau_g$. 
ICRH H-MODES PRODUCED WITH Be-SCREEN ANTENNAS
AND COUPLING-RESISTANCE POSITION FEEDBACK CONTROL

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1. INTRODUCTION: Ion-cyclotron resonance heating (ICRH) H-mode discharges [1,2] in the JET tokamak were first produced using dipole \((0, \pi)\) antenna where beryllium gettering was applied to the nickel antenna screen and the first-wall of the tokamak. These results have now been extended from dipole \((0, \pi)\) to monopole \((0,0)\) phasing of the antenna by the use of new screens made of solid beryllium bars which have reduced all ICRH specific impurity effects to negligible levels. Moreover, with Be-screens the energy confinement time in ICRH H-modes using dipole is better than found earlier with Ni-screens and is now typically 2.5 times Goldston L-mode prediction. But, with monopole the confinement time is typically 35% smaller than the dipole H-mode confinement. These H-modes have a duration of up to 2.8 s and RF power levels of up to 12 MW. ICRH H-modes are generally free of edge-localized modes (ELMs). ICRF heated plasmas often develop long sawtooth-free periods (monster sawtooth) and this feature is retained during H-modes both in monopole and dipole phasing. Further, to maintain the antenna-plasma coupling resistance constant at a desired value during the H-mode a feedback control loop on the radial plasma position has been implemented, for the first time on the JET tokamak, which helps in maintaining a good match between the antenna and plasma and keeps the coupled power constant throughout the ICRH pulse.

2. EXPERIMENTAL CONDITIONS AND RESULTS: The plasma was operated at \(B_T \simeq 2.8\ T\) and \(I_p \simeq 3\ MA\) in the double-null X-point configuration and shaped poloidally to match the antenna profile to obtain good antenna plasma coupling. There are 8 ICRH antennas which are symmetrically distributed around the torus on the low-field side. An ICRH antenna has essentially two radiating elements that are separated toroidally which can be driven in phase (monopole) or out of phase (dipole). ICRH power is delivered to the plasma by the fast magnetosonic wave which is damped mainly via the minority-ion cyclotron damping in a narrow region \((\simeq 30\ cm)\) near the ion-ion hybrid resonance zone.

2.1 Beryllium Antenna Screens: New screens made out of beryllium [3] have replaced the nickel screens on all ICRH antennas in the JET tokamak. The new screen bars are essentially square in cross section, with corners generously chamfered, with a pitch of 29 mm and a gap of 5 mm between the bars. Thus the new screen is much more open than the nickel screen in which successive elements were of inverted-T cross section and plasma particles had practically no direct line-of-sight to the central conductor. Antennas with the new more open Be-screens have had no operational problems. In fact, their construction in beryllium, has reduced the ICRH specific impurities to negligible levels even in the operation with monopole phasing of the antenna. This has allowed to produce ICRH H-modes with monopole which was previously not possible with beryllium gettered nickel screens probably because evaporated beryllium could not reach the screen surface facing the gaps. This was a part of the source of nickel impurity from the screen [4]. Another part of nickel impurity comes from the recirculation (deposition, erosion) of Ni at other first wall surfaces. A comparison of the fraction of radiated power due to nickel impurity with the above two types of screens in H-mode discharges shows that it was about a factor of 4-8 higher with the nickel
Be-gettered screens than with Be-screens at RF power levels of 8-12 MW. At lower powers, the difference was not very significant.

2.2 Coupling Resistance Position Feedback Control: There is a steepening of the edge-density profile at the transition from L to H-phase. The scrape-off length of the edge plasma decreases and the density at the edge falls off more rapidly than in the L-phase. This leads to a sudden decrease of the coupling resistance at the transition and puts a severe burden on the antenna matching system which has to remain tuned to continue to couple power to the plasma even with a significant variation of the load impedance on a rapid time scale (60-80 ms). Often these transients lead to tripping of RF generators resulting in ragged output power trace. Initially this problem was partially mitigated by moving the plasma towards the antenna at the anticipated time of the transition to restore the coupling. However, the problem was resolved satisfactorily only by implementing a radial plasma position feedback control by acting on the current of the vertical field winding. The overall response time is typically 100 ms. The error signal was provided by the difference in the measured coupling resistance and the value desired. The result of such a feed back control is shown in Fig. 1 where the coupling resistance is maintained constant at the desired value of \(3 \pm 0.5 \Omega\) by controlling the plasma position during the entire RF pulse that encompasses the L and H phases of the discharge and the two transitions. The coupled output power was also maintained constant. Note that the generator frequency feedback control which has been implemented for automatic matching on the JET ICRH system is also necessary to compensate for the changes in the reactive part of the loading. Typically \(\Delta f = \pm 20 \text{ kHz}\).

2.3 ICRH H-mode Time Traces: In H-modes produced by ICRH alone with monopole or dipole all characteristics typical of H-mode discharges are found (see time traces of Fig. 2-4). For example in the case of dipole at the transition from L to H phase, one can see from Fig. 3, a drop in \(\Delta x\) emission, increase in the plasma density, a small decrease and then a gradual increase in the radiated power from the plasma, and more importantly, an increase in the stored energy at the transition (at constant power level), and an increase in the D-D reaction rate \((R_{\text{DD}})\). Figure 2 shows a high performance ICRH H-mode discharge obtained with beryllium screen antennas operated in dipole phasing where the energy confinement time \(\tau_\text{E} \approx 2.8\tau_\text{c}\) where \(\tau_\text{c}\) refers to the Goldston [5] L-mode prediction. The DD reaction rate \((R_{\text{DD}})\) of \(5.5 \times 10^{15} /\text{s}\) achieved is about a factor of 10 higher than the value obtained at the same power level in 3 MA limiter discharges. Note that the discharge is ELM free. A long ICRH H-mode discharge with dipole is shown in Fig. 3 in which deuterium gas was pulsed into the top X-point region of the discharge. The \(\Delta x\)-signal was ELM-free up to about 13.5 s during which time the density and the radiated power rose characteristically. However, at this time an ‘event’ was triggered, the cause of which is not yet clear, although there was a break in the rate of change of the X-point gas flow at this time. This event raised the level of \(\Delta x\) slightly (possibly a new divertor state) but the level of the activity on this signal was about an order of magnitude lower than the usual ELM activity present in the ELMs neutral beam H-mode discharges. Concomitant with the start of this new phase, there is a decrease in density and radiated power. This permitted the H-mode to last longer for about 2.8 s albeit with a 20% reduction in energy confinement time but the particle confinement time was reduced by a much larger factor of about 3. This is of importance for the use of H-mode energy confinement in a reactor where simultaneously the particle and impurity species confinement has to be kept low.

The time traces of an H-mode with monopole are shown in Fig. 4. The monopole phasing aids [2] in producing monster sawtooth as can be seen from the electron temperature \((T_{\text{e0}})\) signal that remained sawtooth-free during the H-phase. This discharge was terminated by a disruption at 14.1 s. Notice that the density does not rise sharply and rolls over during the H-phase indicating a low particle confinement time. \(R_{\text{DD}}\) remains steady for about 0.8 s. In this H-mode discharge with monopole, the energy confinement time \(\tau_\text{E} \approx 1.7\tau_\text{c}\). At 2.8 T, the power threshold for transition to an H-mode with monopole is about 8 MW whereas it is about 5 MW with dipole. With beryllium antenna screen, the behaviour of radiated power from low and high-Z impurities in an H-mode produced with monopole is similar to that found in dipole.
2.4 Energy Confinement: A plot of stored energy \( W \) (from diamagnetic-loop measurement) plotted as a function of \( P_r - dW/dt \) for ICRH H-mode discharges at \( I_p = 3.1 \) MA using dipole and monopole phasing (with and without pellets) is shown in Fig. 4. Pellets in these discharges were not intended to produce peaked-density profiles but were used to make a comparison of H-modes discharges in open divertor configuration by gas and pellet fuelling in the edge region. Peaked density profiles pellet-enhanced plasmas (PEP) with ICRH H-modes are presented in another paper [6]. In Fig. 5, the lines drawn on the figure represent a multiple of the Goldston [5] L-mode prediction \( W_g \) for deuterium discharges where we have taken \( I_p = 3.1 \) MA, \( R = 3.1 \) m, \( a = 1.12 \) m and \( \kappa = 1.7 \) as representative values for these discharges. For definition of symbols see [5]. The values plotted were averaged over 0.2 s for the chosen time slices. For discharges with dipole and no pellets, \( 1.8 < W/W_c < 2.8 \) whereas with monopole and no pellets, \( W/W_c \approx 1.7 \) for a loss power of about 8 MW. Shots with pellets used for edge fuelling have a confinement slightly inferior to the ones mentioned above.

3. SUMMARY: H-modes produced by ICRH alone have been extended from dipole \((0, \pi)\) to monopole \((0,0)\) phasing by the use of beryllium screens on ICRH antennas in the JET tokamak. The attainment of H-modes with ICRH was facilitated by the implementation of a feedback control on the plasma position to maintain the coupling resistance constant which helps in matching and keeps the coupled power constant throughout the ICRH pulse. The energy confinement in the dipole phasing has been improved by about 30% compared to the Be-gettered nickel-screen by reducing the nickel-impurity radiation and best value of \( \tau_e \approx 2.8 \tau_0 \). The confinement in the monopole phasing is about 35% smaller than in the dipole phasing. This difference could be understood [7] by the effect of convective cells that exist in the monopole phasing and would decrease the shear in the poloidal rotation in the high confinement layer of the H-mode. ICRH H-modes are generally ELM free and are accompanied by long sawtooth-free periods. Longest duration of H-mode achieved is about 2.8 s. When deuterium gas was puffed into the X-point region, an H-phase of the discharge was found that had practically no ELM activity and only a 20% decrease in the energy confinement but was accompanied roughly by a factor of 3 decrease in the particle confinement time.

ACKNOWLEDGEMENT: We wish to thank our colleagues in the JET team, especially the RF plant team, the tokamak operation team and those operating the diagnostics used in the experiments reported in this paper. We thank P. Noll, A. Sibley and CODAS division for help in implementing the coupling resistance feedback control system.

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FIG. 2. Time traces of a high performance ICRH H-mode with dipole phasing.

FIG. 3. Time traces of an ICRH H-mode with dipole phasing with a gas puff into the top X-point. Note the low particle confinement (LPC) phase. \( < n_e > \) is the volume averaged density.

FIG. 4. An ICRH H-mode with monopole.

FIG. 5. \( W_{\text{Dia}} \) vs \( P_T \cdot dW/dt \) for ICRH H-mode discharges.
EDGE CURRENT DENSITY IN H-MODE DISCHARGES AT JET


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Introduction

A proper determination of MHD equilibria with the appropriate pressure and current density profiles is an essential first step in the analysis of stability, confinement and the plasma wall interaction of shaped tokamak plasmas. The evaluation of confinement and local transport in a shaped plasma depends on an accurate reconstruction of magnetic flux topology over the entire plasma cross-section. The stability of the plasma is a strong function of plasma profiles [1,2] and specifically the stability of the edge region depends on the local edge current density, and current density gradient [3]. The details of the topology of the magnetic flux surface, especially in the region of a magnetic separatrix (X-point), depend on the specifics of profile assumptions made in an equilibrium solver, and the details of the local topology have a profound effect on the local heat flux in the divertor region and the resultant impurity influx into the core plasma [4]. In this paper we report on the reconstruction of JET equilibria using measured pressure profile data. We include finite current density near the boundary and use the infinite domain equilibrium solver EFIT recently modified to include an iron core to calculate JET equilibria [5]. Comparisons are made with the edge current calculated from a local expansion near the X-point and with that obtained from detailed bootstrap calculations. Brief descriptions of the edge current calculations follow.

1. Bootstrap Current

In low collisionality, high β toroidal plasmas, Ohm's law is modified by neoclassical effects to include bootstrap current. To lowest order in the inverse aspect ratio the magnitude of the bootstrap current is proportional to the pressure gradient. In H-mode discharges, edge confinement improves and the pressure gradient near the boundary increases and may approach the first stability limit. In this case the magnitude of the bootstrap current is expected to increase at the boundary causing a broadening of the current density profile, and from simple considerations it can be shown that the maximum bootstrap current near the boundary is approximately

\[ J_{\text{SEP}} = J_{\text{BS}} \left( I_p / A \right) \leq \frac{0.3C_{BS}}{\sqrt{\epsilon} + 0.3C_{BS}} \]

where \( I_p \) is the total plasma current, \( A \) the area, \( \epsilon \) the inverse aspect ratio and \( C_{BS} \) a constant.
Detailed neoclassical calculations of the bootstrap current in the H-mode phase at JET, using experimental temperature and density profiles from Thomson Scattering and ECE measurements, have been carried out [6,7]. The calculation gives the surface averaged bootstrap current density on a poloidal flux surface and shows that $C_{BS}$ is often as high as 2.0, leading to a $J_{SEP N} \approx 50\%$.

2. Local Expansion Method

The X point of a divertor configuration is defined as a point where the poloidal field vanishes. We expand $\psi (R, Z)$ in a Taylor series around the X-point and determine the coefficients of the expansion by fitting to the magnetic field measurements in the vicinity of the X-point. The Grad-Shafranov equation then relates the expansion coefficients to the current density at the X-point. The angle made by the separatrix field lines is determined by the current density at the X-point and is no longer $90^\circ$ if the current density is finite. Since the surface averaged current density on the separatrix has a significant contribution from the X-point, the value of current density at the X-point gives a good approximation to this quantity.

3. MHD Equilibria with Finite Edge Current

The Grad-Shafranov equation is solved while approximately conserving the experimental measurements available. Some of the measurements which may be used in the equilibrium solution are: poloidal flux measurements, poloidal magnetic field measurements, total plasma current, diamagnetic flux, currents in the coil sets, pressure measurements, and internal measurements of the magnetic field (from motional Stark effect or Faraday rotation). The solution arises by minimizing

$$\chi^2 = \sum_{j=1}^{n_M} \left( \frac{M_j - C_j}{\sigma_j} \right)^2,$$

where $C_j$ and $\sigma_j$ are the computed values and the uncertainty of the corresponding measured quantities $M_j$, respectively, and $n_M$ is the total number of data used [8]. If the minimum of $\chi^2$ equals the difference between the number of independent measurements and the number of adjustable parameters, the fit is good and there is no justification for introducing more parameters. We consider a fit to be good if $\chi^2 \leq 60$. In order to determine the surface averaged edge current density, values of this quantity are used as input and the value which gives the minimum $\chi^2$ is taken as the edge current density.

4. Analysis of JET Data

We consider a double null high $\beta$ shot at a time during its H-mode phase. The most important parameters of the discharge are shown in Table 1. $\chi^2$ is given as a function of the normalized edge current in Fig. 1. We see that there is a minimum of $\chi^2$ at a finite value of $J_{SEP N}$ which corresponds to a surface averaged edge current of 0.23MA/m^2. We take $\chi^2$ to be statistically meaningful if
there is a significant drop in the local $\chi^2$ of measurements in the vicinity of the X-points. From Fig. 1 we see that this leads to values of $J_{\text{SEP}}$ with $0.3 < J_{\text{SEP}} < 0.75$, which is consistent with the rough estimation of bootstrap current given previously. The position of an X-point can be sensitive to the existence of edge current and in Fig. 2 we see the variation of X-point position with $J_{\text{SEP}}$. By artificially changing the X-point position in the local expansion technique and examining the $\chi^2$s obtained similar values of $J_{\text{SEP}}$ as found from the full equilibrium solution are obtained. A comparison of the three different methods to calculate edge current is given in Fig. 3, and we see that to within the experimental error there is good agreement. The goodness of fit to the kinetic pressure profiles is given in Fig. 4 where a $\chi^2$ of 4.0 was obtained.

5. Conclusions

Magnetic equilibrium reconstruction usually assumes vanishing current density at the plasma boundary. This assumption conflicts with the existence of steep pressure gradients at the edge during the H-mode phase of tokamak discharges, and with bootstrap calculations based on these pressure gradients which show that a significant amount of edge current may exist. These calculations however suffer from the large uncertainty in determining pressure gradients. We have developed a method of determining MHD equilibria for iron core tokamaks which is fully consistent with the available experimental data including soft X-ray determination of $q(0)$ and finite pressure gradient effects near the plasma boundary leading to finite edge current. This is a necessary first step for a consistent detailed analysis of the physics and stability of the plasma edge. The results obtained show good agreement with the bootstrap calculations, giving confidence in the results previously obtained.

References

Table 1. Major Plasma Parameters for JET Shot 21022 at $t = 51.23$ secs.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Major radius, $R_o$</td>
<td>3.06 m</td>
</tr>
<tr>
<td>Minor radius, $a$</td>
<td>1.08 m</td>
</tr>
<tr>
<td>Elongation, $\kappa$</td>
<td>1.82</td>
</tr>
<tr>
<td>Upper triangularity, $\delta_U$</td>
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<tr>
<td>Lower triangularity, $\delta_L$</td>
<td>0.45</td>
</tr>
<tr>
<td>Toroidal beta, $\bar{\beta}_T$</td>
<td>1.51%</td>
</tr>
<tr>
<td>Poloidal beta, $\bar{\beta}_P$</td>
<td>0.46</td>
</tr>
<tr>
<td>Internal inductance, $\ell_i$</td>
<td>1.36</td>
</tr>
<tr>
<td>$q_{95}$</td>
<td>3.6</td>
</tr>
<tr>
<td>$I_p$</td>
<td>3.11 MA</td>
</tr>
<tr>
<td>$B_T$</td>
<td>2.26 T</td>
</tr>
</tbody>
</table>
Evidence of edge current density from equilibrium Mag probes, Flux loops, coil current, Kin press fitted
21022.05123 q0=0.90, fixed
I_p=3MA, B_t=2.26T, double null divertor

Fig. 1

X-point position varies with edge current

Fig. 2

Edge current calculations

Fig. 3

Fil to kinetic pressure profile
21022.05123 q0=0.90, fixed

Fig. 4
POWER THRESHOLD FOR L-H MODE TRANSITION IN JET

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1. INTRODUCTION

The transition from L to H-mode regime of tokamak confinement during additional heating is generally observed when the power input exceeds a threshold $P_{th}$. The transition is marked by the sudden drop of the $D_0$ line emission, by the onset of edge plasma poloidal rotation $v_p$, by the enhanced radial electric field at the edge and by a suppression of the edge fluctuations. An hysteresis effect is also observed, namely the power required to sustain the H-mode can be lower than $P_{th}$. Therefore most of the theoretical models proposed to explain the transition rely on the identification of a bifurcation in some state variable describing the global equilibrium or the edge transport properties. The bifurcation appears only above a critical value of a control parameter which must be related to the power input.

A statistical analysis of JET H-mode data has been carried out in order to obtain a scaling for the power threshold for the L-H transition. Those thresholds have not only an operational relevance, but they can also help to discriminate between the different theoretical models. Since a variety of spurious effects can prevent the H-mode from taking place at powers higher than the ideal threshold, such as the MHD activity, the details of the wall conditioning, the impurity content, etc, we have included in our database pulses selected with the following criteria:

- power input roughly constant over 1 s (except for a limited number of cases with a power input waveform with multiple steps lasting ~ 0.5 s).
- total radiated power $P_{rad} < 50\% P_{tot}$.
- L-mode pulses with $Z_{eff} < 5$.
- H-mode pulses with no edge-localized modes (ELM), either ‘grassy’ or ‘giant’, for at least 0.5 s after the transition.

The discharges examined, ~500 in total, span the experimental campaigns from 1986 to 1990, with the toroidal field $B$ in the range 1.2 - 3.4 T, and plasma current in the range $I_p = 2.0 - 5.1$ MA. Plasmas limited by magnetic separatrix both in Single (SNX) and Double Null X-point (DNX), and in limiter configuration touching the inner wall have been considered, as well as both forward and reversed toroidal field (TF) direction. The additional heating can be with co-injected 80-140 kV $D^0$ Neutral Beam (NBI), with Ion Cyclotron Resonance Heating (ICRH) or both. Most of the discharges are in Deuterium; no sufficient data have been collected to study a possible mass dependence of $P_{th}$. The measured plasma parameters have been taken immediately before the transition (within 20 ms), or during a quasi-stationary phase of the energy and density evolution in the cases remaining in L-mode.

2. SCALING FOR NBI ONLY H-MODES

The main result is the increase of $P_{th}$ with the toroidal field $B$. The scaling $P_{th} = k B^{1.8}$ fits reasonably well all the available data, with the meaning of a boundary between the L and the H region. This has to be considered as a necessary, rather than sufficient condition, since some discharges remain in L-mode even if they are well inside the H-region.

The coefficient $k$ varies with the magnetic configuration, i.e. $P_{th}$ increases from upper SNX, forward field ($k = 1.5$), to DNX ($k = 1.75$), to upper SNX, reversed field ($k = 2.3$). No good H-modes have been obtained so far in lower SNX, reversed field configuration, due to the unoptimised Be protection tiles. H-modes with the transition ridden with ELMs have a somewhat lower $P_{th}$. Fig. 1 shows $P_{th}$ vs $B$ for DNX discharges at 3 MA.
In contrast with DIII-D results [1], it is also observed that \( P_m \) increases with the plasma current. At \( B = 2.1 \) T, \( P_m \) approximately doubles in going from 2 to 4 MA.

Profile effects are also associated with the transition, determining the equilibrium and stability properties of the H-mode. It is found that broader current profiles, as indicated by the internal inductance \( l_n \), and attributable to finite bootstrap or NBI driven current, have lower \( P_m \) (Fig. 2). On the other hand, no dependence on the average density has been found, again in contrast with DIII-D [1]. With the Beryllium gettered walls, the critical density threshold previously reported is not found anymore.

3. ICRH ONLY H-MODES

The introduction of several technical improvements have made possible to obtain H-modes with ICRH only [2]. In most of the discharges the resonance has been placed on axis. The magnetic configuration is always Double Null X point. A feature common to the ICRH only or combined heating discharges is the multiple-step nature of the transition, visible both from the \( D_\alpha \) signal and the density and energy evolution. The results reported refer to the first one of those steps. Generally at that time the ICRH power input is still in its rise phase.

It is found that \( P_m \) is quite different in the two possible phasings of the antennae (see Fig. 3). In monopole \( P_m = 8.2\) MW, in dipole \( P_m = 5.0\) MW. This effect could be related to the different electric field pattern in front of the antennae. In fact, a convective cell develops in the monopole case, which can affect the boundary flows and the transition itself.

Most of the ICRH only discharges were made at 2.9 T, so that a scaling with \( B \) as in the NBI case is unknown. It has to be noted, though, that \( P_m \) with NBI only at that field in DNX is ~8.5 MW. Combined heating discharges have still \( P_m = 5.0\) MW, which correspond to the \( P_m \) with NBI only at B ≤ 2.0 T.

4. DEPENDENCE ON PLASMA-WALL CLEARANCES

Various systematic scans in the distance between the X point and the vessel protection tiles have been performed. No dependence of \( P_m \) on the X point distance from the vessel protection tiles \( \delta_x \) has been found either in the upper single null, forward TF case and reversed TF (Fig. 4). The SN lower data is insufficient to draw any conclusion on this point.

On the other hand, the distance \( \delta_w \) of the plasma last closed surface from the inner wall has been found to be crucial. In NBI only case with forward TF, at \( \delta_w \geq 10 \) cm, \( P_{nbi} \) is as low as possible, but \( P_{nbi} \) goes up to \( \geq 11 \) MW at \( \delta_w \leq 2 \) cm. Indeed, H-modes obtained in limiter configuration with the plasma curvature matching the inner wall have higher \( P_m \) compared to DN at the same \( B \) [3].

5. COMPARISON WITH THEORETICAL MODELS

Although most of the models which describe the transition are cast in a form that involves edge quantities, some of those can be expressed in terms of the global power input, at least under various assumptions. The Hinton model [4] explains qualitatively the dependence on the direction of the ion flow relative to the X point location discussed in Par. 2 (see also [5]). A preliminary, qualitative agreement with the Hinton scaling \( P_m \propto n_e T_t \) (where the quantities on the r.h.s. must be taken at the edge), using Charge Exchange Spectroscopy (CXS) data, is found as well.

It has been observed in DIII-D that the L-H transition is associated with a spin up of the poloidal velocity \( v_p \) to values of tens of km/s [6]. Most of the more recent theories invoke a bifurcation in \( v_p \) or on the associated radial electric field at the edge, taking place at the transition. A model based on the coupled power and poloidal momentum balance written in nonclassical terms predicts directly the observed increase of \( P_m \) with \( B \) [7]. While the profile of toroidal velocity \( v_p \) is routinely measured at JET by CXS, spectroscopic measurements of \( v_p \) are very preliminary. Therefore indirect measurements of it have been pursued using the magnetic data.

The propagation velocity of \( B_\phi \) is \( v^* = v + v_p \), where \( v = v_{i\phi} + v_{e\phi} \) is the fluid velocity and \( v_p = \nabla p \times B \times \nabla B^2 \) is the diamagnetic electron drift velocity. Using the fact that the propagation is perpendicular to \( B \), and under the assumption of slowly varying pressure gradients \( \nabla p \), the variation of poloidal velocity is given in the cylindrical geometry by \( v_p = v_p [2 \pi v + n v_{e\phi} R] \), where \( r \) is the minor radius of the flux surface resonant at the poloidal mode number \( m \) and toroidal mode number \( n \), R is the major radius, \( v \) is the observed frequency of \( B_\phi \) with that elicty. Fig. 5 shows the time evolution of \( v_p \) obtained as just described, for a discharge with 3.1 T, 4.2 MA, upper SNX and 12.7 MW of NBI. The trace of \( D_\alpha \) highlights the spin up in the poloidal rotation immediately after the transition. Values obtained for the variation in \( v_p \), taking the parallel viscosity into account, are
in rough agreement with the preliminary spectroscopic results [8]. The time scale needed for the 
spin up is of the order of few ion-ion collision times, which is again in agreement with [7].

The Minardi model is based on totally different concepts, namely the vanishing production 
of a suitably defined entropy [9]. At low power the transition threshold for the additional power 
satisfies the relation \( P_{\text{add}} \propto E_{\text{th}} B_r \), where \( E_{\text{th}} \) is the ohmic axial electric field. A manageable way 
to estimate it is by means of \( P_{\text{th}} \propto P_{\text{meas}} \), plotted vs the experimental \( P_{\text{add}} \) in fig.6 (NBI only cases, 
all configurations). The separation of the L and H regions is reasonable.

ACKNOWLEDGMENTS

We wish to thank N Hawkes, J Jacquinot, C Lowry, E Minardi, and G Saibene for useful 
discussions, M Keilhacker, P Thomas and B Tubbing for useful remarks, M Bures for the validation 
of ICRH data and M Johnson for help in the setup of the database.

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conference.
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Fig. 1 - Total power input \( P_{\text{tot}} \) (MW) vs toroidal field \( B_r \) (T) for \( I_p = 3 \) MA discharges with 
Double Null X-point and forward toroidal field. Circles are L-remaining cases and crosses 
cases at the L-H transition. The dotted curve corresponds to \( P_{\text{tot}}(\text{MW}) = 1.75B_r^{-1.5}(T) \), the 
straight line to \( P_{\text{tot}}(\text{MW}) = 4.0B_r(T) - 3.0(\text{MW}) \).

Fig. 2 - \( P_{\text{tot}}(\text{MW}) \) vs internal inductance \( l_i \) for NBI heated only discharges with \( I_p = 3.0 \) MA, at \( B_r = 2.3 \) T and \( 3.3 \) T. + and x stand for L-modes, 
squares and triangles for L-H transition cases. The broken line is drawn to guide the eye only.
Fig. 3 - $P_{\text{tot}}$ (MW) vs plasma-ICRH antennae distance (cm) for ICRH only cases with $B_T=3.0\ T$, $I_T=3.0\ MA$ and Be gettered walls. Monopole and dipole phasing of the antennae are differently marked. L and H-modes are indicated as in fig. 2. The dotted line corresponds to $P_{\text{in}}=8.2\ MW$, the broken one to $P_{\text{in}}=5.0\ MW$.

Fig. 4 - $P_{\text{tot}}$ (MW) vs the upper X-point position $\delta_x$ relative to the tiles for NBI only, upper Single Null, forward field cases ($\delta_x<0$ means X point inside vessel). Open symbols for $B_T=2.3\ T$, full symbols for $B_T=3.0\ T$. L and H-modes are indicated as in fig. 1.

Fig. 5 - $D_e$ trace, frequency of $m=4, n=1$ mode from $B_N$ signal, $v_e$ at $r=3.9\ m$ from CXS and estimated variation of $v_e$ around the L-H transition (see text) for a 13 MW NBI heated discharge.

Fig. 6 - Additional power input $P_{\text{add}}$ (MW) vs $P$ (Minardi) $\propto kP_{\text{ohm}}q_{\text{eq}}$. L and H-modes are indicated as in fig. 1. L-H transition cases are found mainly above the straight dashed lines.
DOES THE ION CONFINEMENT IMPROVE IN ASDEX H-MODE DISCHARGES?

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1. Introduction
There is still uncertainty which transport channels actually are plugged after the L→H transition in the plasma core apart from the obvious transport barrier in the H-mode at the plasma edge. On JET the transport analysis indicates an improvement of the ion confinement alone /1/. On DIII-D the H-mode primarily exists at high densities with a close coupling between ions and electrons, but in the hot ion H-mode a reduction of both transport channels could be derived /2/. On ASDEX we have found that it is the electron transport which is reduced.

In order to explore the modifications of electron (κₐ) and ion (κᵢ) heat diffusivities between the L- and the H-mode in more detail, a comparative study between auxiliary heated discharges with peaked (L-mode with Ctr-NI) and broad (L-mode with Co-NI, H-mode with Co- and Ctr-NI) density profiles was carried out /3,4/. Transport analyses of L-mode discharges with the TRANSP transport code show that the ion thermal transport and the toroidal momentum diffusivity (κᵢ) - closely connected with it, as momentum is exclusively carried by the ions - can be consistently described by a combination of neoclassical and anomalous transport. The latter one is reduced with peaking ion density profiles (κᵢ = Lᵢ/ΔLᵢ < 1) and might be described by ion temperature gradient (ITG) driven transport. The origin of the anomalous electron transport remains still open.

Enhanced ion energy transport from ni-modes would, of course, also be expected to persist in the interior regions of H-regime discharges still having κₑᵢ > 1. And in fact our previous analyses of those discharges have always shown the necessity for an enhancement of κᵢ above the neoclassical ion heat conductivity κ(CH) as derived by F. Hinton et al (including impurities) to match the average ion temperature and the magnetically measured plasma energy. But the formally assumption κᵢ = ακ(CH) with α = 2-3 yields much more peaked Tᵢ-profiles than given by passive and active CX measurements. Having on hand reliable kinetic electron data from a 16 channel YAG-Thomson scattering system, Zₑᵢ measurements from IR bremsstrahlung to obtain improved ion density profiles and ion temperature measurements from active CX neutral energy analyzers (taking into account in gross form beam halo effects and neglecting high energy tail contributions near the plasma edge), further TRANSP analyses of well diagnosed H (with ELM’s) and quiescent H* (without ELM’s) discharges have been done to further address the above posed question.

2. Discharge characteristics
In this paper we present the results of a series of ten nearly identical discharges performed under boronized wall conditions and with a closed divertor (DV-II) at a plasma current Iₛ = 320 kA, a magnetic field Bₜ = 2 T and a line averaged density nₑ of about 3·10¹⁹ m⁻³ in the ohmic phase. Similar results are obtained over a large range of Iₛ (280 - 460 kA) and Bₜ (1.5 - 2.6 T). Fig.1 shows the time evolution of various plasma parameters during discharge # 33141. After beginning of NBI (H⁰ → D+ co-injection with 2 MW) first an L-mode develops (0.992 - 1.136 s), which is followed after a short H-mode phase by an H⁺-mode lasting about 120 ms (from 1.18 to 1.304 s) as can be seen from the H⁺-signal. Improved energy and particle confinement results in rising plasma energy, density and - due to a rise in impurity concentration- radiation and Zₑᵢ(0). Central ion and electron temperatures begin to decline. The energy replacement time τₑᵢ, which has fallen from 80 ms in the ohmic phase to 25 ms in the L-phase increases to 60 ms. The H⁺-mode is terminated by a sudden collapse of energy, nₑ and temperatures, while the central density, Zₑᵢ(0) as well as the to-
tal radiation still increase. The density and impurity peaking is accompanied by sequences of ELM’s, but the collapse is obviously caused by the resulting high central radiation power density of up to 1.5 MW/m³ exceeding the local heating power. The discharge falls back into the L-mode with \( \tau_E = 20 \) ms, recovers, and finally a H-mode with regular ELM’s follows lasting for 400 ms.

Here we concentrate on the time-dependent confinement analysis till the stationary second H-phase, and compare OH, L and H’ transport, the last one - although not stationary - allows confinement analyses without strong perturbations by the ELM’s.

3. Analysis of energy and momentum transport

Fig. 2 compares the measured radial profiles of \( T_i, T_e \) and \( n_e \) (4 channel DCN interferometer, broadband reflectometry) for the L and H’ phases. All three profiles are broader in the H-mode. For the analysis the shown profile fits and two additional ones displaying the reasonably largest and smallest gradients are used. We have gone more to the upper boundary of the temperature measurements, to match on the one hand the magnetically derived energies by the kinetic data within 5 % and, on the other hand, to obtain the lowest limits for the ion transport coefficients in the light of the posed question. The “effective” heat diffusivities \( \chi_i \) and \( \chi_e \) are defined by \( \chi_{i,e} = -q_{i,e} / (k n_{i,e} V T_{i,e}) \) with the conductive heat fluxes \( q_{i,e} \). These heat diffusivities are compared with the neoclassical ion heat conductivity \( \chi_{CH} \) including impurities \( (Z_i=8) \) and with the effective toroidal momentum diffusivity \( \chi_p \) (the toroidal plasma rotation velocity is measured by CX recombination spectroscopy). Combinations of \( T_i \) and \( T_e \) profiles leading to such large i-e energy exchange terms, that \( \chi_i < \chi_{CH}/2 \) or that even the ion heat flows against \( VT_i \) have been discarded.

Going from the L to the H phase \( \chi_e \) is strongly reduced over the whole plasma cross-section, whereas \( \chi_i \) and \( \chi_{CH} \) increase. The anomalous ion heat transport given by the difference \( \chi_i - \chi_{CH} \) is slightly reduced, however, as is the momentum transport \( \chi_p \). The results for the transport coefficients are further confirmed by the time evolution of the same quantities given at \( r = 2a/3 \) in Fig. 3. The strong \( \chi_i \) enhancement above \( \chi_{CH} \) at the beginning of the H’-mode declines later during this phase as does the momentum diffusivity. The momentum confinement time is close to \( \tau_E \) over the whole discharge and improves at the L to H transition therefore, but \( \chi_p \) is not at all reduced (the velocity profile flattens) in correspondence with \( \chi_i \). The anomalous ion and momentum transport are obviously linked together. In the following L-mode phase heat and momentum diffusivities rise again.

The often used ansatz \( \chi_{e} = \chi_{i} = \chi_{p} \) is apparently not appropriate. To demonstrate this further, an analysis with \( \chi_{p} = \chi_{i} \) has been done, yielding much more peaked \( T_i \) profiles (see Fig. 4) than measured (Fig. 3). Additionally \( \chi_i \) is one order of magnitude below its neoclassical value in the plasma centre, and a factor 2 - 5 even over the whole plasma in the H’ phase. As a consequence the time dependences of \( \chi_i \) and \( \chi_{p} \) do not agree at all.

Going from the L- to the H-mode not only \( \chi_i \) but also the ion heat conduction loss power increases, whereas reduced electron heat transport and convective losses are responsible for the confinement improvement (see Fig. 3). Energy convection can drive energy towards the plasma centre in the H’-phase with its very broad and, depending on the divertor and wall conditions - even slightly hollow density profiles.

4. Comparison with theory and conclusions

We can compare the above results for the anomalous transport with fluctuation-based anomalous transport theories. For the ion and momentum transport the connection with the ITG-modes has already been stated in the introduction. In Fig. 3 the \( \eta_i \) and \( \eta_e \) values at \( r = 2a/3 \) are compared with two threshold values for the ITG modes considering long density decay lengths in connection with toroidal effects according to F. Romanelli: \( \eta_i; \theta_1 = 1 + (1 + T_i/T_e)(L_{n_i}/R - 2) \), and with the magnetic shear according to P. Diamond et al: \( \eta_i; \theta_2 = 1 + 2 (1 + T_i/T_e)(L_{n_i}/L_e) \).

A remarkable clamping of the actual \( \eta_i \) values to the thresholds for the onset of the turbulence is found. This is also revealed by the radial profiles of \( \eta_i \) in the different discharge
phases. $\eta_1$ drops from high values in the centre (above 5 in the H*-mode) towards 1 near the boundary as do the $\eta_{1,th}$ values confirming the close relation of the anomalous ion transport with the $\eta_1$-modes. Further, within experimental error bars, our results indicate a similar radial and time behaviour of $\chi_p$ and $\chi_i \chi_{CH}$, but - as already stated for co- and ctr-injection discharges /4/ - $\chi_p$ is larger.

Regarding the electron heat transport the trapped electron driven turbulence predicts at least the increase of $\chi_e$ from the OH- to the L-phase due to the increasing $T_e$, whereas the decreasing $\chi_e$ in the H-mode is hard to explain with the broadening of the profiles alone ($\chi_e \sim T_e^{3/2}/(T_\infty L_ne L_T e)$). Resistive ballooning modes, which might be a possible candidate to explain the observed radial increase of $\chi_e$ towards the boundary not described by the TEM's /5/, have the wrong parameter dependence too.

Finally, this new analysis of the transport properties in the L and H* mode at ASDEX clearly demonstrates again, that the improvement in $\tau_E$ is not in the ion transport channel. Comparing with the results from other devices, we must conclude, that the observed differences are possibly caused by the parameter dependencies of the driving transport processes.

References

Fig.1: Time evolution of various plasma parameters in discharge #33141
(I_D = 320 kA, B_t = 2 T,
2 MW H^0 - D^+, single-null operation with $\Delta z = 1$ cm).
Fig. 2: Radial profiles of $T_e$, $T_i$ and $n_e$ and of diffusivities $\chi_e$, $\chi_i$ and $\chi_q$ and neoclassical ion thermal diffusivity $\chi_{CH}$ in the bulk of the L- and H'-mode plasmas.

Fig. 3: Time dependences of $n_e$, $\eta$ and two threshold values for ITG-modes ($\eta_{th1}$ and $\eta_{th2}$), of the heat conduction losses $q_e$ and $q_i$ and the convection loss $P_{conv}$ and of the diffusivities $\chi_e$, $\chi_i$ and $\chi_q$ and neoclassical ion thermal diffusivity $\chi_{CH}$ at $r = 2a/3$ of the H' discharges.

Fig. 4: Radial profiles of $T_e$ and $T_i$ (calculated under the assumption $\chi_e = \chi_i$) and of diffusivities $\chi_e$ (inferred) $= \chi_i$ and neoclassical ion thermal diffusivity $\chi_{CH}$ in L- and H'-mode plasmas.
LONG PULSE STATIONARY H-MODE WITH ELMS ON ASDEX

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1. Introduction
Extensive H-mode investigations have been made in ASDEX with boronised walls and reduced vacuum conductances between divertor and plasma chamber [1] [2] [3] [4]. Under these conditions also long quasi-steady-state H-mode discharges with optimised ELMs (Edge Localised Modes) have been achieved by careful radial and vertical positioning of the plasma column. An ELM-regulated discharge (H-ELM) is shown in Fig.1: $t_E=44$ ms; $I_p=0.28$ MA; $q_{95}=3$; $\bar{n}_e=3.1*10^{19}$ m$^{-3}$; ELM period =5.5 msec. The duration of the steady H-phase (3.5 s, ca. 80 confinement times) is only limited by the neutral beam pulse length. ELM controlled H-modes are obtained over a fairly wide range of plasma current (270 kA $\leq I_p \leq 440$ kA), injected power (1.6 MW $\leq P_{NI} \leq 2.7$ MW) and magnetic field (1.7 T $\leq B_t \leq 2.7$ T) both for Hydrogen and Deuterium injection into Deuterium plasmas. From the large number of ELMy discharges performed on ASDEX, only H-mode phases with distinct regular ELM events occurring at stable frequencies between 200 Hz and 500 Hz are analysed in this study. In the case of these regular ELMs with constant frequency the impact of the ELM-activity on the confinement can be estimated. Since for the ELM stabilised H-mode radiation losses ($P_{Rad} \leq 30\%$ of input power) can be neglected, the "intrinsic" H-mode confinement can be calculated in a simple way by the correction for the ELM losses, in contrast to the evaluation of the ELM-free H-mode ($H^*$) where corrections for radiation losses and the non-stationary energy content are essential.

In section 2 confinement scalings with and without ELM corrections are given and compared with existing confinement evaluations for $H^*$ and H-ELM discharges. A "subset" of the ELMy discharges where the confinement degrades irreversibly is discussed in section 3.

2. Confinement scaling
Energy loss by ELMs

Within the analysed parameter range no systematic dependence of the ELM-frequency on the injected power and the plasma current can be observed. The energy loss per ELM, however, increases with decreasing ELM-repetition rate. The relative energy losses per ELM have been calculated from slope changes of the plasma energy signal induced by the onset of ELM activity after short ELM free phases. The result is shown in Fig.2 for Hydrogen injection. For Deuterium injection a similar dependence can be measured, but the energy loss by the ELMs tends to be lower. In both cases the averaged energy loss by the ELMs ($\Delta E_{ELM \cdot f_{ELM}}$) increases slightly with the ELM-frequency.

$I_p$-Dependence

The variation of the confinement time with plasma current is shown in Fig.3 for Hydrogen beams. The data points corrected for the ELM losses show that the slight saturation for the higher currents can be attributed at least to some extent to the impact of ELMs causing a higher correction for larger $\tau_E$. Earlier data [5] for short pulses (duration of H-mode phase $\approx 100$ ms)
obtained with stainless steel walls of the plasma chamber are located within the gap to the ELM corrected data points. This may be explained by the fact that these short H-mode phases are not necessarily as stationary as the long ELMy discharges analysed in this paper.

The slope of the linear fit gives the "Quality-Factor" $\tau_I/I_p=120$ ms/kA without ELM correction. The "Quality-Factor" for the ELM corrected data points ($=190$ ms/kA) is in the range given for ELM free discharges [3]. Fig.4 gives the data for deuterium injection. The "Quality Factors" are 0.16 ms/kA and 0.22 ms/kA respectively. Again earlier data [6] for stainless steel plasma chamber walls fit between corrected and uncorrected H-ELM data.

No dependence on $q_a$ can be seen for the "Quality-Factors" in the covered range ($2.3 \leq q_a \leq 3.6$). Therefore the influence of the toroidal field must be weak.

Enhancement factors

Fig.5 shows the confinement times normalised to ITER89-P scaling for Deuterium injection. No definite current dependence can be seen. Included are data derived from ELM free discharges (H*) corrected for the non-steady-state condition ($dE/dt < 0.3 P_{abs}$) and central radiation losses [7]. The agreement between these data and the H-ELM data corrected for ELM losses is remarkable, since intrinsic H-mode confinement has been evaluated in completely different ways starting from two different discharges (H* and H-ELM).

Dependence on Heating Power

Because of the limited range in the injected power no meaningful fit for the dependence of the confinement time on the heating power is possible. A degradation of the confinement with increasing power however can be clearly seen for ELM corrected and uncorrected data.

3. Current profile effects

If the effectiveness of the ELMs in preventing impurities to increase in the plasma center is too low, the confinement degrades with time. This is demonstrated in Fig.6 which compares a stationary discharge (Fig.6a) and a discharge where the confinement degrades approximately 100 ms after the H-mode transition (Fig.6b). The discharge parameter settings are identical; the injected power is slightly different. During the stationary discharge $Z_{eff}$ in the center ($Z_{eff}(0)$), the $Z_{eff}$ profile from Bremsstrahlung and the radiated power do not change. The impurity behavior of the second discharge is different: $Z_{eff}(0)$ rises and the radiated power increases. With the peaking of $Z_{eff}$ the confinement (measured by the poloidal beta) starts to decrease to a lower level. The radiated power to the end of the neutral beam heated phase is only $\approx 20\%$ of the total input power, close to the initial level. The reduced confinement time can therefore not be attributed to higher radiation losses.

The influence of the $Z_{eff}$ profiles on the plasma current profile has been calculated by a current diffusion code [8]. This code uses measured $Z_{eff}$- and $T_e$-profiles to trace the development of the current profile. The result of the calculation is shown as $j(0)$ in Fig.6. The transient peaking of the $Z_{eff}$ profile results in a flattening of the initially peaked current profile. It is interesting to note that the current in the plasma center is reduced on the same time scale as the confinement of the discharge. This coincidence may indicate a close relation between the current profile and the H-mode confinement.

After the confinement degradation the $Z_{eff}$ profile flattens again. However, it has never been observed on ASDEX, that the confinement recovers from this initial degradation even if the following H-phase lasted for seconds. Obviously the transient rise of $Z_{eff}$ in the core triggers a broadening of the plasma current profile and enforces a transition into another equilibrium state. Under good confinement conditions, the equilibrium current profile is peaked. Figure 6 shows the peaking factor of the electron temperature profile $T_e(0)/<T_e>$. In the good case it remains large at about 1.9. In the degraded case the $T_e$ profile broadens irreversibly during the central
$Z_{\text{eff}}$ rise. But even after $Z_{\text{eff}}$ is restored at the original level the enhanced energy losses at the reduced confinement maintain the flat $T_e$ profile and the reduced central density.

4. Conclusion

Confinement scalings have been derived for ASDEX H-mode discharges stabilised by regular ELM of constant frequency. In Deuterium plasmas current scaling gives quality factors of $\tau_E/I_p=120$ ms/kA for Hydrogen injection and $\tau_E/I_p=150$ ms/kA for Deuterium injection. Correction for the energy loss by the regular ELMs leads to H*-mode confinement levels.

Irreversible confinement degradations which can be found in H-mode discharges with ELMs are related to calculated current profile changes.

References


[7] Ryter, F., priv. communication


Fig. 1 Quasi-steady state H-discharge (H-ELM)

Fig. 2 Energy loss by ELMs measured for different ELM frequencies
Fig. 3 Confinement versus plasma current for Hydrogen Injection into Deuterium plasma.

Fig. 4 Confinement versus plasma current for Deuterium injection into Deuterium plasma.

Fig. 5 Confinement time normalised to ITER89-P L-mode scaling.

Fig. 6a Steady-state H-discharge; Ip=320KA, qa=3, ne= 4*10^19 m^-3, PNI=1.9MW

Fig. 6b H-discharge with degrading confinement; parameters same as Fig.6a, PNI=2.4MW
ACHIEVING IMPROVED OHMIC CONFINEMENT VIA IMPURITY INJECTION

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1. Introduction

Improved Ohmic Confinement (IOC) was obtained in ASDEX after a modification of the divertors that allowed a larger (deuterium and impurity) backflow from the divertor chamber [Söldner et al., 1988]. The quality of IOC depended crucially on the wall conditions, i.e. IOC was best for uncovered stainless steel walls and vanished with boronization. Furthermore, IOC was found only in deuterium discharges. These circumstances led to the idea that IOC correlates with the content of light impurities in the plasma. To substantiate this working hypothesis, we present observations in deuterium discharges with boronized wall conditions into which various impurities have been injected with the aim to induce IOC conditions.

The remainder of this contribution is organized according to the following lines: Firstly, the plasma behaviour in typical IOC discharges is characterized. Secondly, injection experiments with the low-Z impurities nitrogen and neon as well as with the high-Z impurities argon and krypton are discussed. Then, we concentrate on optimized neon puffing that yields the best confinement results which are similar to IOC conditions. Finally, these results are compared with experiments in other tokamaks and some conclusions are drawn about the effects of the impurity puffing on both, the central and the edge plasma behaviour.

2. Plasma behaviour in the IOC regime

The IOC regime in ASDEX is characterized by a linear increase of the energy confinement time $\tau_E$ with increasing line-averaged density $\bar{n}_e$ up to the density limit. This differs from the usual behaviour of Saturated Ohmic Confinement (SOC) in ASDEX above $3 \times 10^{19}$ m$^{-3}$ and looks like a continuation of the Linear Ohmic Confinement (LOC) behaviour of low-density plasmas. The transition to the IOC is achieved with either uncoated or carbonized walls when the external fuelling rate is reduced. Simultaneously, the electron density at the separatrix drops which establishes favourable boundary conditions for the peaking of the radial density profile [Bessenrodt-Weberpals et al., 1990]. Moreover, the improved particle confinement is reflected in the particle flux into the SOL which drops by a factor of 2 [Verbeek et al., 1990].

It was striking that the quality of the IOC, measured by the increase of $\tau_E$ and closely correlated with the peaking of the radial density profile, degrades with improved wall conditioning. More precisely, the best results are obtained with uncoated SS walls, only minor improvements result from carbonization, and it seems not possible to approach IOC conditions after boronization of ASDEX. Recently [Bessenrodt-Weberpals et al., 1991], the trigger mechanism for the entrance into the IOC regime was investigated by means of a one-dimensional analytical model for the Scrape-Off Layer (SOL) plasma. This model solves the heat conductivity equation along field lines together with a simple form of the momentum balance [M. Ali Mahdavi et al., 1981]. The results reveal that the electron density and the electron temperature at the separatrix are determined by the power flux into the SOL and onto the target plates. If low-Z impurities are present, which radiate power from the edge of the plasma,
the power influx into the SOL is significantly reduced without detriment to central energy confinement. A reduced power load leads to lower edge densities; this seems to be at the heart of the gradual peaking process of the radial density profile, which is characteristic of IOC.

3. Studies with puffing of different impurities and flux optimization

With the knowledge of the close link between improved confinement and wall conditioning and with the results of the SOL model in mind, we performed a first series of experiments where various impurities are injected into deuterium discharges. These are tailored to include three density plateaus at $n_e = 2.7, 4.0$ and $5.4 \times 10^{19} \text{ m}^{-3}$, i.e. LOC conditions as well as SOC or – if possible – IOC conditions (see figure 1(a)).

![Graphs showing electron density, Ohmic input, SOL load, and main plasma radiation over time.](image)

**Fig. 1:** (a) Temporal behaviour of the line-averaged electron density. — Ohmic power, main plasma radiation and hence SOL load for (b) a typical IOC discharge with uncoated SS walls, (c) a pure deuterium discharge with boronized wall conditions, and (d) optimized neon puffing into a deuterium discharge with boronized wall conditions.

Figure 1 shows also the corresponding Ohmic heating power and radiation loss in the main plasma for (b) a typical IOC discharge with uncoated SS walls and (c) a pure deuterium discharge with boronized wall conditions. Obviously, the radiation loss overcome 50 per cent of the Ohmic heating power in the IOC discharge, whereas roughly 20 per cent is the value for pure deuterium cases. As a consequence, the power load into the SOL is much higher in this case than in the IOC discharge. To enhance
the radiation level, impurities are injected into the plasma; the impurity puffing is limited to the density ramps for the rare gases neon \((Z = 10)\), argon \((Z = 18)\), and krypton \((Z = 36)\); nitrogen \((Z = 7)\) as a non-recycling gas was added during the whole discharge.

![Diagram](image.png)

**Fig. 2:** (a) Power influx into the SOL for pure deuterium discharges and discharges with additional impurities at \(n_e = 2.7, 4.0, 5.4 \times 10^{19} \text{ m}^{-3}\). For comparison, the time-integrated target heat load \((kJ)\) is indicated by dots. — (b) Energy confinement time for the same discharges as in (a).

As indicated in figure 2(a), the total power into the SOL is actually reduced during all three density plateaus with impurity injection as compared to two discharges in pure deuterium, which have been performed before and after the impurity puffing experiments. This finding is corroborated by calorimetric observations which measure the divertor heat load integrated over the whole discharge; they are included in figure 2(a) as additional dots. All values reveal a minimum power load in the case of neon puffing due to the high radiation power in the peripheral plasma for low-Z elements. According to the model results given above, the reduced power load with impurity injection lets the separatrix pressure drop. In fact, the separatrix density is decreasing. Consecutively, the radial density profile in the central plasma is peaking. Such peaked density profiles are believed to establish favourable boundary conditions for the suppression of ion temperature gradient (ITG) driven modes, since the decay lengths of density and temperature are reduced such that the onset of ITG turbulence is hindered. Altogether, the energy confinement time is improved as shown in figure 2(b). Here again, neon puffing yields the best results.

To optimize the neon contribution to the deuterium plasma, we then proceed with a second series of discharges where the concentration and the puffing intervals are varied. Finally, conditions with strong neon puffing between 0.5-0.8 s and 1.5-1.9 s have been achieved which fairly resemble IOC data with uncoated SS walls. This comprehends many aspects. With respect to the power, both improved scenarios are characterized by a rather low target load (see figures 1(b) and (d)). Due to this low energy flux, the triggering of lower separatrix densities and a radially peaked density profile follow. Simultaneously, the electron as well as the neutral density in the divertor chamber are reduced like in the IOC discharge. In addition, the flux of neutrons reaches \(5 \times 10^{10} \text{ neutrons s}^{-1}\) for the neon puffed discharge, which is similar to the IOC value and is an order of magnitude higher than the flux \(3 \times 10^9 \text{ neutrons s}^{-1}\) observed in the pure deuterium case. Putting all pieces together, the confinement times for both particles and energy prolong.
4. Discussion

The success of achieving IOC conditions in ASDEX with optimized neon puffing under otherwise clean machine conditions reminds one of similar experiments in other tokamaks where also improved confinement was established by impurity injection.

This comprehends former studies of Ohmic discharges in PLT [Meservey et al., 1984] and the B-mode in T-10 [Alikaev et al., 1988] where neon puffing yields longer electron energy confinement times than regular discharges. Whereas the PLT authors ascribe their findings to an increase in $Z_{eff}$ and hence to higher Ohmic heating power, the ASDEX results show only a moderate change of the Ohmic heating power but an elevated radiation level. In this context, the impurity puffing experiments during neutral beam injection into ISX-B are also interesting [Lazarus et al., 1985]. The authors start from a degraded confinement after chromium gettering, i.e. clean wall conditioning, and enhance both the particle and the energy confinement if they add neon. This is called the Z-mode and is interpreted as a consequence of a cold plasma mantle built up by radiative cooling. Impurity injection has also been applied to H-mode discharges in D III-D [Luxon et al., 1990] in order to control the divertor target heat flux. In fact, deuterium and nitrogen puffing into the divertor region have been tested; whereas $\tau_e$ decreased only by 5 - 15 per cent, the target heat load is reduced by a factor more than 2.

5. Summary

To investigate the physical processes which happen during Improved Ohmic Confinement (IOC) in ASDEX, impurity injection experiments have been performed. Discharges with optimized neon puffing give the best energy confinement times. The results show that the low power flux into the SOL is the key to reduce the electron density at the separatrix. This drop seems to be the trigger of a change in the main plasma transport and hence leads to the peaking of the density profiles in the main plasma. In fact, the access to IOC can now be understood as being determined by the balance between the heating power and the radiated power loss. Thereby, the impurity content and the recycling conditions play a prominent rôle.

References


I - Introduction

In order to interpret theoretically the H mode it is of crucial importance to further study the fast phenomena occurring at the plasma edge and also to understand the link between edge and bulk phenomena. Several studies showed the relevant role played by the plasma density in the H mode physics /1/, /2/.

We have studied the plasma density and its fluctuations during the H phase in ASDEX with an O~mode broadband microwave reflectometry diagnostic. Combined broadband and fixed frequency measurements, and the use of a novel data analysis technique applied to the broadband data /3/, enabled the detailed study of the density profile development and of the temporal evolution of density fluctuations from the Ohmic to the H phase.

II - Density Profiles and Fluctuations

Plasma density profiles were obtained from the scrape off layer until close to the plasma center (0.4 - 4.5*10^{13}cm^{-3}, 43>r>8cm) and fluctuations were localized in the measured profiles, by studying the disturbed phase shift of the broadband reflectometric signals /3/. Figure 1 shows the development of the plasma density during the H mode along with the fluctuations level in the L and H phases.

OH and L phases: In the ohmic phase fluctuations are present both in the edge plasma (between 35 and 40cm) and in the central region (close to and within the q=1 rational surface). In the intermediate zone (16<r<35cm), MHD modes with low level of density perturbation could be detected close to the expected locations of q=3/2 and q=2. In the course of the L phase, the fluctuations increase; no dramatic changes are observed at the OH-L transition.

L-H transition: At the onset of the H phase, the level of edge fluctuations decreases as fast as the H_{t} radiation drops within the divertor chamber, to levels well below those observed in the OH phase. The edge density gradient steepens; density shoulders were detected ~0.4ms after the L-H transition. In Fig. 1(a) a density shoulder is shown 1.5 ms after the transition. At the edge, a region with reduced fluctuations is formed that extends radially beyond the steep gradient zone. However, within this region two narrow zones of coherent fluctuations remain, spaced 1-2cm apart, and probably located inside and outside the separatrix. The location of the fluctuation zones does not significantly vary during the H-phase, regardless of the great changes in plasma density; however, the width of the zone greatly varies, attaining a maximum shortly after the L-H transition. No dependence of these phenomena on the plasma conditions (B\_t, l\_p, q_a~3.3 is kept constant) was found.

After the rapid formation of the edge shoulder at the L-H transition, the interior plasma profile (r<35 cm) flattens along a much slower time scale, while fluctuations in that region increase to the levels observed in the OH phase. In H regimes with an ELMy phase the flattening of the interior plasma is initiated after the last ELM; in H regimes without ELMs or with a single ELM the flattening occurs some 30-40ms after the L-H transition. The observation of a fast time scale for the density rise at the edge and a slow one for changes in the bulk plasma, suggests that the improvement of the bulk confinement is rather a consequence of the changes in the edge conditions, whose effects slowly propagate inward.
**ELMy Phase:** Fluctuations at the edge in many cases increase to the levels observed in the L-phase, and coherent MHD modes (namely around the surfaces \( q = 2 \) and \( q = 3/2 \)) that had been suppressed after the L-H transition often reappear.

**H⁻ Quiescent Phase:** A large increase of density occurs during the quiescent phase; contrary to expectation, however, large fluctuations can occur inside the edge density shoulder. At the onset of these fluctuations the edge \( T_e \) is reduced and decreases thereafter until the end of the H-phase.

**H⁻ L Back Transition:** Prior to the H⁻ L back transition, fluctuations increase and recover the pattern of the L-phase. A further increase of both incoherent density fluctuations and MHD activity is observed after the back transition. However, in the flat part of the profile a region with a very low level of fluctuations remains for some time (>10ms) after the transition, suggesting that particle confinement is still improved in the interior plasma. This is in agreement with the observed slow time scale for density changes in the interior of the plasma.

**III - MHD Activity and Density Fluctuations**

MHD activity and turbulent fluctuations (in the range 0-150 kHz), were further studied with fixed frequency reflectometry, providing radial wave number integrated measurements \((0 < k < 25cm^{-1})\), with good spatial resolution (<3cm) and high sensitivity (<1.5mm). Three reflectometers were used to probe simultaneously density layers, that are located: (i) in the scrape off layer; (ii) close to and within the separatrix; (iii) in the central region.

1 - Density Fluctuations

An H mode discharge was studied where after the L-H transition an ELMy phase with very frequent ELMy occurred, followed by a quiescent phase (#33144). Edge \((n_{ec} = 0.5*10^{13}cm^{-3})\) fluctuations were measured in the frequency range 0-140kHz. In the ohmic phase, fluctuations exhibit frequency components up to 100kHz. After the NI pulse they transiently decrease (preferably at the highest frequencies) but increase afterwards during the L phase back to the levels observed in OH phase; higher frequency components (up to 140kHz) are then detected. At the L-H transition fluctuations are strongly reduced for frequencies >30kHz. During the ELMy phase typical broadband turbulence is seen (see section 2.1); in between ELMs fluctuations drop abruptly. After the last ELM, coinciding with the appearance of a coherent \( m = 1 \) satellite mode, and with the flattening of the interior plasma (see section II), fluctuations reappear. In contrast to the turbulence observed in the L phase, however, the highest frequency components \((f>80kHz)\) are suppressed. The observed correlations suggest that these high frequency components might play an important role in determining the differences concerning plasma anomalous transport in the OH, L, and H-phases.

2 - MHD activity

2.1 - ELMs

Figure 2 shows the frequency distribution of density fluctuations at the end of the quiescent H-phase, and the \( H_{\perp} \) trace at the divertor, for the previous discharge (#33144). The main features of the ELMs, as observed by the magnetic diagnostics /4/, can be found in the density measurements, of the edge plasma \((n_{ec} = 0.5*10^{13}cm^{-3})\). A precursor appears ~800\( \mu \)s before the first ELM, with frequency \( f \sim 100kHz \). During the ELM, broadband turbulence is seen up to the highest measured frequency (figure 2). This broadband turbulence can also be observed in the inner plasma layers \((n_{ec} = 3 \cdot 4*10^{13}cm^{-3})\), suggesting that ELMs might even affect the core of the plasma.
2.2 - m = 1 activity

The frequency distribution of edge (n_{\text{ec}} = 0.5 \times 10^{13} \text{cm}^{-3}) density perturbations (due mainly to MHD activity), detected during an NI pulse is shown in Fig. 3. The Hα trace, and the frequencies of m=1 satellite modes in the SOL, as measured with Mirnov coils are also presented.

Prior to the L-H transition a high amplitude coherent mode is detected with f ~ 7kHz, that is also seen before the H-L back transition. From profile measurements it can be concluded that this mode is located close and inside the separatrix. These modes have also been detected in L regimes by Mirnov coils /5/.

At the L-H transition fluctuations drop drastically and a so called "m = 1 satellite" mode, /6/, is detected. Its frequency increases from ~14kHz to 24kHz; identical frequencies are observed by Mirnov coils (Fig. 3(c)), and by the FIR scattering diagnostic. The frequencies agree with the central m = 1 activity as seen both by reflectometry and soft X-rays diagnostic, and follow the increase of central toroidal velocity. The m = 1 activity and the coupled mode are suppressed after the first sawtooth observed in the quiescent H phase and are destabilized shortly before the second sawtooth. The frequencies then decrease as a corollary to the decrease of the toroidal plasma rotation; these modes are suppressed after the second sawtooth.

About 40ms after the transition high amplitude narrowband fluctuations are detected around 45kHz, both inside the separatrix and in the interior plasma; these modes are also seen by FIR scattering and by Soft X-ray diagnostics. The density perturbations propagate to the inner plasma layers along with the flattening of the central plasma and the increase of the bulk density, (see section II).

IV - Conclusions

Several aspects of the H regime were studied with an O-mode broadband reflectometry system in the ASDEX plasma. Steep edge density gradient were measured after the L-H transition (<400\mu s). Flattening of the interior plasma occurs along a slower time scale after the ELMy phase or some delay time after the transition (30-40ms). The edge barrier is observed to be formed after the transition. The zone of reduced fluctuations extends radially beyond the steep gradient region and it shrinks when changes in the interior plasma are initiated. These are accompanied by MHD activity and increased fluctuations. High frequency components (>100kHz) of turbulent fluctuations increase from the OH to the L phases and are suppressed during the H phase. They are present, however, for short time intervals during the ELMs. These facts suggest that high frequency components of turbulence might play an important role in determining the plasma transport. The experimental findings showed the strong coupling existing between the edge and central plasma phenomena, that seems to be provided by the m = 1, n = 1 central MHD activity and by the ELMs. These phenomena are correlated with significant changes of the turbulent fluctuations and the density profile.

References

/3/ M.E.Manso et al., IPP - Report, to be published.
/4/ H.Zohm et al., this conference.
Figure Captions

Fig. 1: (a) Development of density profile from the L-phase into a quiescent H-phase; (b) Fluctuation level in the L-phase and H-phase at two different time points.

Fig. 2: (a) Contour plot of the signal power spectrum for the time window 1.35 - 1.365s, showing the precursor and the ELMs; (b) Hα in the divertor.

Fig. 3: (a) Contour plot of frequency distribution during L and H phases; (b) The Hα in the divertor; (c) the mode frequency from magnetic measurements (I_p = 280kA, B_t = 1.75T).
RANDOM COEFFICIENT H MODE SCALINGS

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Abstract: The random coefficient two stage regression algorithm\(^1\) with the collisional Maxwell Vlasov constraint \((C.M.V.)\) is tested and applied to the I.T.E.R. H mode confinement database\(^4\). The resulting scaling is similar to ITER89P with a slightly stronger size dependence.

I. Data Specification: We use the MHD confinement time for JFT2-M, D3D, PDX, and PBX-M, the diamagnetic \(\tau_E\) for JET and the geometric average of the diamagnetic \(\tau_E\) and the MHD equivalent \(\tau_E\) for ASDEX. We find no significant evidence of confinement differences between elmy and elmnfree discharges and therefore combine the elmy and elmnfree discharges.

We assume that the isotope enhancement factor is \(M^{1/2}\). For PDX, the isotope scaling is roughly linear, \(M^{1.0}\), which may be due to poor beam penetration. We therefore exclude all PDX \(H \rightarrow D\) discharges. We force the JFT2-M discharges to scale as \(M^{1/2}\).

We accept all standard constraints of the H mode database group and the divertor pressure ratio constraint for PDX. We require \(\tilde{n}\tau_E/\tilde{\tau} < 0.4\) and restrict to discharges with \(q_{95} \leq 6.5\). To constrain the data, we define an equivalent Ohmic power, \(P_{\text{ohm}}^* = I_p \times 1\text{V volt} \). We restrict the relative Ohmic power by \(P_{\text{ohm}}^*/(P_{\text{abs}} + P_{\text{ohm}}) < 0.4\).

A number of influential outliers strongly impact the \(B_t\) scaling in PDX and ASDEX and were removed. To examine the \(B_t\) dependencies, we regressed the individual tokamaks versus, \(I_p\), \(\tilde{n}\) and \(P\) and then plotted the residuals, \(y_i - \tilde{y}_i(I_p, \tilde{n}, P)\), versus \(B_t\). The residual errors, \(y_i - \tilde{y}_i(I_p, B_t, \tilde{n}, P)\), often depend on the relative time change of the energy, i.e. nonstationary discharges have better confinement.

Thus these restrictions result in a more uniform dataset which better characterises normal H mode discharges (see Table 1,2). Our scalings are closer to a \(L\) mode type scaling than the dataset of \([4]\). Table 3 summarises how each of our constraints has reduced the number of discharges.

The \(B_t\) scaling is poorly determined as a within tokamak covariate in the present database. The RMSE of the \(B_t\) exponent is much broader than the half width predicted by OLS regression. This indicates that at least one of the assumptions of OLS regression, such as the correctness of the model or the independence of the errors, is violated.

The ITER H mode database consists of a large tokamak, JET, a medium tokamak, D3D and four small tokamaks, ASDEX, PDX, PBX, and JFT2-M. Thus the H mode database is more statistically unbalanced than the \(L\) mode database. JET has a lower power density and beta value. It is unclear, if the scalings on JET are systematic differences from smaller devices or are manifestations of tokamak to tokamak differences.

Ordinary least squares regression of our 823 datapoint dataset yields

\[
\tau_E \sim M^{1/2}(R/a)^{-0.217} R^{2.113} \kappa^{-0.379} I_p^{0.20} B_t^{0.611} \tilde{n}^{-0.060} P^{-0.510}.
\]
The $B_t$ scaling of $B_t^{51}$ even exceeds the $B_t^{20}$ exponent of JET. This occurs because the within tokamak $B_t$ scalings are poorly determined and the RMSE in fitting the corrected mean confinement time of the tokamaks to the between tokamak covariates, $R$, $R/a$, and $\kappa$ can be reduced by including $B_t$ as a basically between tokamak covariate. Determining the standard between tokamak dependencies of $R$, $R/a$, $\kappa$ and the overall constant is already poorly conditioned for six tokamaks. Treating $B_t$ as an additional between tokamak covariate is clearly ill posed in the present database.

II. RANDOM COEFFICIENT SCALING: To model this tokamak to tokamak variation, we treat the scaling differences between devices as random variables. The random coefficient (RC) model is applicable when the tokamak to tokamak differences are due to many small factors. The RC within tokamak scalings are the matrix weighted average of the scalings of the individual tokamaks. Thus the $B_t$ scaling will not exceed the maximum scaling observed in any tokamak. We now summarise our random coefficient analysis.

First, for each tokamak, a scaling and covariance is estimated in $I_p$, $B_t$, $\bar{n}$ and $P$. We calculate the empirical mean and covariance of these within tokamak scalings using the Swamyr random coefficient weighting procedure. Second, the mean confinement time of each tokamak is corrected for the within tokamak scalings. The scalings with $R/a$, $\kappa$ and $R$ are estimated by regressing the corrected mean energy times of the tokamaks. The error, $\Sigma_{RC}$, in our estimate, $\hat{\beta}_{RC}$, of the scaling vector is given in [1]. The collisional Maxwell Vlasov (C.M.V.) similarity is then tested and imposed.

We exclude the $B_t$ principal component in JFT2-M, PDX and PBX-M. We also exclude the $I_p$ principal component in PBX-M. We include all the D3D principal components, however the $I_p$, $B_t$, and $\bar{n}$ scalings are strongly coupled. We apply the projection missing value algorithm. Unfortunately, the uncertainty in the $B_t$ scaling will be systematically underestimated since we are unable to compensate for the fewer degrees of freedom in the $B_t$ direction.

In the second stage regression, to determine the $R$, $R/a$ and $\kappa$ scalings, we weight the larger tokamaks, D3D and JET, a factor of two larger than the smaller tokamaks. In the second stage regression, we apply ridge regression with a relative ridge parameter of $\beta_{norm} = 0.005$, a half percent downweighting. This downweighting affects the $R$ exponent predominately.

The resulting scaling resembles the ITER89P scaling except for a stronger size scaling. A principal components analysis of the random coefficients matrix, $\Sigma_{RC}$, reveals that the size scaling is the most poorly determined exponent. The large variance in the $R$ exponent is a consequence of the database only containing one large tokamak.

The H mode data has a much larger component which violates collisional Maxwell Vlasov similarity than the L mode data. The hypothesis that the data can be explained by a dimensionless power law scaling can be rejected with slightly more than a 75% certainty. This conclusion is based on our extremely crude but self consistent modeling of the R.C. covariance. This may be caused by the presence of other hidden variables.

Nevertheless, we determine a C.M.V. constrained scaling within the R.C. model:
\[
\tau_E = 0.07904 M^{1/2} (R/a)^{-2.33} R^{1.877} \kappa^{3.17} I_p^{0.889} B_t^{2.07} \bar{n}^{1.105} P^{-0.486}
\]

(1).

\( L_p \) is in MAmps, \( B_t \) in Teslas, \( P \) in MWatts, and \( \bar{n} \) in \( 10^{19}/m^3 \). We have determined the dimensionless scaling which is closest to the unconstrained R.C. scaling measured in the \( \Xi^{-1}_{RC} \) metric. Since the dominant uncertainty occurs in the \( R \) exponent, our dimensionally constrained scaling differs from the unconstrained scaling by a weaker size scaling. Since \( R \) and \( R/a \) are strongly anticorrelated, the aspect ratio scaling decreases as well.

The aspect ratio scaling uncertainty is relatively small due to the presence of PBX-M. The larger tokamaks, D3D and JET, have small aspect ratios. Thus the negative aspect ratio scaling indicates that D3D and JET have better confinement than a \( (R/a)^0 \) dependence would indicate. If a \( (R/a)^0 \) dependence were required in our scaling, an even stronger size dependence would result.

In accepting the C.M.V. constraint, we not only set the dimensional projection equal to zero, but also eliminate the R.C. variance in the dimensional direction from our uncertainty estimates. The collisional Maxwell Vlasov constraint couples the size and \( B_t \) scalings, reduces the size uncertainty and thus the I.T.E.R. uncertainty more than the B.P.X. uncertainty.

For I.T.E.R., \( (a = 2.15m, R = 6.0m, \kappa = 2.2, I_p = 22MA, B_t = 4.85T, \bar{n} = 14.0 \times 10^{19}, P_{tot} = 160MW) \), the predicted confinement time is 4.65 sec with a statistical uncertainty factor of 32%. For B.P.X., \( (a = 0.8m, R = 2.59m, \kappa = 2.2, I_p = 11.8MA, B_t = 9.0T, \bar{n} = 40 \times 10^{19}, P_{tot} = 80MW) \), the predicted confinement time of 1.22 sec with a statistical uncertainty factor of 26 %. The unaccounted for uncertainties are discussed in Ref. [2].

We close on an optimistic note! When we weight D3D and JET more heavily than in Eq. 1, the resulting scaling yields higher predicted confinement times for both I.T.E.R. and B.P.X..

Acknowledgments: The author thanks Geoff Cordey and the H mode group for compiling the H mode database. The author thanks C. Bolton, J. DeBoo, R. Goldston, O.J.W.F. Kardaun, S.M. Kaye, and D. Post for many useful discussions. Curt Bolton suggested the use of \( P_{\text{chem}}^* \). This work was supported by U.S. Dept. of Energy Grant No. DE-FG02-86ER53223.

REFERENCES

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<th>$B_t$</th>
<th>$\bar{n}$</th>
<th>$P_{tot}$</th>
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<td>4.26±1.00</td>
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Table 1: Database Summary of Selected Discharges

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<th>$P_{tot}$</th>
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<td>PBX-M</td>
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<td>ASDEX</td>
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Table 2a: Individual Tokamak Scalings: Article’s restrictions

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Table 2b: Individual Tokamak Scalings: H mode database group restrictions only

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Table 3: Breakdown of imposed constraints by tokamak
EDGE PLASMA BEHAVIOUR IN OHMIC H-MODE
AND EDGE POLARIZATION IN TUMAN-3


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H-mode was observed in the experiments with different kinds of auxiliary heating [1,2,3,4]. In some cases H-mode was found in ohmic discharges. In DIII-D it was obtained in divertor configuration [5] and in CCT and TEXTOR it was induced by externally imposed radial electric field [6,7]. In our experiments the transition to H-mode was found in ohmically heated discharges in the circular limiter configuration [8,9]. Plasma parameters in TUMAN-3 Ohmic H-mode shots ranged as follows: \( R_o = 0.52-0.53 \, \text{m} \), \( a = 0.21-0.23 \, \text{m} \), \( B_T = 0.45-0.7 \, \text{T} \), \( I = 90-120 \, \text{kA} \), \( q_{\text{pol}}(a) = 2.0-4.0 \), \( \bar{n}_e = (0.6-2.5) \times 10^{19} \, \text{m}^{-3} \), \( T_e = 0.3-0.8 \, \text{keV} \), \( T_i = 0.11-0.17 \, \text{keV} \), \( Z_{\text{eff}} \approx 2 \).

The peripheral plasma properties in Ohmic H-mode were studied in order to compare them with that in the auxiliary heated H-mode plasmas. Investigation of the edge plasma could clarify the theory of the L-H transition and physical mechanism responsible for the confinement improvement [10,11,12]. It seemed to be useful to find the localization of transport reduction zone, to observe plasma turbulence evolution during transition and influence of the edge polarization on the confinement.

In this set of experiments plasma parameters were measured by UHF-interferometer (\( \lambda=2\,\text{mm} \), 10 vertical channels), multielectrode Langmuir probes, tunable UHF-reflectometer [13], \( D_\alpha \) detectors.

EDGE PLASMA PROPERTIES IN OHMIC H-MODE

Ohmic H-mode frequently appeared spontaneously and could be initiated by short increase of the puffing rate [8]. In our experiments we have found that the edge polarization by radial
electric field imposed with electrode also led to H-mode transition. In all these cases the sharp drop of D$_\alpha$ emission and simultaneous density increase were observed indicating H-mode appearance. Occasionally Ohmic H-mode was accompanied by D$_\alpha$ bursts, which could be interpreted as ELMs. In ELMy shots density increase rate was less than in ELM free cases, fig.1. in ELMy and ELM free shots

The most rapid and significant changes were found in the line averaged densities measured along the peripheral chords. Fig.2 shows density evolutions in interferometer channels looking through SOL region ($R_{\text{int}}$=2.3cm) and crossing Last Closed Magnetic Surface ($R_{\text{int}}$=-2.3cm). Density waveforms prove the formation of the steep gradient zone near the LCMS. Obtained using "abelization" procedure radial density profile is presented in fig.3. Langmuire probe measurements provided better space resolution and showed that after transition density reduced in a factor of 3-5, and

The e-folding length didn't change in SOL ($\lambda_{\text{OH}}^\text{SOL}$=1.2 cm.

the values calculated assuming $T_{\text{e}}=30$ eV ), fig.4a. Probe data confirmed formation of the steep density gradient in vicinity of the LCMS, fig.4b. Derived from these data $\lambda_{\text{LCMS}}^\text{OH}$ was equal to 3.8.
Fig. 4a. Density profiles in the SOL region, dashed line represents LCMS projection.

Fig. 4b. Ion saturation current for probe position 5 mm inside and 5 mm outside LCMS.

Fig. 5. Fluctuation driven particle flux in the LCMS vicinity (triple probe).

As it was found in our experiments H-mode transition was accompanied by suppression of the fluctuation-driven particle flux near the plasma boundary [9], fig. 5. The flux was measured using movable triple probe. The reduction of the
density and its fluctuations didn't lead to change of \( \tilde{n} / n_e \). It seems that suppression of \( \Gamma \) after transition was due to decorrelation between \( \tilde{n} \) and \( \tilde{E}_r \). Suppression of the high frequency fluctuations was registered also by reflectometer adjusted to boundary densities, \( f_{\text{UHF}} = 16.5 \text{ GHz} \) (\( n \approx 10^{12} \text{cm}^{-3} \)) [13]. Amplitude of fluctuation \( \langle f_{\tilde{n}} \rangle \) 200 kHz \( \sqrt{\langle U^2 \rangle} \) presented in fig.6. It should be noted that low frequency fluctuations at the boundary and fluctuations in core region didn't change their intensities after transition.

**EDGE POLARIZATION BY BIASED ELECTRODE**

Edge polarization was created by movable electrode with conducring head inserted inside LCMS from the low field side at \( r_e = a_{\text{LCMS}} - \Delta r \), where \( \Delta r = 3 - 4 \) cm. Potential was applied between electrode and vessel walls/limiters. We have found that radial conductivity was significantly lower in the positive potential case than in the negative one. H-mode was initiated using both polarities when voltage increased up to 500 V, which corresponded to radial field close to 150 V/cm. Observed H-mode had the same signatures as Ohmic H-mode. In a case of low densities the H-mode terminated and ordinary ohmic confinement restored after voltage was switched off.

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NOVEL FEATURES OF H-MODE PLASMAS INDUCED BY EDGE POLARIZATION IN TEXTOR

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Abstract: Whereas the energy confinement improvement in H-modes that are characterized by positive radial electric fields in the plasma edge is as good as that of the more usual negative field H-modes, the improvement in particle confinement of main plasma impurity ions is smaller in the former. Detailed measurements are given of the edge electric field, the electron density and temperature and the electrostatic fluctuation level. During the H-mode an instability can grow that strongly affects the plasma profiles without impairing H-mode behaviour.

INTRODUCTION: Earlier work [1,2] described how H-mode behaviour has been achieved in TEXTOR by using an electrode to induce strong radial electric fields Er and poloidal rotation velocities in the edge of ohmic and auxiliary heated plasmas. It was shown that H-modes can be set up irrespective whether the field points outwards or inwards. The latter situation will henceforth be called negative H-mode (or H_-mode) and the former labeled as positive (H_+mode), corresponding to their respective signs of the edge Er. As indicated before [2], H_-conditions are more difficult to establish in our experiment with the result that most of the previous work focussed on the H_+ behaviour. Recently we were able to considerably improve this situation and further our comparison of the performance of both modes. It should be recalled that the traditional H-mode [3] that develops "spontaneously" in many devices is of the H_-type [4].

Here, we first show the strong confinement improvement that can also be achieved in the H_-mode, thus supplementing our previous H_+ characterization [2]. In the second part we show that whereas the energy confinement of H_+ and H_- is comparable, strong differences are seen in the particle confinement of deuterium and of impurities. In the third part, modifications in the edge profiles of ne, Te and Er upon L to H transition are discussed.

The experiments have been performed under the following conditions: deuterium plasmas, Ip = 190 kA, Bt = 2.25 T, R = 1.75 m, a = 0.46 m and boronised walls. The insulated shaft of the biasing electrode [1] penetrates 6 cm beyond the belt limiter ALTII.

NEGATIVE H-MODE: The interrupted lines of Fig. 1 depict the time evolution for a beam-heated discharge (P Ni= 210 kW, from t =1.0 to 1.6 s) of n_e, H_α at limiter and wall and the kinetic energy in the electron component E_e. The beam-phase confinement has L-character with the electron confinement time τE_e decreasing with respect to the ohmic phase from 41 to 27 ms and the particle confinement time τp dropping in about the same proportion. When a 900 V negative electrode bias is superimposed, a clear H-mode develops yielding a significant gain in density and energy (full lines) with τE_e rising to 37 ms. The relative
change in $\tau_p$ is much stronger and reaches about 6. Note that before and after the NBI pulse the ohmic discharge is also in the H-mode.

**COMPARISON BETWEEN THE H$_+$ AND H REGIMES:** In Ref. 2, it was stated that whereas the energy confinement times of H$_+$ and H are comparable, strong differences exist in the deuterium particle confinement. We have continued such investigations and complemented them with impurity studies. To this end a neon gas puff was applied to discharges where stationary polarization conditions were set up, and the time decay of a central NeVIII line was monitored. An example is shown in Fig. 2 which compares two beam-heated discharges with either polarization, resulting in identical values of $\tau_e$. A detailed comparison of energy and particle confinement in these discharges reveals the following features:

(i) $\tau_e$ reaches 43.5 ms in the H$_+$ case (full lines) and 37 ms in the H one (interrupted lines).

(ii) $H_x$ was measured at the limiter, wall and biasing electrode. Changing the relative radial position of these elements allows to estimate their respective contribution to the total recycling flux as 0.55, 0.15 and 0.3. For the discharges shown, almost equal differences are measured at all three locations and yielding a 1.75 times higher total flux in the H$_+$ case.

(iii) The dwell time of the injected neon is clearly also higher in the H$_+$ discharge, as seen from the decay of the NeVIII brilliance traces in Fig. 2.

We can then conclude that the impurity confinement follows the trend previously established for the main ions. At roughly equal improvement in energy confinement, the H$_+$ mode shows significantly less particle confinement improvement. This result is thought to be of prime importance in the light of the recent debate on the suitability of H-mode conditions for the reactor [5].

The H$_+$ regime owes its superiority with respect to impurity behaviour to an apparent competition between (i) the effect of positive electric fields at the plasma edge capable of either expelling the impurities or preventing them from entering, and (ii) the global confinement improvement which is also felt by the impurities once the electric field reaches the conditions yielding L to H transition. This is illustrated in Fig. 3 showing the decay of the injected neon for different positive polarization voltages applied to an ohmic discharge. It is seen that a very strong decay compared to the ohmic discharge can be achieved with the application of $+400$V bias. Note that this voltage is below threshold for transition and the energy and density are identical to the ohmic ones. At $600$ V H-mode conditions are already established but the decay is still stronger than for the pure ohmic regime. At $900$V some mild retention of the neon is finally seen: comparison with the intermediate voltages yields that this condition is almost asymptotic for energy and density gains. For comparison the $-900$V case of Fig. 2 is reproduced.

**EDGE MEASUREMENTS:** Profile measurements were performed to ascertain the edge modifications upon the L to H transitions.

(i) Figure 4a shows the radial electric field (from floating potential measurements $V_f$) before and after transition to the H$_+$mode. An instantaneous increase ($\Delta t < 1$ms) of $E_r$ by a factor of 3-4 is observed. The $E_r$ values found here are significantly higher than those reported previously [2] and the localisation of the maximum field is even more close to the limiter. These effects are thought to be due to differences in edge recycling, as change-over experiments from carbonisation to boronisation have demonstrated. This might also explain the greater ease in achieving the H mode. Spectroscopical Doppler measurements confirm on the one hand the extreme spatial localisation of the ensuing plasma rotation, and on the other hand, by the measured speeds of up to 40 km/s, the observed values of $E_r = 1$V/cm.
FIG. 1: Time evolution of electron energy, $E_e$, $H_\alpha$, emission at limiter and wall, line density and beam power for shots with (full lines) and without (interrupted lines) bias.

FIG. 2: Time evolution of electron energy, bias voltage, beam power, neon brilliance, $H_\alpha$ emission at the wall and line density for discharges with positive (full lines) and negative (interrupted lines) bias.

FIG. 3: Time evolution of neon brilliance normalized to the line density for different bias conditions.

FIG. 4: (a) $E_r$ versus radius during flattop phase before (+400 V) and after transition (+900 V).

(b) Profile of RMS value of fluctuations under same conditions.
(ii) Density profile data were obtained by means of lithium-beam spectroscopy and Langmuir double probe measurements. Figures 5 shows the results in discharges where constant positive bias voltages were applied at 400, 600 and 900 V. The first voltage is below threshold for transition and the second just above. At 900 V at local flattening in the density occurs which is very pronounced in the probe data, which also reveal a strong peak in electron temperature at the same location (not shown here).

(iii) Figure 4b shows the variation with radius of $\langle V_f \rangle$, the RMS value of the floating potential fluctuations, defined as $N^{-1/2} \left( \sum_k \left( V_f (l_k) - \langle V_f \rangle \right)^2 \right)^{1/2}$, where N is the number of time points considered in the averaging process. The fluctuation level peaks in the high field region and increases roughly proportionally to $E_r$. Between 600 and 900 V, a strong increase is observed in the contribution to the fluctuation level coming from the frequency range 50-250 kHz. This high frequency component appears at the same time as the flattening in the density profile, thus suggesting a possible link. A plausible candidate for this instability could be the Kelvin-Helmholtz one [6]. It is remarkable that this strong fluctuation enhancement does not destroy the H-mode behaviour.

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IMPROVED CONFINEMENT REGIMES IN TEXTOR.


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INTRODUCTION: Earlier work [1,2] on TEXTOR has allowed the identification of an operational regime which is accessed with NBI-co and the combinations NBI-co+ICRH and NBI-co-counter. This regime has an energy confinement which is far in excess of the standard L-mode scaling that was the rule on TEXTOR before and still is for the pure NBI-counter and ICRH operations or their combinations [3,1]. The improved L-regime has recently been further investigated through parametric studies over a wide range of plasma currents and densities (200 < Ip < 500 kA, 1.5 < ne < 4 x 10^13 cm^-3). All results pertain to circular discharges (a = 0.46 m, R = 1.75 m) limited by the belt limiter ALT-II with boronized walls. Heating powers up to 2 x 1.8 MW of tangential NBI (co-counter, D^-+D^+) and up to 3.6 MW of ICRH (H/D < 0.1) have been applied.

COMPARISON OF TEXTOR IMPROVED CONFINEMENT REGIMES WITH DIVERTOR H-MODE SCALING. For comparison of the confinement quality with respect to other machines operated with H-mode discharges, we use the recent ITER H90-P scaling expression [4] obtained for ELM-free H-mode divertor discharges:

\[ \tau_{E,ELM-free} = 52 B_0^{0.15} \tau_e^{0.09} I_p^{0.27} P_{tot}^{-0.5} R^{1.82} a^{0.12} K^{0.35} A_i^{0.5} \]  
with the units ms(\tau), T(B), 10^3 cm^-3(\tau_e), MA(I_p), MW(P_{tot}), m(R,a). K is the elongation and A_i the ion atomic mass. For the stationary ELMY discharges the ITER team proposed to use \[ \tau_{E,ELMY} = 0.75 \tau_{E,ELM-free} \].

Our TEXTOR I-data base is subjected to the same selection criteria as those used for establishing Eq (1): the energy is the MHD (or equilibrium) energy E_{equi} and all shots must have

\[ \frac{P_{rad}}{P_{tot}} < 0.6, \quad \frac{dE_{equi}}{dt} \frac{1}{P_{tot}} < 0.3 \quad \text{and} \quad E_{fast}/E_{equi} < 0.4 \]

(P_{rad} is the radiated power, P_{tot} the total power coupled to the plasma and E_{fast} is the energy in the fast particles). In Fig. 1 we compare the experimental values of \( \tau_{equi} = E_{equi}/P_{tot} \) for the improved confinement regimes, with the predictions of \( \tau_{E,ELM-free} \) and \( \tau_{E,ELMY} \). The values of E_{equi} correspond to the maximum \( \dot{\beta} \) obtained during the heating phase and the values of P_{tot} are corrected for the charge exchange and shine-through losses of NBI. No attempt is made to correct P_{tot} for not optimized deposition profiles of NBI or ICRH [5]. The contribution of E_{fast} is obtained from the simultaneous measurement of E_{equi} of the
diamagnetic energy $E_{\text{dia}}$ and of the kinetic energy $E_{\text{kin}}$. $E_{\text{kin}}$ is obtained from $T_e(r)$, $n_e(r)$ and simulation by TRANSP [6] to fit the 2 points of the $T_i$ profile provided by CXES [1].

Figure 1 shows that the I-regime has a parametric dependence which is quite similar to $\tau_{E,\text{ELMY}}$ and $\tau_{E,\text{ELM-free}}$; the TEXTOR data lie on average between these two scalings with a derating of 0.87 with respect to $\tau_{E,\text{ELM-free}}$. Table 1 shows the corresponding derating for $E_{\text{dia}}$ and $E_{\text{kin}}$ at maximum $\beta$ and during long quasi steady state ($\Delta t \geq 20 \tau_e$). Note that to achieve steady state a derating of about 10 % has to be accepted and that under these conditions the thermal energy alone is still somewhat above L-mode.

<table>
<thead>
<tr>
<th></th>
<th>$E_{\text{equi}}$</th>
<th>$E_{\text{dia}}$</th>
<th>$E_{\text{kin}}$</th>
</tr>
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<tbody>
<tr>
<td>Maximum $\beta$</td>
<td>0.87</td>
<td>0.73</td>
<td></td>
</tr>
<tr>
<td>Stationary</td>
<td>0.79</td>
<td>0.69</td>
<td>0.55</td>
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**TABLE 1**: Derating factor with respect to $\tau_{E,\text{ELM-free}}$

**CHARACTERISTICS OF THE IMPROVED REGIME : CURRENT SCALING.** Figure 2 shows for two discharges (with successive NBI-co, NBI-co+counter and NBI-counter heated phases) the evolution of $E_{\text{dia}}$, $E_{\text{equi}}$ and $E_{\text{kin}}$. The central line densities $n_e$ are also given. Both discharges remain in the high confinement regime during the NBI-co phase but shot 2 has an abrupt transition during the co+counter phase to a lower confinement regime, found to scale like the usual L-mode ($= 0.5 \tau_{E,\text{ELM-free}}$; see also Fig. 4). The NBI-counter heated phase pertains to the L-mode as well. The oscillation appearing in discharge 1 during the combined heating phase is not due to sawtooth oscillations, as these are stabilized in this phase, but to a density profile modulation.

These NBI-co+counter heated discharges can be characterized by a strong profile peaking. The central ion temperature in these discharges varies from 2.5 to 4.5 keV depending upon the density. Figure 3 shows the electron density and temperature evolution for a shot with increasing $n_e$ during the combined phase. The increase of $n_e$ results in a rise of the shoulders of the $n_e$ profile leading to a large density gradient near the limiter radius, which is comparable to those observed for H-mode discharges for machines with a similar size and same $n_e$.

In Fig. 4 the black dots show the evolution of $E_{\text{equi}}$ at maximum $\beta$ for the improved regime (NBI-co+cou) at $2.5 < P_{\text{tot}}(\text{MW}) < 3.1$ as a function of the plasma current $I_p$. A detailed analysis shows the existence of a clear energy limit at low $I_p$ corresponding to (curve b) $\beta_p = 1.5$ which is already reached with $P_{\text{tot}} = 1$ MW at $I_p = 200$ kA and 3 MW at 300 kA. At high $I_p$, $E_{\text{equi}}$ scales roughly as $0.87 \tau_{E,\text{ELM-free}} P_{\text{tot}}$ as indicated by curve a. The best stationary toroidal beta values hitherto obtained reach 0.75 %, about 2/3 of the Troyon limit for TEXTOR. The open dots pertain to L-mode discharges and comprise also those discharges reached after the I to L transition. They are 20 to 30 % higher than the Kaye-Goldston (KG) scaling [7], with its recently introduced correction for the ion mass dependence, and might have a somewhat stronger current scaling ($\sim I_p^{1.24}$). Note that the I-mode has an enhancement factor with respect to this KG scaling of 1.8 to 2.3.

**HIGH-TO-LOW TRANSITION.** Discharges with NBI-co+counter and NBI-co+ICRH operation can suddenly transit from high to low confinement. Whereas only a slight MHD precursor activity is observed around the $q = 1$ region, the L-mode shows a strongly enhanced MHD
activity, seen on ECE, HCN and Mirnov coils with \( m = 1 \) and 2 components. The peaked density profile occurring in the I-mode can persist in the low mode but its shape can strongly oscillate. The frequency of occurrence of a I to L transition increases with decreasing \( n_e \) and \( I_p \). There is no strong difference in the MHD activity between the high and low NBI-co discharges: their difference is thought to be linked to recycling [1].

**CONCLUSIONS:** 1) Besides the usual L-mode regime which prevails in many heating scenarios, we have identified an improved regime with a confinement at least as good as that of ELMY H-mode divertor discharges. The I-mode can be produced with NBI-co, NBI-co+ICRH and NBI-co-counter. 2) The latter case shows peaked density profiles quite reminiscent of those of the supershot regime[8]. 3) The confinement in the L- and I-modes improves with \( I_p \). 4) Transitions from I to L have been observed.

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![Fig. 1 Comparison of \( \tau_{equi} \) for the improved I-regime of TEXTOR (with NBI-co, NBI-co+ICRH or NBI-co-counter) with the ITER H90-P scaling.](image)
Fig. 2. Time evolution of $E_{\text{equi}}$, $E_{\text{dia}}$, $E_{\text{kin}}$, $\bar{n}_e$ and NBI power for two NBI-co-counter heated shots: 1/ for a shot in the I-regime (#44565) and 2/ for a shot with a I to L transition (#44566).

Fig. 3. a) Electron density (expressed in $10^{13}\text{cm}^{-3}$) and b) temperature profiles for a NBI-co-counter heated TEXTOR discharge (#42085).

Fig. 4. $E_{\text{equi}}$ versus $I_p$ for the improved I-mode (closed symbols) and L-mode regimes (open symbols) with $2.5 < P_{\text{tot}} < 3.1 \text{ MW}$ and $1.6 < \bar{n}_e < 3.1 \times 10^{13}\text{cm}^{-3}$. Loci of predictions of the ITER H90-P and modified Kaye-Goldston scaling laws together with $\beta_p = 1.5$ are given.
CANONICAL PROFILES TRANSPORT MODEL FOR IMPROVED CONFINEMENT REGIMES IN TOKAMAKS

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ABSTRACT. The transport model of canonical profiles (TMCP) was developed previously for the description of L-mode in tokamaks [1]. Here we extended this model to the improved confinement regimes. The analysis of experimental data from JET and JT-60 permits us to describe naturally the transition to H-regime and in hot ion and/or hot electron modes.

I. THE MODEL. Recently [1] we considered the TMCP for L-mode in tokamaks. The comparison of calculations with the experiments on T-10, JET, ASDEX and TFTR made it possible to construct the expressions for electron and ion heat fluxes without free parameters. The model [1] was based on the idea of heat fluxes increasing with the deviations of the electron and ion temperature profiles $T_k(r)$ from canonical ones $T_{kc}(r)$ ($k' = e, i$). We substitute the fluxes in the 1-D system of transport equations for $T_k$ and poloidal field. The density profile $n(r)$ is assumed to be known from the experiment. Here we extended the model [1] to the improved confinement regimes: H-mode, hot ion mode and regimes with the peaked density profiles.

The analysis of experiments led us to the conclusion that the zones of the canonical profiles action are bounded. At strong deviations of the real profiles from the canonical ones the heat fluxes are determined by the electron and ion pressure profiles $p_k(r)$: if the deviations of $p_k(r)$ from the canonical one $p_{kc}(r)$ exceed over some critical value, then plasma forgets about the canonical profile and the electron and/or ion heat fluxes very decreased.

Let us suppose

$$\frac{p_{kc}(r)}{p_{kc}(0)} = (1 + r^2/a_j^2)^{-2},$$

$$\frac{T_{kc}(r)}{T_{kc}(0)} = (1 + r^2/a_j^2)^{-\gamma_k}, \quad (k = i, e). \quad (1)$$

Here $a_j = a/\sqrt{q_a} - 1$ is a current radius, $\gamma_k \sim 1$ are the parameters to be determined from the comparison of the calculations with the experiment.

Let us take the heat fluxes as

$$\Gamma_k = \Gamma_k^{con} + \Gamma_k^{an} + \Gamma_k^{PC}, \quad (k = i, e), \quad (2)$$

Here $\Gamma_k^{con}$, $\Gamma_k^{an}$ and $\Gamma_k^{PC}$ denote the contributions from the classical, anomalous and particle confinement effects, respectively.
where \( \Gamma_{\text{con}}^{\text{an}} = \frac{5}{2} r n_{\text{k}} T_{\text{k}} \), \( \Gamma^{\text{a}} \) is the particle flux that known from the experiment,

\[
\Gamma_{k}^{\text{an}} = -n \chi_{k}^{\text{an}} (\partial T_{k}/\partial r), \quad \chi_{k}^{\text{an}} = \text{const} = \alpha_{k}^{\text{an}} / n.
\]

The \( \Gamma_{k}^{\text{PC}} \) designates the heat fluxes associated with deviations from the canonical profiles taking into account the forgetting effect:

\[
\Gamma_{k}^{\text{PC}} = -n \chi_{k}^{\text{PC}} (r T_{k} / a^2) z_{T_{k}} F(z_{pk}). \quad (3)
\]

Here [1]:

\[
\chi_{k}^{\text{PC}} = \frac{\kappa_{k}^{\text{PC}}}{n} = \alpha_{k}^{\text{PC}} \left( \frac{a}{R} \right)^{0.75} q(a/2) q(a) \frac{\sqrt{T_{k} (a/4)}}{n}.
\]

\( B \) is a toroidal magnetic field (T), \( T_{k} \) (keV), \( \chi_{k} \) (m²/s),

\[
z_{T_{k}} = \frac{a^2}{r} \frac{\partial}{\partial r} \ln (T_{k}/T_{k_{c}}) \quad \text{and} \quad z_{pk} = \frac{a^2}{r} \frac{\partial}{\partial r} \ln (p_{k}/p_{k_{c}}) \]

are dimensionless deviations \( T_{k} \) and \( p_{k} \) from \( T_{k_{c}} \) and \( p_{k_{c}} \). \( F = \exp(-z_{pk}^2/2 z_{ok}^2) \) is a forgetting factor, \( z_{ok}(r) \) is the width of the canonical profile action zone. The fluxes (2)-(3) contain some parameters. The comparison with the experiments in L-mode gives [1]: \( \chi_{l} = 2/3 + 1, \chi_{e} = 1 + 4/3, \alpha_{e} = 3.5, \alpha_{l} = 5 \). The other parameters: \( \alpha_{k}^{\text{an}}, z_{ok}(0), z_{ok}(a) \) are determined here by comparison of calculations with the experiments in improved confinement regimes.

Let us discuss the conditions of the hot mode occurrence. The flux (3) can be represented as:

\[
\Gamma_{k}^{\text{PC}} = -k_{k}^{\text{PC}} \frac{r T_{k}}{a^2} \varphi_{k}, \quad \varphi_{k} = \varphi_{k} (z_{T_{k}}) = z_{T_{k}} \exp \left( -\frac{(z_{T_{k}} + \frac{z_{nk}}{2})^2}{2 z_{ok}^2} \right).
\]

\[
z_{nk} = \frac{a^2}{r} \frac{\partial}{\partial r} \ln \frac{n_{k}}{n_{k_{c}}}, \quad n_{k_{c}} = p_{k_{c}} / T_{k_{c}}.
\]

Let the heating power \( Q_{k}^{\text{aux}}(r) \) to be localized in the plasma core. The function \( |\varphi_{k} (z_{T_{k}})| \) is limited, so the condition of the hot mode occurrence is

\[
Q_{k}^{\text{aux}}(0) > Q_{k}^{\text{cr}} = \frac{2}{\pi} \max \Gamma_{k}^{\text{PC}} = 1.6 \times 10^{-3} n_{k_{c}}^{\text{PC}} \frac{2 T_{k}}{a^2} \max |\varphi_{k}|.
\]

where \( Q_{k}^{\text{cr}} \) (MW/m³), \( \max |\varphi_{k}| = |\varphi_{k} (z_{k})|, \hat{z}_{k} = -\sqrt{z_{ok}^2 + z_{nk}^2 / 4} - z_{nk} / 2 \). When the density profile is rather flat, \( n(r) = n_{k_{c}}(r) \), we have \( z_{nk} = 0, \hat{z} = -z_{ok} \), \( \max |\varphi_{k}| = z_{ok} \exp(-1/2) \). When the density profile become more peaked, the
value $|z_{nk}|$ is increased too and $\max|\varphi_k|$ quickly decreased. In particular, if $z(0) = -z_{nk}$, $\max|\varphi_k| \approx 0.25z_{nk} \exp(-1/2)$, i.e. $Q_k^{cr}$ decreases in 4 times.

2. THE SIMULATION OF HOT ION MODE IN JET. The experimental data for 6 different regimes are presented in [2]. The comparison with the experiments gives $z_{o1}(0) \approx 1.5+2$, $z_{oe}(0) \approx 5$. Thus the transition in the hot electron mode needs either more peaked density profile, or more specific power input than for ions.

On the plasma edge the electrons and ions are strongly bounded by energy exchange term and we can't distinguish the dimensions of their canonical profiles action zones. Calculations show that $z_{oe} \approx z_{o1} \approx 4$, but the detail description of L-H transition is rather difficult because of absence of the experimental data about neutral flux from the wall.

The calculated (solid lines) and experimental (dashed line) temperature profiles for shot #18757 in L-regime with hot ions (10.5 MW NBI) are shown in Fig. 1. Here $\gamma_1 = \gamma_e = 1$, $z_{o1}(0) = 1.5$, $z_{oe}(0) = 5$. In the plasma core, $r/a < 0.3$, the forgetting of the ion temperature canonical profile occurred. Here the flux $\Gamma_e^{pe}$ is small and it determines by the flux $\Gamma_e^{an}$ with $\alpha_e^{an} = 0.6$.

The Fig. 2 shows the same profiles for shot #19987 in H-mode with hot ions (15 MW NBI). In the calculations $z_{o1}(a) = z_{oe}(a) = 3.5$. Here the forgetting for ions occurred in the core region, $r/a < 0.3$, and near the edge, $r/a > 0.9$. For the electrons it occurred only near the edge in the same zone as for ions. The $T_e$ value is determined here by the flux $\Gamma_e^{an}$ with $\alpha_e^{an} = 0.9$. The Figs. 3 and 4 show the $T_e$, $T_i$ and effective thermal diffusivity, $\chi_k^{eff} = -\Gamma_k/(\partial T_k/\partial r)$ profiles for shot #17749 in L-mode with hot ions and electrons that formed when the density is very peaked and after pellet injection (12 MW ICRH + 5 MW NBI) [3].

The simulation of hot ion mode on JT-60 [4] with founded values of parameters $\gamma_k$, $z_{nk}$ and $\alpha_k^{an}$ does not contradict to the experimental data.

REFERENCES

Fig. 1. Calculated (solid lines) and experimental (dashed line) temperature profiles in JET. Shot #18757, L-mode with hot ions (10.5 MW NBI).

Fig. 2. Calculated (solid lines) and experimental (dashed line) temperature profiles in JET. Shot #19987, H-mode with hot ions (15 MW NBI). The ion thermal diffusivity profiles are also shown.

Fig. 3. Calculated (solid lines) and experimental (dashed line) temperature profiles in JET. Shot #17749, L-mode with hot ions and hot electrons (5 MW NBI + 12 MW ICRH).

Fig. 4. Radial profiles of thermal diffusivity for shot #17749 from Fig. 3.
ON THE POSSIBILITY OF REACHING THE H-MODE FROM INITIAL CONDITIONS

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Recent experimental and theoretical studies have emphasized the role of the radial electric field and of the poloidal rotation of the tokamak plasmas in connection with the L to H-mode transition. It resulted that a fast, sheared poloidal rotation and a negative radial electric field determined this transition by suppressing the drift turbulence.

Starting from the analysis of the evolution equations for these two parameters, we have identified the possibility that a stable improved confinement state could be established by an appropriate preparation of the discharge in the initial phase. At this stage, the study is mainly qualitative and contains simplifying assumptions. The conclusions need to be verified by numerical computations integrated in a transport simulation.

The radial electric field is obtained from Ampère's law which expresses the balance between the plasma polarisation current and the current arising from the particle fluxes:

$$\varepsilon_0 \kappa_1 \frac{\partial E_r}{\partial t} = e \Gamma_e, r - e \Gamma_i, r + D_e \frac{1}{r} \frac{\partial^2 E_r}{\partial r^2}$$

(1)

where $\kappa_1 = \frac{1}{1 - \left(\frac{v_A}{v}\right)^2}$, $v$ is the Alfvén velocity, $\varepsilon_0$ is the vacuum permittivity and $D_e$ is the electric field diffusion coefficient $/1/$. The ion flux $\Gamma_i, r$ contains the nonambipolar components due to charge exchange processes and to the loss cone loss. The electron flux $\Gamma_e, r$ is the anomalous turbulent flux. According to various models $/2,3,4/$, the difference of these fluxes is a nonlinear function of the radial electric field. This equation can be simplified by making the change of variables: $x = r / \Gamma_e$ and $\tau = t / (\varepsilon_0 \kappa_1)$ which, due to the fact that $x \gg r$, allows to extend the space range to the real axis and to keep only the second derivative in the diffusion term:

$$\frac{\partial E_r}{\partial t} = e \Gamma_e, r - e \Gamma_i, r + \frac{\partial^2 E_r}{\partial x^2}$$

(2)

The characteristic time for the variation of $E_r$ resulting from this equation is:

$$\tau_E = \varepsilon_0 \kappa_1 / \left(\frac{\partial E_r}{\partial t} (e \Gamma_e, r - e \Gamma_i, r)\right)$$

(3)

It gives values of the order $10^{-10}$, much shorter than the time scale for the plasma parameters.

In order to study the qualitative behaviour of the solution of this equation for various initial and boundary conditions, we have used the method of traveling fronts which is based on the change of variable $z = x - ct$. The equation (2) becomes:

$$e_z = e_{zz} + c e_z + f(e)$$

(4)
where \( e(z,\tau) = E_z(z+c\tau,\tau) \) and subscripts mean partial derivatives. The limit \( e^\infty(z) \) of the solution \( e(z,\tau) \) corresponding to the initial condition \( e^0(x) \), as \( \tau \to \infty \) (physically, the radial electric field after a time of the order \( \tau_{\infty} \)) is shown to be the solution of the stationary equation:

\[
\phi_{zz} + c \phi_z + f(\phi) = 0
\]  

(5)

which satisfies the following constriction:
- if \( e^0(x) \) is a supersolution (i.e. \( e^0_{xx} + c e^0_x + f(e^0) \leq 0 \)), then \( e^\infty \) is the maximal stationary solution bellow \( e \);
- if \( e^0(x) \) is a subsolution (i.e. \( e^0_{xx} + c e^0_x + f(e^0) \geq 0 \)), then \( e^\infty \) is the minimal stationary solution above \( e \).

A qualitative picture of the set of stationary solutions is obtained by introducing the phase space with the coordinates \( (\phi, \psi) \) and by transforming Eq. (5) into the system:

\[
\begin{align*}
\frac{d\phi}{dz} &= \psi \\
\frac{d\psi}{dz} &= -c \psi - f(\phi)
\end{align*}
\]  

(6)

The topology of the phase space depends on the nonlinear term \( f(\phi) \) (i.e. on the nonambipolar flux) and more specifically on the number of its zeroes. These zeroes represent the equilibrium points of the system (6) \((d\phi/dz=0, d\psi/dz=0 \) ). The set of stationary solutions consists in these points (i.e. uniform radial electric field) and of the shifts of the separatrices corresponding to these points.

As an example, the phase space analysis for the case of a three zero flux as in the model /2/ shows that the following profiles for the asymptotic radial electric field are possible:
- flat profile with \( E = E_1 \);
- pulse shape profile with \( E(-\infty)=E_1, E(\infty)=E_1 \) and \( E \geq E_1 \);
- monotone profile with \( E(-\infty)=E_1, E(\infty)=E_2 \);
- oscillatory profile with \( E(-\infty)=E_1, E(\infty)=E_2 \) which is moving with a positive velocity \( c \);
- monotone decreasing profile with \( E(-\infty)=E, E(\infty)=E_1 \).

Here, \( E_1, E_2, E_3 \) are the zeroes of the function \( f \) and \( |E_2| < |E_1| < |E_3| \).

For each of these stationary solutions there is a domain in the space of functions \( Q_i \), with the property that any function of this domain when taken as initial condition for Eq. (2) evolves to the corresponding stationary solution \( i \). This means that modifying the initial condition for Eq. (2), the solution is stable with the exception of sudden changes (which appear in times of the order \( \tau_{\infty} \)) when passing from one domain \( Q_i \) to another \( Q_{i+1} \). The modification of the initial condition is performed on the long time scale due to the evolution of the poloidal velocity and of the plasma parameters. For a slowly varying radial ion flux, the following relation holds:

\[
E_r = \frac{1}{\frac{\rho}{\rho_0}} \frac{\partial \nu}{\partial r} + \nu \theta E_t
\]  

(7)

In the situations when this slow evolution determines the crossing of the boundary of the appropriate \( Q_i \), the L-H transition appears. However, it is possible in principle to maintain the slow time scale evolution inside the domain which corresponds to the improved confinement.
The evolution of the poloidal rotation velocity of the plasma is determined from the poloidal projection of the ion momentum balance equation:

\[ m_i n_i \frac{\partial v_{\theta}}{\partial t} = (ZeB_t) \Gamma_{i,r} + \frac{1}{r} \frac{\partial}{\partial r} \left( r \mu \frac{\partial v_{\theta}}{\partial r} \right) \]  

(8)

where \( \mu \) is the ion dynamic viscosity. Experimental and theoretical studies have shown that the viscosity coefficient is a nonlinear function of \( v_{\theta} \), having a maximum for a velocity \( v_{M} \). We have chosen an analytic expression reproducing the qualitative behaviour of \( \mu /5/ \):

\[ \mu = \mu_1 \exp(-\alpha(v_{\theta} - v_{M})^2) \]  

(9)

\[ v_{M} = v_{th,i} B_{\theta}/B_t \]

\[ \mu_1 = \frac{v_{th,i}^2}{p_i m_i n_i v_{ii} q} \]

where \( v_{th,i} \) is the ion thermal velocity, \( v_{ii} \) is the ion poloidal gyro-radius, \( v_{ii} \) is the ion-ion collision frequency and \( q \) is the safety factor.

The term containing the ion radial flux \( \Gamma_{i,r} \) acts as a driving force for the rotation since it is determined by independent processes. The absolute value of this driving force depends on the radial electric field which results from Eq.(2).

Eq.(8) being of parabolic type, the poloidal velocity evolves toward a stationary state in which the driving force is balanced by the viscous damping. However, due to the nonlinearity, the asymptotic state is strongly dependent on the initial condition and on the relative magnitude of the driving and damping forces.

Using the notation \( \lambda = ZeB_t \Gamma_{i,r} \) and \( M \), for the initial magnitude of the viscosity term as determined by \( v_{M}(B_t,0) = v_{0} \), the following cases were revealed by the numerical study of Eq.(8):

- 1. \( v_{0} > v_{M} \) and \( \lambda > M \)

The velocity grows and, according to (9), the viscosity decreases progressively. This produces a continuous increase of the poloidal rotation and of its shear which represents a state of improved confinement as demonstrated in /6/;

- 2. \( v_{0} > v_{M} \) and \( \lambda < M \)

The resulting decrease of \( v_{\theta} \) induces an even greater viscosity and a phase of fast decay of the rotation leading to an asymptotic value of \( v_{\theta} \) smaller than \( v_{M} \);

- 3. \( v_{0} < v_{M} \) and \( \lambda > M \)

The velocity increases and there are two possibilities for the further evolution: a) if \( \lambda < M \) (the maximum viscosity force), a state of asymptotic equilibrium with \( v_{\theta} = v_{M} \) will be established and b) if \( \lambda > M \), the poloidal velocity will become greater than \( v_{M} \) in a finite time and will continue to increase leading to an improved confinement;

- 4. \( v_{0} < v_{M} \) and \( \lambda < M \)

The rotation velocity decays at a rate which decreases in time and the velocity shear has small values.
Since $\lambda$ is usually small, the situation 1. is the most promising while the situation 3.b) is difficult to appear with negative $E_r$. We note that these results were obtained for a time independent driving force. When the dependence of the ion flux on the radial electric field is included, fast changes of $\lambda$ are possible (at a crossing of the boundary between two $\Omega$ 's) which could modify the characteristic type of the poloidal rotation evolution.

From this analysis results that one possibility of obtaining a stable improved confinement from early stages of the discharge consists in imposing a fast poloidal rotation of a narrow boundary plasma layer (satisfying the conditions of 1.). This velocity induces a negative radial electric field. At the beginning of the discharge when is not expected to have multiple solution for the total nonambipolar flux, the equation for the radial electric field (2) has an unique class of solutions (i.e. there is only one domain $\Omega$). Along the evolution of the nonambipolar flux, the domain of the initial profiles $\Omega$ splits into several subsets which lead to different asymptotic solutions. In the case of a three zero flux it resulted that there is a monotone decreasing profile with $E < E_0$, which is the attractor for all initial functions with the property $E(-\infty)=E$, and $E_1(\infty)<0$. This means that, the changes in the radial electric field produced by the evolution of the poloidal velocity do not determine jumps of the electric field, i.e., in these conditions, the poloidal rotation and the negative electric field support each other through the nonlinear flux.

Another possibility of obtaining improved confinement is to choose a boundary condition for the electric field at plasma edge (to bias the plasma) which determines stable asymptotic solutions for Eq.(2). This have to be a negative value or a positive value greater than $E_3$ (in this case, the radial electric field inside the plasma would equal $E_3$).

INDEX OF AUTHORS
<table>
<thead>
<tr>
<th>Author</th>
<th>Page</th>
<th>Author</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>Abramov, A.V.</td>
<td>II-93</td>
<td>Axon, K.B.</td>
<td>III-73</td>
</tr>
<tr>
<td>Abramov, V.A.</td>
<td>III-217</td>
<td>Azarenkov, N.A.</td>
<td>IV-105</td>
</tr>
<tr>
<td>Aceto, S.</td>
<td>II-169</td>
<td>Azevedo, M.T.</td>
<td>II-313</td>
</tr>
<tr>
<td>Aceto, S.C.</td>
<td>II-161</td>
<td>Azumi, M.</td>
<td>I-233</td>
</tr>
<tr>
<td>Adams, J.M.</td>
<td>I-21</td>
<td>Bachmann, P.</td>
<td>III-17</td>
</tr>
<tr>
<td></td>
<td>I-45</td>
<td>Baek, W.</td>
<td>III-341</td>
</tr>
<tr>
<td></td>
<td>IV-277</td>
<td>Baek, W.Y.</td>
<td>I-333</td>
</tr>
<tr>
<td></td>
<td>IV-281</td>
<td>Baelmans, T.</td>
<td>IV-369</td>
</tr>
<tr>
<td></td>
<td>III-389</td>
<td>Bagatin, M.</td>
<td>III-197</td>
</tr>
<tr>
<td>Afanasiev, V.I.</td>
<td>I-149</td>
<td>Bagdasarov, A.A.</td>
<td>III-69</td>
</tr>
<tr>
<td>Afanasjev, v.I.</td>
<td>I-209</td>
<td></td>
<td>II-1</td>
</tr>
<tr>
<td>Afanes'ev, V.I.</td>
<td>IV-349</td>
<td>Bakity, F.W.</td>
<td>III-361</td>
</tr>
<tr>
<td>Agostini, E.</td>
<td>I-61</td>
<td>Bakunin, O.G.</td>
<td>II-129</td>
</tr>
<tr>
<td>Airoldi, A.</td>
<td>I-157</td>
<td>Balet, B.</td>
<td>I-9</td>
</tr>
<tr>
<td>Akiyama, H.</td>
<td>I-333</td>
<td></td>
<td>I-37</td>
</tr>
<tr>
<td>Alava, M.J.</td>
<td>III-305</td>
<td></td>
<td>I-41</td>
</tr>
<tr>
<td>Alejaldre, C.</td>
<td>III-125</td>
<td></td>
<td>I-189</td>
</tr>
<tr>
<td>Alexander, K.F.</td>
<td>III-25</td>
<td>Bamford, R.</td>
<td>I-81</td>
</tr>
<tr>
<td>Alkaev, V.V.</td>
<td>III-361</td>
<td>Bank, S.L.</td>
<td>II-61</td>
</tr>
<tr>
<td>Alladio, F.</td>
<td>I-77</td>
<td>Barbato, E.</td>
<td>IV-217</td>
</tr>
<tr>
<td>Almagri, A.</td>
<td>II-117</td>
<td>Barbian, E.P.</td>
<td>III-417</td>
</tr>
<tr>
<td>Alvarez, A.</td>
<td>II-209</td>
<td>Barnsley, R.</td>
<td>IV-257</td>
</tr>
<tr>
<td>Anderson, D.</td>
<td>I-125</td>
<td></td>
<td>III-109</td>
</tr>
<tr>
<td>Ando, A.</td>
<td>IV-101</td>
<td>Barth, C.J.</td>
<td>I-121</td>
</tr>
<tr>
<td>Antoni, V.</td>
<td>I-137</td>
<td>Bartiromo, R.</td>
<td>I-73</td>
</tr>
<tr>
<td>Antonov, N.V.</td>
<td>III-69</td>
<td></td>
<td>III-405</td>
</tr>
<tr>
<td>Antsiferov, P.</td>
<td>II-221</td>
<td></td>
<td>III-417</td>
</tr>
<tr>
<td>Aranchuk, L.E.</td>
<td>III-209</td>
<td>Bartlett, D.</td>
<td>I-1</td>
</tr>
<tr>
<td>Arimoto, H.</td>
<td>I-129</td>
<td></td>
<td>I-13</td>
</tr>
<tr>
<td>Arsenev, A.V.</td>
<td>II-141</td>
<td>Bartlett, D.V.</td>
<td>I-45</td>
</tr>
<tr>
<td>Arsenin, V.V.</td>
<td>II-261</td>
<td></td>
<td>IV-357</td>
</tr>
<tr>
<td>Asakura, N.</td>
<td>II-293</td>
<td>Basilico, F.</td>
<td>IV-1</td>
</tr>
<tr>
<td>ASDEX- and NI-Team</td>
<td>I-97</td>
<td>Basu, J.</td>
<td>III-89</td>
</tr>
<tr>
<td></td>
<td>I-109</td>
<td>Batchelor, D.B.</td>
<td>I-113</td>
</tr>
<tr>
<td></td>
<td>I-385</td>
<td></td>
<td>III-309</td>
</tr>
<tr>
<td></td>
<td>IV-269</td>
<td>Bateman, G.</td>
<td>I-93</td>
</tr>
<tr>
<td></td>
<td>I-105</td>
<td></td>
<td>I-237</td>
</tr>
<tr>
<td></td>
<td>I-117</td>
<td>Batenyuk, A.A.</td>
<td>II-237</td>
</tr>
<tr>
<td></td>
<td>I-217</td>
<td>Batistoni, P.</td>
<td>II-117</td>
</tr>
<tr>
<td></td>
<td>I-221</td>
<td>Baty, H.</td>
<td>II-81</td>
</tr>
<tr>
<td></td>
<td>I-249</td>
<td>Bayley, J.M.</td>
<td>II-325</td>
</tr>
<tr>
<td></td>
<td>I-389</td>
<td>Baylor, L.R.</td>
<td>II-165</td>
</tr>
<tr>
<td></td>
<td>IV-305</td>
<td>Beckstead, J.</td>
<td>II-289</td>
</tr>
<tr>
<td></td>
<td>III-117</td>
<td>Behn, R.</td>
<td>II-37</td>
</tr>
<tr>
<td>Askinasi, L.G.</td>
<td>I-149</td>
<td></td>
<td>IV-289</td>
</tr>
<tr>
<td>Assadi, S.</td>
<td>I-401</td>
<td>Behrisch, R.</td>
<td>III-153</td>
</tr>
<tr>
<td>Astapkovitch, A.M.</td>
<td>II-85</td>
<td>Beidler, C.</td>
<td>II-149</td>
</tr>
<tr>
<td>Auge, N.</td>
<td>IV-353</td>
<td></td>
<td>II-181</td>
</tr>
<tr>
<td>Aumayr, F.</td>
<td>IV-365</td>
<td>Bell, J.D.</td>
<td>II-177</td>
</tr>
<tr>
<td>Austin, M.E.</td>
<td>I-325</td>
<td>Bell, H.G.</td>
<td>III-141</td>
</tr>
<tr>
<td>Name</td>
<td>Page</td>
<td></td>
<td></td>
</tr>
<tr>
<td>------------------------</td>
<td>------</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fahrbach, H.U.</td>
<td>I-221</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fall, T.</td>
<td>IV-297</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fang, Z.S.</td>
<td>III-349</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Paulconer, D.W.</td>
<td>II-13</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Feix, M.</td>
<td>III-297</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Feneberg, W.</td>
<td>IV-153</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Feng, Y.</td>
<td>IV-141</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Ferreira da Cruz, D.</td>
<td>II-121</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Ferrer Roca, Ch.</td>
<td>II-305</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Ferron, J.R.</td>
<td>II-105</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fessey, J.</td>
<td>IV-73</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fessey, N.</td>
<td>I-369</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fielding, S.J.</td>
<td>III-73</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Field, A.R.</td>
<td>III-113</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fijalkow, E.</td>
<td>IV-153</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Filippas, A.V.</td>
<td>I-309</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fishman, H.</td>
<td>I-133</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fishpool, G.</td>
<td>I-41</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fitzpatrick, R.</td>
<td>II-21</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>II-61</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>IV-77</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fletcher, J.D.</td>
<td>I-261</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>I-293</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fogaccia, G.</td>
<td>IV-65</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fonck, R.J.</td>
<td>I-269</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>II-77</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fowler, R.H.</td>
<td>II-157</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fraguas, A.L.</td>
<td>II-125</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Franz, D.</td>
<td>II-221</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fredrickson, E.</td>
<td>I-265</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>II-5</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Freeman, R.</td>
<td>III-369</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Froissard, P.</td>
<td>III-389</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>III-393</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Fuchs, G.</td>
<td>II-17</td>
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