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Dubrovnik, May 16—20, 1988

Editors: S. Pešić, J. Jacquinot

Contributed Papers

Part I

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Contributed Papers
Part I
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Preface

The 15th European Conference on Controlled Fusion and Plasma Heating was organized by the Boris Kidrič Institute of Nuclear Sciences, Belgrade, Yugoslavia, on behalf of the Plasma Physics Division of the European Physical Society (EPS). It was held in Cavtat, near Dubrovnik, from 16 to 22 May 1988.

The 15th Conference concentrates on experimental and theoretical aspects of Plasma Confinement and Plasma Heating. The programme, format and schedule of the meeting was determined by the International Programme Committee, which was appointed by the Plasma Physics Division Board of the EPS. The Programme Committee selected 19 invited and 24 oral contributed papers.

This volume contains all accepted contributed papers received in due time by the Organizers. It is published in the Europhysics Conference Abstracts Series and it follows the rules for publication of the EPS. The 4–page extended abstracts were reproduced photographically using the camera ready manuscripts submitted by the authors who are therefore responsible for the quality of the presentation. Post-deadline papers are not included in this volume.

All invited papers will be published in a special issue of the journal "Plasma Physics and Controlled Fusion". This journal may also publish rapidly an extended version of contributed papers, following an accelerated refereeing procedure.

The Organizers would like to acknowledge the support of the Serbian Science Research Council and the technical assistance of the Publishing Department of the Boris Kidrič Institute of Nuclear Sciences. The Conference has been organized under the general sponsorship of the Union of Yugoslav Societies of Mathematicians, Physicists and Astronomers, member of EPS.

April 1988

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## Contents

### Paper Identification

<table>
<thead>
<tr>
<th>Part I - A. TOKAMAKS</th>
<th>Part II - B. STELLARATORS</th>
</tr>
</thead>
<tbody>
<tr>
<td>A1. Experiments</td>
<td>C. ALTERNATIVE CONFINEMENT SCHEMES</td>
</tr>
<tr>
<td>A2. H Mode</td>
<td>C1. Reversed Field Pinch</td>
</tr>
<tr>
<td>A3. Theory</td>
<td>C2. Other Alternative Magnetic Confinement Schemes</td>
</tr>
<tr>
<td>A4. Sawteeth, Disruptions and Other Related MHD Phenomena</td>
<td></td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>Part III</th>
</tr>
</thead>
<tbody>
<tr>
<td>E3. Lower Hybrid Heating</td>
</tr>
<tr>
<td>E4. Alfvén Wave and Other RF Heating Methods</td>
</tr>
<tr>
<td>F. CURRENT DRIVE AND PROFILE CONTROL</td>
</tr>
<tr>
<td>F1. Lower Hybrid Current Drive</td>
</tr>
<tr>
<td>F2. Other Profile Control Methods</td>
</tr>
<tr>
<td>G. NEUTRAL BEAM INJECTION HEATING</td>
</tr>
<tr>
<td>H. DIAGNOSTICS</td>
</tr>
<tr>
<td>I. BASIC COLLISIONLESS PLASMA PHYSICS</td>
</tr>
<tr>
<td>J. INERTIAL CONFINEMENT PHYSICS</td>
</tr>
</tbody>
</table>

### First Author Index

<table>
<thead>
<tr>
<th>Full Author Index</th>
</tr>
</thead>
<tbody>
<tr>
<td>VI</td>
</tr>
<tr>
<td>VII</td>
</tr>
<tr>
<td>I-1</td>
</tr>
<tr>
<td>I-3</td>
</tr>
<tr>
<td>I-207</td>
</tr>
<tr>
<td>I-255</td>
</tr>
<tr>
<td>I-330</td>
</tr>
<tr>
<td>XXXV</td>
</tr>
<tr>
<td>II-445</td>
</tr>
<tr>
<td>II-531</td>
</tr>
<tr>
<td>II-533</td>
</tr>
<tr>
<td>II-589</td>
</tr>
<tr>
<td>II-649</td>
</tr>
<tr>
<td>II-707</td>
</tr>
<tr>
<td>II-709</td>
</tr>
<tr>
<td>II-807</td>
</tr>
<tr>
<td>XXXV</td>
</tr>
<tr>
<td>III-874</td>
</tr>
<tr>
<td>III-924</td>
</tr>
<tr>
<td>III-985</td>
</tr>
<tr>
<td>III-987</td>
</tr>
<tr>
<td>III-1019</td>
</tr>
<tr>
<td>III-1059</td>
</tr>
<tr>
<td>III-1097</td>
</tr>
<tr>
<td>III-1203</td>
</tr>
<tr>
<td>III-1295</td>
</tr>
<tr>
<td>XXXV</td>
</tr>
<tr>
<td>XLI</td>
</tr>
</tbody>
</table>
Paper Identification

Each paper is identified with a 6 character code printed on the top right corner of the first page. The code $u v wx yz$ has the following structure:

- $u$ -- type of contribution; $u = I, O, P$ for invited paper, oral and poster contributed paper
- $v$ -- the day of event; $v = 0, 1, 2, ..., 9$ for Monday morning, Monday afternoon, Tuesday morning, ..., to Friday afternoon
- $wx$ -- the topic and subtopic of the contribution
- $yz$ -- the poster board number: $yz = 01, 02, ..., 99$

Note that the authors of oral contributed papers are given the possibility to post their contributions. Therefore, the data in the code refer to the corresponding poster presentation.

Example

P 8 E2 47

<table>
<thead>
<tr>
<th>Paper type</th>
<th>Half-day number of event</th>
<th>Topic and Subtopic</th>
<th>Board number</th>
</tr>
</thead>
<tbody>
<tr>
<td>Poster</td>
<td>Friday morning</td>
<td>RF Heating</td>
<td>47</td>
</tr>
<tr>
<td></td>
<td></td>
<td>Electron Cyclotron Heating</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>Heating</td>
<td></td>
</tr>
</tbody>
</table>
Title List of Contributed Papers

A. TOKAMAKS

E1. Experiments

Murmann H., Wanger F. the ASDEX—, N1—, ICRH—teams
The Isotope Dependence of Global Confinement in Ohmically and
Auxiliary Heated ASDEX Plasmas
P 3 A1 01

Gehre O., Gentle K. W., Richards B., Eberhagen A. et al.
Evaluation of Particle Transport from Gas oscillation
Experiments in Ohmic and Neutral Beam Heated ASDEX Plasmas
P 3 A1 02

Numerical and Experimental Investigation of Neutron
Scattering on ASDEX
P 3 A1 03

Momentum Confinement of ASDEX Plasmas during
Co and Counter Neutral Beam Injection
P 3 A1 04

Improved Confinement at High Densities in Ohmically Heated
and Gas Refuelled Divertor Discharges in ASDEX
P 3 A1 05

Gruber O., Wagner F., Kaufmann M., Lackner K., Murmann H. et al.
Influence of Density Profile Shape on Plasma Transport in ASDEX
P 3 A1 06

Gruber O., Kaufmann M., Lackner K., Lang R. S., Mertens V. et al.
Comparison of Confinement in Hydrogen versus Deuterium in
Multi–Pellet Fuelled OH Discharges in ASDEX
P 3 A1 07

Zeoff—Profiles in Different Confinement and Heating Regimes of ASDEX
P 3 A1 08

McCormick K., Eberhagen A., Murmann H. and the ASDEX Team
q—Profile Measurements in the Central Plasma Region of ASDEX
P 3 A1 09
Mertens V., Sandman W., Kaufmann M., Lang R. S., Buchl V. et al.
Improvement of Beam—Heated Discharges by Repetitive Pellet
Fuelling in ASDEX
P 3 A1 10 

Dodel G., Holzhauer E., Niedermeyer H., McCormick K. et al.
Measurements of Density Turbulence with FIR Laser
Scattering in the ASDEX Tokamak
O 3 A1 11 

Egorov S. M., Kuteev B. V., Miroshnikov I. V., Sergeev V. Yu., et al.
Magnetic Field Line Tracing in T—10 Tokamak
P 3 A1 12 

Dnestrovskij Yu. N., Esipchuk Yu. V., Lysenko S. E., Neudatchin S. V. et al.
Electron Temperature Profile Consistency under ECRH in T—10 Tokamak
P 3 A1 13 

Experimental and Numerical Study of Sawteeth on T—10 Tokamak
P 3 A1 15 

Impurity Transport Study in $\beta$ and S Regimes on the T—10
P 3 A1 17 

Belashov V. I., Borisnikov A. V. and Brevnov N. N.
Experimental Study of Some Problems for a Two—Chamber Tokamak
P 3 A1 18 

Small—Scale Plasma Turbulence in the FT—2 Tokamak
P 3 A1 19 

Vinogradov N. I., Izvozhikov A. B., Silin V. P., Urupin S. A. et al.
Turbulent Ion Heating in TUMAN—3 under the Fast Current Ramp
P 3 A1 20 

Bender S. E., Deshko G. N., Izvozhikov A. B., Kaminskij A. O. et al.
Modelling of the Current Density Distributions under the
Different Discharge Scenarios in TUMAN—3 Plasmas
P 3 A1 21 

Measurements of Fast Ion Radial Diffusion in TFTR
P 3 A1 22 

Ion Temperature Profiles and Ion Thermal Confinement in TFTR
P 3 A1 23 

Barnes C. W., Bosch H. S., Nieschmidt E. B., Saito T., Bitter M. et al.
Triton Burnup Studies on TFTR
P 3 A1 24 

Murphy T. J., Barnes C. W., Schmidt G. L., Strachan J. D., Bosch H. S. et al.
Injection of Deuterium Pellets into Post—Neutral—Beam TFTR Plasmas
P 3 A1 25 
Zarnstroff M. C., Bell M. G., Bitter M., Bush C., Fonck R. J. et. al.
Convective Heat Transport in TFTR Supershots
P 3 A 1 26 .......... 1–95

Goldston R., Takase Y., Bell M., Bitter M., Cavallo A. et al.
Low Power Heating Studies on TFTR
P 3 A 1 27 .......... 1–99

Analysis of Rotation Speed Radial Profiles on TFTR
P 3 A 1 28 .......... 1–103

High Frequency Emission from TFTR Plasmas
P 3 A 1 29 .......... 1–107

Bickerton R. J., Apruzzese G., Tanga A., Thomas P. and Wesson J.
Ignition Tokamaks
P 7 A 1 0 1 .......... 1–111

Christiansen J. P., Connor J. W., Cordey J. G., Lauro—Taroni L. et al.
Local Heat Transport in JET Plasmas
P 7 A 1 0 2 .......... 1–115

Campbell D. J., Christiansen J. P., Cordey J. G., Thomas P. R. and Thomsen K.
Global Confinement Characteristics of JET Limiter Plasmas
P 7 A 1 0 3 .......... 1–119

Lomas P., Bhatnagar V., Campbell D., Christiansen J. P., Chui1on P. et al.
High Current Operation in JET
P 7 A 1 0 4 .......... 1–123

Ion Temperature Profiles and Ion Energy Transport in JET during
Additional Heating and H—Modes
P 7 A 1 0 5 .......... 1–127

Sadler G., Jarvis O. N., van Belle P. and Adams J. M.
Diagnosing RF Driven High Energy Minority Tails with γ—Ray and
Neutron Spectroscopy
P 7 A 1 0 6 .......... 1–131

Batistoni P., Argyle J., Conroy S., Gorini G., Huxtable G. et al.
Measurement and Interpretation of Triton Burnup in JET
Deuterium Plasmas
P 7 A 1 0 7 .......... 1–135

Morgan P. D.
The Evolution of Z eff (r) Profiles in JET Plasmas
P 7 A 1 0 8 .......... 1–139

Kupschus P., Cheetham A., Denne B., Gadeberg M., Gowers C. et al.
Multi—Pellet Injection on JET
O 7 A 1 0 9 .......... 1–143
The JET Multipellet Launcher and Fueling of JET Plasmas
by Multipellet Injection
P 7 A1 10

Gondhalekar A., Campbell D., Cheatham A. D., Edwards A. et al.
Simultaneous Measurements of Electron Thermal and Particle
Transport in JET
P 7 A1 11

O'Rourke J., Blum J., Cordey J. G., Edwards A., Gottardi N. et al.
Polarimetric Measurements of the q−Profile
O 7 A1 12

Operation Regime and Confinement Scaling of Neutral Beam
Heated JT−60 Discharges
P 7 A1 13

Yoshida H., Shimizu K., Shirai H., Tobita K., Kusama Y. et al.
Energy Confinement Analysis of Neutral Beam Heated JT−60
Discharges
P 7 A1 14

High Energy Ion Tail Formation and its Behavior in Additionally
Heated JT−60 Plasmas
P 7 A1 15

Okabayashi M., Asakura N., Bell R., Bol K., Ellis R. et al.
Initial Results from the PBX−M Tokamak
P 7 A1 16

Bracco G., Podda S., and Zanza V.
Ion Temperature and Energy Balance in Ohmic FT Discharges
P 7 A1 17

De Angelis R., Bartiromo R., Mazzitelli G. and Tuccillo A. A.
Impurity Confinement in FT
P 7 A1 18

Experimental Observation of Ion−Temperature−Gradient−Driven
Turbulence in the TEXT Tokamak
P 7 A1 19

Kim S. K., Brower D. L., Foster M. S., McCool S. C., Peebles W. A. et al.
Coupling of Particle and Heat Transport Measured via Sawtooth
Induced Pulse Propagation
P 7 A1 20

Heavy Ion Beam Probe Measurements of Space Potential and Electrostatic
Fluctuations in TEXT with a Resonant Magnetic Field
P 7 A1 21
Zurro B., Hidalgo C., Garcia–Castaner B., Pardo C. and TJ–1 Group
Observation of Anomalous Ion Heating in the TJ–1 Tokamak
P 7 A1 22

Hidalgo C., Navarro A. P., Pedrosa M. A. and Rodriguez R.
Fluctuation Studies in the TJ–1 Tokamak
P 7 A1 23

Plasma Disruptions in Tokamak TBR–1
P 7 A1 26

A2. H Mode

Wagner F., Gruber O., Gehre O., Lackner K., Muller E. R. and Stabler A.
The Power Dependence of RF in the H–Mode of ASDEX
P 3 A2 30

Becker G.
Transport Analysis of the L to H Transition in ASDEX
by Computer Simulation
P 3 A2 31

Budny R., Bell M., Bitter M., Bush C., Dylla H. F. et al.
q–Dependent, H–Mode–Like Phenomena in TFTR
P 3 A2 32

Shoji T., Hoshino K., Kasai S., Kawakami T. et al.
Confinement Studies of H–Mode in Divertor / Limiter Discharges on JFT–2M
P 3 A2 33

Neutral Beam Current Driven Operation of the DIII–D Tokamak
P 3 A2 34

H–Mode Study in DIII–D
P 3 A2 35

Keilhacker M., Balet B., Cordey J., Gottardi N., Muir D. et al.
Studies of Energy Transport in JET H–Modes
P 7 A2 28

The JET H–Mode
O 7 A2 29

Profile Behaviour during L and H Phases of JET Discharges
P 7 A2 30

Lazzaro E., Avinash K., Brusati M., Gottardi N., Rimini F. and Smeulders P.
Analysis of Current and Pressure Profiles in JET H–Mode Discharges
P 7 A2 31
Rebut P. H., Watkins M. L. and Lallia P. P.
Spontaneous Transitions in the Temperature of a Tokamak Plasma with Separatrix
P 7 A2 32

Simulation of Soft X-Ray Emissivity during Pellet—Injection and H—Mode in JET
P 7 A2 33

A3. Theory

Morozov D. Kh.
Anomalous Transports and Relaxation of Toroidal Rotation in Plasma
P 3 A3 36

Rebut P. H.
Current Transport in a Chaotic Magnetic Field and Self—Sustainment of Islands
P 3 A3 38

Frigione D.
Density Fluctuations in FT Tokamak
P 3 A3 39

Amein W. H. and Mohamed B. F.
Time and Temperature Dependent Magnetic Diffusion of an Inhomogeneous Plasma
P 3 A3 40

Haas F. A. and Thyagaraja A.
Theoretical Interpretation of Turbulent Fluctuations and Transport in TEXT
P 3 A3 41

Minardi E.
Resistive Bifurcating States Related to Auxiliary Power in a Tokamak
P 3 A3 42

Chang C. T.
A comparison between Predicted and Observed Pellet Penetration Depth in JET Ohmic—Heated Discharges
P 3 A3 44

Lengyel L. L.
Interaction of Cold High—Density Particle Clouds With Magnetically Confined Plasmas
P 3 A3 46

Lalousis P. J. and Lengyel L. L.
Pellet Particle Deposition Profiles with Allowance for Neutral Gas Expansion Effects
P 3 A3 47
Heimsoth A.
Comparison of Selfconsistent $\beta$—Scaling Laws with Experiments
P 7 A3 35

Hender T. C., O'Brien M. R., Riviere A. C., Robinson D. C. and Todd T. N.
Optimising the Thermal $\alpha$—Particle Yield
P 7 A3 36

Bittoni E. and Haegi M.
Preliminary Results on the Alpha Confinement in NET
P 7 A3 37

Maddison G. P., Hastie R. J. and Bishop C. M.
Direct Losses of Alpha Particles in Spin Polarised Plasmas
P 7 A3 38

Marchenko V. S. and Taranov V. B.
Effect of the Lower Hybrid Wave on the Tokamak Drift Modes
P 7 A3 39

Briguglio S., Bishop C. M., Connor J. W., Hastie R. J. and Romanelli F.
Stability of Toroidicity Induced Drift Waves in Divertor Tokamaks
P 7 A3 40

Jarmen A.
Toroidal Ion Temperature Gradient Driven Drift Modes with Dissipative
Trapped Electron Effects
P 7 A3 41

Van Milligen B. Ph. and Lopes Cardozo N. J.
Tokamak Equilibrium Determination through Function Parametrization
P 7 A3 43

Cenacchi G., Coppi B. and Lanzavecchia L.
Poloidal Field System and MHD Equilibria for the IGNITOR—U
Experiment
P 7 A3 44

Maschke E. K. and Morros Tosas J.
Representation of Toroidal MHD and its Application to Nonlinear
Stationary States
P 7 A3 45

A4. Sawteeth, Disruptions and
Other Related MHD Phenomena

Ward D. J., Gill R. D., Morgan P. D. and Wesson J. A.
The Final Phase of JET Disruptions
P 3 A4 48

Jarvis O. N., Sadler G. and Thompson J. L.
Study of Photoneutron Production Accompanying Plasma
Disruptions in JET
P 3 A4 49
<table>
<thead>
<tr>
<th>Title</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>Impurity Transport in JET during H-Mode, Monster Sawteeth, and after Pellet Injection</td>
<td>1–338</td>
</tr>
<tr>
<td>Measurements of “Snakes” Following Multiple Pellet Feuelling of JET</td>
<td>1–342</td>
</tr>
<tr>
<td>Effects of Large Amplitude MHD Activity on Confinement in JET</td>
<td>1–346</td>
</tr>
<tr>
<td>The Sawtooth in JET</td>
<td>1–350</td>
</tr>
<tr>
<td>High Electron and Ion Temperatures Produced in JET by ICRH and Neutral Beam Heating</td>
<td>1–354</td>
</tr>
<tr>
<td>Effect of Sawteeth and Safety Factor q on Confinement during ICRF Heating of JET</td>
<td>1–358</td>
</tr>
<tr>
<td>Magnetic Measurements of the Sawtooth Instability in JET</td>
<td>1–362</td>
</tr>
<tr>
<td>Resistive Ballooning Modes under Plasma Edge Conditions</td>
<td>1–366</td>
</tr>
<tr>
<td>Ballooning Instabilities in Tokamaks with Sheared Toroidal Flows</td>
<td>1–370</td>
</tr>
<tr>
<td>Resistive Ballooning Modes in Different Collisionality Regimes</td>
<td>1–374</td>
</tr>
<tr>
<td>Analysis of Sawtooth Stabilization in JET</td>
<td>1–377</td>
</tr>
<tr>
<td>The Dynamic Behaviour of the Electron Temperature Profile in the TEXTOR Tokamak Plasma</td>
<td>1–381</td>
</tr>
<tr>
<td>Role of Sawtooth Crashes in a Saturation and Collapse in the PBX Tokamak</td>
<td>1–385</td>
</tr>
</tbody>
</table>
Central Electron Power Deposition from dTeo/dt Measurements on TFTR
P 7 A4 50 .......................... 1–389

Lao L. L., Strait E. J., Taylor T. S., Chu M. S., Burrell K. H. et al.
MHD Stability in High $\beta_T$ DIII-D Divertor Discharges
P 7 A4 51 .......................... 1–393

Characteristics of Low–q Disruptions in PBX
P 7 A4 52 .......................... 1–397

Westerhof E. and Goedl1eer W. J.
Transport Code Studies of M=2 Mode Control by Local Electron
Cyclotron Heating in TFR
P 7 A4 53 .......................... 1–401

Simulation of Plasma Control in the TCV Tokamak with High Frequency
Stabilization
P 7 A4 54 .......................... 1–405

Vlad G. and Bondeson A.
Numerical Simulation of Sawtooth Activity in Tokamaks
P 7 A4 55 .......................... 1–409

White R. B., Rutherford P. H., Colesstock P. and Bussac M. N.
Sawtooth Stabilization by Energetic Trapped Particles
P 7 A4 56 .......................... 1–413

Schep T. J., Pegoraro F. and Porcelli F.
Internal Kink Modes in the Ion–Kinetic Regime
P 7 A4 57 .......................... 1–417

De Blank H. J. and Schep T. J.
The m=1 Internal Kink Mode in a Toroidal Plasma with a Flat
q–Profile near $q=1$
P 7 A4 58 .......................... 1–421

Cap F. and Khalil Sh. M.
Eigenvalues of Relaxed Toroidal Plasmas
P 7 A4 60 .......................... 1–425

Bruschi A., De Luca F. and Jacchia A.
Effect of the Electron Energy Transport Coefficient on the Stability of
the Tearing Modes
P 7 A4 61 .......................... 1–429

Edenstrasser J. W. and Hohenauer W. M. M.
Finite–Beta Minimum Energy States Arising from a Multiple Scale Approach to
Taylor's Minimum Energy Principle
P 7 A4 62 .......................... 1–433

Hender T. C., Gimblett C. G. and Robinson D. C.
Mode Locking in the Tokamak and RFP
P 7 A4 64 .......................... 1–437
### B. STELLARATORS

<table>
<thead>
<tr>
<th>Authors</th>
<th>Title</th>
</tr>
</thead>
<tbody>
<tr>
<td>Nave M. F., Lazzaro E., Gowers C., Hirsch K., Hugon M. et al.</td>
<td>Observation of Nonlinear Resistive Mode Structure on JET Temperature Profiles</td>
</tr>
<tr>
<td>Likin K. M., Ochirov B. D. and Skvortsova N. N.</td>
<td>Energy Deposition Profiles of Simulation of the ECRH in the L–2 Stellarator</td>
</tr>
<tr>
<td>Voronov G. S. and Donskaya N. P.</td>
<td>L–2 Stellarator Plasma Rotation in Ohmic Heating Regime</td>
</tr>
<tr>
<td>Alejaldre C., Castejon F., Goldfinger R. C. and Batchelor D. B.</td>
<td>Analysis of ECRH in the TJ–II Flexible Heliac Using Rays</td>
</tr>
<tr>
<td>Shishkin A. A., Bykov V. E., Peletminskaya V. G. and Khodyachikh A. V.</td>
<td>Magnetic Surface Destruction Due to Equilibrium Plasma Currents in Torsatrons</td>
</tr>
</tbody>
</table>
C. ALTERNATIVE CONFINEMENT SCHEMES

C1. Reverse Field Pinch

Ion Heating Studies in the ZT-40M Reversed Field Pinch
O1 C1 13

Evans D. E. and Tsui H. Y. W.
Comparison of the Theory of Fluctuation-Driven Diffusion with Experimental Observations from HBTX
P1 C1 14
Shinohara S.  
Analysis of Linear Current Distribution and Induced Error Field in RFP Device  
P 1 C1 15  

Fluctuation Measurements in Edge Plasma of the REPUTE-1 Reversed Field Pinch  
P 1 C1 16  

Fujisawa A., Nagayama Y., Yamagishi K., Toyama H. and Miyamoto K.  
Ion Temperature Measurements of REPUTE-1 RFP  
P 1 C1 17  

Antoni V., Bagatin M., Baseggio E., Bassan M., Buffa A. et al.  
Recent Results from the ETA—BETA II RFP Experiment  
P 1 C1 18  

Merlin D., Ortolani S., Paccagnella R. and Scapin M.  
Linear MHD Stability Properties of RFP Configurations  
P 1 C1 19  

Carraro L., Ortolani S. and Puiatti M. E.  
Impurity Diffusion in RFP Plasmas  
P 1 C1 20  

Brandt S z., Jerzykiewicz A., Kociecka K. and Nawrot W.  
Investigations of Breakdown between Plasma—Focus Electrodes  
P 1 C1 21  

Ortolani S., Schnack D. D., Harned D. S. and Ho Y. L.  
Three Dimensional Resistive MHD Modeling of the RFX Reversed Field Pinch Experiment  
P 8 C1 11  

Taylor P., Greenfield C., La Haye R., Ortolani S., Schaffer M. and Tamano T.  
Sustainment of Reversed Field Pinch Plasmas in OHTE  
O 8 C1 12  

Carolan P. G., Lazaros A., Long J. W. and Rusbridge M. G.  
Ion Power Balance Model for Reversed Field Pinch Plasmas  
P 8 C1 13  

Hayden R. J. and Alper B.  
Coherent Oscillations in HBTX1B Reversed Field Pinch Plasmas  
P 8 C1 14  

Tsui H. Y. W. and Evans D. E.  
Fluctuation—Driven Diffusion and Heating in the Reversed Field Pinch  
P 8 C1 15
C2. Other Alternative Magnetic Confinement Schemes

Drozdov V. V. and Martynov A. A.
Calculating the MHD Equilibrium and Stability of a Plasma with Anisotropic Pressure in Axisymmetric Open Traps
P 1 C2 24

Particle Balance Studies by Spectroscopic Method on the Tandem Mirror GAMMA 10
P 1 C2 25

Arsenin V. V.
Average Minimum—B in an Axisymmetric Steeply—Curved Mirror
P 1 C2 26

Zhao Hua and Yang Si—ze
ECRH Trapping of High Energy Gyrated Electron Beam in a Magnetic Mirror
P 1 C2 28

Lehnert B.
Lower and Upper Limits of the Pinch Radius in Extrap
P 8 C2 16

Jin Li
Breakdown in the Toroidal Extrap Experiment
P 8 C2 18

Bortolotti A., Brzosko J. S., Mezzetti F., Nardi V., Powell C. et al.
Heavy Ion Fusion in a Dense Pinch with Enhanced Compression, Ion Acceleration and Trapping
P 8 C2 19

Sinman S. and Sinman A.
Initial Results from Drift Wave Scheme in a Compact Toroid
P 8 C2 20

Sinman A. and Sinman S.
Heating and Confinement in SK/CG—1 Spheromak
P 8 C2 21

Sugisaki K.
Toroidal Z Pinch Experiment
P 8 C2 22

Recent Results of the Toroidal Screw Pinch SPICA II
P 8 C2 23

Składnik—Sadowska E., Baranowski J. and Sadowski M.
Investigation of Convergent Deuteron Beams within a Penetrable Electrode System
P 8 C2 24
Eggen J. B. M. M. and Schuurman W.
Stability of Extended Taylor States in a Weakly Resistive, Cylindrical Finite β Plasma to Helical Perturbations
P 8 C 2 2 5

Bruhns H., Raupp G., Steiger J. and Brendel R.
Magnetic Properties of the Spherical Torus
O 8 C 2 2 6

Studies of the Small Tight Aspect Ratio Torus Concept
P 8 C 2 2 7

D. PLASMA EDGE PHYSICS

Edge Fluctuation Measurements during X—Point Plasmas in JET
O 3 D 6 2

Harbour P. J., de Kock L., Clement S., Erents S. K., Gottardi N. et al.
The Role of the Scrape—off Layer in X—Point Discharges in JET
O 3 D 6 3

Brinkshulte H., Clement S., Coad J. P., de Kock L., Erents S. K. et al.
Plasma Edge Effects during Additional Heating in JET with Belt Limiter Configuration
P 3 D 6 4

Analysis of Thermographic Measurements on the Toroidal Pump—Limiter ALT—II
P 3 D 6 5

Goebel D. M., Corbett W. J., Conn R. W., Dippel K. H. and Finken K. H.
Edge Plasma Characteristics in TEXTOR with the ALT—II Toroidal Belt Pump Limiter
P 3 D 6 6

Plasma Boundary Studies in DITE with ECRH and a Pump—Limiter
P 3 D 6 7

Tokar M. Z.
Tokamak Edge Plasma Transition to the State with Detachment from Limiter
P 3 D 6 8

Study of Edge Plasma Parameters under Ohmic Heating and ECRH on T—10 Tokamak
O 7 D 6 5
Rubel M., Bergsaker H., Emmoth B., Waelbroeck F., Wienhold P. and Winter J.  
The Influence of Limitor Configuration on the Impurity Fluxes in the  
Scrape-off Layer in TEXTOR  
P 7 D 66 

Martin Y. and Hollenstein Ch.  
Influence of the Alfvén Wave Spectrum on the Scrape-off Layer of the  
TCA Tokamak  
P 7 D 67 

Wu C. H., Davis J. W. and Haasz A. A.  
The Formation of Methane by the Interaction of Very Low Energy Hydrogen  
Ions with Graphite  
P 7 D 68 

Martinelli A. P., Taglauer E. and ASDEX Team  
Impurity Flux onto the Divertor Plates of ASDEX  
P 7 D 69 

Zurro B. and TJ-I Group  
Study of Edge Ion Thermal Asymmetries in the TJ-I Tokamak  
P 7 D 70 

Behaviour of Edge Plasma under ICRH in TO-2 Tokamak  
P 7 D 71 

E. RF HEATING

E1. Ion Cyclotron Heating

Tibone F., Evrard M. P., Bhatnagar V., Campbell D. J., Cordey J. G. et al.  
Predictions for ICRF Power Deposition in JET and Modulation  
Experiments during Sawtooth—Free Periods  
P 1 E 1 30 

Bures M., Bhatnagar V., Corti S., Devillers G., Denne B. et al.  
Role of Antenna Screen Angle during ICRF Heating Experiments in JET  
P 1 E 1 31 

Identification of Radial and Toroidal Eigenmodes in the Coupling of the Well  
Defined k1 Spectrum of the New JET ICRH Antennas  
P 1 E 1 32 

Cottrell G. A., Sadler G., van Belle P., Campbell D. J., Cordey J. G. et al.  
Study of ICRF Driven Fusion Reactivity  
O 1 E 1 33 

Hellsten T. and Core W. G. F.  
Resonant Ion Diffusion in ICRF Heated Tokamak Plasmas  
P 1 E 1 34
Kaufman A. N. and Ye H.
Analytic Theory of Absorption, Conversion and Reflection of the Fast Magneto-Sonic Wave at the Second-Harmonic Layer
P 1 E1 35

Davydova T. A. and Lashkin V. M.
Parametric Instabilities of an Inhomogeneous Plasma near Ion-Ion Hybrid Resonance under Ion Cyclotron Heating
P 1 E1 36

Plasma Heating and Quasi-Steady-State Current Drive at Half-Integer Ion Cyclotron Frequency Harmonics in the Toroidal Omega Device
P 1 E1 37

Grekov D. L., Pyatak A. I. and Carter M. D.
Ion Trapping Effect on Cyclotron Fast Magnetosonic Wave Absorption in a Tokamak
P 1 E1 38

Longinov A. V. and Lukinov V. A.
ICRF Antenna System for Exciting Slow Waves in a Plasma
P 1 E1 39

Longinov A. V., Pavlov S. S. and Stepanov K. N.
ICRF Heating Method Using Two-Species Ion Admixture
P 1 E1 40

Krucken T. and Brambilla M.
Applications of the 3-Dim ICRH Global Wave Code FISIC and Comparison with Other Models
P 1 E1 41

Puri S.
Particle Acceleration Near the Faraday Shield via Cyclotron Harmonic Interaction during ICRF Plasma Heating
P 1 E1 42

Sauter O. and Vaclavik J.
Integro-Differential Equation Approach to Electrostatic Wave Problems in ICRF
P 1 E1 43

Parametric Decay in the Edge Plasma of ASDEX during Fast Wave Heating in the Ion Cyclotron Frequency Range
O 8 E i 29

Second Harmonic ICRF Experiment with Ohmic and Strong NBI-Heated Plasmas in JT-60
O 8 E i 30

The ICRF Antennas for TFTR
P 8 E1 31
Comparison of ICRH Heating Scenarios and Antenna Configurations  
in TEXTOR  
P 8 E 32  
II–774

Observations of Harmonics and Parametric Decay Instabilities during  
ICRF Heating on TEXTOR  
P 8 E 33  
II–778

Loninov A. V., Pavlov S. S. and Chmyga A. A.  
On a Possibility of Realizing a High–Power RF Heavy Minority Ion  
Heating of a Plasma  
P 8 E 34  
II–783

Scharer J. and Sund R.  
ICRF Full Wave Field Solutions and Absorption for D–T and D–3He Scenarios  
P 8 E 35  
II–787

Cattanei G. and Murphy A. B.  
Ion Cyclotron Minority Heating of a Two–ion Component Toroidal  
Plasma  
P 8 E 36  
II–791

Analysis of Plasma Coupling with the Prototype DIII–D ICRF Antenna  
P 8 E 37  
II–795

Goedbloed J. P. and D’Ippolito D. A.  
RF Stabilization of External Kink Modes  
O 8 E 38  
II–799

Lam N. T., Lee J. L., Scharer J. and Jost B.  
Analysis and Simulation Measurements of ICRF Waveguide Coupling to  
Divertor Tokamaks  
P 8 E 1 94  
II–803

E2. Electron Cyclotron Heating

Hugill J., Ashraf M., Cox M., Deliyanakis N., Lean H. et al.  
Transport Studies in the DITE Tokamak with Modulated ECRH  
P 1 E 2 45  
II–807

Downshifted Electron Cyclotron Heating Experiments in a near Thermal  
Plasma  
P 1 E 2 46  
II–811

Mantica P., Argenti L., Cirant S., Hugill J. and Millar W.  
Experimental Investigation of Magnetic Field Oscillations on DITE Tokamak  
P 1 E 2 47  
II–815

Electron Cyclotron Resonance Heating on CT–6B Tokamak  
P 1 E 2 48  
II–819
Efficiency of Electron Cyclotron Heating in FT-1 Tokamak  
P 1 E2 50  
II–823

Low–Power ECH Results in TEXT  
P 1 E2 51  
II–827

Cirant S., Argenti L., Cima G., Mantica P., Maroli C. and Petrillo V.  
Preionization and Start–up Experiments with ECRH on THOR Tokamak  
P 1 E2 52  
II–831

Wang Z., Jian G. and Wang E. Y.  
Hot Electron Ring Formation in ECR Heated Plasma  
P 1 E2 53  
II–835

Lazarev V. B.  
Method of Determination of ECR Emission Polarization Characteristics in Tokamak Conditions  
P 1 E2 89  
II–839

Kasparek W., Muller G. A., Schuller P. G., Thumm M. and Erckmann V.  
Performance of the 70 GHz/1 MW Long–Pulse ECRH System on the Advanced Stellarator W VII–AS  
P 8 E2 40  
II–843

Jory H., Felch K., Huey H. and Jongewaard E.  
Millimeter–Wave Gyrotrons for ECRH  
P 8 E2 41  
II–847

Giruzzi G.  
Kinetic Effects on Electron Cyclotron Emission during Electron Cyclotron Heating in Tokamaks  
P 8 E2 42  
II–850

Kriven ski V.  
Quasi–Linear Evolution of the Wave–Damping during High Power Electron Cyclotron Heating  
P 8 E2 43  
II–854

Pesic S.  
Second Electron Cyclotron Harmonic Absorption in the Presence of a Superthermal Tail  
P 8 E2 44  
II–858

Castejon F. and Alejandre C.  
Quasi Electrostatic Branch of X–Mode: a Theoretical Study  
P 8 E2 45  
II–862

Pozzoli R.  
Absorption and Propagation of ECH Pulses in the Presence of Strongly Distorted Electron Distributions  
P 8 E2 46  
II–866

Moser F. and Rauchle E.  
Dispersion and Absorption of Electron Cyclotron Waves in Anisotropic, Relativistic Plasmas  
P 8 E2 47  
II–870
E3. Lower Hybrid Heating

LH Power Absorption and Energy Confinement during Combined Lower Hybrid and NBI Heating on JT 60
P 8 E3 48 ........ III-874

Alladio F., Barbato E., Bardotti G., Bartiromo R., Bracco G. et al.
Lower Hybrid Experiments at 8 GHz in FT
O 8 E3 49 ........ III-878

Cardinali A., Cesario R. and Paoletti F.
Lower Hybrid Parametric Instabilities in the FT Plasma
P 8 E3 52 ........ III-892

Cesario R., Mc Williams R. and Pericoli—Ridolfini V.
Interaction of the Lower Hybrid Pump Wave with FT Edge Plasma
P 8 E3 53 ........ III-896

Lower Hybrid Ion Heating in FT—2 Tokamak
P 8 E3 54 ........ III-900

Lower Hybrid Wave Absorption Studies in FT—2 Tokamak
P 8 E3 55 ........ III-904

Baranov Yu. F., Dyachenko V. V., Larionov M. M., Levin L. S. et al.
Ion Heating in the Tokamak FT—1 at Frequences \( \omega_d < \omega_{\text{LH}} \)
P 8 E3 56 ........ III-908

McCune E. W.
High Efficiency Klystrons for Lower Hybrid Heating Applications
P 8 E3 57 ........ III-912

Shukla P. K., Pavlenko V. N. and Panchenko V. G.
The Saturation of Purely Growing Instability Due to Parametric Excitation of Convective Cells in Plasma
P 8 E3 59 ........ III-916

Pan C. H. and Qiu X. M.
Low—Hybrid Wave Propagation Affected by a Random Medium Layer
P 8 E3 60 ........ III-920

E4. Alfven Wave and Other RF Heating Methods

Joye B., Lister J. B. and Moret J.—M.
Effects of the Alfven Wave Heating on the TCA Plasma Studied by the Dinamical Response
O 1 E4 54 ........ III-924

Joye B., Lister J. B. and Ryter F.
Shafranov Parameter Limits for Ohmic and RF Heated Plasmas in TCA
P 1 E4 55 ........ III-928
De Chabrier A., Duval B. P., Lister J. B. Monpean F. J. and Moret J. M.
Ion Temperature Evolution during Alfvén Wave Heating in TCA
P 1 E4 56

Borg G. G., Howling A. A., Joye B., Lister J. B., Ryter F. and Weisen H.
Kinetic and Current Profile Effects of Alfvén Waves in the TCA Tokamak
P 1 E4 57

Ballico M. J., Brennan M. H., Cross R. C., Lehane J. A. and Sawley M. L.
Alfvén Wave Heating Studies in the TORTUS Tokamak
P 1 E4 58

Dmitrieva M. V., Ivanov A. A., Sidorova A. V., Tishkin V. F. et al.
Twodimensional Computation on Alfvén Heating of a Toroidal Plasma
P 1 E4 59

Diver D. A. and Laing E. W.
Alfvén Resonance Absorption in a Magnetofluid
P 1 E4 60

Puri S.
Alfvén Wave Heating of Toroidal Plasmas with Non-Circular Cross Sections
P 1 E4 62

Borg G. G., Knight A. J., Lister J. B., Appert K. and Vaclavik J.
Alfvén Wave Coupling in Large Tokamaks
P 1 E4 63

Cross R. C.
Propagation of a Magnetically Guided Alfvén Beam in the Edge Plasma
P 1 E4 64

Elfimov A. G.
Plasma Toroidicity Effects on Alfvén Resonances
P 1 E4 65

Brambilla M. and Krucken T.
On the Local Power Absorption of HF Waves in Hot Inhomogeneous Plasmas
P 8 E4 71

Dendy R. O. and Lashmore-Davies C. N.
A Gyrokinetic Description of Cyclotron Resonance Absorption in Toroidal Plasmas
P 8 E4 72

Cardinali A., Lontano M. and Sergeev A. M.
Dynamical Self-Focusing of the High-Power FEL Radiation in a Magnetized Plasma
P 8 E4 73

Murphy A. B.
Surface Waves in a Two-Ion Species Plasma with Finite Edge Density
P 8 E4 74
F. CURRENT DRIVE AND PROFILE CONTROL

F1. Lower Hybrid Current Drive

Leuterer F., Soldner F., Yoshioka K., Okazaki T. and Fujisawa N.
Lower Hybrid Current Drive Efficiency in ASDEX
P 8 F1 62

Briffod G.
Evaluation of the Current Profile in L. H. C. D. Tokamaks
P 8 F1 64

Moreau D., Rax J. M. and Samain A.
Lower Hybrid Wave Stochasticity in Tokamaks: a Universal Mechanism for
Bridging the $n_{||}$ Spectral Gap
O 8 F1 65

Belikov V. S., Kolesnichenko Ya. I. and Plotnik I. S.
Current Drive by LH Waves with the Wide Spectrum
P 8 F1 66

Neudatchin S. V. and Pereverzev G. V.
Numerical Simulation of Current Drive by Lower-Hybrid Waves in T-7
Tokamak
P 8 F1 67

Ray Tracing Studies for the Lower Hybrid Experiments in JET
P 8 F1 68

Barbato E., Cardinali A. and Romanelli F.
Propagation and Absorption of LH Waves in Presence of MHD Turbulence
P 8 F1 69

Jiang T. W., Liu Y. X., Wu G. P. and Zhang X. L.
Multijunction Grill and its Application on Lower Hybrid Current Drive
Experiments
P 8 F1 70

F2. Other Profile Control Methods

Potapenko I. F., Efimov A. G. and Sidorov V. P.
Electrical Field Effect on Alfven Driving Currents
P 1 F2 67

Dudok de Wit T., Howling A. A., Joye B. and Lister J. B.
Alfven Wave Heating and its Effect on the Tokamak Current Profile
P 1 F2 68

Vdovin V. L.
Current Drive by ICRF Waves in Tokamaks
P 1 F2 69
Electron Cyclotron Current Drive Experiments in the WT--3 Tokamak
P 1 F2 70

Bornatici M. and Pieruccini M.
Electron Cyclotron Current Drive: Theoretical Considerations
P 1 F2 71

Giruzzi G.
Optimizing Current Drive by Electron Cyclotron Waves in the Presence of
Trapped Particles
P 1 F2 72

Farina D. and Finardi S.
Propagation, Absorption and Current Generation by EC Waves in the LH
Current Drive Regime
P 1 F2 73

Heikkinen J. A., Karttunen S. J. and Salomaa R. R. E.
Current Drive in Tokamak Plasmas by Beating of High Frequency Waves
P 1 F2 74

Kishimoto Y., Takizuka T., Yamagiwa M., Itoh S. I. and Itoh K.
Effect of Electron Spatial Diffusion on Current Drive
P 1 F2 75

Devoto R. S., Tani K. and Azumi M.
Computation of Self—Consistent 2—D MHD Neutral—Beam and Bootstrap
Currents in Elongated Plasmas
P 1 F2 76

G. NEUTRAL BEAM INJECTION HEATING

Hawkes N. C., von Hellermann M., Boileau A., Horton L., Kaline E. et al.
Profiles of Toroidal Plasma Rotation
P 3 G 70

Transport Studies of High Density Ohmically Heated Plasmas and High
Power Neutral Beam Heated Plasmas on JT—60
O 3 G 71

Carlson A., Buechli K., Gehre O., Kaufmann M., Lang R. S. et al.
Mass Loss with Pellet Refuelling on ASDEX during Neutral Injection
Heating
P 3 G 72

Current Drive and Heating Systems Based on High—Energy (1— to 3—MeV)
Negative Ions Beams
P 3 G 73

Fumelli M., Jequier F. and Pamela J.
First Experimental Results of Energy Recovery on the Tore Supra Neutral
Beam Injector Prototype
P 3 G 92
Feneberg W. and Hellberg M. A.
Transport in an Ergodic Magnetic Field with Ambipolar Electric Field Effects
P 7 G 73

Miljević V. I.
Large Area 4 cm Hollow Anode Ion Source
P 7 G 74

Čadež I., Hall R. I., Landau M., Pichou F., Popović D. and Schermann C.
Determination of Vibrational and Rotational State Population
in Hydrogen by Dissociative Electron Attachment
P 7 G 75

Archipov N. I., Zhitlukhin A. M., Safronov V. M., Sidney V. V. and Skvortsov Yu. V.
Electrodynamic Accelerators Use for High Temperature Plasma Production
P 7 G 90

H. DIAGNOSTICS

Bowden M. D., Brand G. F., Falconer I. S., Fekete P. W., James B. W. et al.
Scattering of Millimetre–Submillimetre Waves from the Tortus Tokamak
Plasma
P 3 H 74

Navarro A. P., Anabitarte E., Alejaldre C. and Castejon F.
A Microwave Reflectometer for the TJ–II Flexible Heliac
P 3 H 75

Tartari U. and Lontano M.
Investigations on ECRH via 140 GHz Collective Scattering in a Tokamak
Plasma
P 3 H 77

Bornatici M., Ruffina U. and Spada M.
Approximate Formulas for Electron Cyclotron Emission at High
Temperatures (50–500 keV)
P 3 H 78

Reflectometry on JET
P 3 H 79

Bartlett D. V., Campbell D. J., Costley A. E., Gottardi N., Gowers C. W. et al.
Integrated Electron Temperature and Density Measurements on JET
P 3 H 80

Kuttel O.
Measurements of Density Fluctuations Using a Homodyne Small Angle Scattering
Technique with a Simplified Wave Vector Selection
P 3 H 81

Manso M., Serra F., Mata J., Borroso J., Comprido J. et al.
A Microwave Reflectometric System for the ASDEX Tokamak
P 3 H 82
Kritz A. H. and Fisch N. J.
Sensitivity of Transient Synchrotron Radiation to Tokamak Plasma Parameters
P 3 H 83

Airoldi A., Orefice A. and Ramponi G.
Polarization Change of Electromagnetic Waves Passing through Toroidal Sheared Plasmas
P 3 H 84

Schild P. and Cottrel G. A.
Ion Cyclotron Emission Measurements on JET
P 3 H 87

Giannone L., Holzhauer E. and Gernhardt J.
Radial Decay of Broadband Magnetic Fluctuations in ASDEX
O 3 H 88

Neudatchin S. V.
Versatile Technique of Finding a Local Dynamic Value of Electron Heat Conduction Coefficient Local Dynamic, \( v_e^H \) from Experimental Data
P 3 H 89

Kasperczuk A., Miklaszewski R., Paduch M., Tomaszewski K. et al.
Plasma Sheath Structure in the PF--150 Plasma--Focus Device
P 3 H 90

Two--Dimensional Optical Tomography of Impurities in the FT--2 Tokamak
P 7 H 76

Barbian E. P., Van Blokland A. A. E., Donné A. J. H. and Van der Ven H. V.
The Applicability of Rutherford Scattering Ti--Measurements at Medium-- and Large--Sized Experiments
P 7 H 77

Tabares F. L.
Applications of Resonant Multiphoton Ionization of Atoms to Fusion Plasma Research
P 7 H 78

Lakićević I., Mucha Z., Hintz E., Samm U. and Uhlenbusch J.
High Resolved Time Measurement of the Light Density Distribution from the TEXTOR Main Limiter Cross--Section
P 7 H 79

Gott Yu. V. and Shurygin V. A.
Measurements of X--Ray Radiation from Plasma by a Photoelectron Method
P 7 H 80

Morsi H. W., Behringer K., Denne B., Kallne E. Rupprecht G. et al.
Results on JET Plasma and Impurity Behaviour Based on Measurements of Radial Profiles in the Soft X--Ray Region
P 7 H 82

Krause H., Kornherr M., ASDEX Team and NI Team
High Resolution Sparse Channel Tomography for Slowly Varying Rotating SXR Profiles
P 7 H 83
Orsitto F. P. and Buratti P.
Collective Thomson Scattering for Alpha Particles Diagnostics in Tokamaks
P 7 H 84

Nagatsu M., Peebles W. A. and Luhmann Jr. N. C.
Current Profile Determination via Polarimetry in High Density, High
Field Tokamaks
P 7 H 85

Hubner K., Batzner R., Bomba B., Rapp H., Herrmann W. et al.
Ion Temperature Determination from Neutron Rate during Neutral
Injection in ASDEX
P 7 H 86

Separate Measurement of Particle and Radiation Losses by Using Time–of–
Flight Type Neutral Particle Energy Analyzer
P 7 H 87

Von Hellermann M., Summers H. and Boileau A.
Investigation of Slowing–Down and Thermalized Alpha Particles by
Charge Exchange Recombination Spectroscopy — a Feasibility Study
P 7 H 88

I. BASIC COLLISIONLESS PLASMA PHYSICS

Shukla P. K. and Stenflo L.
Linear And Nonlinear Coupled Alfvén–Varma Modes in Inhomogeneous
Plasmas
P 11 77

Pavlenko V. P., Petviashvili V. I. and Taranov V. B.
Numerical Simulation of the Evolution of Flute Vortices
P 11 78

Fasoli A., Fontanesi M., Galassi A., Longari C. and Sindoni E.
Electrostatic Ion Cyclotron Waves in a Steady – State Toroidal Plasma
P 11 79

Fasoli A., Galassi A., Longari C., Maroli C. and Petrillo V.
Electrostatic Dispersion Relation in the Ion Cyclotron Regime
P 11 80

Van Niekerk E. and Krumm P.
EIC Properties as a Function of the Current Channel Diameter
P 11 81

Nakah R. and Misguich J. H.
Onset of Chaotic Diffusion in Dynamical Guiding Centers Systems with
More than 2 Electrostatic Waves
P 11 82

Pfirsch D.
Negative Energy Waves in the Framework of Vlasov–Maxwell Theory
P 11 83
Sitenko A. G. and Sosenko P. P.
Nonlinear Generation of Large-Scale Magnetic Fields in Plasmas
P 1 1 8 4

Hansen F. R., Knorr G., Lynov J. P., Pecseli H. L. and Juul Rasmussen J.
Finite Larmor Radius Effects on Particle Diffusion in a Turbulent Plasma
P 1 1 8 5

Lehnert B. and Scheffel J.
On Large Debye Distance Effects in a Fully Ionized Plasma
P 1 1 8 6

Kotelnikov V. A., Nikolaev F. A. and Gurina T. A.
Distribution Functions of Charged Particles in a Time-Dependent Magnetic Field Near the Charged Surface
P 1 1 8 8

Bornatici M. and Chiozzi G.
A New Representation of the Relativistic Dielectric Tensor for a Magnetized Plasma
P 8 1 7 5

Bornatici M., Ruffina U. and Westerhof E.
Fundamental Harmonic Electron Cyclotron Emission for Hot, Loss-Cone Type Distributions
P 8 1 7 6

Amein W. H. and Mohamed B. F.
Wave-Wave Interaction of Hot Collisionless Plasma
P 8 1 7 8

Zaki N. G. and Amein W. H.
Heating of the Plasma by Incident Electron Beam
P 8 1 7 9

Cade V. M. and Okretić V. K.
Leakage of MHD Surface Waves in Stratified Media
P 8 1 8 0

Guha S. and Bose M.
Electron Acoustic and Lower Hybrid Drift Dissipative Instabilities in Multi-Ion Species Plasmas
P 8 1 8 1

El Ashry M. Y. and Papuashvili N. A.
Modulation Instability of Electron Helicon in a Magnetized Collisional Plasma
P 8 1 8 2

Masoud M. M., Soliman H. M. and El-Khalafawy T. A.
Magnetic Reconnection and Instabilities in Coaxial Discharge
P 8 1 8 3

Čerček M. and Jelić N.
Experiments on Double Layers
P 8 1 8 4
Skorić M. M. and Kono M.
Ponderomotive versus Linear Fluid Response in a Magnetized Fusion Plasma
P 8 I 85

Hadžievski Lj. and Skorić M. M.
A Numerical Study of a Stability of Upper--Hybrid Solitons
P 8 I 86

De Angelis U., Jovanović D. and Shukla P. K.
Upper--Hybrid Solitary Vortices
P 8 I 95

J. INERTIAL CONFINEMENT PHYSICS

Filyukov A. A.
Binary Shock Wave Formation Structure in Hydrogen Plasma
P 8 J 92

Borodziuk S., Kostecki J. and Marczak J.
Laser Simulation of Impact of Particles and Foil Acceleration
P 8 J 93

Gribkov V. A., Nikulin V. Ya., Lebedev P. N. and Zmiëvskaya G. I.
The Investigation of the Nonequilibrium Processes within the Laser Produced Streams
P 8 J 96
Tokamaks
The Isotope Dependence of Global Confinement in Ohmically and Auxiliary Heated ASDEX Plasmas

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Introduction.

The investigation of the confinement properties of a tokamak plasma, and its dependence on external as well as intrinsic parameters has been a topic since the very beginning of fusion research and most of the observed effects could be explained by theory. A crucial but still unsolved problem are radial ion and electron energy transport. The energy confinement time has been found to depend on parameters as plasma size, -current, -density etc. and all dependencies are interlinked in a complicated way.

The topic of this paper deals with the dependence of the total energy confinement on the ion mass. A comparison of many discharges in hydrogen (H$_2$) and deuterium (D$_2$) with ohmic (OH) or additional neutral beam heating (NI) reveals a remarkable improvement of energy confinement in the case of deuterium; this effect has been observed on many different machines so far.

Ohmic heating.

The confinement of ASDEX OH - Plasmas and its dependence on local as well as global parameters has been studied extensively /1/ for hydrogen discharges, and various scaling laws have been given. The general behavior since then has not changed although the divertor chamber has been rebuilt completely to enable long pulse additional heating and the recycling properties have changed significantly. In fig.1 the energy confinement time \( t_e \) is plotted versus line averaged density and shows the familiar linear rise for small densities and saturation at about 50 ms as stated before; evidently there is also a dependence on the safety factor \( q_a \), indicating the influence of \( T_e \) and \( n_e \) profile shapes on energy transport, which will have to be studied in a separate work.

With deuterium as working gas the situation does not change in principle; the profiles remain unaltered, but the central temperature becomes higher with the same input power applied - thus increasing the global energy confinement time. This is demonstrated in fig. 2, which shows a series of 10 discharges in which the machine was switched from hydrogen to deuterium (H$_2$ gas puffing up to shot 22473, D$_2$ thereafter). During the discharge the line averaged density has been stepped up, thus increasing the ratio \( x = D_2/H_2 \) also step by step, due to the decreasing admixture of hydrogen recycling from the walls during the discharge. So, with increasing shotnumber the hydrogen will eventually disappear from the vacuum vessel and the ratio \( x \) will approach one when the transition to deuterium is completed.

The loop voltage \( U_l \) rises during the discharge when the density is ramped up, but on average it drops from shot to shot during the \( H \rightarrow D \) transition as a consequence of higher electron temperature and electrical conductivity respectively (fig. 3); with \( I_p \) being kept constant at 0.32 MA the ohmic heating power drops for that reason during the transition; the total energy content, represented by \( \beta_p \) (diamagnetic loop) increases, however, and it can be concluded that the energy confinement in a deuterium plasma is better than in a hydrogen plasma (fig.3 and 4). It has to be noted, that this improved confinement does obviously exist already at low densities in the linear range where the confinement is exclusively dominated by electron transport. In the saturation region deuterium plasmas exhibit the same qualitative levelling off behavior as in the \( H_2 \)-case, but the saturation level is much higher than in the \( H_2 \)-case. (s.fig. 2)
The transition from ohmic to neutral beam heating

Additional heating affects plasma confinement dramatically. This was observed for hydrogen and deuterium as well and documented on ASDEX /2/ as well as many other machines. The beneficial effect of improved confinement for deuterium, however, is conserved. This is demonstrated for a series of discharges with neutral beam heating. Four scenarios are compared: H⁺, and D⁺ as target plasmas and injection of a H⁰ or D⁰ neutral beams. Fig. 5 a,b show a comparison of the confinement development for H⁰ injection into H⁺ and D⁺ target plasmas for various heating powers. τₑ exhibits the familiar deterioration with increasing power and saturation with high density, but there is a clear beneficial effect with deuterium as working gas.

In a comparison of all four heating cases, mentioned above, with discharges of comparable line averaged densities, the electron temperature profiles do not differ relatively in shape, which is revealed in fig. 6, where the radial profiles for all four cases are shown absolutely and normalized at half radius; the only decisive difference is within the confinement zone, which is governed by internal mode activity, sawtoothing, etc. The edge density is higher, however, for deuterium discharges.

The general observation is the decrease of τₑ with increasing total heating power as shown in fig. 5c. The average confinement is best in the case with deuterium injection into a D₂ target plasma and becomes worse, the more hydrogen is involved. The two cases H⁰ → D⁺ and D⁰ → H⁺ seem to have similar characteristics as far as confinement is concerned. The differences of confinement under OH conditions remain with beam heating.

In the high confinement regime (H-regime) during additional heating above the threshold power, the energy confinement time rises again and τₑ values corresponding to those during ohmic heating alone are recovered on ASDEX. This behavior is observable for hydrogen and deuterium as well. In fig. 5c a few examples are shown, representing additional heated plasmas after its transition to the H regime for three cases: H⁰ → H⁺, H⁰ → D⁺ and D⁰ → D⁺ injection. An extensive study of the parameter dependence of τₑ in steady state H-regime discharges, however, is a still matter of forthcoming experimental work.

Summary.

The global energy confinement is dependent on the mass of the plasma ions. This effect has been investigated on ASDEX by comparison of many different discharge types using hydrogen and deuterium as working gas. All of the relevant physical quantities that describe the quality of a plasma in terms of good confinement, as electron- and ion temperature, density, particle and energy confinement times, heat conductivities etc. do change dramatically, when the discharge conditions are altered. Interlinking dependencies, as the well known ohmically constraint obstruct the interpretation of measured data, and unique statements can not always be derived. The empirically found isotope effect, however, seems to be an invariant effect of very fundamental importance. As the gyro radius increases with the ion mass, the transport across the magnetic field is expected to enhance in the case of deuterium; this can certainly not explain the observed isotope effect. In particular the isotope dependence of τₑ is also present in regimes dominated by electron transport. Other, probably more subtle explanations have to be found.

References:
  /1/ O. Klüber, H.Murmann IPP report III/72 -198
  /2/ A.Stäbler, F.Wagner 4th Int. Sympos. of Heating, Rome 1984
Fig. 1. Energy confinement time $\tau_E$ vs. line averaged density $n_e$. Hydrogen, ohmic heating, divertor plasma.
- $2 < q_a < 3$; • $4 < q_a < 5$;

Fig. 2. Energy confinement time $\tau_E$ vs. $n_e$ for the transition of H$^+$ to D$^+$ for an ohmically heated plasma within 10 shots.
Open symbols: last hydrogen discharge and beginning of transition
Solid symbols: 5th to 10th shot in D$_2$

Fig. 3a. Loop voltage $U_1$ vs. time for two ohmically heated discharges, while the line averaged density is being stepped up. Transition from H$_2$ to D$_2$ within 10 discharges:
- Hydrogen (shot 22474)
- Deuterium (shot 22483)
$I_p = 320$ kA; $B_t = 2.2$ T for all discharges.

Fig. 3b. Beta poloidal $\beta_p$ vs. time for the same shots as in fig. 3a. During the transition to deuterium the thermal energy rises although the input power is reduced.

Fig. 4.: (shot 22475-22483): line averaged density and plasma current vs time.
Fig. 6: Thomson scattering measurement of electron temperature $T_e$ and density $n_e$ for four neutral injection heating scenarios:
- injection of $H^0$ into a $H^+$ plasma,
- injection of $D^0$ into a $H^+$ plasma,
- injection of $H^0$ into a $D^+$ plasma,
- injection of $D^0$ into a $D^+$ plasma;
(6a) $T_e$ of a flux tube versus its radius, (6b) $T_e$ normalized at half radius,
(6c) $n_e$ vs. flux tube radius, (6d) $n_e$ normalized at half radius.

Fig. 5: Energy confinement time $\tau_E$ vs. line averaged density $n_e$ for ohmic heating and various neutral injection powers:
(a) hydrogen injection into a hydrogen target plasma ; □ = OH ; ▬ =2 sources
(b) Hydrogen injection into a deuterium target plasma; ● = 1 source ; ○ =4 sources
(c) $\tau_E$ versus the total heating power (OH and NI); symbols as in fig. 6;
transition to the H-regime: + $H^0$->$H^+$; × $H^0$->$D^+$; $D^0$->$D^+$;
EVALUATION OF PARTICLE TRANSPORT FROM GASOSCILLATION EXPERIMENTS IN OHMIC AND NEUTRAL BEAM HEATED ASDEX PLASMAS

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INTRODUCTION:
Gasoscillation experiments were performed at the divertor tokamak ASDEX to evaluate particle transport under different discharge conditions. These experiments followed a method described in /1/. In order to induce density perturbations in the plasma, the external gas feed was modulated with a sinusoidal wave form of chosen frequency (5, 10 or 20 Hz) and the density modulation at the four horizontal chords (r = 0, 10, 21, and 30 cm) of the ASDEX HCH-laser interferometer was observed. A typical result of this measurement is shown in Fig. 1. The precision of the technique arises from the extraction of the complex Fourier amplitudes from the full wave form over several periods. The amplitude and phase have an accuracy typical of a signal averaged over several hundred milliseconds.

DETAILS OF THE ANALYSIS
The transport parameters are determined by picking functional forms for the diffusion coefficient $D(r)$ and the inward convection velocity $V(r)$, which is a major improvement to /1/, where only constant $D$ and $V$ were considered. The particle transport equation is solved and the chord integrals are performed, adjusting the free parameters to give the best least-squares fit to the experimental points. Conceptually, the system can be idealized as a boundary-value problem in which the outermost channel gives the edge density and the inner channels are computed from it, although the actual analysis is more complete with a model for the source layer and plasma beyond the separatrix radius. Furthermore, the experimental data are normalized to the centre channel for comparison with the calculations because the centre channel is most accurate, thus propagating the minimum error.

Mathematically, there are three complex numbers (amplitude and phase of the density perturbation) to be fitted and generally three free parameters (two in the density and one in $V$) to be adjusted for each fit type. The allowed functional forms for $D$ and $V$ are tabulated below; they may be combined in various ways.

\[
D \quad V
\]

\[
D(1+b(r/a)) \quad V(r/a)
\]

\[
D(1+b(r/a)^2) \quad V((r/a)^2)
\]

\[
D(r<b), \quad V((r/a)^3)
\]

\[
D(r>b) \quad V((r/a)(1-r/a))
\]
All solutions are strongly constrained by the transport equation and chord integration.

No choice of free parameters will be able to fit random numbers for the experimental data, similarly, even infinite freedom of \( V(r) \) could not provide acceptable fits if \( V = 0 \) is imposed.

Questions of accuracy, uniqueness, and robustness must be answered semi-empirically. A measurement of the accuracy of the inferences may be obtained by measuring values of the amplitude and phase in different but apparently identical shots, or by repeating the measurements at different times in a long plateau or at different modulation frequencies. Comparison of the transport coefficients evaluated from several data sets of that kind provides one estimate of the accuracy of the coefficients, which can only be regarded as being determined to within the observed variation.

The quantitative evaluation of the error is based on the "SSQ", the sum of the squares of the differences between observed and computed complex amplitudes at the different radial positions. It determines the choices of functional form and values of the constants and the uniqueness of these choices. When this parameter is used, the coefficients for each data set and fit type are well determined: Convective velocities differing by 20% and diffusion coefficients differing by 10% produce distinctly poorer fits. However, values obtained for different modulation frequencies or shots at the same nominal plasma parameters often differ by more than that. A total error of perhaps 30% in \( V \) and 20% in \( D \) should be used in assigning transport coefficients to discharge conditions.

**TRANSPORT RESULTS**

The following discharge conditions were analyzed and the results are summarized in the table below. The functional dependences of the transport coefficients are those which fit the experimental data best. For the diffusion coefficient the most common one is a two-step \( D \), having one \( D \) \((r<16 \text{ cm})\) and another \( D \) \((r>16 \text{ cm})\); the two values are listed in order, \( D \) always being larger towards the periphery.
Convective velocities of the form $V(r/a)^n$ are considered. The coefficient $V$ and the exponent $n$ are tabulated. The ohmic hydrogen discharges are best fit with a $(r/a)^3$ convective velocity, the results are less clear for deuterium. In cases where no clear distinction could be made, the simple \((r/a)\) form is listed. With the application of injection power the \((r/a)\) form is better than the other radial dependences for $V$.

**CONCLUSIONS**

The analysis of a density scan in hydrogen indicates a decrease in central $D$ of almost an order of magnitude, when $\bar{n}_e$ is increased from $1 \times 10^{13}$ cm$^{-3}$ to $-2.3 \times 10^{13}$ cm$^{-3}$. In parallel the energy confinement time $\tau_E$ linearly increases in this range from -20 msec to a saturation value of 53 msec (Fig. 2).
seems to be closely connected to an improvement in central particle confinement.

The changes in transport near the density limit and with neutral injection are significant and characteristic of the phenomena.

The behaviour near the density limit can best be described as a decrease in $D$ near the centre, which has the effect of improving the central particle confinement. The decrease is quite marked and the results may well be consistent with neo-classical $D$ on axis. Comparable results have been found for TEXT and a similar behaviour is also known for impurities from ISX-B /2/. With carbonized wall the central $D$ rather shows a flat response over a wide density range up to the density limit.

The effect of neutral injection is to increase the rate of diffusion in relation to the ohmic case for the same isotope. No significant effect on convection is to be seen but the effect on $D$ is pronounced even at low powers and becomes stronger (up to a factor of three) with increasing NI-power. Only L-type cases with sawteeth during injection could be measured and for these no difference of co- versus counter-injection is found. Quantitatively, the increase in particle diffusivity seems to be even greater than the increase in thermal diffusivity associated with the L-mode.

NUMERICAL AND EXPERIMENTAL INVESTIGATION OF NEUTRON SCATTERING ON ASDEX

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In the past we used the VINIA-3DAMC and NEPMC software [1] to treat the birth, migration and detection of neutrons for the ASDEX facility and compared the results with nuclear emulsion measurements. The agreement between numerical and experimental results was very good, in particular for neutron energies above 2 MeV, despite the fact that we have used a rather simplified ASDEX structure and some approximations in the description of the plasma neutron source.

The VINIA-3DAMC software was improved in order to treat more complex geometric structures, so that a more realistic ASDEX model is now being used. The plasma neutron source was improved, too. Furthermore, the NEPMC software was extended in order to allow arbitrary positions for the nuclear emulsion. Collimators viewing tangentially or in arbitrary directions to the plasma can thus also be interpreted now.

ASDEX model

In our new ASDEX model we distinguish six different groups of the structural components according to their significance for the scattering of neutrons. The model is described in detail in a forthcoming IPP report [2].

Group 1 consists of the core (central screw and wooden core), the 16 toroidal field coils, the ohmic coils OHI to OH8, and the vertical field coils V1 to V4 together with the central multipole correction coils MC1. The remaining coils and the support structure are neglected, mainly because they contribute only small masses in relation to the vessel and the multipole coils. Furthermore, for most of the measurements they are located outside the aperture fields. The vessel (vacuum chamber and its thermal insulation) forms group 2 and the divertor group 3. Group 4 consists of the torus hall, its foundations, walls, roof, and the air. As yet the roof has not been installed at ASDEX, and so for the calculations presented here we filled the roof volume with air.

Group 5 gives the necessary details of the ports. Here we take into account only the port for the YAG light scattering system, near which most of our neutron diagnostics is at present located. Finally, group 6 contains all the details of the nuclear emulsion equipment, i.e. their supports and collimators.

Plasma neutron source

In the VINIA software, the neutron birth points are determined stochastically by reproducing the measured plasma data. To do so, the NR software [3] is used to calculate the neutron emission profile from the measured density and temperature profiles, the Shafranov shift, and the neutron rate. Only discharges with H0-injection into
deuterium plasmas are considered here. It is thus assumed that the ion and electron temperatures have the same radial profiles and time dependences. The ion temperature and the rotation velocity were determined by nuclear emulsion measurements [4]. Owing to the integration over different neutron emission angles the emulsion measurement delivers a mean value for the rotation velocity and there is no simple way to determine the velocity on axis. Therefore we used this mean value in the VINIA calculations as input for the central velocity. The rotation velocity profile is assumed to be parabolic.

The angle of emission and the energy of the neutrons are determined with the rotation velocity at the place of birth being taken into account. For H\textsuperscript{0}-injection we have purely thermonuclear production [4] and in the frame of the rotating plasma the neutron energy distribution is a Gaussian with a half-width determined by the local ion temperature.

The calculations presented here were done for two different shot series, no. 16744-16748 for collimators 3 and 4 (tangential, and antiparallel and parallel to the direction of injection, respectively) and no. 18949-18959 for position 5 (uncollimated, radially directed nuclear emulsion). The corresponding results of the nuclear emulsion measurements are discussed in [4].

The nuclear emulsions were exposed for the whole duration of the discharges. They thus integrate over the time history of the neutron production. In the VINIA calculation this could be simulated by using appropriate time intervals during which the plasma parameters and therefore the neutron production properties do not change considerably and by creating a number of neutrons, for each interval, proportional to the different total neutron yields. An example of this procedure is given in Figure 1 for the shot series no. 16744-16748. Figure 1 shows the time evolution of the neutron rate (mean value for the five discharges). We considered four time intervals, indicated by the vertical lines, using the plasma parameters corresponding to the times indicated by the dots.

Results of VINIA calculations

We present here the results of VINIA test runs for the two collimators 3 and 4 (2830 neutrons each) and the unshielded position 5 (4000 neutrons). So far we looked for qualitative tendencies only and limited our calculations to small numbers of neutrons.

Figure 2 shows the VINIA calculation of the spectral neutron fluences arriving at the emulsions from the full solid angle. The essentially higher contribution of collided fluence at the unshielded emulsion is obvious. The relative shift of the main line from the emitted neutrons in collimators 3 and 4 is caused by the plasma rotation.

Fig.1: Time dependence of neutron rate and intervals for VINIA calculations

Fig.2: VINIA calculations of the spectral neutron fluence arriving at the emulsion
Results

Table 1 compares the emitted and collided contributions normalized to the number of emitted neutrons for the full energy interval considered (1 to 3 MeV) at each of the three positions. For positions 3 and 4 the values without collimator in place are also given.

Table 1  Neutron fluences per emitted neutron in $10^{-7}$ cm$^{-2}$

<table>
<thead>
<tr>
<th>Pol.</th>
<th>Coll. 3</th>
<th>Coll. 4</th>
<th>Pos. 3</th>
<th>Pos. 4</th>
<th>Pos. 5</th>
</tr>
</thead>
<tbody>
<tr>
<td>emitted</td>
<td>0.56± 6%</td>
<td>0.68± 6%</td>
<td>2.28± 5%</td>
<td>2.30± 5%</td>
<td>3.78± 3%</td>
</tr>
<tr>
<td>collided</td>
<td>0.49±52%</td>
<td>0.10±22%</td>
<td>4.22±21%</td>
<td>2.83±12%</td>
<td>6.82±11%</td>
</tr>
<tr>
<td>ratio c/e</td>
<td>0.88</td>
<td>0.15</td>
<td>1.85</td>
<td>1.23</td>
<td>1.80</td>
</tr>
</tbody>
</table>

Position 5 is near the quartz window and positions 3 and 4 are near the toroidal field coils, and the reduction of the fluences at the unshielded positions 3 and 4 compared with position 5 is mainly an effect of shielding by the field coils. The big scatter of the results for the collided fluence inside the collimators is caused by collisions in the collimator material. Because these collisions are rare, their treatment needs an essentially higher number of neutron histories. This is also evident also from Table 2, where we give the contributions of the different groups of constituents of the ASDEX facility to the collided fluence.

Table 2  Collided neutron fluences per emitted neutron in $10^{-7}$ cm$^{-2}$

<table>
<thead>
<tr>
<th>Group</th>
<th>Regions</th>
<th>Coll. 3</th>
<th>Coll. 4</th>
<th>Pos. 5</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>core, coils</td>
<td>0.0030</td>
<td>0.0053</td>
<td>0.0054</td>
</tr>
<tr>
<td>2</td>
<td>vessel</td>
<td>0.0075</td>
<td>0.0111</td>
<td>0.0087</td>
</tr>
<tr>
<td>3</td>
<td>divertor</td>
<td>0.0021</td>
<td>0.0013</td>
<td>0.0004</td>
</tr>
<tr>
<td>4</td>
<td>hall</td>
<td>0.0048</td>
<td>0.0018</td>
<td>0.0061</td>
</tr>
<tr>
<td>5</td>
<td>quartz window</td>
<td>0.0014</td>
<td>0.0121</td>
<td>0.0007</td>
</tr>
<tr>
<td>6</td>
<td>rest of port 1</td>
<td>0.0009</td>
<td>0.0107</td>
<td>0.0041</td>
</tr>
<tr>
<td>7</td>
<td>coll. 3, 4</td>
<td>0.0108</td>
<td>0.0307</td>
<td>0.2079</td>
</tr>
<tr>
<td>total</td>
<td></td>
<td>0.0305</td>
<td>0.0730</td>
<td>0.2333</td>
</tr>
</tbody>
</table>

Results of NEPMC calculations and measurements

The NEPMC software is used to simulate the response of the nuclear emulsion [1], taking the VINIA results as input. Figure 3 gives the NEPMC results for collimators 3 and 4 compared with the experimental results. In the calculated spectra the line shift caused by plasma rotation is smaller than in the measured spectra by a factor of about 1.9. Hence the true central rotation velocity was about $4 \times 10^7$ cm/sec. A more accurate determination would be possible by detailed VINIA calculations.

Fig.3: NEPMC results for collimators 3 and 4 and experimental results
The collided fluence is reduced by the NEPMC response calculations. This needs further investigation; it may be caused by the directional response of emulsions.

Table 3 compares the normalized calculated and measured neutron fluences for the two energy intervals 1 to 2 and 2 to 3 MeV separately. Good agreement is obtained for the high-energy range, but for the low-energy range the statistical error in the calculation is too large for a comparison.

Table 3: Calculated and measured normalized neutron fluences \([10^{-7} \text{ cm}^3]\)

<table>
<thead>
<tr>
<th></th>
<th>VINIA calculation</th>
<th></th>
<th>Emulsion measurement</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>1 to 2</td>
<td>2 to 3</td>
<td>1 to 2</td>
</tr>
<tr>
<td>Coll. 3</td>
<td>0.24</td>
<td>0.81</td>
<td>0.22</td>
</tr>
<tr>
<td>Coll. 4</td>
<td>0.03</td>
<td>0.74</td>
<td>0.23</td>
</tr>
<tr>
<td>mean value</td>
<td>0.14</td>
<td>0.77</td>
<td>0.23</td>
</tr>
</tbody>
</table>

Figure 4 gives the results of NEPMC calculations for the proton and neutron energy spectra compared with the experimental results. For an uncollimated nuclear emulsion measurement, the neutron energy spectrum is determined by differentiating the proton energy spectrum. The agreement between numerical and experimental spectra is good; unfortunately owing to limitations during scanning of this plate the measured proton energy spectrum is meaningless for energies below 2.2 MeV and therefore a determination of the absolute values is not possible.

The low-energy wing of the measured neutron energy spectrum is broader than the numerical spectrum. This is probably caused by an underestimation of the measuring errors for the track length in the NEPMC software.

Fig.4: NEPMC results for position 5 compared with experimental results

References
[3] K. Hübner, e.a., this conference
MOMENTUM CONFINEMENT OF ASDEX PLASMAS DURING CO AND COUNTER NEUTRAL BEAM INJECTION


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I. Introduction

Toroidal rotation velocities of ASDEX plasmas of up to $2 \times 10^7 \text{cm} \text{s}^{-1}$ have been measured under a variety of conditions and for both co- and counter neutral beam injection. The velocities were obtained from the Doppler shifts of O VIII 2976 Å and/or C VI 3434 Å excited by charge exchange recombination (CXR). Up to five lines of sight have been used (table I), each of which intersects the axis corresponding to a source in the north-west neutral injection beam line. The main objective has been to compare the global momentum confinement with the energy confinement of ASDEX plasmas for different plasma parameters and heating scenarios.

II. Estimates of Momentum Confinement Time

The experimental global momentum confinement time was estimated from $\tau_\phi = L/\Gamma$ where $L$ is the total toroidal angular momentum of the plasma and $\Gamma$ the momentum input from the beams. $L \simeq 2\pi R^2 m_p \int_0^\infty V_\phi(r) n_c(r) 2\pi r dr$ where $f$ is the ratio of the average atomic mass to that of a pure hydrogen plasma at the same $n_c$. With the assumption that only carbon and oxygen impurities, present in roughly equal densities, contribute to $Z_{\text{eff}}$, $f \simeq \frac{1+2g}{1+g} + 0.16 \left(\frac{Z_{\text{eff}}-1}{1+g}\right)$.

The ratio $g = n_D/n_H$ was typically rather small ($\approx 1$) in the $H_\alpha \rightarrow D_+ \beta$ plasmas and was estimated to about $\pm 50\%$ from neutron measurements. $Z_{\text{eff}}$ was obtained from infra-red continuum measurements. Errors in $g$ and $Z_{\text{eff}}$ do not affect $\tau_\phi$ too critically and these two quantities were assumed to be independent of radius. The main uncertainty in $\tau_\phi$ arises from that in $V_\phi(r)$.

One problem with measuring $V_\phi$ with the observed CXR lines is that they are also excited in the plasma edge region even without NI. It has been possible to investigate the importance of edge excited lines, with plasmas heated by the SE beam which is not seen directly by any of the lines of sight in table I. Interference from the edge excited transition together with low signal to noise for the inner channels, due to beam penetration limitations ($E_\phi = 41$ keV), give rise to problems in measuring $V_\phi$ at small radii, particularly for C VI 3434Å. However, at least for $r \geq 14\text{cm}$, $V_\phi$ is believed to be accurate to within a systematic uncertainty $\pm 10^6 \text{cm} \text{s}^{-1}$ with statistical errors of a similar size. The estimated errors in $\tau_\phi$ are $\approx \pm 15\%$ statistical with a further $\approx \pm 20\%$ systematic.

Apart from neoclassical theory, which gives estimates of $\tau_\phi \approx 10^3 \times$ higher than experiment, the only theory which allows a quantitative estimate of $\tau_\phi$ is the gyroviscous theory of Stacey et al. [1]. Here $\tau_\phi \approx 2R^2 B_0(Z_{\text{eff}})/\langle T_i \rangle$, where the brackets imply
Rotation measurements for a co-injection (L-mode) run at fixed power and different momentum confinement results have been obtained for H₀ injection into mainly D⁺ but also some H⁺ plasmas with co- (L-mode) and counter injection. A study of H-mode discharges has also started. Unlike 17¢, [1], 17¢ shows no clear field dependence (Fig. 3), due particularly to a strong density dependence of Zₑff, which falls off with increasing input power for TFTR plasmas in terms of increasing viscous dissipation at higher rotation velocities [3]. In this case £ = £₀/(1 + £₀ Vₑ² fₚ / R² ω), where ω is the ion gyrofrequency and fₚ is a profile factor.

The observed density dependence of V₁ for ASDEX plasma in principle provides a way of testing this theory at constant input power because in the ohmic (“reference”) phase V₁ < V₁(NI) and £₀ ≈ £₀(NI), which is essentially independent of nₑ (Fig. 4). However, using experimental values for Zₑff, £₀ and V₁, £₀ is also essentially independent of nₑ, though with a reduction from the ohmic heating case about 2× lower than observed (Fig. 4). This is again due to the strong nₑ dependence of Zₑff, which influences the effective £.

Significantly different from the co-injection results are those obtained with counter injection. Here the velocities rise to values more than 2× those obtained at the same density with co-injection (Fig. 1) leading to values of £₁ of up to 90 ms, similar in magnitude to £₁ (Fig. 2). In these discharges, the increased energy and momentum confinement is accompanied by improved particle confinement which leads to a steep rise in nₑ with time. The corresponding impurity accumulation, which occurs for light
as well as heavy impurities, leads to large axial radiation losses and termination of the discharges by major disruption. Gyroviscous theory gives too small a $\tau_\phi$ because $Z_{eff}$ in the counter-NI case is not much larger than in the co-NI case while $T_i$ is approximately the same. The diffusive model gives $\tau_\phi \approx \tau_E$, as observed.

Co-injection discharges where the density is ramped up with time, as well as pellet injected discharges, (both of which can reach similar densities to the counter-injection cases) attain similar velocities to those in steady state co-discharges, well below the corresponding counter-injection velocities (Fig. 1).

H-mode discharges (single-null) with a high ELM frequency ($\approx 400 H_e$) show no significant differences in velocity to corresponding L-mode (double-null) discharges except at the relatively large radius of 35 cm. Here a distinctive increase of $\approx 2 \times$ in $V_\phi$ is seen which disappears when the discharge goes back into the L-mode. Such an increase in $V_\phi$ could be interpreted as a decrease in edge momentum diffusivity which accompanies a decrease in edge electron thermal diffusivity and particle diffusion. The lack of significant increase in global $\tau_\phi$ for such discharges is also to be compared with a lack of increase in $\tau_E$ compared with the L-mode.

IV. Summary

In the results studied to date $\tau_\phi$ is always about the same size as $\tau_E$. In two scenarios where improved energy confinement is seen relative to the L-mode, improved momentum confinement is also found (and indeed improved particle confinement). There is a global improvement with counter injection but only an edge improvement for the H-mode with high ELM frequency. Gyroviscous theory gives agreement with experiment within about a factor of 2 though systematic differences are apparent. A simple diffusion model also gives satisfactory agreement.

References


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Table I: Details of lines of sight used for CXR spectroscopy
(The separatrix is at $r = 40$ cms and the limiter typically at $r = 46.0$ cms.)
Fig. 1

$V_\phi (24 \text{cm})$ versus $\bar{n}_e$ for counter injection (triangles) and co-injection (circles) with $P_{NI} = 1.3 \text{MW}, I_p = 380 \text{kA}$. The symbols $\blacktriangle, \triangle$ (counter-NI) and $\varnothing, \Theta$ (co-NI) each refer to one discharge, with the arrows indicating the time sequence. The symbols $\bullet$ each correspond to an average of several measurements taken during the density plateau of a discharge.

Fig. 3

Comparison of $\tau_\phi$ with $\tau_E$ and $\tau_{EG}$ as a function of $\bar{n}_e$ for co- and counter injection.

Fig. 4

Test of the velocity dependence of $\tau_\phi$ from gyroviscous theory. $\tau_{EG}$, $\tau_E$ are the experimental confinement times for ohmic and beam heated plasmas respectively and $\tau_{EG}$ the values given by gyroviscous theory with experimental $V_\phi$, $\Gamma_{\phi\phi}$ and $Z_{eff}$.
IMPROVED CONFINEMENT AT HIGH DENSITIES IN OHMICALLY HEATED AND GAS REFUELLED DIVERTOR DISCHARGES IN ASDEX

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New regime of improved ohmic confinement in gas refuelled discharges.
In ohmically heated and gas refuelled divertor discharges with deuterium in ASDEX, a bifurcation of the energy confinement time \(\tau_E^*\) is observed at densities above \(n_e = 3.8 \times 10^{13} \text{ cm}^{-3}\) (\(I_p = 390 \text{ kA}; B_i = 2.17 \text{ T}\) ) (Fig. 1; \(\tau_E^* = W / P_{OH}\) ). Discharges at stationary, feed-back controlled \(\tilde{n}_e\) in the original ASDEX divertor configuration (DV-I) and discharges with strong gas puffing (\(\tilde{n}_e\) ramp) in the new divertor configuration (DV-II) / I / show the usual \(\tau_E^*\) saturation at \(\tau_E^* = 85 \text{ ms}\) with increasing \(\tilde{n}_e\) and are in the saturated ohmic confinement (SOC) regime. Now, deuterium discharges at stationary \(\tilde{n}_e\) with a reduced gas puff rate in the new DV-II configuration yield considerably improved \(\tau_E^*\) values at high \(\tilde{n}_e\) (with a present maximum of \(\tau_E^* = 110 \text{ ms}\) ). The plasma is in a new improved ohmic confinement (IOC) regime, indicated by essential changes in the \(\tilde{n}_e\) scaling of \(T_{e,0}^*\), \(T_{i,0}^\ast\), \(Z_{eff}\) of the \(\tilde{n}_e\) profile shape and of the particle transport (see below). So far, the IOC regime is obtained only in deuterium. The energy confinement time of hydrogen discharges, needing more gas puffing than respective deuterium discharges because of their poorer particle confinement, still saturates at \(\tau_E^* = 55 \text{ ms}\). All discharges discussed in this paper have been performed with stainless-steel walls of the vacuum vessel without carbonization or gettering.

In the modified ASDEX divertor configuration DV-II, the surface areas in the divertor chambers pumping atomic hydrogen are much smaller and the conductance values of the bypasses from the divertor to the main chamber are much larger than in the old DV-I configuration. As a consequence, in DV-II the external gas puff rate is roughly halved and, secondly, only about 1/4 of the gas backflow from the divertor into the main chamber is guided through the divertor slits to the stagnation points but 4/5 reach the main chamber in a more uniform way through the bypasses, whereas this fraction was negligible in DV-I. Our experimental results thus suggest that gas refuelling and recycling affect the plasma edge properties in a way that is crucial for the global plasma confinement.

Plasma parameter evolution during transition to improved confinement.
Fig. 2 presents the evolution of various plasma parameters during conversion from the SOC to the IOC regime in the course of a discharge. The change in confinement is initiated by the preprogrammed reduction of the external gas feed at \(t = 1.17 \text{ s}\) in order to establish a \(\tilde{n}_e\) plateau. The power losses into the divertor (see RADDIV signal) drop immediately accompanied by a decay of \(n_e\) in the plasma periphery and within the scrape-off layer. About 65 ms later, 40 ms prior to the \(\tilde{n}_e\) plateau, the increase of the energy content of the plasma (\(\beta_p\)) is already clearly recognizable and it continues for \(280 \text{ ms}\) till the \(\beta_p\) plateau is attained. The improving confinement causes higher central temperatures \(T_{e,0}^*\) and \(T_{i,0}^\ast\)

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at reduced loop voltage and ohmic heating power. Simultaneously with the rise of $\beta_p$ and the peaking of the $n_e$ profile (Fig. 3), the volume integrated radiation power losses within the outer half of plasma minor radius (RAD$_{edge}$) decrease and, reversely, those within the inner plasma half (RAD$_{center}$) increase (see also radiation profiles in Fig. 5). RAD$_{edge}$ is predominated by line radiation of the low-Z impurities oxygen and carbon and is, therefore, time-correlated with the O VI line intensity. RAD$_{center}$ is mainly due to the metals iron and (only in DV-II configuration) copper, and has a similar time behaviour as the soft X-ray signal. The sawtooth pattern of the discharge changes after gas puffing is diminished at $t = 1.17$ s. From $t = 1.51$ s, after a 70 ms interval with stronger impurity accumulation and weaker sawtooth but enhanced MHD activity, sawteeth develop with longer periods and larger amplitudes than in the SOC regime ($t < 1.170$ s). At the same time, the $\beta_p$ plateau is reached. RAD$_{center}$ does not stop to grow but levels off about 120 ms later.

**Peaking of the electron density profile at improved confinement.**

Plasma regimes with improved confinement, resembling the IOC regime, are achieved in ohmic discharges with pellet refuelling / 2 / and in auxiliarily heated discharges with neutral-beam injection in counter direction / 3 /. The last-mentioned regimes are always associated with strongly peaked $n_e$ profiles. During the SOC-to-IOC transition in one discharge, the $n_e$ profile peaks gradually, whereas the $T_e$ profile remains constant in shape though $T_{e,0}$ increases moderately (Fig. 3). Correspondingly, the bifurcation of $\tau_E^*$ on the $\tilde{n}_e$ scale is connected with a bifurcation of $Q_T = n_{e,0} / < n_e >$, a quantity representative of the degree of $n_e$ profile peaking (Fig. 4): In the IOC regime, the $n_e$ profile becomes more and more peaked with growing $\tilde{n}_e$ as well as $\tau_E^*$ and, contrarily, in the SOC regime the $n_e$ profile broadens somewhat at larger $\tilde{n}_e$. The respective profile shape factor $Q_T$ for the electron temperature does not vary with $\tilde{n}_e$, despite the fact that $T_{e,0}$ falls from 1920 eV to 880 eV over the entire density range.

**Impurity accumulation during improved confinement.**

The central peaking of the radiation power profile in the IOC regime (Fig. 5) is the result of an impurity accumulation at the plasma centre. The peaking of the $n_e$ profile accelerates the central radiation growth but is too small to explain it fully. $Z_{eff}$ rises from 1.6 to 2.7 and the $Z_{eff}$ profile slightly peaks at the plasma centre during the SOC-to-IOC transition in shot #22990 / 4 /, which also points towards impurity accumulation. Injection of titanium by laser blow-off technique proves that the impurity confinement time of the plasma core (0-to-30 cm) is enlarged from 55 ms in the SOC regime to 185 ms in the IOC regime. Preliminary impurity transport code analysis using spectroscopic data reveals that a reduction of the diffusion coefficient is mainly responsible for the improved transport properties, whereas the inward drift velocity is scarcely changed. Both impurity accumulation and high central radiation losses delay the peaking of the current density profile and contribute thus to prolong the sawtooth period, an effect which again facilitates impurity accumulation.

**No plasma detachment in both confinement regimes.**

Bifurcating ohmic $\tau_E^*$ behaviour has been found, too, in the limiter tokamak TEXTOR / 5 /. Detachment, i.e. plasma shrinkage due to extremely high edge radiation power losses has been quoted as cause of the deteriorated confinement. In ASDEX, detachment does definitely not occur in the SOC regime. Even at the highest densities of the SOC regime in DV-I configuration, the total radiation power losses of the main chamber plasma (RAD$_{center}$ + RAD$_{edge}$) are only about 33% of the ohmic power input and the power flow accounted for into the divertor (RAD$_{DIV}$ + DEP) amounts to about 60% of the heating
power (Fig. 6 (a)). In the SOC regime, both the electron density and temperature profile do not show any shrinking at the edge. Just the reverse is observed, at improved confinement in the IOC regime the \( n_e \) profile is more centrally peaked and the edge density is reduced. In the IOC regime, also both \( \text{RAD}_{\text{center}} \) and \( \text{RAD}_{\text{edge}} \) are larger and the energy losses into the divertor (\( \text{RAD}_{\text{DIV}} \)) are smaller (compare Figs. 6 (a) and (b)), a fact which is in contrast with TEXTOR results. It should be stressed, however, that there is also no detachment in the IOC regime. The SOC regime in DV-II demonstrates that \( n_e \) profiles broader than in the IOC regime well exist at \( \text{RAD}_{\text{edge}} \) values higher than in the IOC regime (compare Figs. 6 (b) and (c)).

**References**

/1/ H. Niedermeyer, et al., this conference.
/2/ M. Kaufmann, et al., acc. for publ. in Nucl. Fus.; O. Gruber, et al., this conference.

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**Fig. 1** (above) : Energy confinement time \( \tau_E \) versus line averaged density \( \bar{n}_e \) in ohmic divertor discharges (\( I_p = 390 \) kA; \( B_t = 2.17 \) T)

(A) \( \triangle \) \( \text{D}_2 \)-gas; DV-II; \( \bar{n}_e \)-plateau

(B) + \( \text{D}_2 \)-gas; DV-II; \( n_e \)-ramp

(C) \( \square \) \( \text{D}_2 \)-gas; DV-I; \( \bar{n}_e \)-plateau

(D) \( \star \) \( \text{H}_2 \)-gas; DV-II; \( \bar{n}_e \)-plateau

**Fig. 2** (right) : Time evolution of various plasma parameters during transition from \( \tau_E \) saturation (SOC) to improved \( \tau_E \) regime (IOC) - note the increase in \( \beta_p \).

Radiation power losses of the torus plasma are divided into central (\( \text{RAD}_{\text{center}} \)) and edge part (\( \text{RAD}_{\text{edge}} \)). The volume power losses of the divertor plasma (\( \text{RAD}_{\text{DIV}} \)) are mainly due to neutral particle losses.

(\( \text{D}_2 \)-gas; DV-II; \( I_p = 390 \) kA; \( B_t = 2.17 \) T)
Fig. 3: Radial profiles of electron density \( n_e \) and temperature \( T_e \) at two discrete times representative of the two kinds of ohmic confinement regime during one discharge (for plasma parameters, see Fig. 2).

Fig. 4: Plasma density \( \bar{n}_e \) dependence of profile shape factors of electron temperature \( Q_T \) and density \( Q_n \) for both cases of ohmic confinement regime (same discharges and symbols as in Fig. 1).

Fig. 5: Evolution of radiation power density profile during transition from SOC to IOC regime (same discharge as in Fig. 2).

Fig. 6 (right):
(a) SOC regime in \( \bar{n}_e \)-plateau of DV-I (discharge parameter set C of Fig. 1)
(b) IOC regime in \( \bar{n}_e \)-plateau of DV-II (discharge parameter set A of Fig. 1)
(c) SOC regime in DV-II during strong gas puffing (parameter set B of Fig. 1)
DEP denotes the power deposition on the divertor neutralizer plates.
INFLUENCE OF DENSITY PROFILE SHAPE ON PLASMA TRANSPORT IN ASDEX


1. Introduction
In ASDEX divertor discharges we observe - besides the H-mode - four regimes of long lasting improved energy and particle confinement compared with gas-fuelled (GP) ohmic (OH) and co-beam heated plasmas, namely OH and beam heated pellet refuelled discharges [1, 2], OH discharges without sawteeth, discharges with neutral injection in the counter direction (ctr-NI) [3] and ohmic discharges near the density limit with reduced gas-puff (GP) [4]. All these discharges have strongly peaked density profiles with ratios of the axial (n_e(0)) to the volume averaged (< n_e >) electron density of up to 2.5 compared with 1.4 ± 1.6 of the other discharge types. During the density increase to line-averaged values (n_e) of up to 1.2·10^{20} m^{-3} the electron temperature (T_e) profile shape stays nearly self-similar depending only on q_a. Good particle confinement and high n_e(0) lead to high-Z impurity accumulation and high central radiation losses, which can dominate the central power balance and can give rise to hollow T_e-profiles and internal disruptions in a final phase. The tendency toward the observed reduced sawtooth activity may be a further result. Z_{eff} increases owing to light impurities.

According to theory peaked density profiles may lead to reduced anomalous heat transport if the threshold condition of the trapped electron or ion temperature gradient modes, namely \eta_{e,i} = d\ln T_{e,i}/d\ln n_e = L_{n_e}/L_{T_e,i} \geq \eta_c \approx 1.5, is not violated [5]. The measured \eta \approx 1.5, is not violated [5]. The measured \eta_e and - as T_i \approx T_e at the high n_e's - \eta_i values of the high confinement discharges indeed decrease to values below 1 over a large part of the plasma radius (see Fig. 3a, [1]). An alternative explanation follows the profile consistency model. According to it, T_e(r)/T_e(0) in the bulk of the plasma is fixed by stability restrictions, whereas the normalization constant is determined by local transport processes in a near-boundary zone. The roll-over of \tau_E(n_e) in GP OH discharges follows then from a decrease of T_e(0) with density and the concomitant increase in ohmic dissipation, whereas the high confinement discharges gain in \tau_E from a more favourable weighting of T_e(r)/T_e(0) with the peaked density profiles.

In this paper we try to identify the dominating energy loss channels using the TRANSP analysis code and measured plasma parameters: n_e, T_e, T_i and radiation profiles and global parameters (loop voltage, Z_{eff}, \beta_{pol}).

2. Gas-fuelled OH and co-beam heated discharges
In ASDEX we observe "broad" density profiles with n_e(0)/ < n_e > \approx 1.4 ± 1.6 and \eta_{e,i} > 1 at all radii in gas-fuelled ohmic and co-beam heated plasmas at q-values around 3. OH discharges show the roll-over from the linear dependence \tau_E \sim n_e to a saturated \tau_E regime beyond n_e \approx 3 \cdot 10^{19} m^{-3}. Confinement is degraded in the additionally heated L-mode plasmas and improves again in the H-mode even at high
heating powers. In all three regimes a reduction of $\tau_E$ of hydrogen ($H^+$) discharges in comparison with deuterium ($D^+$) discharges holds. To describe the observed confinement we have to add to an anomalous electron heat conduction channel and to the neo-classical ion energy losses as given by Chang-Hinton, $\chi_{CH}$, an additional heat conductivity contribution causing for instance the saturation of ohmic confinement at high densities. CX measurements of $T_i$, measured $\beta_p$ and neutron productions can be described consistently when we assume enhanced ion losses with an enhancement factor of $\chi_i = (3/4)\chi_{CH}$ over the neoclassical value.

This brings low and high density OH results into line with a $\chi_e(OH) \sim 1/(nT_eq)$ [1], dominating at low densities the power balance. We have started to simulate the ion losses by a $\chi_i = \chi_{CH} + \chi_{\eta_i}$ [5], with $\chi_{\eta_i} = 0$ for $\eta_i \leq 1$ and fully developed for $\eta_i \geq 1.8$ including an enhanced threshold for the long density decay length ($L_{n_i}$) region. Figure 1 shows for a GP ohmic discharge that $\chi_{\eta_i}$ yields obviously the necessary $\chi_i$ enhancement, and $\eta_i$, which is smaller than $\eta_e$, is clamped to values between 1 and 2. Electron and ion heat conduction losses ($P_{ee}, P_{ci}$) are about the same at this medium density. The $\eta$-modes have, however, in their present theoretical form the wrong dependence on the ion mass $A_i$ ($\chi_{\eta}$ is increasing with $A_i$) to explain the observed isotope dependence of $\tau_E$. This discrepancy may be explained by the more peaked density and broader $T_i$ profiles of the $D^+$ plasmas compared with those of the $H^+$ plasmas at nearly the same $T_e$ shape yielding $\eta_i(D^+) < \eta_i(H^+)$ and $\chi_{\eta_i}(D^+) < \chi_{\eta_i}(H^+)$. In L-mode discharges $P_{ee}$ exceeds $P_{ci}$ and $\chi_e(L) > \chi_e(OH)$ holds at the same $n_e$ and increased temperatures. This can be seen by comparing Fig. 1 and Fig. 2, which shows the analysis results for a $H^+$ beam heated $D^+$ discharge. Only an $\chi_i > \chi_{CH}$ can explain the $T_i$ measurements, whereas a $\chi_i = \chi_{CH}$ would yield too high $T_i$ values.

3. Pellet fuelled discharges
Ohmic and co-beam heated pellet-fuelled discharges with strong density profile peaking exhibit a confinement improvement compared with GP fuelled discharges (doubling of $\tau_E$) [1, 2]. In the OH pellet discharges with a density peaking of $n_e(0)/n_e > 2.5$ the reason for the roll-over of $\tau_E$ is removed and the effect causing the ion transport enhancement has to be quenched: $\chi_i > \chi_{CH}$ would require an electron heat transport against $\nabla T_e$ in order to satisfy the power balance. With $\chi_i = \chi_{CH}$ during these pellet phases, however, a $\chi_e \sim 1/(nT_eq)$ at fixed $q$ results again. Global confinement is then governed by the electron heat transport. The confinement times for $D^+$ ($\leq 110$ ms) exceed those for $H^+$ ($\leq 160$ ms) considerably and a $\chi_e \sim A_i^{-\alpha}$ with $\alpha = 0.3 \div 0.7$ can be deduced [1].

4. Ctr-beam heated discharges
Ctr-injection in ASDEX leads to a doubling of $\Delta \beta_p$ due to NI compared with a comparable co-injection discharge and an improvement of $\tau_E$ up to 80 ms [3]. Confinement is gradually improving along with a continuously peaking of $n_e$ yielding $n_e(0)$ above $1 \cdot 10^{20} m^{-3}$ and a peaking factor $n_e(0)/n_e > 1.9$. The $T_e$ profile shape changes mainly in the central part, where it becomes hollow due to increasing radiation losses. $\eta_e$ and the calculated $\eta_i$ values are below 1 over 2/3 of the plasma radius (see Fig. 3a).
Again as in the pellet discharges $\chi_i = \chi_{CH}$ has to be assumed: with $\chi_i > \chi_{CH}$ the energy content of the plasma is underestimated even in the extreme of no additional electron transport. In this situation the neoclassical ion losses dominate over the electron heat losses ($T_i \approx T_e$). $\chi_e$ (see Fig. 3b) is strongly decreased compared with that of the co-NI case shown in Fig. 2, having nearly the same plasma current, but a somewhat higher heating power. Along with the improvement of the energy and particle confinement also the one of momentum is observed to increase with ctr-NI. The plasma rotation velocity is measured outside $a/2$; it increases throughout the ctr-beam phase up to $v_\phi(\frac{a}{2}) \approx 1.5 \cdot 10^8 m/s$. Assuming $v_\phi \approx (1 - r^2/a^2)^{1.3}$ the momentum confinement time at the end of the ctr-phase is $r_\phi = 90 ms$ and the momentum diffusivity is comparable to the electron heat diffusivity.

5. Summary
There are 4 regimes with peaked density profiles at ASDEX: pellet refuelled and ctr-NI discharges, OH discharges without sawteeth and those with reduced GP. The reason for the development of the peaked density profiles may be quite different and is not yet understood in all cases. But all regions have improved confinement which is partly offset by core radiation. Transport analysis - only performed for the first two regimes up to now - reveals that the ion transport has to be reduced in comparison to the broad density profile cases (OH-saturation and co-NI L and H-mode). $\eta_i$-modes may explain this result. Interestingly, in the small tokamak Pulsator the $\eta_i$-values are below 1 and the ion transport was consistently observed to be neoclassical.

References

Fig. 1: Radial profiles for a OH discharge ($I_p = 380 kA, B_t = 2.2 T; n_e = 4 \cdot 10^{19} m^{-3}$); a) $T_e(0)$ and $n_e(0)$ from Thomson scattering, $\eta_e$ and $T_i$ (calc. using $\chi_i = \chi_{NC} + \chi_{\eta_i}$); b) transport coefficients from TRANSP analysis.

Fig. 2: Radial profiles for a co-NI L-mode discharge ($I_p = 440 kA; B_t = 2.3 T; n_e = 4.5 \cdot 10^{19} m^{-3}, P_{NI} = 1.35 MW$); a) $T_e, T_i(0)$ from pass. CX-meas., $\eta_e, T_i$ (using $\chi_i = \chi_{NC} + \chi_{\eta_i}$) and $T_i$ ($\chi_i = \chi_{NC}$); b) transport coefficients from TRANSP analysis.
Fig. 3: Radial profiles for a ctr-NI high confinement discharge ($I_p = 420\,kA$, $B_t = 2.7$, $n_e = 7 \cdot 10^{19}\,m^{-3}$, $P_{NI} = 0.9\,MW$); a) $T_e(o), n_e(o)$ and $\eta_e$; b) heating ($P_{bi}, P_{be}, P_{OH}$) and loss ($P_{ce}, P_{ci}, P_{conv}$ and $P_{rad}$) power fluxes.

c) $x_i = x_{CH}$ and $x_e$ from TrANSP analysis.
COMPARISON OF CONFINEMENT IN HYDROGEN VERSUS DEUTERIUM IN MULTI-PELLET FUELED OH DISCHARGES IN ASDEX

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1. Introduction
In earlier investigations with multiple deuterium (D) pellet injection into ohmically heated D discharges (ASDEX tokamak, carbonized walls, \( I_p = 380 \text{ kA}, B_t = 2.2 \text{T}, q^* = 2.7 \)) a long-lasting period of peaked electron density profiles was obtained and the energy confinement time \( (r_E) \) was increased by a factor of nearly two in relation to discharges with gas puffing (GP) only (see Fig. 1). The analysis of energy confinement [1] showed that during this high confinement phase ion energy transport is the dominant energy loss mechanism in the center \( (r/a < 0.6 \text{ cm}) \) and revealed that the ion energy transport is neoclassical. Therefore, further improvement could be expected by changing from deuterium to hydrogen (H) if this situation were to hold. To investigate this question, the centrifugal pellet injector was modified to inject larger H and D pellets with \( 1.7 \times 10^{20} \text{ atoms each} \), compared to \( 5 \times 10^{19} \text{ atoms} \) in the small D pellets. At the same velocity of 600 m/s the typical penetration depths increased from \( 20 \div 25 \text{ cm} \) (small pellets) to \( 30 \div 35 \text{ cm} \) (large pellets). We describe first the density build-up and particle transport and then we compare the bulk plasma energy transport of both H and D pellet refuelled discharges.

2. Density build-up and bulk plasma particle transport
With both D and H pellet injection a strong increase of the line-averaged \( (n_e \) up to \( 1.2 \times 10^{20} \text{m}^{-3} \)) and volume-averaged \( (n_e > \) up to \( 8.5 \times 10^{20} \text{m}^{-3} \)) densities could be obtained. This is accompanied by a strong peaking of the density profiles with \( n_e(o)/n_e > \)-values up to 2.5 compared with 1.3 \( \div \) 1.5 for the GP phase. During this peaking the electron temperature \( (T_e) \) drops but the \( T_e \)-profile shape stays nearly self-similar \( (T_e(o)/T_e > \approx 1.9) \). This is shown in Fig. 2 where the time development of \( n_e \) and \( T_e \) are given for a discharge with large H pellets measured by Thomson scattering with a 16 channel YAG laser system, providing data very 17 ms. The measured data points are smoothed between adjacent pellet pop times and extrapolated to them. The peaked density profiles remain nearly stationary for times up to 230 ms (D) and 150 ms (H), respectively.

These strong density increase and profile peaking originate from 3 effects. Firstly, a fast transport process (time scale \( < 1 \text{ms} \)) can produce both a decrease of \( T_e(o) \)- keeping the temperature profile self-similar - and an increase of \( n_e(o) \) immediately after the pellet injection despite of the penetration depth being smaller than the plasma radius \( a \). Secondly, reduced sawtooth activity leads to a \( n_e \)-profile peaking due to the omission of an instantaneous particle outward flow during a sawtooth disruption. From the nearly
stationary and source-free density profiles of GP D plasmas one can deduce the ratio of an inward velocity $V_D$ and the particle diffusion coefficient $D_p$ by $V_D/D_p = -dn/dr/n$ which is increasing from $V_D/D_p = r/a^2$ (time averaged) with sawteeth to $3.5 r/a^2$ in a sawtooth free discharge. Only D discharges tend to loose sawteeth. Finally, reduced $D_p$ or increased $V_D$ values are observed during and after the pellet injection phases, where a $V_D/D_p$ up to $5.5 r/a^2$ can yield a further $n_e$ profile peaking [1]. With large H and D pellets the first effect seems to dominate during the peaking phase, but is certainly subsidized by the larger penetration (see Fig. 2a). The last two effects are the predominant cause for the peaking in the case of small D pellets and for the nearly stationary post pellet phases [1].

During the density build-up central radiation increases drastically. This results from the high $n_e(o)$ and from the increase of highly ionized metal impurities. While the total losses remain below 40% of the ohmic input power, the radiation loss dominates the central power balance and internal disruptions occur strongly reducing energy and particle content. $Z_{eff}$ obtained from absolute bremsstrahlung measurements and the loop voltage using neoclassical resistivity shows no peaking (the contribution due to the metal impurities is less than 0.1).

3. Energy confinement

H and D discharges with GP fuelling only show the roll-over from the linear dependence $\tau_E \sim \bar{n}_e$ to a saturated $\tau_E$ regime beyond $\bar{n}_e \approx 3 \times 10^{19} m^{-3}$. In the saturation region with “broad” density profiles $(n_e(o)/ < n_e >= 1.3 \div 1.5)$ the known reduction of energy confinement of H discharges $(\tau_E \approx 60 \div 70 ms)$ in relation to D discharges $(\tau_E \approx 80 \div 90 ms)$ holds. In the pellet discharges with successful density build-up and peaking the $\tau_E$'s could be improved to 110 ms (H) and 160 ms (D) at the same discharge conditions (see Fig. 1).

This suggests that the pellets remove the reason for the roll-over which can be explained by two alternatives. A profile consistency picture with fixed $T_e(r)/T_e(o)$ depending only on $q$ yields a $\tau_E$ saturation in GP discharges from a decrease of $T_e(o)$ and an increase of ohmic dissipation, whereas the pellet discharges gain in $\tau_E$ due to a more favourable weighting of $T_e(r)/T_e(o)$ with the high $n_e(o)$ values [1]. A “local” model for conductive energy transport was investigated with the TRANSP code. To describe the $\tau_E$ saturation we have to add to the neo-classical ion energy losses as given by CHANG-HINTON, $\chi_{CH}$, and to an anomalous electron heat conductivity, $\chi_e$, an additional heat conduction channel causing the $\tau_E$ saturation. CX measurements of $T_i$, measured $\beta_p$ and neutron productions guide as to describe this loss channel as an ion loss resulting in $\chi_i \approx (3 \div 4) \times \chi_{CH}$ and a $\chi_e \sim 1/(n_e T_e q)$. During the discharge phases with strongly peaked density profiles the effect causing the ion transport enhancement has to be quenched: $\chi_i > \chi_{CH}$ would require an electron heat transport toward the plasma center against $\nabla T_e$ in order to satisfy the power balance. With neo-classical ion losses $(\chi_i = \chi_{CH})$ during this pellet phases, however, GP and pellet results can be brought well into line with a continuous explanation of electron losses through $\chi_e \sim 1/(n_e T_e)$(fixed $q$).

The Alcator team has suggested first to attribute this additional ion loss to the ion temperature gradient mode triggered when the criterion $n_i = dlnT_i/dlnn_i = L_{ni}/L_{TI} > 1.5$ is fulfilled [2]. Figure 3a shows that due to the $n_e$ peaking $n_i \approx n_e (T_i \approx T_e$ at high $\bar{n}_e)$
decreases from above 1.8 at all radii in the GP phase to values below 1 over a large part of the plasma cross-section. Using a $\chi_i = \chi_{CH} + \chi_{\eta i}$ with $\chi_{\eta i} = 0$ for $\eta_i < 1$ and fully developed for $\eta_i \geq 1.8$ [3] the time development given in Fig. 3c for the $i$ and $e$ heat diffusivities at $r/a = 2/3$ of a H pellet discharge are obtained. While $\chi_e$ is only slightly decreasing with time, $\chi_i$ drops immediately after the first pellet from $(3/4)\chi_{NC}$ to $\chi_{NC}$. This model assumption is further supported by two facts. $\tau_E$ is at once improved when already the first pellet reduces $\eta_i$ to 1 (as is the case given in Fig. 3), but is only gradually rising when the $n_e$-peaking and $\eta_i$ reduction occur more slowly which happens especially with small pellets. Secondly the confinement improvement is only marginal when the $\eta_i < 1$ region is small.

The neoclassical ion loss in the high confinement pellet regime accounts still for the major part of the total non-radiated power flow within half the plasma radius. But the global energy confinement is governed by the anomalous electron heat transport in the outer plasma region with $\chi_e > \chi_i$. As $\chi_{NC}$ is the upper bound for $\chi_e$ one can calculate the radial dependence of $\chi_e$. Figure 4 compares $\chi_{CH}$ and $\chi_e$ both for a H and D post pellet phase at the same $q^* = 2.7$, having about equal $T_e$ profiles and the same $n_e$ profile shape. The $\chi_i$'s differ by the $\chi_{NC} \sim \sqrt{A_i}$ dependence (ion mass $A_i$), whereas $\chi_e(H)$ clearly exceeds $\chi_e(D)$. Taking into account the higher density of the D discharge a $\chi_e \sim A_i^{-\alpha}$ with $\alpha = 0.3 \div 0.7$ can be extracted; explaining also the inferior energy confinement times of H pellet discharges.

4. Summary

H and D ohmic heated pellet discharges show a density build-up (beyond the GP density limit) with strongly peaked $n_e$ profiles ($n_e(0) < n_e \leq 2.5$). They show in parallel a remarkable improvement of energy confinement by nearly a factor of two. With these peaked density profiles $\chi_i$ is restricted to $\chi_{CH}$. This allows a determination of $\chi_e \sim A_i^{-\alpha}$ ($\alpha \approx 0.5$) dominating the global energy confinement and reducing $\tau_E(D) \approx 160$ ms to $\tau_E(H) \approx 110$ ms.

References


Figure Captions

Fig. 1: Energy confinement time $\tau_E$ vs. $n_e$ for GP (points) and pellet fuelled hydrogen (H) and deuterium (D) single ohmic discharges.

Fig. 2: Axial ($r=0$) and volume averaged ($\langle \rangle$) $n_e$ and $T_e$ vs. time for a hydrogen pellet discharge (points are measured by Thomson scattering).

Fig. 3: a) $\eta_e = d\ln T_e/d\ln n_e$ at three radial positions,
   b) $\tau_E$, electron ($\tau_{Ee}$) and ion ($\tau_{Ei}$) confinement times and
   c) $\chi_e, \chi_i$ and $\chi_{CH}$ at $r/a = 2/3$ as a function of time for the discharge of Fig. 2.

Fig. 4: Radial dependence of $\chi_e$ and $\chi_i$ during post-pellet phase of a H and D ohmic discharge.
Fig. 1

\[ \tau_e \] [ms]

\[ \tilde{n}_e \] \left[ 10^{19} \text{ m}^{-3} \right]

Fig. 2

\[ \eta_e = \frac{L_{\tilde{n}_e}}{L_{T_e}} \]

\[ \tau \] [s]

Fig. 3

\[ X \] \left[ \text{m}^2 \text{s}^{-1} \right]

\[ \left( \frac{r^2}{a} \right)^{-3} \]

Fig. 4

\[ X \] \left[ \text{m}^2 \text{s}^{-1} \right]

\[ X_{CH} \]

\[ X_{D^+} \]

\[ X_{H^+} \]

\[ X_{e} \]
**Ze**ff-PROFILES IN DIFFERENT CONFINEMENT-AND HEATING REGIMES OF ASDEX

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**Introduction**

Ze**ff** is an important quantity in fusion research. It should ideally be as small as possible to minimize both fuel dilution and radiation losses. Further, radial profiles of Zeff, together with T_e(r), are needed to derive j(r) used in stability analysis. We report measurements of Zeff profiles for a variety of ASDEX plasmas: ohmic (H and D; gas and pellet fuelling), and co and counter neutral beam heated plasmas.

**Zeff measurements**

Ze**ff** profiles are measured from the intensity of plasma radiation in the near infrared. In this wavelength region recombination and line radiation are generally unimportant and the radiation is mainly bremsstrahlung. The intensity profiles are measured with the ASDEX 16 channel YAG laser Thomson scattering system (Fig. 1) /Ref. 1/. Each of the 16 detection boxes contains 3 broadband interference filters set at wavelengths of 850, 950 and 1000 nm. This provides a means of checking the \( \lambda^{-2} \) scaling of bremsstrahlung intensity and gives an approximate measure of the absence of line radiation. Up to the silicon avalanche diodes used as detectors, the system is identical to that used for Thomson scattering. For the Zeff measurements, however, a DC-coupled circuit with a bandwidth of 50 kHz is used. The sample rate is typically 200 Hz. Comparison of the Abel-inverted radiation with hydrogen bremsstrahlung determined from simultaneously measured n_e and T_e profiles results in Zeff profiles from +20 cm to -38.5 cm with a spatial resolution of about 3 cm. Under some conditions comparisons have been made with radially averaged Zeff values from plasma resistivity and from sawtooth analysis. The latter method is based on the electron power balance in ohmic discharges, in the plasma centre, during the period of recovery after the sawtooth crash /ref. 2/. In addition, Zeff has been estimated from impurity densities derived by two methods. Firstly, absolutely measured VUV line intensities of the four dominant impurities (Fe, Cu, O, C) are compared with those calculated by a time dependent impurity transport code using measured n_e and T_e profiles.
From this comparison an independent density profile is obtained for each impurity species, allowing its $Z_{\text{eff}}$ contribution to be calculated. Secondly, the densities of the two light impurity species (0, C) can be also derived from charge exchange recombination spectroscopy (CXRS) resulting in an independent measurement of their $Z_{\text{eff}}$ contribution. These spectroscopic $Z_{\text{eff}}$ are generally less accurate than bremsstrahlung $Z_{\text{eff}}$ but the resulting impurity densities yield considerable insight into the $Z_{\text{eff}}$ behaviour reported in this paper.

Results and Discussion

Axial values of $Z_{\text{eff}}$ as a function of $n_e$ for steady state ohmic and co-NI heated discharges are shown in Figs. 2 and 3. Normally, $Z_{\text{eff}}$ is found to decrease with increasing $n_e$, to be higher in $D_+$ than in $H_+$ plasmas, and higher in NI than in ohmic plasmas. The VUV $Z_{\text{eff}}$ are in good agreement with the bremsstrahlung $Z_{\text{eff}}$, as are the sawtooth $Z_{\text{eff}}$ for OH discharges in $H_+$. The CXRS $Z_{\text{eff}}$ must be normalized to one point from the bremsstrahlung measurements because optical components of the CXRS diagnostic, inside the vacuum chamber, have been inaccessible to direct calibration. The CXRS $Z_{\text{eff}}$ then show the same trends with $n_e$.

The above behaviour can be understood from the spectroscopic results. In ohmically and NI heated discharges O and C are the dominant light impurities with O giving the largest contribution to $Z_{\text{eff}}$. The VUV results for a transition from OH → co-NI heated plasma are shown in Fig. 4 for $H_+$ and Fig. 5 for $D_+$. The dominant heavy impurity in the $H_+$ discharge is Cu. It originates from the divertor plates and is more important at lower densities and in NI than in OH discharges. The densities of C and O (measured by VUV and CXRS) do not change strongly with $n_e$ so it is this fact, together with the decrease of Cu, that explains the decrease in $Z_{\text{eff}}$ with increasing $n_e$. Fe is unimportant in $H_+$ discharges due to carbonization, but gives a significant contribution to $Z_{\text{eff}}$ in $D_+$ discharges where the walls were not carbonized (Figs. 3 and 5).

Particularly interesting are the $Z_{\text{eff}}$ results obtained in three rather different scenarios in which improved energy confinement is found in ASDEX. These are OH discharges in $D_+$ with reduced gaspuff (Fig. 2) /Ref. 3/, pellet fuelled OH discharges (Fig. 6) /Ref. 4/, and counter-NI heated discharges (Fig. 7) /Ref. 5/.

In these cases there is a reversal of the trend of $Z_{\text{eff}}$ decreasing with increasing $n_e$ which shows that the improved energy confinement is accompanied by improved particle confinement. The strong accumulation of light impurities in counter NI discharges can be seen from the $O^{8+}$ (Fig. 7) and $C^{6+}$ densities. The relatively slow initial decrease of $n_O^{8+}$ and $n_C^{6+}$ leads to a decrease of $Z_{\text{eff}}$ as in the normal co-NI case. However, the impurity density increase is much steeper later on, so that $Z_{\text{eff}}$ increases during the phase when the global energy confinement is also increasing. For counter-NI the neoclassical $Z_{\text{eff}}$ is in good agreement with the bremsstrahlung $Z_{\text{eff}}$, and this is also true for the OH pellet case (Fig. 6). In the improved confinement OH regime $Z_{\text{eff}}(r)$ shows slight peaking at the plasma centre at later times (Fig. 8) but this is not seen in the pellet (Fig. 9) or counter-NI discharges.
Summary

The present measurements have shown a wide variety of behaviour of $Z_{\text{eff}}$ in ASDEX plasmas and $Z_{\text{eff}}$ values ranging from 1.2 to -5 have been found. In particular, for three different improved energy confinement regimes, the $Z_{\text{eff}}$ results indicate correspondingly improved particle confinement.

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Ref. 5 O. Gehre et al, Phys. Rev. Lett., to be published

Fig. 2: $Z_{\text{eff}}$ as a function of $n_e$ for OH-plasmas in hydrogen (H) and deuterium (D).

Fig. 3: $Z_{\text{eff}}$ as a function of $n_e$ for CO-NI heated plasmas in H and D.

Fig. 4: bremsstrahlung $Z_{\text{eff}}$ and the VUV imp. contributions during a transition from OH+co-NI in a hydrogen plasma.

Fig. 5: bremsstrahlung $Z_{\text{eff}}$ and the VUV imp. contributions during a transition from OH+co-NI in a deuterium plasma.
Fig. 6: Time dependence of $n_e$ (above) and $Z_{eff}$ (below) in an OH pellet refuelled discharge in hydrogen.

Fig. 7: Time dependence of $n_e$ and $n_0^{8+}$ (above) and $Z_{eff}$ (below) in a counter NI heated discharge in deuterium.

Fig. 8: Radial profiles of $Z_{eff}$ for OH discharges in deuterium. (a) "normal" regime, (b) improved confinement regime ($n_e=5.6 \times 10^{13}$ cm$^{-3}$, see Fig. 2).

Fig. 9: Radial profile of $Z_{eff}$ for a pellet fuelled OH discharge in H (see Fig. 6).
q-PROFILE MEASUREMENTS IN THE CENTRAL PLASMA REGION OF ASDEX

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Abstract: Using the Lithium-Beam-Spectroscopy (LBS) technique /1/ q-profile measurements have been carried out in the central region of diverted, ohmic ASDEX plasmas over a range \( q(a) \approx 2.9-5.3 \) for \( n_e \approx 1 \times 10^{13} \text{cm}^{-3} \). The experimental results are consistent with a flat \( q(r) \) profile having \( q(0) \approx 0.06 \) and exhibiting a slight tendency towards lower \( q(0) \) for decreasing \( q(a) \).

Motivation: Previous LBS measurements /2/ of \( q(0) \) in ASDEX for \( q(a) \approx 3.3 \) gave \( q(0) \approx 1 \) during ohmic heating. On TEXTOR \( q(0) < 1 \) is measured over the range \( q(a) = 2.1-6.3 \) /3/, whereas on TEXT \( q(0) \) is determined to be related to \( q(a) \), having \( q(0) > 1 \) for large \( q(a) \) and vice versa.

In order to provide a broader data base for comparison with these findings, the LBS technique is used to investigate \( q(0) \) for a variety of \( q(a) \).

Experiment: The \( q(a) \) scan is performed in four discharge series whose parameters are listed in Table 1 along with the deduced \( q(0) \) values from neoclassical resistivity (NR) and LBS, as well as the sawtooth (ST) inversion radius \( r_{ST} \) from LBS, NR (here \( r_{q=1} \) is equated to \( r_{ST} \)) and electron-cyclotron-emission (ECE). Three series have two \( I_p \) plateaus each, thereby enabling the relative \( \Delta q(r) \) to be monitored more precisely since each plateau has a common measurement base line. The fourth series features a radial displacement of the plasma (R=166.5-163 cm) over 0.8 sec to achieve a moderate radial scan within one shot. Except for \( q(a) = 5.3 \), where it is not certain, ST are present in all series with \( \Delta T_{e0}/T_{e0} \approx 5-8\% \).

Series #1 is with hydrogen, the others with deuterium. Based on NR, \( Z_{eff} \approx 4-6 \).

The density \( n_e \) of \( \approx 1 \times 10^{13} \text{cm}^{-3} \), plasma currents and toroidal fields chosen represent a balance between optimizing the LBS signal (low \( n_e \), high \( B_T \)) and the hindering of runaways (high \( n_e \), low \( I_p \)). Based on the gathered experience a larger \( q(a) \) range than covered here is accessible.

Results: The LBS technique measures the local magnetic field pitch angle \( \theta_p = \tan^{-1}(B_P/B_T) \). Figure 1 illustrates the noise level and temporal behavior of \( \theta_p(t) \) for series #1 at several values of the flux-surface radius \( r_f \). The corresponding pitch angle \( \theta^0_p \) profiles, adjusted to cylindrical geometry, are plotted in Fig. 4a; the values represent averages over 200 ms around the given time points, hence ST activity is completely averaged out. The representative error bar reflects the base line uncertainty (direction of \( B_T \)) - accounting for \(-2/3\) of the total - and of \( \theta^0_p \) itself.
Straight-line fits to the $\delta_p$ data in the central region, corresponding to a constant $q(=r_F/R_\tan \delta_p^o)$, yield $q(o)=-1.06$ and 0.95 for $t=1$ and 1.8 sec, respectively. Although neither Thomson scattering nor ECE profiles are available for corroborative evidence, the $r_F(q=1, t=1.8)$ ~15.2 cm value matches the $r_F(q=1)\geq 1+a/q(a)$ (1+40.2/2.93=14.7 cm) scaling of the other series. By the same account, $r_F(q=1)$ for $t=1$ sec should be ~8.6 cm, meaning that only two measurement points are inside the expected $q=1$ surface, which is not adequate to characterize $q(o)$.

The 2nd series was plagued with x-rays to the extent that no reasonably reliable $\theta_p$ measurement was possible. Nevertheless, the series serves to show that the $r_F(q=1)$ radius derived using neoclassical resistivity, under the assumption of a flat electric field- and $Z_{eff}$-profile, corresponds well with $r_{ST}$.

For the 3rd series, x-ray development again disturbed the diagnostic going into the second $I_p$ plateau. Fig. 3 illustrates that up to this point, $\theta_p$ changes little in the inner region. The $\delta_p^o (r_F, t=1.0)$ plot of Fig. 4b yields $q(r_F<9 \text{ cm})=1.01\pm 0.12$. For comparison, $\delta_p^o$ and $q$ derived from neoclassical and Spitzer conductivity are also presented.

The typical broadening of the $Te$ profile in response to a decrease in $q(a)$ is demonstrated in Fig. 3a. Figure 3b serves both to convey an impression of the error bars on the $Te$-based calculation of $\delta_p^o$ over the series, as well as to confirm that a perceptible change in $\theta_p$ vs $q(a)$ is to be expected only for $r_F>10$ cm.

The Fourth series produced only a very marginal scan in $r_F$ of ~0.6-1.1 cm due to the fact that $\theta_p$ is interrogated along a line inclined about 59° to the midplane. Nevertheless, $\delta_p^o (r_F>4 \text{ cm})$ in Fig. 4c is well documented: $q(r_F<11.2 \text{ cm})=1.0\pm 0.05$; $r_{ST}=12.1$ cm and $r_F(q=1)=12.9$ cm for neoclassical resistivity.

Discussion: Within this limited data base, varying $q(a)$ over 5.3-29 has ostensibly altered $q(o)$ from ~1.06 to 0.95. However, not enough radial points were present within the potential $q=1$ surface to convincingly describe $q(o)$ for $q(a)=5.3$, and at $q(a)=2.9$ the uncertainty in $\delta_p^o$ encloses $q(o)=1$. One might fault the linear $\delta_p^o$ data fit within $q=1$; notwithstanding, a close examination of such furnishes no compelling motivation to introduce any other algorithm in the central region. Within any one series, local excursions of $q(r)$ from the indicated value cannot be precluded, but may be regarded as unlikely when considering the overall direction of the results.

The experimental data are not consistent with the neoclassical prediction that $q$ continues to monotonically decrease within $r_{q=1}$. On the other hand, the better agreement with the Spitzer profile is probably specious: Quite systematically - over a wider parameter range than presented here - Spitzer resistivity fails to correctly give $r_{ST}$ and for low-$q(a)$, sawtooth discharges it often yields $q(o)>1$. In short, it is necessary to invoke neoclassical effects to approximately describe the experimental situation. The discrepancies within $r_{q=1}$ may be at least partially attributed to two effects: a) Due to the absence of points near the plasma center and the fit function chosen, the $Te$ profile is taken to be more peaked than in fact. (See Fig. 3a), b) The calculations assume a uniform E-field over the plasma cross section, which is known to be invalid in the presence of ST /5/.
Fig. 1: Temporal behavior of the measured pitch angle $\theta_p$, $n_e$ and $I_p$ for series #1.

Fig. 3: (a) Te profiles from the YAG Thomson scattering system at $t=1.0$ and 1.8 sec for series #3, (b) Pitch angle profiles derived from Te profiles averaged over the series, assuming neoclassical resistivity.
Table 1: Experimental Parameters and Results

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<th>(I_p) (kA)</th>
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<th>(q(o))</th>
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<th>ECE</th>
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<td>1.00±0.05</td>
<td>0.62</td>
<td>12.2</td>
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Fig. 4: Pitch angle profiles vs the flux surface radius \(r_f\) for: (a) first series, at \(t=1.0, 1.8\) sec near the end of each \(I_p\) plateau. The corresponding \(q(r_f)\) dependence derived from \(q(r_f)=r_f/R_t\eta\) is also given. (b) Third series, during the first plateau at \(t=1.0\) sec; Comparison among the measured \(\Theta_p^d\) points and \(\Theta_p^d\) profiles based on neoclassical and Spitzer resistivity and associated \(q\)-profiles. The arrow indicates \(r_{ST}\) from ECE. (c) Fourth series: • R-166.5, \(\Delta R-163.9\) cm.

References:

IMPROVEMENT OF BEAM-HEATED DISCHARGES BY REPETITIVE PELLET FUELLING IN ASDEX

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Introduction: Experiments with repetitive pellet injection were performed on the ASDEX tokamak to study additionally heated high-density divertor discharges using pellets of different sizes as a fuel source. This paper reports pellet-refuelled L- and H-mode deuterium discharges. In preceding investigations with multiple pellet injection, especially in ohmic heated discharges, a long-lasting period of peaked electron density profile and considerably improved global energy confinement were obtained. In recent experimentation L-mode plasmas could be created which show advanced plasma performance in relation to the common gas fuelled and neutral-beam-heated (NBI) discharges. These large pellet refuelled plasmas were characterized by moderately centrally peaked electron density profiles, high edge recycling, reduced sawtooth activity, central impurity radiation and significantly improved energy confinement. At a heating power of 2.5 MW the H-regime could be attained together with injection of small pellets.

Experimental Parameters: ASDEX is a tokamak with a major radius of 1.65 m and minor radius of 0.40 m. The poloidal plasma cross section exhibits nearly circular shape. The discharge parameters were typically \( B_t = 2.2T, I_p = 380 \text{kA} \) and \( q_a = 2.7 \). The line-averaged target electron density ranged from \( \bar{n}_e \sim 1 \times 10^{13} \text{cm}^{-3} \) to \( \bar{n}_e \sim 8 \times 10^{13} \text{cm}^{-3} \). The pellets with about \( 4.5 \times 10^{19} \) or alternatively \( 1.5 \times 10^{20} \) deuterium atoms each were accelerated by a centrifuge to a velocity of approximately \( 600 \text{m/s} \) [1] and yield penetration depths of roughly half the plasma radius. Normally up to 40 pellets were injected with a repetition rate of 30 ms. The additional gas puffing is not reduced during the pellet injection. Pellet ablation and penetration were monitored by photodiodes with \( D_\alpha/H_\alpha \) line filter and plasma photography. The power of the hydrogen beams is scanned from 0.35 MW to 3 MW in co-direction. When large pellets were applied, the heating power was technically limited to 1.35 MW. In typical cases of good confinement ASDEX was carbonized.

Experimental Observations: In earlier investigations combining injection of small pellets and strong NBI no remarkable density build-up could be produced [2]. Up to 70% of the injected pellet mass was missing in the discharge and nearly all the ablated mass left the plasma between two pellets (fig. 1a). The sawtooth activity increased and the energy confinement time \( (\tau_E = W_p/(P_{heat} - dW_p/dt)) \) degraded to values comparable to the gas puff case. The electron density and temperature profiles behaved also like in the gas puff L-regime. In recent experimentation with increased pellet size
and limited NBI power it was possible to improve significantly the plasma performance of the L-mode. The operational density range is extended to $n_e = 1.3 \times 10^{14} \text{cm}^{-3}$. During the heating phase the plasma stored energy from beta measurements increased by a factor of about 2 (fig. 1b) although still up to 50% of the measured pellet mass is lost during injection [3]. The situation with respect to the pellet ablation and penetration depths ($\sim 22 \text{ cm}$) corresponded to the ohmically heated discharges. Typically the sawteeth continued through the pellet injection but the period increased and the sawteeth lock to the pellets. When the pellet injection repetition rate is enlarged close to the sawtooth period, the sawtooth activity could nearly be suppressed (in OH discharges the sawtooth dynamic can vanish completely). Under these conditions maximum Murakami parameters ($M = n_e R / B_i$) of $10 \times 10^{16} \text{m}^{-2} \cdot \text{T}^{-1}$ were achieved. In parallel, the energy confinement time improved by approximately 40%. A weak density dependence as in the gas puff case and no saturation of the energy confinement were observed in the explored density range (fig. 2).

Pellet injection is able to generate strongly peaked electron density profiles $(n_e(0)/\bar{n}_e \approx 2)$. This effect indicates a change of the particle transport properties of the discharge, as seen with pellet injection into ohmic discharges. The electron temperature profile shape, on the other hand, exhibits no remarkable change compared to the pre-pellet phase. With increasing beam heating power, the electron density profile peaking in the L-mode becomes less prominent (fig. 3), and the energy confinement degrades (fig. 4) [4]. The confinement degrades much faster than the density profile peaking. At 1.35 MW neutral beam power the profile peaking is close to the standard gas refuelled L-mode discharge.

In the first H-regime experiments together with injection of small pellets (penetration depth of pellet $\sim 12 \text{ cm}$) the plasma performance is very similar to the gas puff case: density build-up takes place typically for the H-phase even without gas puffing and in between the pellet cycles. Starting at $\bar{n}_e = 8 \times 10^{13} \text{cm}^{-3}$, the density could be increased to $\bar{n}_e = 1.2 \times 10^{14} \text{cm}^{-3}$ by pellets and the intrinsic H-properties. The electron density profile showed the typical H-type shoulder and no pronounced profile peaking ($n_e(0)/\bar{n}_e = 1.25$). There was no sawtooth activity. Typical values of the energy confinement time were 70 ms at $\bar{n}_e = 5 \times 10^{13} \text{cm}^{-3}$ and 40 ms at the maximum density $\bar{n}_e = 1.2 \times 10^{14} \text{cm}^{-3}$.

Strong accumulation of high-Z impurities and central radiation (when sawteeth could be suppressed) were found in L- and H-shots with successful density build-up. The discharges often terminated through radiation collapse, in particular when Kr is puffed into the discharge to smother the sawteeth. Absolute bremsstrahlung measurements demonstrated that $Z_{eff}$ stays nearly constant at $\sim 1.5$ in most of the plasma cross-section during pellet injection [5], indicating that there is no low-Z accumulation.

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Figure Captions:

1: The density build-up of two pellet-refuelled and NBI-heated discharges is shown. In conjunction with the successful density build-up of discharge #21427 the global energy confinement time increases during pellet injection to ~ 65 ms. The increase of the diamagnetic beta (dashed-dotted curve) is also shown. When small pellets are injected (#18913), neither a high density nor improved confinement is attained.

2: Global energy confinement time as a function of the line-averaged electron density of pellet-refuelled (a) and gas-puff-refuelled (b) deuterium discharges at different NBI heating powers. The confinement time and electron density of pellet-refuelled discharges reach values which are considerably higher than those of the standard gas-puff case.

3: Peaking factor (ratio of the peak electron density to the volume-averaged electron density) as a function of the total plasma heating power $P_{\text{tot}}$ of pellet-refuelled discharges. This power is the sum of the ohmic input power $P_{\text{OH}}$ and the absorbed neutral-beam power: $P_{\text{tot}} = P_{\text{OH}} + 0.9 \times P_{\text{NBI}}$. The peaking factor at $P_{\text{tot}} = 0.5$ MW corresponds to pure ohmically heated discharges with injection of small pellets. When NBI power is applied, only discharges refuelled via large pellets are considered. The peaking factors at $P_{\text{tot}} = 0.5$ MW and 1.0 MW are determined during stationary density phases. The other values are ascertained close to stationary density conditions.

4: Global energy confinement time as a function of the total heating power $P_{\text{tot}}$ of pellet-refuelled discharges (see also caption of figure 3). All energy confinement times are determined at a line-averaged electron density of $\bar{n}_e = 1 \times 10^{14} \text{cm}^{-3}$. 

![Figure 1](image-url)
Fig. 2

Fig. 3

Fig. 4
MEASUREMENTS OF DENSITY TURBULENCE WITH FIR LASER
SCATTERING IN THE ASDEX TOKAMAK

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Previous measurements of density turbulence in ASDEX [1], [2] were limited to a wavenumber range \( k_\perp > 5 \text{ cm}^{-1} \) and to distances \( \leq 21 \text{ cm} \) of the measuring chords from the plasma centre. After modification of the scattering system the important range of lower \( k_\perp \) and the region near the separatrix are now accessible. First results of the investigations are presented.

**Scattering system:**

The scattering system using a 100 mW, 119 µm CW CH\textsubscript{3}OH laser and homodyne detection with a Schottky diode is shown in a schematic view in Fig.1.

![Schematic view of the optical setup](image)

**Fig. 1** Left: Schematic of the optical setup. Right: Poloidal section of ASDEX showing the beam paths inside the plasma vessel for the different chords and the window array. The scattered beams are indicated schematically by dashed lines.

The parameters of the scattering experiment are as follows:

- Beam waist in the plasma: \( w_0 = 1.6 \text{ cm} \)
- Wavenumber range: \( 2.5 \text{ cm}^{-1} \leq k_\perp \leq 25 \text{ cm}^{-1} \)
- Wavenumber resolution: \( \pm 1.25 \text{ cm}^{-1} \)
- Spatial resolution: \( \pm 1.6 \text{ cm} \) perpendicular to beam, chord averaged along line of sight

Accessible chords (distance measured from plasma center):
- Horizontal: 0 cm, 10.5 cm, 25 cm; vertical: 33 cm, 39.5 cm

Frequency analysis: spectrum analyser with fixed channels and continuous frequency sweeps in plateau phases

Wavenumber scan: possible within one plasma shot.

Spatial scan: different measuring chords can be chosen from shot to shot.
Ohmic discharges: Evidence for driftwave nature of the turbulence.

The following findings on ASDEX are consistent with the assumption of density gradient driven driftwave turbulence:

a) The rms value of the frequency integrated scattered power scales linearly with the mean electron density if the relative density profiles remain fairly similar. This was established for $n_e < 5 \times 10^{13} \text{ cm}^{-3}$ (where $T_e \propto n_n$ in ASDEX) in the important $k_L$ range and in different chords.

b) In the central chord which sees primarily poloidally propagating fluctuations a maximum of the scattered power is observed around $-100$ kHz in the dominant $k_L$ range. This is on the order of the diamagnetic drift frequency evaluated in the gradient region of the discharge. In the outer vertical chord which sees predominantly radially propagating fluctuations the frequency spectra are significantly narrower (Fig. 2).

c) The maximum of the frequency integrated $k_L$ spectrum shifts towards lower $k_L$ with increasing $T_e$ (Fig. 3). A value $k_L^{\text{max}} \cdot 0.3$ is inferred from the "cold" shots in Fig. 3. It should be kept in mind, however, that the shape of $P_S(k_L)$ and the shape of the fluctuation spectrum $n_e(k_L)$ are not identical. The latter results from an integration over all spatial Fourier components while $P_S(k_L)$ contains only those components which are selected by the scattering geometry.

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Fig. 2: Frequency spectra of scattered power in the central horizontal chord and outer vertical chord (D+ plasma)

Fig. 3: Wavenumber spectra in "hot" and "cold" ohmic hydrogen plasmas. The densities in the center are $1.75 \times 10^{13} \text{ cm}^{-3}$ and $4.8 \times 10^{13} \text{ cm}^{-3}$, respectively. Signals are normalized to maximum value. Chord 10.5 cm.

Fig. 4: Change of wavenumber spectra with gas filling. Note: Same vertical scale for both curves.
d) The maximum of the frequency integrated $k_L$ spectrum shifts towards lower $k_L$ and its value increases when the gas filling is changed from pure hydrogen to a $1:1$ mixture of hydrogen and deuterium at constant electron density (Fig. 4).

e) The frequency and wavenumber integrated scattered power decreases with increasing toroidal magnetic field at constant plasma current.

A detailed analysis is beyond the scope of this summary. It should be noted that unambiguous experimental tests of theoretical models for drift wave turbulence would require pure one-parameter scans which in general are difficult to realize.

L-phase with neutral beam injection (NI):

NI heating produces complex changes of the density turbulence. A dramatic broadening of the frequency spectra with respect to the ohmic phase occurs as is illustrated in Fig. 5. The temporal development of the frequency integrated scattering signal was recorded at the transition from an ohmic to an L-phase, the density profiles remaining essentially unchanged (Fig. 6). The $k_L$ spectrum is shifted towards longer wavelengths with NI, and the time behaviour is different for different $k_L$. It is an open question whether the change in the spectra is solely due to the increase in $T_e$ or due to a change in nature of the turbulence induced by NI.

As can be seen in Fig. 6 sawtooth activity during NI strongly affects the scattering signals at low $k_L$. A clear difference in arrival time and shape of sawteeth on the scattering signals is observed in the different chords (not shown in the figure).

Fig. 5: Change of frequency spectra with NI-heating (L-shot; chord 0 cm)  
Fig. 6: Change of frequency integrated scattering signals for different $k_L$ in a series of identical L-shots.
Transition from the L into the burst-free H-phase.

By shifting the horizontal plasma position by a few cm the region inside and outside the separatrix can be scanned with the outer vertical chord. In a series of L-H transitions the behaviour of the fluctuations was investigated at distances of -0.5 cm outside and -1 cm inside the separatrix. In both cases the scattering signals decrease sharply at the time of the transition (Fig. 7). This behaviour cannot be explained by a drop in density and/or decrease in the density gradient length, as can be seen from Fig. 8. It shows the density profiles obtained from the Li beam probe before and after the transition as well as the position and radial extent of the FIR laser beam with respect to the separatrix radius $r_s$.

Further investigations are necessary in order to obtain a conclusive picture of the fluctuation behaviour at the L-H transition.

![Fig.7](image1.png)

**Fig.7:** Density fluctuation signals measured close to the separatrix in different frequency channels at the transition from an L-phase into a burstfree H-phase (Outer vertical chord; lower traces: $D_x$ monitor and line electron density).

![Fig.8](image2.png)

**Fig.8:** Electron density profiles in the separatrix region before and shortly after the L-H transition. The position and width ($1/e^2$ of intensity) of the laser beam relative to the separatrix are marked with vertical lines.

References

MAGNETIC FIELD LINE TRACING IN T-10 TOKAMAK

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There are several reasons rousing interest for observations of ablation clouds near pellets injected in tokamak plasma. Firstly, the clouds elongate in the magnetic field line direction [1, 2] and this can be used for studying magnetic fields in tokamaks. Secondly, it is necessary to know light emissivity and particle density distributions in the cloud for examination and advancement of pellet ablation models [3, 4]. Finally, in order to indicate the runaway electrons in plasma and to determine an asymmetry of electron velocity distribution function one can use the observation of pellet trajectory [5].

The results of fast photography of pellets injected into the T-10 plasma are presented in this paper. Some new information (in comparison with [2]) was obtained using an automatic image processing system including the AMD-1 microdensitometer. It made possible the determination of the rotational transform angle $\Theta (r)$ with accuracy of $5 - 10$ mrad, scale length in poloidal ($l_\perp$) and toroidal ($l_{||}$) directions and the deflection of the electron velocity distribution function from Maxwellian.

In experiments carbon pellets with diameter $d_p = 300 - 400$ $\mu$m and velocity $v_p \sim 140$ m/s were injected during a stationary phase of the discharge in the direction of plasma core with the angle $\alpha = 30^\circ$ relative to an equatorial plane. Photographs of the ablation clouds were taken in spectral range 390-600 nm with a period of 25 - 50 $\mu$s using the VSK-5 camera.

Fig. 1 presents a typical densitogrames of the ablation cloud. Fig. 2 shows the cloud light distributions in toroidal
Z (solid) and poloidal Y (dashed) directions. It is seen from Fig. 2 that the emissivity spatial decay in poloidal direction is close to exponential. The radial profiles of scale length $l_{\perp}(r)$, $l_{\parallel}(r)$ calculated from $e$-decrease of radiation in the middle part of the curves are shown in Fig. 3. The value of $l_{\perp} \sim 1 - 1.5$ mm changes weakly during pellet ablation and is close to ionization length of carbon atoms $l_{\perp} = v_s/(n_e \langle \delta v \rangle_{\text{ion}})$ that leave the pellet surface heated to temperature $T_s = 4500 - 5000$ K [3] with the velocity $v_s = \sqrt{\gamma T_s/m_e}$. One can easily see it using Lotz's ionization coefficient data [6] ($\langle \delta v \rangle_{\text{ion}} \sim 10^{-7}$ cm$^3$/s for $T_e \geq 100$ eV). The calculated values $l_{\perp}$ for density profile $n_e(r) = 3 \times 10^{13} \left[ 1 - (r/38)^2 \right]$ cm$^{-3}$ are presented on Fig. 3 by crosses. The nearly proportional decrease of $l_{\perp}$ with increasing of plasma density $n_e$ was observed in impulses with different density. The data on $l_{\perp}$ are in agreement with the carbon ablation model assumptions [3]. The $l_{\parallel}$ dependence on minor radius qualitatively corresponds to ablation rate $\dot{N}(r)$ (see Fig. 3). But, there is still no clear understanding of the mechanism responsible for the light distribution along a magnetic field line. Perhaps, the growth of $l_{\parallel}$ with $\dot{N}$ may be explained by carbon ions light emission. However, this process can not describe the decrease of $l_{\parallel}$ at the final evaporation stage. Another cause of such behaviour of $l_{\parallel}(r)$ may be a radiation emitted by carbon atoms originating due to recombination from ions of cold and dense secondary plasma flowing away from the pellet along magnetic field line. It is necessary to note that the last mechanism can explain qualitative agreement in behaviour of $l_{\parallel}(r)$ and $\dot{N}(r)$.

A cloud elongation observed made it possible to determine the rotational angle $\theta = B_{\text{pol}}/B_{\text{tor}}$. One can see from Fig. 1 that the sign of angle $\theta$ changes when the toroidal magnetic field direction changes to the opposite. The direc-
tion of the magnetic field was supposed to be the same as the cloud axis with minimum moment of inertia. We supposed a constancy of optical density as well. The value and accuracy of $\Theta (r)$ was calculated using the least squares method. The absolute magnitude of $\Theta$ in the limiter region was obtained by fitting it with calculated value. The profile $\Theta (r)$ for regime with low safety factor $q(a_1) = 2$ ($I_p = 285$ kA, $B_{tor} = 2$ T, $n_e = 2.10^{13}$ cm$^{-3}$, $a_1 = 28$ cm) is shown on Fig.4. The highest accuracy takes place in the phase of intensive ablation when the elongation is large enough.

The fact that the carbon pellet evaporation begins at $r/a_1 > 1$ indicates the shift of plasma column up to 5 cm in the outer direction. In contradiction to early data for $q(a_1) = 3,5$ [2] the presented profile $\Theta (r)$ shows that a valuable part of plasma current flows in the peripherical region. The current density distribution $j(r)$ is consistent enough with profile $j(\varphi) = j(0) [1 - (\varphi/a_1)^2]$, where $\varphi$ - magnetic surface radius.

The ultrafast photography allowed to determine that the ablation is more intensive from the electron current side of pellet. The corresponding pellet deflection in toroidal direction ($\sim$5 cm) points at this fact. The value ($v_s \sim 2.10^5$ cm/s) of velocity of sound is known well enough. So we were able to estimate absolute values of heat flow $Q_f$ transported by asymmetrical part of electron distribution function more correctly than Andersen [5]. In the discharge with parameters $I_p = 230$ kA, $B_{tor} = 2$ T, $n_e = 10^{13}$ cm$^{-3}$, $a_1 = 34$ cm the $Q_f$ value in the vicinity $r \sim 20$ cm reaches $3.10^5$ J/cm$^2$s which makes up $10 - 15\%$ of the full heat flow estimated using carbon ablation rate [3].
\[ S_{\text{hot} = 43116}: I_p = 230 \text{ kA}, B_{\text{tor}} = 2 \text{T} \]
\[ n_e = 10^{13} \text{ cm}^{-3}; a_L = 34 \text{ cm} \]

\[ S_{\text{hot} = 43196}: I_p = 285 \text{ kA}, B_{\text{tor}} = 2.1 \text{T}; \]
\[ n_e = 2 \times 10^{13} \text{ cm}^{-3}; a_L = 28 \text{ cm} \]

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ELECTRON TEMPERATURE PROFILE CONSISTENCY UNDER ECRH IN T-10 TOKAMAK
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Abstract. A particle and energy flux model based on the idea of canonic pressure or electron temperature profiles is proposed. The calculations for the T-10 tokamak have been done, It is possible to represent reasonably the on-axis and mixed ECRH, heat wave propagation and the possible transition from the L-mode to the H-one with a given model.

1. The experiments on on-axis and off-axis ECRH on T-10 show a strong dependence between the effective local heat conduction coefficient, \( \mathcal{L}_{\text{eff}} = \frac{1}{r} \frac{\partial}{\partial r} \left( \frac{1}{r} \mathcal{L} \right) \), and the profile of deposited ECRH power (\( \Gamma \) and \( T \) are the heat flux and electron temperature, respectively) \([1,2]\). A simple quasilinear model based on the inclusion of "hard" terms representing large parameters into the equations is adopted in a given paper for the description of profile consistency. These terms are different from zero only at deviation of the profiles from the canonical ones.

2. Model. A set of equations for energy and particle balance has the form

\( \frac{\partial n}{\partial t} = - \frac{1}{\varepsilon} \frac{\partial}{\partial r} \left( \varepsilon \frac{\partial n}{\partial r} \right) + P_n \) \hspace{1cm} (1)

\( \frac{3}{2} \frac{\partial}{\partial t} \left( n T_e \right) = - \frac{1}{\varepsilon} \frac{\partial}{\partial r} \left( \varepsilon \frac{\partial T_e}{\partial r} \right) + Q_{\text{OH}} + Q_{\text{EC}} + Q_{\epsilon i} + Q_{\text{rad}} \) \hspace{1cm} (2)

\( \frac{\partial}{\partial t} \left( \frac{\partial n}{\partial r} \right) = \frac{1}{2} \frac{\partial}{\partial r} \left( \frac{1}{\varepsilon} \frac{\partial}{\partial r} \left( \frac{\partial n}{\partial r} \right) \right) \) \hspace{1cm} (3)

Here \( n(n,r,t) \) is the plasma density, \( \mathcal{L} = 1/q, Q_{\text{OH}} \) is the Ohmic power, \( Q_{\epsilon i} \) is the power of exchange between ions and electrons. The ion temperature, \( T_i(n,r,t) \) the ECRH power \( Q_{\text{EC}}(n,r,t) \) and the radiation power \( Q_{\text{rad}}(n,r,t) \) are assumed to be known from the experiment.
Let us represent the particle and heat fluxes, $\Gamma_n$ and $\Gamma_T$, as sums:

$$\Gamma = \Gamma_n + \Gamma_T, \quad \Gamma_n = \Gamma_{n,\text{no}} + \Gamma_{n,\text{an}}, \quad \Gamma_T = \Gamma_{T,\text{no}} + \Gamma_{T,\text{pc}}$$  \hspace{1cm} (4)$$

Here the subscripts "no" and "an" designate neoclassical and anomalous fluxes. We assume:

$$\Gamma_{n,\text{no}} = - D_{n,\text{no}} \left( \frac{\partial \rho_n}{\partial \tau} \right) T_e, \quad \Gamma_{n,\text{an}} = - D_{n,\text{an}} \left( \frac{\partial \rho_n}{\partial \tau} \right) T_e, \quad \Gamma_{T,\text{no}} = - D_{T,\text{no}} \left( \frac{\partial \rho_T}{\partial \tau} \right) T_e,$$

where $D_{n,\text{no}} = D_{n,\text{an}} = D_{T,\text{no}} = D_{T,\text{pc}}$ are constant factors of the order of unity, $\Gamma$ in (6) is the average density times $10^{13}$ cm$^{-3}$. The form factor $F(r)$ ($0 < F(r) < 1$) represents a change in hardness across the plasma cross-section. The parameter $K_p = K_p(r)$ is chosen so that $\Gamma_n$ and $\Gamma_T$ become zero at the canonical pressure profile. Choosing a canonical profile, according to B.B. Kadomtsev [3], one has

$$r_p = r_0 \left( \frac{1}{c_T^2} \right), \quad c_T = \sqrt{c_n c_e} - q,$$  \hspace{1cm} (6)

Instead of (2)-(6) we can use a model based on a current canonical profile. In this case the canonical profile $T_{e0}(r)$ exists for $T_e(r)$ and we can write

$$\Gamma_{T,\text{pc}} = - \chi T_{n,\text{no}} \left( \frac{\partial \rho_T}{\partial \tau} \right) T_e, \quad \Gamma_{T,\text{an}} = \chi T_{n,\text{an}} \left( \frac{\partial \rho_T}{\partial \tau} \right) T_e,$$

where $\chi = \left( \frac{3 \rho_n^2}{1 + \rho_n^2} \right)$. So a more simple model consists of equations (2)-(3), (7) [4]. In this case, the density, $n(r,t)$, should also be taken from the experiment.

3. NCR-heating. The experimental results (solid lines) and the calculations based on the model (2)-(3), (7) (dashed lines) for two shots with NCRH are given in Fig. 1. Here $I=200$ kA, $n_0 = 3 \times 10^{13}$ cm$^{-3}$. In the shot No. 45439 the on-axis heating was performed (3 gyrotrons with total power, $P_{3G}=630$ kW). In the shot No. 45443 one gyrotron heated plasma on the axis (200 kW)
and three gyrotrons, at the radius 17 cm (650 kW). The power deposited within the region $r < 12$ cm varies more than three times. Nevertheless, (Fig. 1) $T_e$ and $\frac{\partial n_e}{\partial z}$ have a slight difference. The $\bar{\gamma}_c$ -flux structure (7) and a great value of hardness parameter $\bar{\gamma}_c$ allow one to represent reasonably the experiment. The profiles $H_{\text{eff}}(r)$ in these shots are given in Fig.2. $H_{\text{eff}}$ at the centre varies almost by the factor of four.

4. Transition to the H-mode. In the examples under consideration $T_e(a)$ is small so $\left| \frac{\partial n_e}{\partial z} \right| > \bar{\gamma}_c T_e$ at the plasma edge and the relative deviation of $T_e(r)$ from $T_{e0}(r)$ is great. As a result, $H_{\text{eff}}$ has a strong rise, approaching the plasma edge (Fig.2). Such a behaviour of $H_{\text{eff}}$ is characteristic for the L-mode. The model (2)-(3),(7) allows one to represent the transition to the H-mode in the following way. Let us have an opportunity to increase the edge temperature, $T_e(a)$. In this case, the profile $T_e(r)$ approaches the canonic one, and $H_{\text{eff}}$ is reduced. The dependences of $H_{\text{eff}}$ on the radius at various edge temperatures, $T_e(a)$, for OH ($\bar{n} = 1.9$) and for ECRH ($\bar{n} = 2.4$, on-axis heating at $P_{\text{ECR}} = 2$ MW) in the plasma with $I_0 = 200$ kA are shown in Fig.3. One can see an abrupt deep in $H_{\text{eff}}$ across the whole plasma column, when $T_e(a)$ rises. Such a behaviour of the plasma corresponds to the transition to the H-mode.

5. Propagation of heat waves. The heat wave profile has a strong difference from the canonic one. Therefore $H_{\text{eff}}$ for a heat wave can noticeably exceed $H_{\text{eff}}$ for a smooth, quasi-stationary $T_e(r)$ - distribution. Let us consider the simulation of a heat wave emerging under MHD - mixing with a set (2)-(3),(7) for the regime with the parameters: $I = 175$ kA, $\bar{n} = 3$, $a = 34$ cm, $\bar{n}_a \sim 20$. The experimental and calculated dependences of electron temperature, $T_e(r,t)$, at $r = 6.5$ cm under OH, are given in Fig.4. The curve (1) corresponds to the experiment, other curves are calcula-
An experimental $\chi^{\text{PB}}$, found from the local energy balance is used to build-up the curve (2). The curves (3) and (4) are obtained with the model (7) at $\alpha = 0.5$ and $\alpha = 1$, respectively, $F = 1$. To make the model more precise we assume $F = F_0(r)$, where $F(r) = 0$ at $r < 6$ cm, $0 < F(r) < 1$ at $6$ cm $< r < 8$ cm, $F(r) = 1$ at $r > 8$ cm (dashed line). In this case, $\chi_{\text{eff}} \approx 5 \chi^{\text{PB}} \approx 2 \times 10^{17}$ cm$^{-1}$s$^{-1}$ for a heat wave.

EXPERIMENTAL AND NUMERICAL STUDY OF SAWTEETH ON T-10

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ABSTRACT. The most probable mechanism of sawteeth on T-10 is ergodization of magnetic field.

Results of experimental study of sawteeth on different tokamaks are dramatically different and are often in contradiction with one another (see for example References in /1/). One possible reason is simultaneous excitation of different instabilities during a crash, one of which plays the main role. The leading role of a particular mode depends on the current distribution (the shear) in the plasma core, according to modern theories.

At low shear in the core, \( r < r_1 \), \( r_1 \) being radius of the \( q=1 \) surface, \( \Delta q = 1 - q(0) \approx 0.05 \) the ideal internal mode is excited; this is the case for large tokamaks. At higher shear (0.8 \( \leq q(0) \leq 1 \)), the internal resistive mode \( m=1, n=1 \) is dominant and is localized near the separatrix of a magnetic island. At \( q(0) \leq 0.8 \), stochasticization occurs across the whole cross-section of the \( m=1, n=1 \) island and the sawtooth itself is the process of magnetic field ergodization /2-7/.

In all these cases the energy source of the instability is energy of the poloidal magnetic field. The question is: what serves as a trigger for the crash? For T-10 /1/, the value of \( \gamma_T = \frac{d \ln T_e}{dr} \) at \( r = r_s \) plays such a role, \( r_s \) being the phase inversion radius. As \( \gamma_T \sim \gamma_j (= \frac{d \ln j}{dr} \) at \( r_s \) during the stationary stage, one can suppose that \( \gamma_j \) at the \( q=1 \) surface is the critical one for the crash start (radius of the \( q=1 \) surface is believed to be equal to \( r_s \)).

The machines of T-10 size, such as TFR and ORMAK, have the same \( \gamma_T \) immediately before the crash (Fig.1), \( \gamma_T \) being an extremely weak function of discharge conditions. Change in the tokamak size leads to change in \( \gamma \text{crit} \), which is lower for
larger machines (TFTR and JET). It is important that shear in
the core, \( \sigma_0 \), is lower in these machines than in T-10. This
allows us to suppose that \( \gamma_{\text{crit}} \) decreases with decreasing \( \sigma_0 \).
This is confirmed by results from TEXTOR /5/.

Experiments on T-10 have been carried out to examine the
possibility for transition from one mechanism of the internal
disruption to another by changing \( \sigma_0 \). Variation in the \( \sigma_0 \)
value has been obtained by changing the position of the movable
graphite limiter with \( q_L < 0.0 \). Decreasing the \( a_L \) value without
changing other parameters has led to a sharpening of the \( T_e \)
profile and to an increase of the \( r_S \) value (Fig. 2, \( B_t = 2.6 \) T,
\( n_e = 3.10^{13} \) cm\(^{-3} \), \( I_p, kA: 305 (\circ), 260 (\square), \) and \( 240 (\Delta) \)).
The \( r_S \) increase points to peaking of the \( j \) profile. In the
305 kA regime the shear at the \( r_S \) surface was about 0.4 \( (q(0) = 0.78) \) when calculated from the \( T_e \) profile at \( a_L = 32 - 34.5 \)
cm (conductivity was supposed to be of Spitzer type, \( Z_{e'}(r) \)
and \( E(r) \) to be constant). Variation of shear has been calcu-
lated from the \( r_S \) variation under the supposition that \( r_S = r_1 \)
and that distribution of the "additional" current is homogene-
ous within the \( r_S \) surface. The \( s_0 \) variation was about 20% when
\( a_L \) was decreased from 34.5 to 28 cm (Fig. 2, points a, b, and
c). One can see on Fig. 2 that the temporal evolution of saw-
teeth did not change with such a variation (see also Table 1
where \( T_S \) is the sawtooth period). The characteristic time of
the crash, \( t_{\text{crash}} \), also did not change within the experimental
accuracy (\( \delta t = 0.1 \) ms). However, this result does not mean
that the transition is not possible, because the change in
shear may be not large enough.

High shear, which is believed to be the case for T-10 pla-
smacore, allows us to suppose that the stochasticity mecha-
nism of /2/ is responsible for the crash in T-10. Relatively
high depth of ripples in T-10 ( \( \sim 1\% \) at \( r_S \)) results in exten-
tion of the stochasticity region near the separatrix of the
\( m=1, n=1 \) island and provides the stochasticity mechanism in
T-10 /2/. This is foundation for our numerical model of saw-
teeth /6/.
In the simulation based on the Dnestrovskii-Kostomarov transport code and the Lichtenberg stochasticity theory /2/, we assumed a fast increase in the heat conduction coefficient $K_e$ within the $r_S$ surface during the crash /6/. In addition to simulation of high-q regimes with $q_L = 4 - 6$ /6/, we presented now the results of simulation of the $q_L = 2$ regime (Fig.3). Rectangles on Fig.3 illustrate the process of propagation of perturbed region with $K_e = K_e^0 + K_e^{turb}$, $K_e^{turb}$ being of order $10^3 K_e$. The simulation reproduces well the step-like character of the perturbation propagation towards the center with velocity about $5 \times 10^5$ cm.s$^{-1}$ and enhanced velocity of the heat pulse propagation beyond the $r_S$ surface. Both the $r_S$ value and period of sawtooth also agree well with the experimental ones. The $j$ profile change is extremely small in the simulation which agrees with unchanging position of the $m=1$, $n=1$ island during the crash in the T-10 experiment.

In summary: (i) internal disruption in the T-10 plasma can be described by the stochasticity model of /2/; (ii) the $\gamma_T$ value before the crash depends on tokamak size and probably on shear in plasma core; (iii) small variation in the shear value ($\sim 20\%$) on T-10 does not result in any significant change of the sawtooth features.

**TABLE 1. PARAMETERS OF SAWTEETH AT DIFFERENT $a_L$.**

<table>
<thead>
<tr>
<th>$a_L$, cm</th>
<th>28</th>
<th>30</th>
<th>32 - 34.5</th>
</tr>
</thead>
<tbody>
<tr>
<td>$r_S$, cm</td>
<td>9.9</td>
<td>9</td>
<td>8.2</td>
</tr>
<tr>
<td>$\tau_S$, ms</td>
<td>11 - 12</td>
<td>11 - 12</td>
<td>11 - 12</td>
</tr>
<tr>
<td>$t_{crash}$, ms</td>
<td>0.32</td>
<td>0.40</td>
<td>0.33</td>
</tr>
</tbody>
</table>

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IMPUlRY TRANSPORT STUDY IN B AND S REGIMES ON THE T-10 TOKAMAK

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A given paper describes the study of plasma parameters in the regime with enhanced impurity confinement (B-regime) and in the regime with reduced impurity confinement (S-regime) in the T-10 tokamak [1]. Transition from the S-regime to the B-regime was performed by switching a working gas puff valve off. The T-10 chamber was carbonized a few times in the process of these experiments. In operation after carbonization, the chamber wall was saturated with neutral atoms of hydrogen and deuterium; the main plasma impurity was carbon [2]. Using different programming regimes of the working gas puff valve in shots between carbonizations, we were able to carry out our experiments with wide variations in the influx of neutral atoms from the chamber walls and from gas valve.

Changing the influx of neutral atoms in different ways, we obtained several modifications of discharge evolution in time at the same integral plasma parameters: \( I_p = 200 \text{ kA} \), \( r_a = 32 \text{ cm} \), \( \bar{n}_e = 3.5 \times 10^{-13} \text{ cm}^{-3} \), \( B_T = 30 \text{ kG} \) (see Fig. 1). After switching the gas puff valve off, the influx of neutrals abruptly drops and the discharge transits to the B-phase. The transition is accompanied by a 4-5 times rise in the SXR-signal. This rise is produced besides by a rise in \( T_e(0) \), by an increase in the life-times of impurities (confirmed by the decay time of injected potassium: \( \tau_k \sim 100 \text{ ms} \) (see Fig. 1)). In a given example after the gas puff cut-off, the influx of neutrals was retained at a low level so that \( \bar{n}_e \) = const. If the neutral influx from the gas valve or from the chamber walls remains rather high, the discharge will transit to the S-regime. In this case, the SXR-signal rises insignificantly and the life-time of impurities is short (\( \tau_k \sim 50 \text{ ms} \) (given by the dashed line in Fig. 1a).

The further evolution of the discharge after transition to the B-phase depends on the programming of the gas-puff (Fig. 1b-c). In the case of an insignificant influx of neutral atoms from the chamber walls and completely-closed gas puff valve, \( \bar{n}_e \) drops (Fig. 1b), and the SXR-signal is reduced due to this drop. The operation continues in the B-phase. However, if the valve is switched on again in the B-phase (Fig. 1c), \( \bar{n}_e \) rises, and an additional high neutral influx provides the conditions for the transition to the S-phase. The process is accompanied by a 2.5 - time drop in the SXR-signal. A drop in \( T_e(0) \) and a rise in \( n_e(0) \) in this case cannot completely explain the
change in the SXR-signal. It means that there is a reduction in the concentration of impurities in the plasma. Thus, the repeated switching of the valve on provided the transition of the discharge to the S-phase.

The behaviour of $D_\alpha$-line intensity near the chamber wall and near the limiter in the case of switching the gas valve off is of a weak neutral influx from the chamber walls (Fig. 1b) is given in Fig. 2. After gas puff cut-off, the discharge transitions to the B-phase and then a gradual plasma column expansion for about 120 ms occurs up to a time, when the plasma edge touches the limiter. At this time the influx of neutral atoms from the chamber walls and from the limiter steeply rises; however, $n_e$ and $\int n_e \, dx$ drop, as seen in Fig. 2. Such a change in the integral of particles in the plasma column can be explained only by, at least a factor of 2 reduction in the life-time of particles $\tau_p$. Thus, one can see that the plasma column evolution after switching the gas valve off is non-stationary, and it is difficult to obtain the stationary B-regime by variation in the neutral influx.

An important role of the neutral atomic influx in obtaining the B-regime is confirmed by a series of ten successive shots in the tokamak. The measurements were done after the ECRH experiments at $n_e \sim 2 \times 10^{13}\text{cm}^{-3}$, as a result of which the chamber walls were purified of hydrogen and deuterium atoms. The neutral influx was at a low level, and the gas valve switching-off provided the transition to the B-regime (Fig. 1a, solid line). In the process of further tokamak operation at a higher $n_e = 3.5 \times 10^{13}\text{cm}^{-3}$ without ECRH the walls were saturated with atoms of working gas from shot to shot, and the level of neutral influx was raised. The growth of periphery plasma density and the drop of central temperature also took place. The level of neutral influx gradually became so high that the transition to the B-regime became impossible (Fig. 1a, dashed line) the evolution of the $n_e(r)$ - profile is shown in Fig. 3a. One can see that the density profile for the shot in the B-regime is more peaked than that in the S-regime. The density profiles for discharge with Ne-puffing and without it are shown in Fig. 3b. In the case of Ne-puffing, the transition to the B-regime is observed. One can see that a change in the profile is close to that at the gas valve switch-off.

The effect of transition from the B-regime to the S-regime can be achieved by a pulsed gas puff at the plasma edge. An example is shown in Fig. 4: a single puff of He into the discharge in the B-regime results in an abrupt transition to the S-phase. This transition is accompanied by rearrangement of the
A reduction in the parameter $\lambda_{ne}$ undergoes the most significant change in the range $10 \text{ cm} < r < 20 \text{ cm}$ and as a result the parameter $\zeta_e$ is changed from $\eta_e = 3$ in the $S$-regime to $\zeta_e = 1.6$ in the $B$-regime. One should note that $\zeta_e \sim 2$ in the $B$-regime. Thus the data confirm the conclusion that an increase in the neutral influx provides a drop in $T_e$ at the plasma edge and a rise in $n_e$.

In this case, the parameters $\zeta_e$ and $\zeta_r$ rise, and the conditions for the emergence of instability are produced [3]. The process is accompanied by a reduction in the life-time of impurities in the plasma.

The study of $\zeta_e$ in different plasma regimes has shown that not only $\zeta_e$ can be widely changed, but the amount of injected impurity reaching the plasma centre in changed as well. Two regimes with very similar integral parameters are shown in Fig.6. In this case, $\eta_e$ slightly rises due to the gas influxes from the valve in the regime $\text{dashed line}$, in the other one the valve is switched off $\text{solid line}$. The $B$-regime is realized in both cases, but when a KCl-pellet is injected, and increment in the $\text{SIR-signal}$ in the first discharge in eight-time smaller than that in the second one, although the estimates using the bolometric signal and the loop voltage give a difference not greater than a factor of two in the pellet size. One can conclude that the discharge conditions have permitted impurities to reach the centre to a greater extent in the second case than that in the first one.

Thus, the variation in the influx of neutral atoms results in a change in the plasma parameters at the plasma periphery. A reduction in the influx is accompanied by a rise in the temperature, $T_e$, and by a drop in $n_e$ at the plasma edge. An improvement in the life-time of impurities and particles in the plasma column takes place. The observed phenomenon can be induced by the instability damping produced by a high influx of neutrals, as shown in TEXT [3]. Similar changed in transport processes are observed in the analogous rearrangements of $n_e(r)$ and $T_e(r)$-profiles in other facilities [4]. Providing different conditions at the plasma column periphery by variations in the gas puffing, one can change the amount of impurities reaching the plasma column centre in the process of pellet injection.

REFERENCES

EXPERIMENTAL STUDY OF SOME PROBLEMS FOR A TWO-CHAMBER TOKAMAK

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I. The problems related with quenching of a fusion reaction and with removal of reaction products from the operating chamber have not been solved yet. According to the existing ideas, this phase of the reactor operation will occupy no less than 15% of the fusion reaction burn time [1]. This will result in large temperature variations in the first wall of the reactor and in a noticeable reduction in its service life. A break between shots can be reduced if one uses a scheme of operation for the tokamak-reactor proposed in [2]. After burning the necessary portion of fuel out, the plasma, as a whole, is rapidly ejected into an auxiliary chamber with the developed internal surface and with the high power pumping out. The chambers are connected with each other through a narrow toroidal slot for reducing a back gas flow. As a result of plasma ejection the operating chamber will be ready for igniting a new discharge and then the cycles will be repeated.

2. The passage of a plasma column through a narrow toroidal slot has been studied on the T-13 tokamak. The necessary condition for the plasma passage without destruction of the magnetic plasma column structure is the requirement: $\tau_{\text{motion}} = \frac{2a}{V} < \tau_{SK}^{p} \tau_{SK}^{w}$ where $V$ is the speed of plasma column motion along the major radius, $\tau_{SK}^{p}$ is the plasma skin time, $\tau_{SK}^{w}$ is the skin time of the slot walls. Under our conditions this requirement is satisfied with the speed $V > 10^4$ cm/s. One cannot provide such a rapid plasma motion in T-13 by technical reasons, changing the currents in the poloidal coils only. Therefore we have used a horizontal plasma column instabi-
luty in the confining magnetic field with a high decay index,
\[ n = - \frac{R}{B} \frac{dB}{dR} > \frac{3}{2} \]
, in the region of toroidal transition.

A partition splitting the chamber into two vessels was installed (R=58 cm) in the T-13 tokamak (Fig. 1). The partition includes four semirings with triangular cross-sections extended in the toroidal direction. Each semiring (made of aluminium) is isolated from the chamber and from other semirings. A width of the slot between the upper and lower semirings was adjusted in a range 2-7 cm.

The experiments were done with the following parameters: maximal plasma current \( I_p = 15 \text{kA} \), toroidal magnetic field at the radius \( R_0 = 40 \text{cm} \) \( B = 10 \text{kG} \), major radius was changed in the range \( R = 48-65 \text{cm} \) in motion, plasma density \( n \sim 10^{13} \text{cm}^{-3} \). The minor radius of the plasma column, \( a = 7.5 \text{cm} \), was preset with rail limiters.

3. The discharge parameters at the passage through the slot with a relative width, \( \delta = \frac{h}{2a} \times 100\% = 40\% \), where \( h \) is the distance between the upper and lower semirings (Fig. 1), are shown in Fig. 2. The speed of motion in the region of plasma column formation is \( v = 10^3 \text{cm/s} \), in the slot \( v = 3 \times 10^4 \text{cm/s} \). Spectroscopic measurements and the data from the Langmuir probes confirm the readings of magnetic probes about the plasma positions along the major radius.

The motion through the slot is accompanied by a some flux of cold deuterium atoms from the walls of the slot. It is confirmed by a rise in the intensity of \( D_a \) -line and by an increase in the atomic charge-exchange flux from the plasma.

A relative change in the plasma current in its motion through the slot, differently wide, is shown in Fig. 3. A change in the plasma current in its motion within the T-13 chamber without partition is also shown there.

The losses of charged particles to the partition are estimated by a saturated ion current from one semiring. The flux of charged particles in motion through the slot is reduced the
stronger, the higher its speed (Fig. 4). The power lost by the plasma through this channel does not exceed 10% of the Joule power of heating. \( P = \gamma J_{\text{sat}} T_e \), where \( \gamma \) is the heat transmission rate, \( J_{\text{sat}} \) is the saturated ion current from semirings, \( T_e \) is the electron temperature at the plasma column edge.

A width of the slot determines the nature of the plasma column passage through it. When \( \delta < 30\% \), there is no passage, the plasma column is destructed in front of the partition. When \( 30\% \leq \delta \leq 40\% \), the plasma passes through the slot, but the discharge is often disrupted that results in a slight deceleration of the column and in the loss of a plasma energy fraction. (\( \sim 15\% \)). When \( \delta > 40\% \), as a rule, the plasma column passes the slot without disruption.

Thus, on the basis of the experiments made one can conclude that under definite conditions it is possible to provide the plasma ejection through the toroidal slot with a width essentially-less than the plasma column diameter.

References


Figure captions

Fig. 1. Geometry of the experiment and location of diagnostics in T-13.[3]

Fig. 2. Discharge parameters in motion of the plasma through the slot \( \delta = 40\% \) (\( J_p \) is the plasma current, \( R \) is the major radius, \( U_l \) is the loop voltage, \( I \) is the intensity of the OV-line at the radius \( R = 58 \text{ cm} \), \( J_{\text{probe}} \) is the current from the Langmuir probe at the centre of transition, \( X\)-ray is the hard X-ray intensity.

Fig. 3. A change in the plasma current at the passage through the slot with relatively-different width.

Fig. 4. Saturated current from one semiring in motion of the plasma column with different speed.
SMALL-SCALE PLASMA TURBULENCE IN THE FT-2 TOKAMAK


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Small-scale plasma turbulence in tokamaks is of substantial interest for it can be responsible for anomalous energy transport [1]. It has been reported already [2] of investigations of the small-scale density fluctuations in the FT-2 tokamak during the current rise phase via CO₂ laser scattering. This report presents some new data on plasma density and magnetic field fluctuations during the flat current phase of the discharge in this machine.

The FT-2 tokamak has major radius of 55 cm and limiter radius of 8 cm. The investigations were performed in the ohmic heating regime: B₉ = 2.2-2.4 T, Iₚ = 30-40 kA. The plasma density, electron and ion temperatures on the plasma axis were estimated to be (1-3) · 10¹³ cm⁻³, 300-500 eV and 100-150 eV respectively. A high-power pulsed CO₂ laser was used as probing beam source providing peak power of 25-30 kW and pulse duration of 5 μs. That guaranteed the sensitivity enough to analyse the plasma density fluctuations of relatively low intensity in 2-20 MHz frequency range with wave vectors of 18-90 cm⁻¹. The plasma column was probed along several vertical chords of its small section plane.

The density fluctuations single-shot spectrum within the wavenumber resolution band Δkₚ=5 cm⁻¹ presents a number of well defined frequencies with Δω<<ω (see Fig.1) as it was previously determined for the current rise phase [2]. The spectra peak frequencies were not reproduced from one discharge to another. However, the peaked nature of the spectra remained quite appreciable after the averaging over 6-15 tokamak shots. The fact is that...
"quasilinear modes". Further on the broad frequency range being observed for a fixed \( k_L \) was probably due to simultaneous scattering from different magnetic surfaces. On approaching the plasma edge a rise in the scattered signal was obtained with no fluctuations being found for the chords close to the plasma column centre where scattering from strictly poloidal perturbations can be observed \((I_p=45 \text{ kA}, B_t=2.4 \text{ T})\). In order to clarify the spectra transformation an additional averaging over several frequency bands of 2 MHz width was performed. In Fig. 2 the isolines of such averaged intensity plotted against \( \omega \) and \( k_L \) are compared for chords \( x=+5, +6, +7 \text{ cm} \) \((x\) being a probing beam spacing from the plasma column axis). The spectra for +6 and +7 chords depict substantial similarity while for \( x=+5 \text{ cm} \) the spectra tend towards smaller \( k_L \). This spectra shape evolution observed previously in the current rise phase (see [2]) seemed to be due to essentially anisotropic structure of plasma density fluctuations with the poloidal correlation length exceeding the radial one localized near the plasma periphery. The spectra transformation in the flat current phase of the discharge as well as the absence of scattering from the central chords can be explained in the same way.

The density fluctuation level is estimated to be \( \delta n/n \approx 10^{-3} \). The poloidal phase velocity exceeded the electron drift one: \( \omega/k \approx (3-5)v_{de} \).

The density fluctuations spectra being averaged over 10 shots for probing along symmetric chords +7 and -7 cm are shown in Fig. 3. The fact that scattering is observed both from the low and high field side of the discharge proves that the microinstability is connected not with toroidal effects only.

A high frequency magnetic probe (a small loop about \((2\times10)\text{mm}^2\)) was placed in the equatorial plane in the limiter shadow close to the scattering region. Besides the intensive
magnetic field fluctuations in 20–50 kHz range the perturbations with \( f = 0.2–6 \text{ MHz} \) were found at \( I_p > 40 \text{ kA} \). The density and magnetic field fluctuations spectra are compared in Fig. 1. The coincidence of a number of frequency peaks can be indicated. This effect disappeared with chord displacement towards the plasma column centre and \( k_\perp \) increasing over \( \sim 50 \text{ cm}^{-1} \). The radial decay of the magnetic field perturbations intensity are presented in Fig. 4. Assuming the field associated with fluctuating plasma currents decays like a multipolar vacuum field [3] one can conclude the \( B \) signal for fixed frequency (see Fig. 4) may originate from a sum over different \( m \)-modes. For example the higher mode for \( f = 2 \text{ MHz} \) has \( m = 50 \) \( (k = 7 \text{ cm}^{-1}) \) and such density fluctuations may be indicated in scattering experiment at the plasma periphery [2]. The relative intensity of the magnetic field oscillations being extrapolated from the limiter shadow to the hot region was near \( 10^{-4} \).

According to the estimates the observed microinstability cannot be explained by electron dissipation processes (microtearing or rippling modes). It is more likely to be connected with plasma rotation in the radial electric field. That fits the localization of density fluctuations near the plasma edge and their anisotropy due to the rotation shear [4]. A nonuniform impurity distribution could affect the instability development also.

References

Fig. 1. The density (1) and magnetic field (2) fluctuations single-shot spectra. \( k_\perp = 25 \text{ cm}^{-1}, X = +7 \text{ cm} \).

Fig. 2. Isolines of frequency smoothed value of power spectra.

Fig. 3. The averaging over 10 shots density fluctuations spectra for probing chords +7 (1) and -7 (2) cm. \( k_\perp = 45 \text{ cm}^{-1} \)

Fig. 4. The radial decay of the magnetic field fluctuations intensity. Solid curves - calculations for a sum over two \( m \)-modes.

---

**Fig. 1.**

\[
\tilde{n} \propto 10^{7} \text{ cm}^{-3/2} \text{ MHz}^{-1/2}
\]

**Fig. 2.**

\[
\tilde{B} \propto 10^{-6} \text{ T} \cdot \text{ MHz}^{-1/2}
\]

**Fig. 3.**

**Fig. 4.**
TURBULENT ION HEATING IN TUMAN-3 UNDER THE FAST CURRENT RAMP.


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The possible role of ion-sound turbulence (IST) in fast current penetration into the plasma and plasma heating in tokamaks TORTUR and TUMAN-3 was mentioned in Refs. [1,2]. Here we analyzed experimental data on ion heating under the fast current ramp in TUMAN-3 [3]. For explanation of observed effects the mechanism of IST is evoked.

The evolution of different plasma parameters is shown in Fig.1. Conventionally the discharge might be divided into 3 stages: I – ohmic heating with small current (40-45 kA); II – fast current ramp phase up to 85-100 kA with \( I_p = 12.5 \) MA/s; III – phase with a large current. In the first phase plasma loop was slowly shifted from the center to the outer chamber wall for about 5 cm. The plasma with \( a \sim 18 \) cm was formed (Fig. 2,3) and after that fast current ramp was arranged with a fast (during 1-1.5 ms) shifting to the chamber center.

In the second stage of the discharge fast ion heating was observed. Central ion temperature increased from 100-150 eV to 200-300 eV during 3 ms (Fig.4). Characteristic two-fold increase time of ion temperature was 2-3 ms. The increase of charge-exchange atomic fluxes with the superthermal energies was observed simultaneously with the ion temperature increase.
The appearance of the fast particles was observed in a number of cases just after the beginning of current ramp. Sharply peaked ion temperature profile corresponds to this time (≈3 ms from the beginning of current ramp), as it is shown on Fig. 6.

It's impossible to give account for phenomena described above by energy transfer from electrons to ions due to coloumb collisions, because of long characteristic time of that processes ~10 ms. It's known at the same time, that the IST can provide plasma ion heating during the shorter time.

The ion heating with the classical ion-electron times, took place when the hard X-rays and synchrotron emission increase was registrated during current ramp probably connected with the existence of fast electron beam.

The experimentally measured dependencies of ion energy distributions allow us to draw two conclusions. Firstly it is possible to speak about hot ion tail formation in the energy region $\varepsilon > Z T_e$ (or $v > v_s$ ) with the high effective temperature $T_h \sim Z T_e$. Secondly it is possible to speak about significant increase of number of ions in the region $\varepsilon > Z T_e$ during current ramp (~4 ms). It is possible to explain such peculiarities of the ion distribution in frames of IST theory by the following way. The ion heating in the velocity region $v > v_s$ up to $T_h \sim Z T_e$ is accounted for by their interaction with the ion sound oscillations. The increase of fast ion number in it's turn may be the result of bulk ion heating with $v \sim v_T$, determined by induced sound scattering and caused by this heating ion influx in the region $\varepsilon > Z T_e$. Characteristic time of the heating $\tau \sim T_i / Z E v_s \sim 2-3$ ms is in good agreement with the experimentally observed times of the resonant ions number increase.

For the excitation of ion-sound instability in plasma it's necessary the electric field to be greater threshold value $E_{th}[4]$. Our case with $Z T_e / T_i \ll 6, E_{th} = 3,5$ mV/cm corresponds approximately to the nonisothermicity observed before current ramp. The electric field outside the plasma column becomes
greater than the threshold value ($E > 10$ mV/cm) during current ramp and ion-sound instability develops at the plasma periphery.

It should be noted, that the IST-theory may be applied in the case of plasma with nonmagnetized electrons. In our case plasma discharges were arranged in the magnetic field $\sim 5$ kGs, where $\omega_L/\Omega_e \sim 1 - 2$, i.e. electrons were not magnetized practically.

As the ion-sound instability causes ion turbulent heating, it also has to decrease plasma conductivity in $\sim 10$ times. As a result the penetration time of the electric field in plasma has to decrease correspondently. In such a way anomalously fast current penetration in the plasma center may be accounted for by the excitation of IST. The fast current penetration is proved in our case by the absence of the electron temperature skinning (Fig. 2).

It should be noted, that the experiments were carried out, where the beginning of the shifting coincided with current ramp, occurred in advance or after it. In all these cases the time difference between the beginning of plasma shifting and current ramp $\Delta t = t_R - t_I$ was within $1.5 - 3$ ms. The fast ion heating at the different values of parameter $\Delta t$ and in the discharges with plasma shifting to the chamber center from the inner wall was observed. These facts allow us to draw a conclusion, that seemingly column shifting didn't provide ion heating directly.

References.
Fig. 1  a) Discharge current and loop voltage, b) plasma shifting along the large radius and mean electron density.

Fig. 2 Evolution of the electron temperature distribution over small radius.

Fig. 3 Evolution of the electron density distribution over large radius.

Fig. 4 Time evolution of the charge-exchange atomic flux with the energy 430 eV and ion temperature (points denote values averaged over 1 ms).

Fig. 5 Charge-exchange atomic spectra at the different time.

Fig. 6 Evolution of the ion temperature distribution over small radius.

+ Time is counted from the beginning of the current ramp.
MODELLING OF THE CURRENT DENSITY DISTRIBUTIONS UNDER THE DIFFERENT DISCHARGE SCENARIOUS IN TUMAN-3 PLASMAS.


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A large number of papers is devoted to the study of the current penetration into the tokamak plasmas and experimental measurement of current density evolution under the fast current ramp ($I > 10^7$ A/s) with diagnostic techniques capable to measure distribution $j(r)$ [1,2] is of special interest. Investigations of fast current ramp on TUMAN-3 in the last years [3,4] were interpreted in terms of convective models of current penetration and neoclassical values of plasma conductivity, assumptions were made also about the possible role of ion-sound turbulence and production of the accelerated electrons [5] in the fast current penetration to the center of plasma loop.

In this paper we analyzed one of the typical discharges in tokamak TUMAN-3. On Fig.1 main parameters of the discharge are presented: a) B - magnitude of the toroidal magnetic field at $R=55$ cm, $I_p$ - total plasma current, $V_1$ - loop voltage, $\Phi_p$ - toroidal flux, produced by the current, $R_p$ - position of the plasma axis, b) $n_e$ - mean plasma density, c) ECE, HXR - intensity of synchrotron and hard X-ray radiation. Between the shots the shape of the last closed flux surface was re-
constructed from the magnetic diagnostic data (Fig. 2a). The analysis of the data shows that during the displacement of plasma loop inside the torus $R = 5 \text{ cm}$ current ramp rate reaches the value $I_p \approx 1.25 \times 10^7 \text{ A/s}$. The evolution of the distribution of safety factor $q(r)$ is shown on Fig. 2b. During whole current ramp phase $q_L$ value at the last closed flux surface doesn't change substantially ($3.3 \leq q_L \leq 3.6$). Close to the axis value of $q$ during the current ramp phase is $\approx 1.0$. In the quasistationary phase of the discharge due to shrinking of the current channel safety factor decreases up to $q \approx 0.6$. Inspite of low $q(0)$ values the relaxation oscillations in the discharge are absent. We must mention that $q < 1$ values near the magnetic axis were measured also earlier [2]. The measurement of the main plasma characteristics made in the discharge makes it possible to emphasize the following features of fast current ramp ($I_p > 1.2 \times 10^7 \text{ A/s}$) in the fully ionized TUMAN-3 plasmas: a) the delay of $T_e$ increase is seen on the time behavior of the electron temperature $T_e$ in the center of discharge in the quasistationary phase ($t > 31 \text{ ms}$) temperature distribution remains peaked, b) the current ramp rate $I_p \approx 1.25 \times 10^7 \text{ A/s}$ doesn't cause significant skin current distribution, c) in the quasistationary phase of the discharge shrinking of the current channel takes place without relaxation oscillations and cancelled simultaneously with the decreasing of the MHD oscillations level, d) the current ramp phase might be followed by substantial increase of synchrotron and hard X-ray radiation (examples for different discharges of such type (1 and 1') are shown on Fig. 1c).

On Fig. 3 the experimental and calculated values of $j(0)$ are shown (1 - model with the electron heat conductivity anomaly, 2 - model with the changing of $Z_{eff}$, 3 - model with fast electrons generation). The transport model of the discharge, contains one dimensional time dependent equations for current density $j$, plasma density $n$, ion $T_i$ and electron $T_e$ temperatures, was used. As the input data measured values of
integral parameters $B$, $I_p$, $U_1$, $R_p$, $a_p$ (radius of the last closed magnetic surface) were used. Runaway electrons current generation rate [6] was calculated with $Z_{eff} = 2.5$, taking into account finite orbit losses and conductivity anomaly. Runaway electrons velocity was assumed equal the light velocity. The development of the model can include the finite life time of runaway electrons and finite value of their velocity. It is necessary to mention that similar approach (without taking into account the spatial distributions) was used in Ref. [7]. On Fig. 4 the experimental and calculated current density distributions in the current ramp phase are presented. The measured fast current density increase in the center of the discharge can be satisfactorily explained in frames of the described model by runaway current generation. In these discharges also was detected three-fold increase of X-ray spectra hardness (from $\sim 100$ kev to $\sim 300$ kev). We must also mentione that same results might be obtained under the different scenarious of the discharge: $\Delta R = 0$ and $\Delta R < 0$.

References.
5. N. I. Vinogradov et al., preprint FTI-1177 (in Russian), 1987
Fig. 1 Evolution of the discharge parameters
a) $B$, $I_p$, $V_p$, $\Phi_p$, $R_p$;
b) $n_e$; c) $I_{ECE}$, $I_{HXR}$.

Fig. 2 a) position of the last closed flux surface, b) distribution of safety factor at different time.

Fig. 3 Evolution of current density $j(0)(a)$ (experiment and calculation), rate of runaway current generation $S_r(0)(b)$ and runaway current density $j_r(0)(c)$.

Fig. 4 Evolution of current density distribution over the plasma radius a) experiment, b) calculation:

$\jd$ - conductivity current (- - -)

$\jfr$ - runaway current (---)
MEASUREMENTS OF FAST ION RADIAL DIFFUSION IN TFTR

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The effective confinement of energetic ions is an important requirement for efficient plasma heating by neutral beams, by most forms of RF, and by charged fusion products. Anomalous radial diffusion of the fast ions as they slow down would increase fast-ion losses and decrease the power delivered to the thermal plasma. We report here the first local measurements of fast-ion radial diffusion, obtained in TFTR by injecting the most tangential beam so as to miss the magnetic axis and to create a hollow annular ring of fast ions in the region \( r > 0.35 \) m. The experimental configuration is illustrated in Figure 1. A set of two mass- and energy-resolving neutral particle analyzers (NPAs) scanned horizontally in the torus midplane near the inner edge of the annular ring (tangency radius \( R_T \approx R_0 - 0.35 \) m) to measure the density of fast ions which diffused into the hollow core as they slowed down. As expected, the fast-ion density increased abruptly at the edge of the hollow core. Fast-ion diffusion would cause this edge to become more diffuse at lower energies, since the lower-energy ions have had more time to move radially in the presence of turbulence. We obtain an upper bound on the magnitude of fast-ion radial transport (diffusivity \( D < 0.05 \) m\(^2\)/s) by comparing the charge-exchange measurements with Fokker-Planck calculations which incorporate an ad-hoc level of radial diffusion.

The experiment was conducted during a sequence of 27 nominally identical \( B_b = 4.9 \) Tesla, \( I_F = 1.2 \) MA, \( q_\phi = 5.0 \) deuterium discharges. The plasmas were in contact with the inner toroidal belt limiter (IBL), which defined a minor radius \( a = 0.71 \) m and a major radius \( R_0 = 2.36 \) m. A single neutral beam source with \( P_b = 2.0 \) MW, \( E_b = 95 \) KeV, was injected in the horizontal midplane from 4.0 \(- 4.5 \) seconds at a tangency radius of \( 2.84 \) m. Beam injection increased the line-averaged electron density \( n_e \) from its ohmic value of \( 1.19 \times 10^{19} \) m\(^{-3} \) to \( 1.35 \times 10^{19} \) m\(^{-3} \), and increased the central ion temperature from \( 3.0 \) to \( 4.5 \) KeV. During injection, the discharges were sawtoothing with a period \( \tau_{\text{st}} = 22 \) ms, amplitude \( \Delta T_e(0) = 0.25 \) KeV, and inversion radius about \( 13 \) cm. There was no discernable mode activity from Mirnov coils or soft x-ray detectors. The measured plasma profiles and calculated beam deposition profile are shown in Fig. 1b. At the minimum minor radius of beam deposition \( r \approx 0.35 \) m, \( T_e = 2.5 \) KeV, corresponding to a critical energy \( E_c = 47 \) KeV, so the fast ions slow down predominantly on electrons between \( E_b \) and \( E_b/2 \). The slowing-down time \( (E_b \to E_b/2) \) at \( r = 0.35 \) m in this plasma is \( 75 \) ms, so the fast ions would diffuse about \( 27 \) cm for \( D = 0.5 \) m\(^2\)/s (using \( \Delta r \sim \sqrt{2D\Delta t} \)). These plasmas were heavily contaminated by carbon: the \( Z_{\text{eff}} \) deduced from tangential visible bremsstrahlung emission was \( 5.8 \pm 0.5 \), with a contribution \( Z_{\text{eff}}(\text{metals}) = 0.3-0.4 \) measured by soft x-ray pulse-height analysis. The gross plasma conditions remained remarkably constant throughout the scan, including the line-averaged electron density...
which varied by less than ±2.5%. The edge neutral density as measured by a radially-viewing Hα array decreased systematically by a factor 4-5 during the course of the shot sequence (these plasmas resemble standard TFTR “conditioning” shots used to degas the IBL [1] and reduce recycling). To compensate for this neutral density variation, which changed the target density for fast-ion charge exchange, the measured CX flux for each discharge was divided by the measured Hα line emission.

The measured CX energy spectra along two sightlines, \( R_T = 2.07 \text{ m} \) and \( 2.25 \text{ m} \), are shown in Figure 2. Along sightlines with \( R_T < 2.15 \text{ m} \), the CX signal comes almost entirely from ions in the “inner” part of the fast-ion annular ring, \( R \approx R_\alpha - 0.35 \text{ m} \). Along the sightlines with \( R_T > 2.15 \text{ m} \), the CX signal comes almost entirely from ions in the “outer” part of the fast-ion annular ring which have pitch-angle scattered significantly prior to undergoing charge exchange. There is excellent agreement in the shape and relative magnitude of the energy spectra between the measurements and calculations by the Fokker-Planck code FPFF [2].

Figure 3 plots the CX flux (arbitrary units, normalized to \( H_\alpha \) emission intensity) as a function of sightline tangency radius \( R_T \) for two energies: 95 KeV, corresponding to the beam injection energy, and 57 KeV, which is slightly above the beam half-energy component. The two large peaks (at \( R_T = 1.98 \) and \( 2.90 \text{ m} \)) correspond to direct views of the inner and outer parts of the fast-ion ring, with no pitch-angle scattering. Of greatest interest to the study of fast-ion radial transport is the CX signal in the region \( 1.9 \leq R_T \leq 2.2 \text{ m} \). This signal originates from nearly parallel-going ions (\( \psi_j/\nu \approx 1 \)) at the inner edge of the fast-ion ring. At the injection energy, the gradient of CX flux in this region is determined by the beam deposition profile. At the lower energy, \( E \geq E_b/2 \), any radial diffusion experienced by the beam ions as they slow down would push them into the previously hollow core, or equivalently to larger major radius; such ions would
provide CX signals at larger $R_T$ (since $R_T = R \times v_\parallel/v \approx R$), and would thus tend to flatten the observed gradient. The measured fast-ion gradient is in fact about the same at $E = E_b$ and $E = E_b/2$, which provides qualitative confirmation that fast-ion diffusion is small. To quantitatively evaluate the magnitude of radial transport implied by these measurements, we compare the variation of CX flux versus sightline $R_T$ with predictions of the FPPRF code including radial diffusion.

The code (FPPRF) solves the standard Fokker-Planck equation in toroidal geometry including a new term, $\frac{1}{r J} \frac{\partial}{\partial r} (r D \frac{\partial J}{\partial r})$ to represent radial diffusion, where $D(E, r)$ is the fast-ion radial diffusivity. The code alternates between velocity space transport where $r=$constant, and real space transport (diffusion) where $E$ and $\mu$ are conserved. Particle orbits are computed for concentric circular cross-section, but the plasma is allowed to have a Shafranov shift ($\approx 5.5$ cm in this experiment) in the calculations of beam deposition and simulated CX flux. The code uses the measured $Z_{eff}$ and the measured profiles of $T_e(r), T_i(r),$ and $n_e(r)$. Because the neutral source is strongly poloidally asymmetric for plasmas resting on the IBL, a 3-D Monte-Carlo neutrals code [3] calculated the neutral density used in the simulation of the CX flux.

To obtain the most unambiguous comparison of the Fokker-Planck calculated flux versus $R_T$ with experiment, we must subtract the small contribution to the CX flux originating from the "outer" ring. This is easily accomplished within the FPPRF simulation by setting the target neutral density to zero for $R > 2.4$ m. To make the same adjustment for the data, we fit a smooth curve through the measured flux from $R_t \approx 2.4$ to 2.7 m (which can only represent pitch-angle scattered ions from the outer ring) and subtract

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Figure 2: (a) Flux versus energy measured by a single NPA along two sightlines and the CX flux calculated by FPPRF assuming $D(r, E) = 0$. (b) Measured CX count rate versus tangency radius (normalized to $H_n$, emission intensity) and the Fokker-Planck calculation (using $D = 0$), showing peaks from both the outside and inside edges of the fast-ion annular ring. For clarity, the data and simulation for $E = 56$ keV have been shifted down by a factor 100. The open and shaded symbols differentiate measurements by the two NPAs. Up-arrows identify NPA measurements under near-saturated conditions which thus represent a lower bound on the actual flux.
from the measured flux in the important region $R_T = 1.9 - 2.4$ m. The result is shown in Figure 3. Notice that for the first factor $\approx 10$ decrease in CX signal strength beyond the normalization point at $R_T = 2.05$ m, the data are in best agreement with the $D = 0$ simulation. At lower signal levels, the data analysis becomes more sensitive to the subtraction algorithm and to the relative calibration of the two NPAs (such systematic errors are not reflected in the error bars).

The detailed comparison near the fast-ion inner edge indicates an upper bound on $D$ of 0.05 m$^2$/s, which is much smaller than the ion and electron thermal diffusivities at $r = 0.35$ m ($\chi_i = 0.7$ m$^2$/s, $\chi_e = 0.45$ m$^2$/s, inferred from the measured temperature profiles). We also remark that measurements of the electron density profile evolution following a perturbative gas puff ($\Delta n/n < 5\%$) indicate that the particle diffusivity in ohmic discharges ($I_p = 1.4$ MA, $\bar{n}_e = 1.2 \times 10^{19}$ m$^{-3}$, $B_t = 4.8$ T) is 0.4 - 1.0 m$^2$/s[4]. These trends are consistent with results of theoretical studies of test particles in a turbulent plasma [5] which indicate that the magnitude of radial transport of energetic particles should be significantly less than that of the thermal population. Our conclusions regarding the magnitude of radial diffusion are applicable for co-passing ions in the low-density, very weakly heated regime; results on radial diffusion of energetic counter-passing ions with and without strong central co-heating power will be reported elsewhere.

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ION TEMPERATURE PROFILES AND ION THERMAL CONFINEMENT IN TFTR*

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Introduction
During the 1987 experimental run period on TFTR, extensive efforts were made to expand our capability to measure the plasma ion temperature, especially in the energetic ion, or "supershod", regime. These efforts were motivated by the desire to conclusively verify the high ion temperatures (>20 keV) reported earlier and, with the measurement of ion temperature and toroidal rotation radial profiles, study the details of ion thermal confinement in these plasmas. Indeed, the TFTR supershot operation regime offers an attractive environment for ion power balance studies since the high electron temperatures attained here can result in almost negligible ion-electron coupling. The ion power losses are strongly dominated by ion thermal conduction over the range 0.2 < r/a < 0.8, which allows an extension of the study of ion thermal diffusivity, such as that done on the D-III tokamak, to hot plasmas in larger devices. The influence of plasma rotation on ion power balance is also a topic of intense interest since the neutron production of the supershot plasmas has been found to depend strongly on the degree to which the neutral beam injection is balanced with respect to applied torque to the plasma.

Very High Ion Temperature Measurements on TFTR
A multi-spatial and -spectral imaging change-exchange recombination spectroscopy diagnostic (CHERS) became fully operational during 1987, and it provided radially resolved measurements of $T_i$ (R,t) and $u_\phi$ (R,t) for a wide range of discharge parameters. Both the heating neutral beams and a dedicated diagnostic neutral beam provided the fast doping neutrals needed for the charge exchange excitation of visible lines from a variety of impurity ions. Up to 12 positions in the plasma can be measured simultaneously, and the spatial separation of the measured location is 10 cm for a single shot, but can be reduced to 5 cm by changing the field of view between shots. The spatial resolution of the measurements is typically 3 cm. For most of the measurements reported here, the CVII 5292 Å (n = 8-7) transition was used.

In general, there is reasonable agreement between the central values of $T_i$ and $u_\phi$ obtained in the supershot regime from the CHERS diagnostic and central $T_i$ and $u_\phi$ values obtained with the standard crystal X-ray doppler broadening measurements of the Fe XXV K-alpha line. For example, Fig. 1 shows a comparison of the central C$^{6+}$ ion temperature and the values derived from the Fe XXV K-alpha measurements for a series of balanced injection supershot plasmas. Overall, the CHERS measurements are ~ 5-10% higher than the K-alpha values, which is consistent with the facts that the CHERS measurements are highly localized while the K-alpha...
lines are excited over a small range of the plasma interior where $T_e$ is high, and that the ion temperature profiles tend to be very peaked (more than the $T_e$ profile) in these plasmas. Similar comparisons between CHERS and the K-alpha measurements have verified the attainment of toroidal rotation speeds up to $8 \times 10^5$ m/sec when unbalanced neutral injection is used in these low-density target plasmas. The analysis of these measured rotation profiles is reported in a companion paper at this conference.

The measured profiles have been analyzed with both the SNAP and TRANSP kinetic analysis codes, which use, among other parameters, the measured $T_i(r)$, $T_e(r)$, and $\nu_0(r)$ profiles as input. The steady-state SNAP code has been used in either an analysis mode, where the measured $T_i(r)$ profiles are used as input to derive an ion thermal conductivity, $\chi_i(r)$, from the data, or in a predictive mode wherein a given model is chosen for $\chi_i(r)$ and the resulting calculated $T_i(r)$ profile is compared to the measurements. Both balanced injection supershots with negligible rotation and purely Co-injected plasmas with high rotation speeds have been studied to explore the effects of rotation on the ion power balance. In general, ion temperature profiles are relatively peaked for the balanced injection cases, while they are broader for the lower-confinement Co-injection cases with strong toroidal rotation.

**Balanced Injection "Supershots"**

Comparison of the peaked $T_i(r)$ profiles for the balanced injection cases with model transport calculations show that the ion thermal conduction is strongly non-neoclassical. For example, Fig. 2 shows a comparison of the measured $T_i(r)$ profile for a high temperature supershot with the values calculated under the assumptions that $\chi_i = 1 \times \chi_i^{\text{neo}}$, where $\chi_i^{\text{neo}}$ is the standard Chang-Hinton neoclassical ion thermal diffusivity, and that $\chi_i(r) = 1.5 \times \chi_e(r)$ where $\chi_e(r)$ is the electron thermal conductivity derived from the measured $T_e(r)$ profiles. The measured $T_i(r)$ profile is clearly much narrower than that expected from neoclassical theory, while $\chi_i \propto \chi_e$ produces a much closer approximation to the measured values. Indeed, the value of $\chi_i(r)$ derived from the data in the analysis mode of the SNAP code reveals that $\chi_i(r)$ has a magnitude and radial dependence comparable to that of $\chi_e(r)$. It is of interest to note that a convective ion power flow given by $P_{\text{i,conv}} = 3 T_i \Gamma_i / 2$ (where $\Gamma_i$ is the thermal ion particle flux), is necessary to prevent unrealistic negative thermal conductivities in the analysis mode of SNAP modelling or provide reasonable agreement between the measured and calculated central values of $T_i$ in the predictive mode.

Examination of several balanced injection supershots reveals that, within a considerable range of scatter, $\chi_i(r) = 1-2 \times \chi_e(r)$ provides a crude but reasonable approximation for $\chi_i(r)$, especially in the $r = a/2$ confinement region. This also appears to be true for a few higher current L-mode discharges which have been analyzed but not discussed here.

The ion power balance derived from the kinetic analysis confirm that electron-ion coupling is negligible over the entire profile and the ion power losses are strongly dominated by thermal conduction over most of the plasma cross section, except for the central $r/a \leq 0.2$ region, where
conductive and convective losses of equal magnitude (cf. Fig. 3).

**Unbalanced Injection with Strong Toroidal Rotation**

It has been observed that plasma rotation has a direct correlation to plasma performance in the supershot regime. The overall neutron yield is lowered by the reduced relative energy between the fast ions and the rotating plasma, and rotation has a strong effect on the plasma power deposition and balance. The reduced relative speeds between the plasma and neutral beam particles results in reduced beam penetration to the plasma core and reduced beam energy deposited in the plasma in its rotating frame. This results in a broadened power deposition with consequently less central electron heating. Likewise, the power deposition to the ions is much less peaked in radius than for the balanced injection case. The viscous damping of the plasma rotation provides an additional source of heating power to the ions which, since it is strongest where the $u_\phi(r)$ gradient is largest, tends to peak off-axis.

These combined effects of rotation lead to the broader $T_i(r)$ profiles observed in unbalanced injection cases, and, in extreme cases, can even result in flat or hollow $T_i(r)$ profiles, as shown in Fig. 4. Here, full Co-injection of 11 MW gives rise to a peaked $u_\phi(r)$ profile with $u_\phi(0) \approx 8 \times 10^5$ m/sec for a standard 0.9 MA discharge with $<N_e> = 1.8 \times 10^{19}$ m$^{-3}$. As with the balanced injection case, the ion thermal conductivity $\chi_i$ is strongly non-neoclassical and $\chi_i \sim \chi_e$ plus the effects of ion viscous heating reasonably reproduces the measured $T_i(r)$ profile.

The ion power balance analysis of this rotating plasma (Fig. 5) indicates that convection dominates the core ($r/a \leq 0.3$) power flow, while conduction is more important in the range $0.3 \leq r/a \leq 0.8$. However, both convection and electron-ion coupling are non-negligible ion thermal loss mechanisms for $r/a \geq 0.5$.

The role of rotation is evident in a comparison of the input power densities calculated for the balanced and pure Co-injected discharges discussed here. The balanced injection discharges have a very peaked heating profile both for the ions and electrons, resulting in very high central temperatures, while the unbalanced injection case shows a broad heating profile with relatively low core ion and electron heating, and significant off-axis ion heating due to rotation damping.

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References
Fig. 1. Comparison of central ion temperatures from charge exchange recombination spectroscopy of C+6 and Fe K-alpha doppler broadening.

Fig. 2. Measured ion temperature profile for a nonrotating plasma compared with calculations assuming chi-i(r) = 1.5 x chi-e(r) and chi-i(r) = chi-neoclassical.

Fig. 3. Volume integrated power balance for ions in a balanced injection superheat. Pinj = 13 MW.

Fig. 4. Measured ion temperature profile for a rotating plasma compared with calculations assuming chi-i = 0.5 x chi-e with and without ion heating from rotational viscous damping included.

Fig. 5. Volume integrated power balance for ions in a rotating plasma with unbalanced injection. Pinj = 11 MW.
Triton Burnup Studies on TFTR*  
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The 1-MeV tritons produced by the $d(d,p)t$ fusion reaction in nearly equal numbers to the 2.5-MeV neutrons from the $d(d,n)^3$He reaction have gyroradii and slowing-down times similar to those of the 3.5-MeV $\alpha$ particles to be produced in DT experiments. The confined tritons can slow down and a fraction can react (burn up) with the plasma deuterons to produce 14-MeV neutrons. The ratio of 14-MeV to 2.5-MeV neutron production (the burnup) provides an indication of the fusion product confinement and slowing down.

Activation Foils  The fluence of neutrons is measured by threshold reactions in an activation foil system. Direct calibration and analysis has revealed that near the foils the machine structure is much thicker on average than previously used in neutron transport analysis[1]. The 2.5-MeV $d(d,n)^3$He neutrons are measured by the $^{115}In(n,n')^{116}In$ reaction (effective energy threshold about 0.8 MeV)[2]. The foil system has been calibrated and the fluence observed to be reduced (relative to a source in vacuum) by a factor of about 0.45-0.32, apparently due to attenuation and scattering[3]. The $d(i,n)\alpha$ neutrons are measured by $^{63}Cu(n,2n)^{62}Cu$ activation (threshold of 11 MeV, calibrated reduction of fluence by factor of 0.16-0.36) or $^{27}Al(n,\alpha)^{24}Na$ activation (threshold of 6 MeV, calibrated reduction of fluence by a factor of 0.29-0.37)[3]. The first number in the range of fluence reduction is the observed reduction seen in the data (limited in toroidal angle coverage), while the second number is the estimated reduction from models of machine structure attenuation. Each number in the range is itself uncertain due to calibration accuracy or model assumptions by factors of 25% (for the copper foil measurements) to 50%. The fluence from TFTR plasmas observed by the In foils is proportional to the epithermal neutron fission detector system yield to within 10% over a range of $10^2$. The calibrated yield of 2.5-MeV neutrons measured by the activation foils is 1.6-2.2 (each number with 50% uncertainties) times the yield from the epithermal neutron fission detector system[4].

The time-integrated triton burnup can be measured by taking the ratio of the yield of 14-MeV neutrons measured by the foils divided by the 2.5-MeV yield from the fission detector system. Figure 1 shows the de-
The single discharge. and of foils instead of factor-
factor of I
Figure
tron when born. (The observed reduction in fluence from the calibration has been used to provide the vertical scale. Using the modeled attenuation factors, or using the increased yield from the indium foils instead of the fission detector yield, would reduce the magnitude of the measured burnup.) The copper data is from single discharges, while the aluminum foils require accumulating over 4-20 discharges. Between 0.9 and 1.4 MA the measured burnup increases with the duration of the NBI (more closely achieves steady-state). For similar beam duration the burnup increases with current, and perhaps relatively more than expected from just the increase in the neo-classical confined fraction of the tritons when born. (The confined fraction of tritons decreases from unity at 1.5 MA and above to around 50% at 0.6 MA.) The magnitude of the burnup and the increase with current and beam duration are similar to those seen by the extrapolated fractional burnup method[5], although now the entire duration of the discharge is measured, and a wider range of currents is studied.

**TIMEEV Time-Dependent Calculation**

Since the burnup magnitude is time dependent (Fig. 1), a time-dependent burnup code[6, 7, 5] is used to analyze the expected burnup from TFTR plasmas. The code assumes classical slowing down at the radius of birth of the triton and no effects from pitch-angle scattering or energy diffusion. For TFTR the uncertainties in the input data are much greater than the expected error from these assumptions. Experimental data for each discharge are used for the time and spatial dependence of the electron temperature (from ECE radiometer and Michelson interferometer, and Thomson scattering) and the electron density (from the Multichannel InfraRed Interferometer [MIRI] and Thomson scattering).
The deuterium density is calculated using the $Z_{\text{eff}}$ from visible bremsstrahlung with estimates of the elemental contributions obtained from X-ray pulse height analysis. The uncertainty in the calculated classical burnup is about 50% due to uncertainties in the input data, primarily in the deuterium density (from the $Z_{\text{eff}}$, its radial profile, and its elemental composition), in the triton source profile, and in the magnitude of $T_r$. The magnitude and scaling of the burnup on TFTR are consistent with being classical (Figure 2), since there are about 50% uncertainties in the calculated burnup values, and factor-of-two errors in the absolute measured burnup. Any increase in the burnup with $I_p$ over and above the effect of improving neoclassical confined fraction appears explainable by slight increases in $T_{\text{ef}}$ and decreases in $Z_{\text{eff}}$ with increasing current for this data set.

Time-Dependent Measurements

Time-dependent 14-MeV neutron measurements have now been made throughout the duration of TFTR discharges. The 14-MeV neutron emission was measured by a NE213 proton recoil spectrometer with pulse-shape discrimination to eliminate $\gamma$ rays (Figure 3). The NE213 system can simultaneously measure both the 14-MeV and 2.5-MeV neutron flux, although limited in dynamic range and 14-MeV count rate. The absolute calibration of the NE213 was done relative to the epithermal neutron fission detector system. For the 14-MeV response this was done by comparing 25 “supershots” and finding discharges with maximum NE213 counts for minimum fission counts for intervals 0.6-2.0 seconds after the NBI pulse. For those discharges it was then assumed that all the neutrons this late in time were 14 MeV[5]. This relative

Figure 2: Comparison of measured time-integrated burnup with that calculated by the TIMEEV code. Results from both copper and aluminum foils are shown (some from similar sequences of discharges) as well as the value from the NE213 system.

Figure 3: Time-dependent triton burnup measurements averaged for a sequence of eight TFTR discharges using a NE213 proton recoil spectrometer. Comparison is made to the TIMEEV code. The large error bar on the TIMEEV calculation represents the uncertainty due to the $Z_{\text{eff}}$, for which there was no visible bremsstrahlung measurement for these discharges.
Calibration is within the 40% uncertainties of the absolute calibration done later, although intervening small movement of the collimator makes direct use of the absolute calibration undesirable.

The 14-MeV neutron emission on TFTR has also been measured in Silicon Surface Barrier Diodes (SBD) using the \( ^{28}\text{Si}(n,\alpha)^{25}\text{Mg} \) and \( ^{28}\text{Si}(n,p)^{28}\text{Al} \) threshold reactions. Figure 4 shows a comparison of the 14-MeV count rate and the time-dependent burnup from the TIMEEV code (the magnitude of the SBD count rate has been normalized to the theoretical calculation; the actual magnitude is within the factor-of-two uncertainty of the expected count rate). For the NE213 data (Fig. 3) the measured 14-MeV neutron rate appears marginally high during the triton buildup stage, and somewhat low during the decay compared with the classical prediction. This difference may be due to unresolved variations in \( Z_{\text{eff}} \), or it could be due to slowing down that is more rapid than classical or to beam deuteron interactions with thermal tritons. For the SBD data (Fig. 4) the agreement with the calculated classical burnup is good. The differences between Figs. 3 and 4 may be due to differences in the discharges.

Conclusions Time-integrated and time-dependent measurements of triton burnup have been made on TFTR plasmas. Within the factor-of-two experimental uncertainties and the 50% calculational uncertainties the observed triton burnup cannot be distinguished from "classical", that is neoclassical prompt loss of the tritons and classical triton slowing-down over a range of current from 0.5–1.8 MA.

INJECTION OF DEUTERIUM PELLETS INTO POST-NEUTRAL-BEAM TFTR PLASMAS

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Introduction—After neutral beam injection on the Tokamak Fusion Test Reactor (TFTR), the neutron signal[1] is observed to decay with two time scales. The first decay constant is due to the slowing of the deuterium beam ions.[2] The second, slower-decaying component of the signal has been identified with the burnup of fusion-produced tritons.[3]

In order to confirm that the neutron emission late in time is dominated by the burnup of fast tritons, deuterium pellets were injected into post-neutral-beam TFTR plasmas. Late in time, it is expected that all the injected deuterium ions will have thermalized leaving only fusion-produced tritons as a source of beam-target reactions. By injecting a pellet into this plasma, the beam-target neutrons can be differentiated from thermonuclear neutrons. If the reactions are beam-target in nature, an immediate rise in neutron emission will be seen due to reactions between the beam ions and the pellet material. Pellet injection into a Maxwellian plasma will not lead to a rise in neutron emission on the timescale of pellet ablation since the fusion reactions in Maxwellian plasmas occur between ions in the high-energy tail of the distribution and not with bulk ions, as are added by the injection of a pellet.[4] This has been shown by pellet injection into ohmic and beam-heated plasmas.[5]

The quality of triton confinement is important for predicting the single-particle confinement of alpha particles in the DT phase of TFTR operation[6] since the alpha particles must be confined for a time comparable to their slowing-down time for alpha-particle heating to be effective. This experiment provides further supporting evidence that the reactions at late times for these discharges were dominated by reactions between fusion-produced tritons and the deuterium plasma. Some fraction of these tritons remain well confined as they slow down in the plasma for at least 1.5 s after the beams (the time when the pellets were injected). It also suggests that TFTR can achieve an alpha-storage mode in which the density would be quickly raised by injection of a pellet or by a large gas puff into the post-neutral-beam phase of a TFTR DT plasma. In this way, the heating rate of the alpha particles can be varied.[7]

Experiment—This experiment was conducted on TFTR (R = 2.46 m, a = 0.81 m, \( B_T = 4.8 \) T, \( I_p = 0.9 \) MA for these discharges) with neutral beam powers in the range 6–12 MW injected in the direction of the plasma current. The current chosen is that for which it has been found[8] that for available beam power the DD reaction rate and, thus, the number of fusion-produced tritons in the plasma are maximized. Peak neutron source strengths were approximately \( 2 \times 10^{18} \) s\(^{-1}\). Pre-pellet line-integrated densities rose from

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1 \times 10^{19} \text{ m}^{-2} \text{ during the ohmic phase to three times that during the beam injection phase. The beams were on for 1.5 s and pellets were injected during eight shots with the injection time varying from 0.25 to 1.5 s after beam turn-off. One of the shots received two pellets. The cylindrical pellets[9] that were injected were 3.5 mm in diameter by 3.5 mm long, each being made up of approximately 1.8 \times 10^{21} \text{ deuterium atoms.}

Prior to pellet injection, the central electron temperature, as measured by electron cyclotron emission, was typically in the range of 4–5 keV. After the injection of a pellet, the central temperature dropped to between 1 and 3 keV, the higher values occurring for earlier injection. Approximately 10–20 ms after the injection of the pellet, a sawtooth occurred resulting in an additional sudden drop in the central temperature of 0.3–1.0 keV. Electron temperatures ranged from 150 eV to 2.5 keV immediately after the sawtooth, again the higher values occurring for earlier pellet-injection times. The central electron density, as measured using Thomson scattering and Abel-inverted far-infrared interferometry, was typically 1–1.5 \times 10^{19} \text{ m}^{-3} \text{ prior to pellet injection, and } 7–10 \times 10^{19} \text{ m}^{-3} \text{ after injection. The far-infrared system shows that the deposition of pellet material yields a hollow density profile immediately after injection, but the chords are too widely spaced to give detailed profiles near the center of the discharge.}

**Neutron Signals During Pellet Injection**—The neutron fluctuation system on TFTR[10] consists of an array of scintillators connected by fiber-optic light guides to a set of photomultiplier tubes in the basement. The system operates in current mode with a time resolution of 40 \mu s.

At the time of pellet injection, a spike is seen (Fig. 1&2) by the detectors at the same toroidal location as the pellet injector, but not by a detector displaced toroidally by 108°. The spike has a full width at half maximum of approximately 100 \mu s. The average delay between the first spike and the rise of the neutron signal on the toroidally-displaced detector was 140 \mu s. This corresponds to a velocity for the pellet-injected ions to become toroidally symmetric of 3.3 \times 10^4 \text{ m s}^{-1} \text{ and, if this is identified with the ion sound speed,[5] to an electron temperature of 23 eV. After 700 \mu s, the signals on the detectors near the pellet injector and toroidally displaced from the injector coalesce indicating that the deuterium density has become toroidally isotropic. The fast rise in the neutron emission with pellet injection demonstrates that the neutron emission late in time has a significant beam-target component. Following this spike there is a period of decay followed by a period of nearly constant emission which lasted between 7 and 20 ms. For injection earlier than 0.75 s, the period after the initial decay has a slightly negative slope. This nearly-flat period is terminated, for injection earlier than 0.75 s after the beams (Fig. 1), by a sudden decrease in emission and, for later injection (Fig. 2), by a sudden increase in emission (approximately 50% increase) and a quick decay. Following this quick decay is a much slower decay which decreases the emission level by a factor of 100.

At the time of injection, the epithermal neutron system[1] (Fig. 3) also measures a spike of neutron emission and has the same general characteristics as the neutron fluctuation signal, but with a slower time response.

As the neutron signal moves from the period in which DD reactions are believed to dominate, the ratio of the neutron fluctuation signal to the epithermal neutron signal from $^{235}\text{U}$ fission detectors changes by about a factor of two. We interpret this as due to the larger amount of energy deposited in the scintillator when higher energy neutrons dominate the total neutron emission. This effect...
A comparison was also done during the period of enhanced neutron emission following pellet injection. The ratio of the neutron fluctuation and epithermal neutron signals increases with later pellet injection times, indicating that the neutron emission is becoming more dominated by 14-MeV neutrons later in time. This is consistent with the energy-resolving NE-213 system [12] which, despite marginal statistics, indicates that the enhancement of neutron emission is due predominantly to 14-MeV neutrons late in time.

Conclusions—This work contributes more experimental evidence that some fraction of the fusion-produced tritons remain well confined for at least 1.5 s after the turn-off of the beams. For these shots, we interpret the signals to mean that the dominant contribution to neutron production earlier than about 0.75 s after the beams is from the d(d,n)He reaction and, for times later than about 0.75 s, the d(t,n)He reaction is important. [3] The enhanced emission period after pellet injection is terminated by a sawtooth which causes rapid cooling of the central plasma and subsequent slowing of the fast ions. If alpha particles slow down classically, as some fraction of the tritons did in this experiment, alpha heating experiments in near-Q = 1 DT plasmas, such as the alpha-storage mode,[7] should be possible.

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Figures

**FIG. 1.** Neutron signals obtained with the neutron fluctuation system during the injection of a pellet measured with detectors in the same toroidal location as the pellet injector and displaced 108° from the pellet injector. The pellet was injected 0.495 s after the turn-off of the beams. Note that one curve has been offset by a factor of two for clarity.

**FIG. 2.** Neutron signals obtained with the neutron fluctuation system during the injection of a pellet measured with detectors in the same toroidal location as the pellet injector and displaced 108° from the pellet injector. The pellet was injected 0.745 s after the turn-off of the beams. Note that one curve has been offset by a factor of two for clarity.

**FIG. 3.** Neutron signals obtained with the epithermal neutron system during the injection of a pellet at 4.745 s into the post-neutral-beam phase of the TFTR discharge. The slow increase at 4.8 s is due to ohmic heating of the dense, post-pellet plasma.

**FIG. 4.** Measured neutron decay time constant \( \left( \frac{1}{\tau_n} \right)^{-1} \) at the end of the enhanced emission period vs. \( T_{e}^{3/2}/n_{e} \) using the central electron temperature and density from Thomson scattering. The line is the theoretical decay time.
CONVECTIVE HEAT TRANSPORT IN TFTR SUPERSHOTS


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Introduction: The total radial heat flow $Q_a$ and particle flow $\Gamma_a$ for each plasma species can be experimentally determined by transport analysis. If we define $\alpha_a \equiv Q_a / \Gamma_a T_a$, we can test its value against theoretical transport models. For example, the minimum value of $\alpha$ in an ideal gas is $\frac{n}{2}$, while $\alpha_0$ is expected to be very large for "magnetic flutter"-driven transport. $\alpha$ can theoretically have other minimum values due to off diagonal terms in the transport matrix (e.g. heat flows driven by $\nabla n_a$). In particular, for electrostatic turbulence the minimum value of $\alpha$ is $\frac{3}{2}$ [1]. Thus determining the value of $\alpha$ or placing bounds on its value is of fundamental interest in understanding tokamak anomalous transport.

The TFTR high-confinement neutral-beam heated plasmas ("Supershots") [2] have very large central particle loss rates, resulting in convectively dominated central energy transport, and flattened core $T_e$ profiles [3]. Thus, evaluation of the ratios $Q_e / T_e T_a$ and $Q_i / T_i T_i$ in the center of these plasmas yields relevant upper bounds on convective heat transport for each species. For all such plasmas, it is found that $\alpha_e$ in the core region is strictly less than $\frac{5}{2}$, and is typically $\leq 2$. For balanced co- and counter-beam injection, $\alpha_i$ is $\sim \frac{5}{2}$. However, for highly-rotating plasmas with co-only injection, $\alpha_i \leq 2$.

Observations: TFTR is heated by co- and counter-tangential neutral-beams, providing $\sim 20$ MW of $\sim 100$ keV D$^0$ for up to 2 seconds. The deuterium plasmas discussed here are of the enhanced confinement "supershott" type, and have $I_p = 0.8 - 1.2$ MA, $B_T = 4.8 - 5.2$ T, $R = 2.47$ m, and $a = 0.82$ m. With neutral-beam heating power $P_B$ from 10 to 15 MW and near balanced co- and counter-injection they typically have $T_e(0) \approx 8$ keV, $T_i(0) = 20 - 30$ keV, $Z_{eff} = 2 - 4$, moderate toroidal rotation velocity $\nu_\phi \sim 10^5$ m/sec, and very peaked density profiles $n_e(0)/n_e \sim 2.5$, $n_i(0) \sim 6 \times 10^{19}$ m$^{-3}$. These plasmas are fueled only by the neutral beam injection and limiter recycling. No external gas source is supplied. These plasmas are in the collisionless regime with $\nu_{ei} < 5 \times 10^{-3}$, $\nu_{ce} < 5 \times 10^{-2}$ at $r = 0.3$ m. The measured diamagnetic $\beta p_\perp \sim 2$ and Shafranov shift $\sim 0.35$ m (from Thomson scattering density and temperature profiles). Co-only injection into similar plasmas produces very high toroidal rotation velocities (up to $10^6$ m/sec) but lower $n_e(0) \sim 3.2 \times 10^{19}$ m$^{-3}$ and $T_e(0) \sim 6$ keV. Both types of discharges have large amounts of non-ohmic current [4], thus the toroidal loop voltage is very small ($\sim 0.1$ V). While sawteeth are present before and after beam injection, they are suppressed during injection.

The $T_e$ and $n_e$ profiles are measured by Thomson scattering. The $T_e$ profile is also measured by first harmonic ECE radiometry. The $T_i$ and $\nu_\phi$ profiles are measured by charge-exchange recombination spectroscopy [5] using a diagnostic neutral beam or a heating beam. Central values of $T_i$ and $\nu_\phi$ are also measured spectroscopically, using Doppler broadening of Ni XXVII $K_\alpha$ lines. Central $Z_{eff}$ and metallic impurity densities are determined from tangential and radial measurements of visible bremsstrahlung emission and radial x-ray spectroscopy. $Z_{eff}$ is assumed to
be uniform. The particle source rate from recycling on the toroidally symmetric carbon inner limiter is determined from poloidally resolved absolute measurements of $D_\alpha$ emission. The radiated power loss is measured by horizontal and vertical bolometer arrays, and is negligible in the central region of interest here.

**Analysis:** The radial particle and energy flows have been analyzed in these plasmas by the time-independent 1-D transport analysis code SNAP. These plasmas are in equilibrium (for co-injection) or near equilibrium ($S_e(0)/n_e(0) \approx 20$ for the balanced injection cases) at the time of analysis. The code assumes that the flux surfaces are circular, though subjected to the Shafranov shift. The horizontally measured profiles of $n_e$ and $T_e$ are mapped onto the calculated internal flux surfaces by shifting the respective iso-contours to the flux surfaces with the same horizontal minor diameter. This shift is typically $< 3$ cm, indicating good agreement between the calculated and experimental flux surfaces.

Neutral-beam deposition is calculated including the finite size and divergence of the source, and includes deposition on all plasma particles except non-thermal beam ions. Thermalization is calculated in the rotating frame of the thermal ions using an analytic solution to the Fokker-Planck equation [6]. The neutral particle density and temperature profile is calculated using a version of ANTIC [7]. The calculated particle and energy source profiles, Fig. 1, are extremely peaked due to the tangential injection and the low edge density which allows good beam penetration. The “heating effectiveness” parameter [8] $\eta \sim 0.85$ for these discharges. The cross-field particle fluxes are calculated using the continuity equation in equilibrium, the measured $n_e$ profile and $Z_{eff}$, and the calculated beam and neutral particle ionization rates. The local beam-particle source is larger than the recycling source inside $\frac{3}{2}a$. Similarly, the radial heat flows $Q$ are calculated for each species from the energy balance equations, including classical electron-ion coupling. Typical ion power balances are shown in Ref. [5]. These simulations are in good agreement with magnetic measurements of stored energy. The results obtained are in agreement with the analysis by TRANSP of a subset of the plasmas discussed. TRANSP [9] is a $1\frac{1}{2}$-D time-dependent transport analysis code with Monte-Carlo beam simulation including beam-beam deposition and beam charge-exchange loss recapture.

In the co-injection cases the large central $\nu_\phi$ implies that in the plasma’s rotating frame the beam energy is sharply reduced. In the most severe cases both the $E_{b0}/3$ and $E_{b0}/2$ beam energy-species have less energy than $\frac{3}{2}T_i(0)$ in the rotating frame, and thus cool the plasma! It is interesting that the convective term in the torque balance [10] for these plasmas is not large. This is due to the fact that $\nu_\phi$ is always smaller than the injection velocity for any of the beam ions and the assumption that the torque convective multiplier equals unity.

**Results:** Fig. 2 shows the central profile of the ratio $\alpha_\alpha$ for both electrons and ions for a number of plasmas. In all cases, the central electron transport falls below $\alpha_e = \frac{5}{2}$. For co-only injection, some plasmas are nominally inconsistent with $\alpha_e \geq \frac{3}{2}$, though for most the error bars encompass this value in all but the central 0.05 m.

For balanced injection, the ion transport is compatible with $\alpha_i \geq \frac{5}{2}$. However, for co-only injection, $\alpha_i$ is less than $\frac{5}{2}$. For the lowest current co-injection plasmas studied ($I_P = 0.9$ MA) $\alpha_i \sim 1$. However, in these cases the measured $T_i$ profiles are hollow [5], apparently due to the viscous dissipation of the large $\nabla \nu_\phi$. Thus, the central $q_i$ should be carrying heat inwards.
o pp o sing the co nvectiv e loss , and $\alpha_i$ is not a meaningful bound on convective transport. It is still found that $q_i/\nabla T_i > 0$ unless the convective heat flow is assumed $< 2\Gamma_i T_i$, and thus this appears to be a general bound for co-only injected plasmas.

**Conclusion:** The central electron transport in TFTR supershots and related plasmas with co-only injection is clearly inconsistent with $\alpha_e \geq \frac{5}{2}$. This is also true of the ion transport with co-only injection. These results suggest the strong role of off-diagonal transport terms in these plasmas, and appear to rule out stochastic magnetic fields ("magnetic flutter") as the explanation for the transport observed. The difference between the bounds on $\alpha_i$ during balanced- and co-only injection may be evidence of a large (inward) off-diagonal transport term related to toroidal rotation.

We are grateful for discussions with J.D. Callen, H.P. Furth, and D. Meade. This work was supported by the US DOE, contract number DE-AC02-76-CHO-3073.

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Figure 1: (a) Particle and (b) energy source profiles for a typical supershot with balanced injection, $I_p = 0.9$ MA, and $P_{B_e} = 13.6$ MW. The source profiles for co-only injection have similar shapes, but lower magnitude due to the strong toroidal rotation.
Figure 2: Central profiles of $\alpha_a$ for various plasmas: (A) $\alpha_e$ for balanced injection, (B) $\alpha_i$ for balanced injection, (C) $\alpha_e$ for co-only injection, and (D) $\alpha_i$ for co-only injection. The solid lines are discharges with $I_P = 0.9$ MA, short dashes for $I_P = 1.0$ MA, and long dashes for $I_P = 1.2$ MA. $P_B$ ranges from 10.6 MW to 14.6 MW. The typical error bars shown are estimated by varying the input kinetic profiles within the error bars of the individual measurements.
LOW POWER HEATING STUDIES ON TFTR


Introduction

Experiments have been performed on TFTR to study the transport physics of the transition from ohmic heating to low power neutral beam injection. An ohmic density scan from $n_e = 1.2 \times 10^{19}$ m$^{-3}$ to $2.9 \times 10^{19}$ m$^{-3}$ was performed in a plasma with $R_0 = 2.36m$, $a = 0.71m$, $I_p = 1.2$ MA, $q_\psi = 5$. At the high and low density ends of the scan, $\approx 2$ MW of neutral beam power was co-injected into the center of the plasma. By this means we were able to produce ohmic and auxiliary-heated plasmas at the same densities both below, and close to, the "knee" in ohmic $\tau_E$ vs. $n_e$. Electron temperature and density profiles were measured using Thomson scattering at ohmic equilibrium, and at the end of the 0.5s neutral beam pulse. $Z_{eff}$ and metallic impurity content were determined by visible bremsstrahlung and X-ray PHA. Absolute $H_{\alpha}$ data provided a measurement of $\tau_p(a)$, and so a calibration of the absolute neutral density within the plasma. Ion temperature and plasma rotation profiles were measured by means of charge-exchange recombination spectroscopy (CHERS), using a diagnostic neutral beam, for four conditions during this scan: low and high density, ohmic and beam-heated.

Thermal Confinement Results

Figure 1 shows the thermal energy confinement results, where $\tau_{E_{th}}$ indicates ion + electron stored energy ($W_{th} = W_e + W_i$), divided by the total bulk plasma heating power ($P_{heat} = P_{oh} + P_{bi} + P_{th} + P_{be}$). The squares indicate the 4 discharges for which CHERS data were employed. For the remaining discharges the ion temperature was calculated assuming $\chi_i(r) \propto \chi_e(r)$ as an approximation to the CHERS results (see below and ref. 1). The proportionality constant $\chi_i/\chi_e$ was adjusted for each discharge so as to match the measured neutron emission. The ion temperature profile was flattened within the calculated $q=1$ surface by setting $\chi_i$ to a high value in the region of $q<1$. Note that the ohmic heating data cover the range from well below the confinement "knee" to just above it. The modest density dependence of confinement with low-power auxiliary heating that is observed here is exactly what is expected on the basis of the "sum of squares" formulation of the transition to L-Mode. We can also examine this data, with the perspective of "incremental confinement time" and evaluate a thermal $\tau_{E_{inc}} = \Delta W_{th}/\Delta P_{heat}$ at fixed $n_e$. We find $\tau_{E_{inc}} = 82$ msec for the lower density, and 88 msec for the higher density. The fast ion storage is greater, however, at low density as would be expected, resulting in 35% greater $\Delta \beta_p^{dia}$ for the low density case.

Heat Pulse Propagation

Callen et al. have connected $\tau_{E_{inc}}$ during high-power auxiliary heating with a $\chi_{flux}$, defined by the ratio of incremental heat flow to incremental $n_e VT_e$. They further hypothesized that this $\chi_{flux}$ could be identified with $\chi_{e_{HPP}}$, determined from the analysis of sawtooth heat-pulse propagation. We have therefore analyzed the heat-pulse propagation in these discharges, using data from a 20-channel grating polychromator with 200 usec time resolution. We find (figure 2) that the time-to-peak $\chi_{e_{HPP}}$ depends strongly on plasma density, but is independent of heating. Over the density range where
\(\tau_{\text{Pinc}}\) changes very little we see that \(\chi_{e}^{\text{HPP}}\) from the time-to-peak analysis falls by a factor of 1.7, contradicting the hypothesis of ref. 4. In these cases, as is sometimes observed on TFTR, the \(\chi\)'s determined from time-to-peak analysis are much lower than those deduced from phase-shift analysis\(^5\). Under these circumstances the basic hypothesis that sawtooth heat-pulse propagation is due to simple heat diffusion is itself doubtful.

**Transport Analysis**

CHERS ion temperature data are available for four discharge conditions in this scan, as indicated in figure 1. Figure 3 shows \(T_{i}(r)\) and \(T_{e}(r)\) for the low density beam-heated case. In both beam-heated cases ion-electron coupling plays a minor role in the power balances, so \(\chi_{e}\) and \(\chi_{i}\) can be reasonably well distinguished, and the estimated 1-\(\sigma\) error bar on the \(\chi\)'s (averaged over \(\pm 0.1\) m) at \(r=2a/3\) is a factor of 1.5. In the ohmic cases \(T_{i0}\) is well below \(T_{e0}\), but the accumulated errors in \(T_{e}, T_{i},\) and \(\eta_{i}\) result in a factor of 2 or more uncertainty in \(\chi_{i}\), and somewhat less than a factor of 2 uncertainty in \(\chi_{e}\).

\(\tau_{p}\) from analysis of the \(H_{\alpha}\) data was constant over the ohmic density scan, at about 0.11 seconds. In the high density beam-heated case, however, \(\tau_{p}\) fell to 0.06 seconds. The low density case employed a well-degassed limiter (as for TFTR supershots) so \(H_{\alpha}\) emission was reduced by an order of magnitude compared to the higher density beam-heated case. The deduced \(\tau_{p}\) rose to 0.2 seconds during NBI, presumably due to the strongly centrally peaked particle fueling. We have assumed that convective heat flow for both ions and electrons = 1.5nT\(v_{r}\) (as discussed in ref. 1). Analysis of the high density CHERS cases, both ohmic and beam heated, indicates that ion power flow at the edge of the plasma is dominated by convection and charge-exchange. By contrast the well-conditioned low density cases show conduction dominating the ion power balance up to the edge of the plasma.

In all cases transport analysis using the CHERS data indicates that \(\chi_{i}\) has a similar profile shape to \(\chi_{e}\). In the two beam-heated cases we also find that \(\chi_{\phi}\), the calculated radial diffusivity of toroidal angular momentum, is close in value and profile to \(\chi_{i}\). Figure 2 shows \(\chi_{e}\) and \(\chi_{i}\) at 2a/3 for the four cases with CHERS data. \(\chi_{e}\) and \(\chi_{i}\) are also plotted for the other cases, assuming \(\chi_{i}(r)\propto\chi_{e}(r)\) as described above. In ohmic conditions we see that \(\chi_{e}\) drops steadily with increasing density, while \(\chi_{i}\), despite its considerable error bars, can be seen to rise very slightly. At the low density end we see that beam heating causes a severe degradation in \(\chi_{i}\) at 2a/3, but little change in \(\chi_{e}\). At the high density end both \(\chi_{e}\) and \(\chi_{i}\) appear to degrade, and as a result \(\chi_{e}\) during beam injection is found to be independent of density. In all cases \(\eta_{i}\) over the confinement region is found to be in the range 1.5 - 2.5, with the highest value occurring in the low density ohmic case.

**Conclusions**

A number of conclusions can be drawn from this data. First, as noted above, we see that \(\chi^{\text{HPP}}_{e}\) from time-to-peak analysis of sawtooth heat-pulse propagation does not scale with density in the same way as \(\tau_{\text{Pinc}}\). Thus, at least for this low-power TFTR data, a simple connection between thermal \(\tau_{\text{Pinc}}\) and \(\chi^{\text{HPP}}_{e}\) seems difficult to support.

Next we note that both ion and electron thermal transport are subject to degradation with auxiliary heating, and that their ratio during heating is similar to what is observed at high density in ohmic heating. This supports the hypothesis of ref. 2 that these two regimes are physically related. The present results bridge and confirm previous experimental results from Alcator C, in the high density ohmic regime\(^7\), and from D-III\(^8\) with NBI.
The most interesting observations, however, are somewhat indirect. We note that with high power beam heating, at high density, $\chi_e$, $\chi_i$, and $\chi_1$ are similar in magnitude and radial profile shape. This is qualitatively what is expected from electrostatic turbulence, and quantitatively what is predicted for the ratio $\chi_1/\chi_1$ in the specific case of ion-temperature-gradient turbulence. On the other hand, at low density with ohmic heating, and consequently at low $B_p$, we find that $\chi_e$ is larger than $\chi_i$, and scales differently with density. Global particle confinement appears to follow the ion trend, rather than the electron trend. (Perhaps the variation of neutral penetration with density obscures a more electron-like behavior. There is strong evidence, however, from Alcator C that impurity confinement time, $\tau_i$, also scales differently from $\tau_{\text{Ee}}$. In particular $\tau_i$ rises with increasing $I_p$ and with the mass of the background gas.) Qualitatively the separation of electron heat transport scaling from particle and ion heat transport scaling is what would be expected if the electron heat flow were responding to fine-scale electromagnetic turbulence, to which the ions were insensitive. This would cause the slower turbulent diffusion of the ions to determine both the ion thermal diffusivity and the particle transport. Thus what emerges naturally from this data is a picture in which electromagnetic turbulence ("magnetic braiding") with a $1/n_e$ confinement scaling dominates electron thermal transport at low $B_p$, and some form of electrostatic turbulence with a different, L-mode-like, confinement scaling dominates ion thermal transport and particle confinement in all regimes. As $B_p$ grows, the electrostatic turbulence increases until it drives all transport channels, electrons and ions.

Alternatively the low $B_p$ transport mechanism could also be electrostatic in nature, but coupled more strongly to electron than to ion heat flow, while the high $B_p$ mechanism perhaps coupled more strongly to ion heat flow. Such a picture can naturally arise from drift wave theory. In this view the ion transport's independence of density reflects the transition from one transport mechanism to the other, at lower $B_p$ than for the electrons. To differentiate these two pictures, a study of the scaling of $\chi_e$, $\chi_i$, and $\tau_i$ with current, at very low density, is required. If $\chi_i$ and $\tau_i$ rise with plasma current, one would be strongly inclined to believe that the ions always respond with L-mode-like scaling, and that this relatively simple behavior is generally obscured by an additional electron thermal transport mechanism, which dominates global confinement at low density and low $B_p$.

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Fig. 2. '*'s - $\chi_e^{HPP}$, OH; stars - $\chi_e^{HPP}$, NBI; '+'s - $\chi_i$ from OH power balance, neutron analysis; open squares - $\chi_e$ from OH power balance, CHERS analysis; open circles - $\chi_i$ from OH power balance, CHERS analysis; triangles - $\chi_e$ from NBI power balance, CHERS analysis; diamonds - $\chi_i$ from NBI power balance, CHERS analysis; all power balance $\chi$'s evaluated at $r=2a/3$.

Fig. 1. Thermal confinement time vs. line average density. Open squares indicate CHERS data, '+'s indicate $\chi_i$ or $\chi_e$, calibrated to neutron emission.

Fig. 3. '+'s - Te(R) from Thomson scattering, solid boxes - Ti(R) from CHERS. Low density NBI Radiation (cm) Temperature (keV)
Measurements of the toroidal rotation speed in a number of tokamaks have shown the momentum confinement time during neutral beam injection to be anomalously short, by more than an order of magnitude, compared to predictions of neoclassical theory. These observations contrast with the traditional experimental assessment of ion energy transport, which until recently held that the transport was within a factor 1-4 of neoclassical calculations. However, measurements of the $T_i(R)$ profile in Doublet-III [1] and more recently in TFTR [2] with both balanced and unbalanced neutral injection, now suggest that the ion thermal diffusivity, $\chi_i(\tau)$, is considerably larger than $\chi_{i}^{\text{neo}}$ and is comparable to $\chi_e$. Furthermore, measurements of the $v_\phi(R)$ profile in Doublet-III [3] during neutral injection into divertor discharges have demonstrated a striking correlation between global momentum and energy confinement.

In this paper we describe radial transport analysis of recent rotation profile measurements in TFTR under several beam injection configurations: low and high power co-injection, partially balanced injection, and edge-heating. The $v_\phi(R)$ profile measurements were obtained simultaneously with $T_i(R)$ profile measurements using a charge-exchange recombination spectrometer (CHERS) viewing across a radially injected diagnostic neutral beam. We find that the measured velocity and temperature profiles are consistent with a momentum diffusivity, $\chi_\phi$, that has a similar magnitude and similar radial dependence to the inferred $\chi_i$: $\chi_\phi(\tau) \geq \chi_i(\tau)$. This result suggests that a common mechanism, such as ion-temperature-gradient-driven turbulence [4], is spoiling confinement of both ion momentum and energy.

Four of the discharge conditions studied were members of a high power, low-density plasma current scan, $I_p = 0.7 - 1.2$ MA. Near the end of the 2-second heating pulse, plasma conditions were $B_\phi = 4.75$ T, $R = 2.46$ m, $a = 0.81$ m, $q_\phi = 6.5 - 9.0$, $T_e(0) = 4.6 - 6.8$ KeV (increasing with $I_p$), and $T_i(0) = 20 - 25$ KeV. Other plasma conditions are presented in Table I. A marked transient in central rotation speed was recorded at the lower plasma currents ($I_p \leq 0.9$ MA) by a horizontally viewing x-ray crystal spectrometer (XCS) at the start of injection, reaching $v_\phi \geq 1.3 \times 10^6$ m/s within 200 ms, then decaying to a common equilibrium value of $7 - 8 \times 10^5$ m/s at the end of injection. The CHERS velocity profile data were acquired near the end of injection with the exception of the 0.9 MA discharge, for which the data were acquired at 650 ms into the beam heating phase. The steady-state plasma density increased with plasma current, thus providing the largest beam torque per particle for discharges with $I_p = 0.7$ MA. Figure 1a shows that the velocities obtained in
the $I_p$ scan, the highest achieved to date on TFTR, are not exceptional when compared to the database of other TFTR unidirectional injection discharges. The shots with lower plasma current experienced higher $m/n=3/2$ MHD activity [5], which may be responsible for the apparent saturation of central velocity with torque per particle in Figure 1a.

<table>
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<th>$P_b$ (MW)</th>
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<th>$V_\phi$ ($10^8$ m/s)</th>
<th>$\tau_\phi$</th>
<th>$\tau_E$</th>
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<td>20</td>
<td>-</td>
<td>-</td>
<td>4.2</td>
<td>0.7</td>
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</table>

Table 1: Summary of momentum transport analysis. Units: $I_p$ (MA); $P_b$ (MW); $n_e$ ($10^{19}$ m$^{-3}$); $V_\phi$ ($10^8$ m/s); $\tau_\phi$, $\tau_E$ (ms); and $\chi_\phi$ (m$^2$/s). Central velocities were measured by CHERS. The “total $\tau_{\phi E}$” column includes a calculated beam contribution to plasma stored momentum. The last column represents the variation in the ratio $\chi_{\phi E}/\chi_\phi$ over the region $0.3 \leq r \leq 0.6$ m.

The high-power discharge with partially balanced injection was somewhat unusual in that it exhibited the characteristics of TFTR “supershots” [6] ($T_e(0) = 7.5$ KeV, $T_i(0) = 26$ KeV, broad $T_e(r)$, narrow $n_e(r)$) with incomplete balance of the co- and counter beam power; the balance parameter $(\Gamma_{co} - \Gamma_{cntr})/(\Gamma_{co} + \Gamma_{cntr})$ was 0.37. The edge-heating discharge (described more fully in reference [7]) and low-power discharge are similar modest-current, low density shots which differ primarily in the radial deposition profile of beam power and torque; the deposition profiles are centrally-peaked for the low-power discharge and hollow for the edge-heating discharge.

Figure 1b illustrates typical measured velocity profiles during high-power neutral injection. The CHERS velocity analysis does not presently include corrections for the energy-dependent excitation rate coefficients, which causes a $\sim 7-10\%$ overestimate of the rotation speed for data acquired with the diagnostic neutral beam as a doping source, and a $\sim 7-10\%$ underestimate for data acquired with the heating beam.[8] The high-power profiles are very narrow, falling as $\sim (1 - (r/a)^3)^4$ over the inner 40 cm of minor radius. The velocity profiles are broader for the low-power discharges, assuming a roughly parabolic-squared profile shape for centrally-weighted heating and a roughly parabolic profile shape during edge-heating.

The steady-state momentum balance is evaluated by the 1-D radial transport code SNAP. SNAP first maps the measured $v_\phi(R)$, $n_e(R)$, $T_e(R)$, and $T_i(R)$ onto a minor radius grid using a shifted-circle equilibrium with a Shafranov shift of about 25 cm for the high-power co-injection discharges. The fast-ion energy distribution, and power and torque delivery to plasma ions and electrons, is calculated by a moments solution to the Fokker-Planck equation [9] in the rotating plasma frame. The radial variation of $\chi_\phi(r)$ is
Figure 1: (a) Rotation speeds as measured by XCS for discharges in $I_p$ scan (labelled by current), the supershot with partially balanced injection ("SS"), the low-power co-injection discharge ("LP"), and the edge-heating discharge ("E"). The database of 1987 TFTR unidirectional injection discharges with $0.7 \leq I_p \leq 1.3$ MA, $R = 2.46$ m is shown for comparison. Open squares represent co-injection and solid triangles represent counter-injection. Neglect of Shafranov shift and profile effects in the analysis of XCS data underestimates central $v_\phi$ by $\sim$20%. (b) Measured rotation speed profile in $I_p$ scan at 1.0 MA (solid). Shaded region represents the total range of rotation speeds measured in other $I_p$ scan discharges at 0.7, 0.9, and 1.2 MA. Open circles represent the rotation speed in the partially balanced (co/counter) supershot, measured with a heating beam as the doping source.

Figure 2 shows the rotation speed for $I_p = 1.0$ MA, $P_b = 10.7$ MW, mapped onto minor
Figure 2: (a) Measured rotation speed at \( I_p = 1.0 \) MA, \( P_b = 10.7 \) MW, mapped onto minor radius. Calculated torque deposition: \( T_{\text{tot}} = \) total torque (density) from beams, \( T_{\text{bi}} = \) collisional torque to ions, \( T_{\text{be}} = \) collisional torque to electrons, \( T_{\text{th}} = \) beam thermalization torque. (b) Inferred momentum and energy diffusivities for this discharge.

radius, the calculated beam torque densities, and the inferred \( \chi_{\phi} \) and \( \chi_i \) as a function of minor radius. The other high-power co-injection discharges exhibit the same features as shown in Figure 2: the inferred \( \chi_{\phi}(r) \) increases monotonically and strongly with radius, and outside \( r \geq 0.25 \) m, \( \chi_{\phi}(r) \) and \( \chi_i(r) \) have similar magnitudes and similar variations with radius. Despite the fact that both diffusivities are deduced from measurements of a gradient (\( v_{\phi} \) or \( T_i \)), the variation in the ratio \( \chi_{\phi}/\chi_i \) is remarkably small, generally in the range 1-2 (Table I). This apparent correlation of the ion momentum and thermal diffusivities suggests that both are being driven by a common mechanism. The global momentum confinement time is typically \(~1/2\) of the global thermal energy confinement time (excluding the significant calculated beam contributions), roughly consistent with measurements in Doublet-III limiter discharges [1].

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HIGH FREQUENCY EMission FROM TFTR PLASMAS

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Introduction  The study of high-frequency emission from tokamak plasmas is of interest because of the potential use of the observed spectra as a diagnostic of beam ions or fusion products, and because of the possible relation of spectral features to MHD activity. Previous work using probes on the PDX tokamak during neutral beam injection revealed an RF emission spectrum composed of a sequence of harmonically-related peaks occurring at multiples of an ion cyclotron frequency evaluated at the outer plasma edge [1]. The emission was correlated, in some cases, with observations of fast ion losses and with MHD phenomena. RF emission from ohmic plasmas has been observed on JET, using ICRF antennas as receivers, and is thought to be driven by fast ions arising from fusion products [2].

A probe was installed in TFTR near the end of the 1987 run period for investigations of high frequency magnetic fluctuations. The wide bandwidth of the system (1 - 500 MHz) allowed extension of the study of magnetic turbulence to higher frequencies as well as the observation of line emission at high harmonics of ion cyclotron frequencies. This paper presents initial results from the diagnostic.

Experimental Arrangement  The magnetic probe consisted of a two-turn, balanced, shielded loop sensor (Fig. 1). The opposite orientations of the two loops allowed improvement of the electrostatic rejection of the probe by differentially subtracting the common mode signal. The frequency response of the probe was adequately modeled with a simple equivalent lumped-element circuit and the model was verified experimentally.

The probe was installed in an outer midplane port and was oriented to measure $\vec{B}_\phi$. For the data presented in this paper, the sensor loops were located inside the vacuum vessel, 0.10 m past the wall, at a poloidal angle of 18.6° above the midplane. Probe-plasma separation in various discharges ranged from 0.42 to 0.71 m. Signals from the probe were processed near the tokamak with low-noise amplifiers and programmable attenuators. The resulting RF signal was transmitted a distance of 350 m to the TFTR control room via a wideband analog optical link. A swept analyzer was used to obtain frequency spectra and could be triggered at rates of up to 40 Hz. Tests performed with the probe withdrawn into its port tube verified the lack of signal contamination from non-plasma sources.

Low-frequency Fluctuations (1-10 MHz)  At frequencies below ~10 MHz, a continuum of emission was generally observed during the entire tokamak discharge. Calibrated frequency spectra taken at the same time during two successive ohmic discharges that differed in plasma current are shown in Fig. 2. Analyzer bandwidths ($\Delta S$) for data in this section were 30 or 100 kHz, and the spectra are normalized by $\sqrt{\Delta S}$. Plasma parameters common to both shots were: $B_{TF} = 3.7$ T (toroidal field at the vessel center), $R_p = 2.46$ m (plasma major radius), $a_p = 0.81$ m (plasma minor radius), and $n_e \ell \approx 3.5 \times 10^{19}$ m\(^{-2}\) (line-integrated electron density), while the plasma currents ($I_p$) were 1.0 MA and 1.4 MA. The shapes of the two spectra were nearly the same: a region of slowly decreasing amplitude up to about 5 MHz, followed by a more rapid decrease up to 7 MHz, and an indication of a flattening of the spectrum above 7 MHz. The signals from ohmic discharges were generally masked by the background noise level of the system at frequencies above 10 MHz. The amplitudes of the two spectra differed by a constant multiplicative factor up to about 7 MHz. The size of this factor varied from 1.4-1.6, and these values were similar to the ratio of the two plasma currents (1.4). The data presented here show that, within a limited range, the dependence of the spectral amplitude of the low-frequency continuum emission on plasma current was compatible with a linear relation.

The effect of deuterium neutral beam injection on the low frequency continuum is shown in Fig. 3 for four injected powers: $P_b = 0.0, 2.0, 8.1$, and 10.7 MW. The spectra shown were taken 1.0 s after the start of a neutral beam pulse of 2.0 s duration. Plasma parameters for these shots were: $B_{TF} = 4.4$ T, $I_p = 1.2$ MA, $R_p = 2.46$ m, $\alpha_p = 0.81$ m. Line-integrated electron densities for the shots were 1.4, 2.1, 3.0, and $3.3 \times 10^{19}$ m\(^{-2}\), respectively. The amplitude of the spectrum increased with beam power, although
the nearly identical spectra at $P_b = 8.1$ and 10.7 MW suggest a saturation. The spectral amplitude increased, for the highest power, by a factor of $\sim 6$ over the ohmic levels in the region below 5 MHz.

The effect of a change in plasma density alone was also investigated. By varying the gas puffing in two successive ohmic discharges, the plasma density was changed by a factor of nearly two ($n_{e,e} = 1.9, 3.6 \times 10^{19} \text{ m}^{-3}$) while other parameters remained constant. The spectra observed during the two shots were nearly identical, indicating that the spectral change seen during beam injection was not an effect of the increased density.

**High-frequency Emission during Beam Injection (10-500 MHz)**

A harmonic sequence of narrow emission peaks, sometimes extending to hundreds of MHz, was observed during deuterium neutral beam injection into deuterium plasmas. The exact form of the emission depended strongly on plasma conditions and upon the particular complement of beam sources firing, but the peak-to-peak frequency spacing, $\Delta f$, did not. Indeed, as first observed by Buchenauer [1], the frequency spacing during beam injection was compatible with ion cyclotron emission from a region near the outermost plasma edge ($R_{edge} = R_p + \alpha_p$).

Emission spectra from 1 - 500 MHz are shown in Fig. 4a for a shot in which only a single beam source (2.2 MW) was used and in Fig. 4b for a similar shot with two active beam sources (2.2 MW each) from different beam lines. Beam injection was in the direction of the plasma current at tangency radii of 2.5 and 2.6 m. Discharge parameters were: $B_{TF} = 4.7 \text{ T}, I_p = 1.0 \text{ MA}, n_{e,e} = 1.6 - 1.8 \times 10^{19} \text{ m}^{-3}, R_p = 2.31 \text{ m}, \alpha_p = 0.66 \text{ m}$. The spectra were taken 400 ms after the start of the beam injection. Note that the analyzer bandwidth for spectra presented in this and the following section was 300 kHz.

Other beam sources produced other complements of peaks. There was no clear relation between the location or angle of the beam path with respect to the probe and the spectrum produced. However, the regular spacing between peaks was observed to scale directly with $B_{TF}$ and inversely with $R_{edge}$. Figure 5 displays schematically a cross section of the torus, showing the plasma boundaries (from magnetics) for two typical shots and the positions (to scale) of the deuterium and tritium ion cyclotron resonance layers ($\omega = \Omega_p$, $\Omega_T$) calculated using $\Delta f$ and $B_{TF}$ for each shot. In both cases, the calculated position of the $\Omega_p$ resonance was very close to the outermost plasma edge, even though the plasma boundary moved by nearly 0.3 m.

During one experimental run where a fast analyzer sweep was used, emission power from spectral peaks was observed to decay (by a factor of $1/e$) in less than 25 msec after termination of the beam, while the 2.5 MeV neutron flux decayed in about 100 msec. This result suggests that emission peaks seen during beam injection are not related to fusion product populations.

While the heating neutral beams injected only deuterium, the diagnostic neutral beam injected hydrogen. An emission spectrum from a helium discharge into which the low-power ($\sim 100 \text{ kW}$) beam was firing is shown in Fig. 6 (the injection angle was nearly perpendicular to the toroidal direction). The spectrum is similar to that obtained with the heating beams, but now the frequency spacing between the peaks (46.2 MHz) is close to the proton cyclotron frequency evaluated at the outermost plasma edge (45.4 MHz). Since $\Omega_H = 2\Omega_{He}$, it appears that emission from beam injected discharges arises from an interaction of the injected species at the outermost plasma edge.

There was no clear relation between beam power and the magnitude of the emission peaks (cf. Figs. 4 and 6). The widths of the emission peaks, $\delta f/f$, were typically 0.02-0.05 and increased with rising density. At high beam powers, the baseline of the spectrum rose by 10-15 dB. This could be due to a broadening and overlapping of the bases of the harmonically related peaks, or due to very broad emission of other origin. In some discharges, large broadening on the high-frequency side of the fifth harmonic peak was regularly observed together with an absence of distinct peaks above that frequency.

**High-frequency Emission during Ohmic Discharges (10-500 MHz)**

Weak peaks in the emission spectra were seen in deuterium ohmic discharges only at high density ($n_{e,e} > 2 \times 10^{18} \text{ m}^{-3}$). The spectra for two discharges with different densities ($n_{e,e} = 1.9, 3.5 \times 10^{19} \text{ m}^{-3}$) are shown in Fig. 7. For these shots, $B_{TF} = 4.4 \text{ T}, I_p = 1.2 \text{ MA}, R_p = 2.36 \text{ m},$ and $\alpha_p = 0.71 \text{ m}$. The peak emission power was only $\sim 10 \text{ dB}$ above the thermal noise of the electronics, and some structure may have been hidden beneath the noise floor. The two clear peaks that emerged at high density were much broader ($\delta f/f \sim 0.1 - 0.2$) than those seen during beam injection. The frequency spacing of these peaks also scaled linearly with $B_{TF}$, but the magnitude of the spacing ($\sim 60 \text{ MHz for the case shown}$) corresponded to the cyclotron frequency of protons ($\sim 58 \text{ MHz}$) evaluated at the outermost plasma edge. This result may indicate the involvement of fusion products, which could provide a source of fast (3.0 MeV) protons.
The ratio of the emission peak widths for the beam-injected and ohmic cases is consistent with the ratio of Doppler widths expected for 100 keV beam deuterons and 3.0 MeV protons.

**Conclusions**  The spectrum of high frequency (1-500 MHz) magnetic fluctuations has been investigated with a probe in the TFTR tokamak for ohmic and neutral beam heated plasmas. A continuum of emission was found at frequencies below ~10 MHz, and the amplitude of these fluctuations increased with beam power or plasma current but not with density. This emission could be related to the broadband MHD turbulence seen in TFTR at frequencies below 1 MHz [3]. A sequence of harmonically related peaks was observed at high frequencies during neutral beam injection. Varying $B_{TF}$, $R_p + a_p$, and the beam species demonstrated that the spacing between peaks corresponded to the local cyclotron frequency of the injected beam species at the outermost plasma edge. Broad, weak, high-frequency emission peaks were observed in high density ohmic deuterium discharges, and the peak spacing also scaled linearly with $B_{TF}$ but corresponded to the proton cyclotron frequency at the outer edge. The above results differ from observations on JET-2, where emission from ohmic and hydrogen neutral beam heated deuterium discharges was seen at harmonics of the deuteron cyclotron frequency.

At the fusion reactivity levels of present TFTR plasmas, it appears that possible contribution to the RF emission peaks from fusion products is obscured by the larger signals produced by the beam ions themselves. During the upcoming D-T experiments, however, it is possible that contributions from the significantly enhanced fusion product populations may dominate the RF emission spectrum.

**Acknowledgement**  This work was supported by US DOE Contract No. DE-AC02-76-CHO-3073.

**References**


**Figure Captions**
1. Balanced magnetic loop probe.
2. Low-frequency spectra for two ohmic discharges differing in plasma current.
3. Dependence of low-frequency spectra on neutral beam power.
4. High frequency emission spectra during neutral beam injection using (a) one and (b) two sources.
5. Location of plasma boundary and deuteron and triton cyclotron resonance layers calculated from $\Delta f$ and $B_{TF}$ for small and large plasmas.
6. High frequency emission spectra with neutral (hydrogen) beam injection and helium target plasma.
7. Emission spectra during ohmic discharges for $n_e \ell = 1.9 \times 10^{18}$ m$^{-2}$ and $3.5 \times 10^{19}$ m$^{-2}$.

**Figure 1** Stainless Wire
0.5 mm dia.
Crushed Alumina
Stainless Tube 1.7 mm o.d.

**Figure 2**

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**Ohmic Plasmas**
Ignition Conditions

For a D-T plasma the condition for thermonuclear ignition with \( T_e = T_i \) is

\[
\bar{n}_D \bar{T}_i \tau_e = 5 \times 10^{21} \text{ m}^{-3} \text{ keVs} \quad (1)
\]

This is valid for the range \( 7 < T_i < 20 \text{ keV} \) in which the \( \bar{G} \bar{T} \) product for thermonuclear reactions is proportional to \( T_i^2 \). Condition (1) is insensitive to the pressure profile.

Applying energy balance arguments the condition (1) can be rewritten as,

\[
P \tau_e^2 \frac{f V}{V_p} = 2.4 \quad (2)
\]

where \( P(\text{MW}) \) is the total input power to the plasma, \( \tau_e(s) \) the global energy confinement time, \( V(\text{m}^3) \) the plasma volume, \( f \) the correction due to depletion by impurities and \( V_p \) the ratio of central to volume average plasma pressure.

Confinement Time

The JET data for limiter-bounded discharges is well described by the L-mode scaling law [1], namely

\[
\tau_e = 3.7 \times 10^{-2} I_p P^{-1/2} R^{1.38} (R/a)^{0.37} K^{1/4} \quad (3)
\]

where \( I_p(\text{MA}) \) is the plasma current, \( K \) the plasma elongation, \( R(\text{m}) \) and \( a(\text{m}) \) the major and minor plasma radii. A study of JET and other data shows that the scaling with \( I_p \), \( P \) and \( R \) is well established while the scaling with \( R/a \) and \( K \) is doubtful.

For H-mode data in JET the simplest representation is that \( \tau_e \) is twice the L-mode value. (JET Team, 1987)

Depletion Factor

Assuming that the impurities in the central plasma are fully ionised and serve only to deplete the reactive fuel density, then

\[
f = \left( \frac{Z_i - Z_{\text{eff}}}{Z_i - 1} \right)^2 \left( \frac{2Z_i}{2Z_i + 1 - Z_{\text{eff}}} \right)^2 \quad (4)
\]

where \( Z_i \) is the average charge of the impurities and \( Z_{\text{eff}} \) the effective
charge. For JET the dominant impurities in the centre are carbon and oxygen, so if we take a mean \( Z_i = 7 \) then for a plausible \( Z_{\text{eff}} = 2 \), we have \( f = 0.8 \). In JET the strongest dependence of \( Z_{\text{eff}} \) is on density [2]. An ignition experiment must operate with \( Z_f < 2 \), which from JET data requires a plasma density greater than \( 5 \times 10^{19} \text{m}^{-3} \).

**Pressure Ratio**

The pressure ratio \( \Gamma_p \) versus the safety factor \( q \) for JET limiter discharges varies linearly from 2.5 at \( q = 2 \), to 4.3 for \( q = 4 \). It then flattens out and \( \Gamma_p = 4.5 \) at \( q = 11 \). There is \( \pm 25\% \) scatter on \( \Gamma_p \) due to sawtooth relaxations. The so-called "monster" sawtooth discharges are at the top of the observed band whilst \( \text{H-modes} \) are at the bottom. From (2) and (3) it is clear that to maximise the performance the ratio \( (\Gamma_p/q)^2 \) that should be maximised. This occurs at the lowest \( q \) at which the discharge can operate, the gain from higher current overwhelming the lower value of \( \Gamma_p \).

**Plasma Parameters**

Substituting from (3) into (2) the condition for ignition is

\[
f \Gamma_p \gamma_g^2 I^2 R^{2.5} a^{-2.74} = 3.46 \times 10^4
\]

where \( \gamma_g \) is the ratio of the confinement time to the \( \text{L-mode} \) value.

The current \( I \) is related to the machine parameters by the expression,

\[
I = \frac{5 a^2 B}{K q_p} \left( \frac{1 + K^2}{2} \right) \left( 1 - \frac{a^2}{R_k^2} \right)^{-2}
\]

where \( q_p \) is the Shafranov q.

**Enhanced Tokamak**

We now estimate the parameters of a JET-like device upgraded to reach ignition conditions based on the JET experimental results as described above. Such an apparatus would permit the early study of a wide range of physics questions related to the achievement and the behaviour of an ignited plasma. In contrast to CIT this Enhanced Tokamak would be on the direct path to an eventual tokamak reactor, ie having moderate magnetic field strength and large physical size.

No significant engineering design work has yet been done on this device but a representative set of parameters is

\( B = 5 \text{T}, R = 4 \text{m}, a = 1.4 \text{m}, K = 2, \gamma = 3, f = 0.8, q_p = 2.5 \) giving \( I_p = 15.9 \text{MA} \) and for ignition \( \gamma_g = 2.2 \). These are a plausible set of parameters in the light of the JET experimental results.

The values of the \( \alpha \)-power and the energy confinement time at ignition are determined by the density. Combining the \( \text{L-mode} \) scaling with \( T \sim 10 \text{keV} \) leads to the formula for the \( \alpha \)-power [3].

\[
P_\alpha = \frac{30 M^2 a^{0.74} q_{\text{CYL}}^2}{\gamma_g^2 R^{1.5} K}
\]
where \( q_{\text{CYL}} = \frac{5}{2} \frac{a^2}{R} \frac{B}{T} \) and \( M = \frac{n \cdot R}{B} \times 10^{-19} \)

from the JET experience a realistic value for \( M \) when operating with high input power at low \( q \) is \( n = 6 \). With the parameters given earlier, \( n = 7.5 \times 10^{13} \) and \( P = 42 \text{MW} \) while from (3), \( T_e = 2.7 \text{s} \). With these parameters \( \beta = 2.4% \) compared with the Troyon limiting value \( \beta_c = 6.4% \) (\( \beta_c = 2.8 \text{I/BAZ} \)).

Comparison with other Devices

The table shows the main parameters of JET, Enhanced Tokamak, CIT and extended NET.

<table>
<thead>
<tr>
<th>Apparatus</th>
<th>R (m)</th>
<th>a (m)</th>
<th>K (T)</th>
<th>I (MA)</th>
<th>( q_{\text{CYL}} )</th>
<th>( q_r )</th>
<th>( \gamma_r ) for ignition</th>
</tr>
</thead>
<tbody>
<tr>
<td>JET</td>
<td>3.5</td>
<td>1.2</td>
<td>1.6</td>
<td>3.5</td>
<td>7</td>
<td>1.9</td>
<td>3</td>
</tr>
<tr>
<td>Enhanced Tokamak</td>
<td>4.0</td>
<td>1.4</td>
<td>2.0</td>
<td>5.0</td>
<td>16</td>
<td>1.5</td>
<td>2.5</td>
</tr>
<tr>
<td>CIT</td>
<td>1.8</td>
<td>0.5</td>
<td>2.0</td>
<td>10.0</td>
<td>9.0</td>
<td>1.5</td>
<td>2.4</td>
</tr>
<tr>
<td>NET-E</td>
<td>5.4</td>
<td>1.7</td>
<td>2.2</td>
<td>4.8</td>
<td>14.8</td>
<td>1.8</td>
<td>3</td>
</tr>
</tbody>
</table>

The figures are given for \( Z_{\text{eff}} = 1 \) for the sake of comparison. Evidently Enhanced Tokamak is superior to CIT and similar to NET-E in performance.

Zero-Dimensional Code

So far we have used very simple arguments which lead to ignition independent of the plasma density provided only that the ion temperature lies in the range of 10-20keV. The density was fixed by appealing to the experimental experience of the Murakami parameter.

Steady state and dynamic calculations have been carried out using a more complex model in which,
1. The proper variation of reaction rate with temperature is included over the entire ion temperature range.
2. The confinement time is determined by the combination of the ohmic (neo-Alcator) value \( \tau_{\Omega} \) and the L-mode \( \tau_{\text{EL}} \) so that
\[
\tau_e = \left( \frac{1}{2} \tau_{\Omega} + \frac{1}{2} \tau_{\text{EL}} \right)^{-0.5}
\]

This has the effect of reducing the confinement time at low density.
3. The density and temperature are assumed to have parabolic profiles giving \( \gamma_{\text{C}} = 3 \).
4. The full energy balance equation is used, ie \( P_{\text{OH}} + P_{\text{AUX}} = P_{\text{COND}} + P_{\text{br}} \)

where \( P_{\text{OH}} \) is the ohmic power input, \( P_{\text{AUX}} \) the auxiliary heating, \( P_{\text{COND}} \) the conduction loss and \( P_{\text{br}} \) the Bremsstrahlung loss. This code has been run for the parameters considered earlier (including \( \gamma_{\text{C}} = 2.2 \)). The ignition curves in the \( n \) vs \( T \) plane for \( P_{\text{AUX}} = 0 \) are shown in Figure 1. The ignition domain, shaded in the figure is bounded by upper and lower temperatures and by a minimum density determined by the density dependence of the ohmic confinement time. If a weaker dependence is chosen, as for example derived from JET ohmic data [4] then this minimum density would be much lower.

Analysis of the results with finite \( P_{\text{AUX}} \) shows that a heating power of order 30MW would be sufficient to take the plasma to ignition.
Machine Parameters

The outline of the machine design can be discussed in terms of the parameters:

- $R_0 =$ geometric major radius of torus
- $\varepsilon_3 =$ clearance between plasma boundary and toroidal field coil
- $\varepsilon_2 =$ thickness of toroidal field coil
- $\varepsilon_1 =$ thickness of primary winding

Then the maximum toroidal field in the system $B_{\text{MAX}}$ is

$$B_{\text{MAX}} = \frac{B_0 R_0}{(R_0 - a - \varepsilon_3)}$$

Taking $B_0 = 5T$, $R = 4m$, $a = 1.4m$ and $\varepsilon_3 = 0.2m$ gives $B_{\text{MAX}} = 8.3T$ (cf 7T in present JET). With $\varepsilon_2 = 0.5m$, $\varepsilon_1 = 0.5m$ and the same $B_{\text{MAX}}$ in the primary winding we find

$$\text{Volt seconds} = 2\pi(R_0 - (a + \varepsilon_3 + \varepsilon_2 + \varepsilon_1))^2 B_{\text{MAX}} = 142$$

This is an adequate number of volt seconds.

The current required in the toroidal field coils to produce 5T at 4m radius is 100MA. The cross-sectional area of the coils in the mid-plane is 6.75m$^2$ giving a mean current density of 1.5kA/cm$^2$. Scaling from JET the resistive power dissipation would be 600MW. The flat-top pulse length would be similar to that of JET, namely 20s.

References

LOCAL HEAT TRANSPORT IN JET PLASMAS


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INTRODUCTION

Two complementary approaches to the study of heat transport in JET plasmas are described in this paper. In the first approach, 1-D simulations have been performed using two different forms for the heat fluxes which are based on global scaling laws, have the scale invariance of theoretical models and represent the two generic models for the heat fluxes[1] capable of reproducing the main features of the plasma response to additional heating. Even though the mechanisms for confinement degradation are different, both forms can reproduce a wide range of data from JET and other machines. It is found with both models that the data can only be matched if the electron and ion transport coefficients have the same form and approximately the same magnitude.

In the other approach which will be described here the parametric dependence of the heat flux on local quantities has been obtained by fitting the data with a form with the scale invariance of the non-linear gyrokinetic equation of Frieman and Chen[2]. A good fit is obtained to a set of data obtained from neutral beam heated discharges with \( I_p = 1–5 \, \text{MA} \).

SIMULATIONS USING DIFFUSIVITIES WITH MODEL SCALE INVARIANCE

An extensive programme of simulation of JET discharges has confirmed the result[3] that models of the electron heat loss based on the dissipative trapped electron mode fail to reproduce \( T_e \) profiles. A similar conclusion has been reached concerning the Pogutse/Parail[4] and \( \eta_e \) models[5] by a direct comparison of the electron conductivity with the experimental values.

Two local transport models have been proposed at JET which provide the degree of profile resilience and confinement degradation required by the experimental data, are related to scaling laws for \( \tau_E \) that fit the results from several devices and are sufficiently complete to be used in transport codes.

The first of these models[6] produces confinement degradation through a transport threshold at a critical value of temperature gradient. For simulations we have used
electron and ion heat fluxes
\[ q_{eRL} = \frac{-n_e x_{eRL}(\nabla T_e - \nabla c)}{\nabla c} - n_e x_{e1R} \nabla T_e \quad \text{for } \nabla T_e > \nabla c T_e \]
\[ q_{iRL} = \frac{K_{iRL} n_i \nabla T_i / n_e \nabla T_e (T_i / T_e)^{1/2}} {\text{otherwise}} \]
where
\[ x_{eRL} = \frac{x_{e0R}(T_e q / B_T R_i^{1/2} T_i)(1 + 2 \ln(n_e) / \ln(T_e))(C_1 + R|q||q|^{-1})}{\text{where}} \]
\[ V_{cT} = \frac{\alpha(E / B_T / n_e \nabla T_e)^{1/2} B_T / q} {\text{and}} \]
If SI units are used with keV for temperatures, then \( \alpha \approx 6, C_1 = 2, x_{e0R} \approx 20 \) and \( x_{e1R} \approx 0.1 \) give a good match to the data. In order to reproduce \( T_i0 \), the total energy content and the NPA ion temperature profile, where available, it is found that \( K_{iRL} \approx 1 \) is required.

The second model is based on a resistive MHD scaling law for \( \tau_e \) and the confinement degradation occurs because of the nonlinear dependence of the heat flux on the pressure gradient. We have used
\[ q_{ePT} = \frac{-n_e x_{ePT} \nabla T_e}{\nabla T_e} \]
\[ q_{iPT} = \frac{-n_i x_{iPT} \nabla T_i}{\nabla T_i} \]
\[ x_{ePT} = \frac{C_{PT}/\mu_0 (\mu_0 / 2 n_A / q R)^{1/2} (\nabla p / \nabla c)} {\text{where}} \]
where \( \sigma \) is the Spitzer electrical conductivity, \( v_A \) the Alfvén speed and \( V_{cP} \) the threshold pressure gradient for ideal ballooning instability. \( C_{PT} \approx 0.035 \) has been found to reproduce a wide range of data.

Both the above models have been used in a simple 1-D time independent code which is used for parameter scans. This code solves the ion and electron temperature equations and models the ICH and NBH fast ions. The heating deposition, electron, ion and current density profiles are represented by analytic forms. For the dynamics of particular discharges a full time dependent 1\( ^{1/2} \)D transport/equilibrium code JETTO is used. JET data is used for some quantities such as the density and radiation profiles.

Some uncertainty remains in the values of \( x_{e0R} \) and \( C_{PT} \). This largely because of the insensitivity of the calculated quantities such as temperature and density to these parameters due to non-linear effects and because of experimental errors. For some conditions the transport inside \( q=1 \) must be modified. The results of the simulation of a range of JET discharges such as the hot electron, hot ion and off-axis heating will be presented. Figures 1 and 2 show the model electron temperature profiles with the ECE data for 2MA/2.5T on- and off-axis ICH discharges.

**EVALUATION OF THE THERMAL DIFFUSIVITY FROM HEAT FLUX DATA**

The local heat flux in JET NBH discharges has been studied with the aim of determining values of local heat diffusivity[8]. No attempt is made to separate the total flux into an ion and electron heat flux since no ion temperature data is available. Similarly no allowance is made for the radiation loss channel; the latter can for some JET pulses be shown to be small compared with the total flux at plasma radii less than
0.85—0.95. Thus the total flux $q_{th}$ includes the Ohmic and NBH power inputs.

The data on $q_{th}$ from 350 observations is analysed in terms of models $q_f$ and it has been shown that\cite{7,8}

$$q_f = -en_xV_T - II$$

represents the best fit to the data. The temperature gradient $V_T$ is determined from ECE and the density is measured by interferometer; the heat pinch is II. Depending on how the data is selected various estimates of the heat diffusivity $\chi$ can be arrived at. One definite pattern emerges from the data: $\chi$ shows a pronounced dependence upon the current $I_p(x)$ enclosed within a surface $x$; $\chi \approx 1/I_p(x)$.

By applying the similarity techniques of Connor and Taylor\cite{9} to the gyrokinetic equation of Frieman and Chen\cite{2}, assuming that small space and short time scale turbulence is responsible for local transport, the diffusivity is constrained to the form

$$\chi = (V_p^2/L) F(\nu L/\nu, \beta)$$

where $\rho$ is the Larmor radius, $L$ is some scale length, $\nu$ a collision frequency, $v$ a thermal velocity and $\beta$ the plasma beta. The function $F$ can be arbitrary. In practice we have employed

$$\chi = C \left[ \frac{T_0^3}{L_p} \right]^{\frac{1}{2}} \left[ \frac{n L}{T_e} \right] \left[ \frac{\alpha}{\nu} \right] \left[ \frac{T_P}{L} \right] \gamma$$

and apply non-linear regression techniques to the data on $q$, $n$, $T$ etc.. It is found that $C=10 \pm 0.4$, $\alpha=0.4 \pm 0.02$, $\gamma=-0.42 \pm 0.02$ and $II=2.6 \pm 0.5$ represent the optimum values. This essentially implies that $\chi \approx 1/I_p(x)$. It should be stressed that there is a 32% scatter in the fit of $q_f$ to the data $q_{th}$. A plot of the best fit against the data is shown in figure 3 and the scatter is evident.

Many approximations are made and the data contains scatter not only from the measurements but also from sawtooth oscillations. To improve on the latter we are repeating the above study for 150 JET ICH discharges. These discharges have been selected as they do not have sawtooth oscillations ("monster sawteeth"). The results from the analysis will be discussed.

REFERENCES

[8] J.P.Christiansen et al., JET Preprint JET—P(87)48 (to be published)
FIGURES 1 AND 2: The $T_e$ and $T_i$ profiles obtained using the two transport models compared with the experimental $T_e$ data from ECE. Both discharges were at 2MA/2.5T with 7MW of ICH power. Figure 1 shows the result of heating on-axis whilst figure 2 corresponds to the ICH being absorbed at the half minor radius.

FIGURE 3: The experimental value of the heat flux $q_{th}$ versus the fitted value $q_f = -\chi n V T_e - II$ where $\chi$ is constrained to a form compatible with the scale invariance of the non-linear gyrokinetic equation.
GLOBAL CONFINEMENT CHARACTERISTICS OF JET LIMITER PLASMAS

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Abstract

Data from a wide variety of plasma pulses on JET (aux. heating, current, field, minority species, plasma shape, etc) are analysed in order to assess the characteristics of global confinement. The scaling of confinement in ohmically and auxiliary heated discharges is examined. The ohmic confinement in the present new JET configuration (Belt Limiter) is essentially the same as previously. Confinement in auxiliary heated discharges shows presently a slight improvement since 1986. Both ohmic and non-ohmic data is used in a set of confinement time regression analyses and certain constraints derived from theory are imposed.

1. Ohmic Heating in JET

Previous work [1] on confinement in JET ohmic discharges was found to be best described by an empirical scaling law

\[ \langle T_e \rangle \sim B_\phi^a n^b q^c K^d \]  

where \( T_e \) is electron temperature, \( B_\phi \) toroidal field, \( n \) is electron density, while \( q \) is safety factor of a plasma configuration with elongation \( K \); the brackets \( <> \) denote a volume average. The scaling parameters \( a, b, c, d \) were for 1984-85 data 1.8, -0.6, -1.0, 0.8 respectively. Eq.(1) is a consequence of Ohm's Law and similar expressions can be found for \( Z_{eff} \) and \( T_E \), the global energy confinement time. The ohmic data obtained from experiments on JET in 1986, 87 and 88 is still described by scaling laws of the type (1). Figure 1 shows a plot of \( \langle T_e \rangle \) measured by ECE versus the expression (1). The fitting parameters \( a, b, c, d \) are now found to be 1.3, -0.55, -1.08, 0.68 respectively; the change in the exponent \( a \) is caused by a greater variety of \( B_\phi \) values during 1986-88; in 1985 \( B_\phi \) was mostly 3.4T. The data in Figure 1 includes 1984-85-86 data with toroidal limiters and 1987-88 data with belt limiters. The introduction of belt limiters have resulted in generally higher values of elongation \( K \). The full data set (1984-88) is now more evenly balanced as regards variations in \( K \) and \( B_\phi \) and this increases the level of confidence in the scaling law (1).

2. Decoupling of Ions from Electrons

The auxiliary heating systems (NBI and ICRH) on JET give access to two regimes in which the ion and electron temperatures are decoupled for periods longer than the global energy confinement time \( \tau_E \). Neutral beam
heated discharges following extensive helium discharge cleaning in 1986 yielded low densities (inner wall pumping) and subsequent hot ion modes [2] for which $T_i \approx 2T_e$. During 1987-88 ICRH discharges with hydrogen minority and up to 16 MW of power have produced $T_e \approx 2T_i$. Figure 2 shows that substantial temperature differences $T_i - T_e$ occur only when the ratio $\tau_{eq}/\tau_2$ exceeds 1; this result demonstrates that energy equipartition in JET is classical and is characterised by the time constant $\tau_{eq}$. It is worth noting that the confinement time $\tau_2$ for the data in Figure 2 is largely independent of the sign of $T_i - T_e$ (but not the magnitude). Since $T_i - T_e = T_i$ or $T_e$, the independence of $\tau_2$ upon the sign of $T_i - T_e$ provides some constraints on future Tokamak confinement theories: neither of the heat diffusivities $\chi_e$ or $\chi_i$ can have a strong temperature dependence.

3. Fast Ions from ICRH

It is possible to establish trends in JET data from power-density scans in ICRH heated discharges with fixed current-field (2 MA, 2.2 T). The energy content $W_{dia}$ measured by the diamagnetic loop estimates the total energy $W_i$ (thermal and non-thermal).

At constant ICRH power $W_{dia}$ decreases as $n$ increases, while the energy content $W_{kin}$ does not decrease with $n$. $W_{kin}$ estimates the total thermal energy from ECE, interferometer and X-ray crystal spectrometer data. Because of this density dependence we can test the data on $W_{dia}$ against the fast energy content predicted by theory, i.e.,

$$W_i = \frac{2}{3} (W_{dia} - W_{kin}) = P_{RF} \frac{\tau_s}{2}$$

where $\tau_s$ denotes the classical slowing down time $\tau_s = 0.075T_e^{3/2}/n$ (for $He^3$ minority). The data values in Figure 3 of $\frac{2}{3} (W_{dia} - W_{kin})$ do exhibit the dependence (2) upon $P_{RF}$ and $\tau_s (T_e, n)$. The scatter arises from taking differences between measurements which include both random and systematic errors. Indeed to achieve the level of agreement depicted in Figure 3 it has been necessary to reduce the ECE measurements systematically by 20%; this reduction brings together the ECE and LIDAR measurements of electron temperature. The value of $\tau_s$ is based on an average over a fraction of the plasma volume.

4. On-Off Axis ICRH Heating

In a sequence of 10 ICRH pulses the resonance positions for each of the 8 JET antennae has been moved from on-axis ($R = 3$ m) to off-axis ($R = 3.5$ m). The confinement properties of plasmas with different ICRH power deposition profiles can be described by an offset-linear law

$$W = W(0) + (\tau P)_{on-axis} + (\tau P)_{off-axis}$$

The incremental confinement time $\tau$ and power level $P$ associated with on-off axis heating allow for a range of increments. Figure 4 shows measurements of $W$ plotted against the values predicted by (3). Eq.(3) gives a very good
description of the data from the 10 ICRH pulses which feature 100% off-axis as well as 100% on-axis heating; it is based on the transport models of
[3]. The $\tau_{\text{on-axis}} = 0.21s$ and $\tau_{\text{off-axis}} = 0.15s$ showing that the confinement depends on the location of the heating power as shown in [3].

5. Confinement Scaling

Data obtained from JET experiments during 1984-1988 (approximately 3000 pulses) is employed in regression analyses as described earlier [1,2]. An empirical expression for the confinement time $\tau_E$ like

$$\tau_E = 0.3 \Phi_{\text{pol}}^{0.1} n^{0.2} P^{-0.5}$$

emerges as the best fit to the entire data. In Eq.(4) the units are $I$(MA), $n(10^{19} \text{m}^{-3})$ and power $P$ in MW. The constraints imposed by applying the Connor-Taylor scale invariance techniques to the gyro-kinetic equation derived in [4] result in the scaling law

$$\tau_E = \rho^2 \frac{v}{L} F (\frac{vL}{v}, \beta)$$

(5)

$\rho$ is Larmor radius, $v$ is thermal velocity, $L$ is a length scale, $v$ is a collision frequency and $F$ is an arbitrary function. If the latter is represented by a power law expression such that

$$\tau_E = \rho^2 \frac{v}{L} (\frac{vL}{v})^\alpha \beta$$

(6)

then a wide range of scaling laws are possible depending upon the choice of $\rho$, $v$, $L$, $\alpha$, $\beta$; eg. $\rho_e$ or $\rho_p$, $v_e$ or $v_i$, $L$ or $V_n/n$, $v = v_e$ etc. Eq.(6) provides for a good test against JET data since $vL/v$ chosen as naZe/Te varies by a factor 20 while $\beta$ chosen as $nT_e/I_{\Phi}^2$ varies by a factor 10. An a priori choice of $\alpha$ and $\beta$ establishes a set of scaling laws: $\alpha = \frac{1}{2}$, $\beta = -\frac{1}{2}$ yields the offset linear law of [2] and gives a good fit; so too do electrostatic models with $\beta = 0$: $\alpha = 0$, $\beta = -1$ (T-11 scaling) and $\alpha = 1$, $\beta = -1$ (resistive fluid turbulence) do not fit the data. Applying regression analysis one finds; $\alpha = 0.3$, $\beta = 0.1$ from the best fit correspond to

$$\tau_E = 0.22 \Phi_{\text{pol}}^{1.05} n^{0.3} P^{-0.5} B_{\Phi}^{0.15}$$

(7)

We notice that (7) is essentially the empirical fit (4).

Acknowledgement

The authors gratefully acknowledge the skill and assistance provided by the ICRH and diagnostic teams on JET.

References
Fig. 1 $\langle T_e \rangle$ vs values predicted by eq (1)

Fig. 2 Temperature differences in decoupled plasma regimes

predicted value

Fig. 3 Fast ion energy vs Power $T_{ave}/2$ in ICRH discharges

Fig. 4 Total energy vs the scaling law (3)
HIGH CURRENT OPERATION IN JET


ABSTRACT. Plasma currents of 6MA have been obtained in JET with a limiter configuration and 4.0MA have been obtained with a magnetic separatrix which was well separated from the limiters. A 4.5MA X-point discharge has been produced, but the separatrix was close to the belt limiter. These have been made possible by improvements to the primary configuration which reduce the stray field at breakdown, increase the volts-second capability and enhance the shaping effect due to the primary. Scenarios have been tested which demonstrate the feasibility of 7MA limiter operation and 5MA X-point operation.

The physics issues to be discussed are magnetic configuration, breakdown, current penetration, stability and volts-seconds consumption. The confinement properties of limiter and X-points will be compared.

1. HIGH CURRENT LIMITER OPERATION

1.1 Breakdown. Reliable breakdown can be achieved with \( V \approx 10-30 \) Volts. At maximum premagnetisation, where the stray field is \(-2E-2\) Tesla, prompt breakdown is obtained provided the fields are tuned to give an hexapolar null in the centre of the chamber. In this case the volts seconds loss at breakdown is less than 1 volt-sec. Careful preprogramming of the vertical field is required until the plasma fields dominate the stray fields.

1.2 Current Penetration and Stability. The empirical stability diagram for JET [1] is shown in Fig.1 with the trajectory of a 6MA discharge superimposed. It is necessary to programme the current ramp rate, elongation, minor radius, gas feed and toroidal field so as to avoid the lower boundary where the rotating MHD which arises at rational \( q = 3 \) reaches sufficient amplitude to lock. These locked modes are particularly dangerous because they often persist from early times in the discharge and later grow to cause disruptions at high current. In the example of Fig.1 a current ramp as slow as 0.25MA/sec was necessary to allow the current to penetrate and thereby reach 6MA at \( q = 3 \). Although the rotating modes are clearly visible these do not lock. At constant \( q < 3 \), larger current ramps are possible. Thus by simultaneous ramp of plasma current and toroidal field the discharge shown could safely be extended from 6MA/2.8Tesla to 7MA/3.4Tesla.

1.3 Volts-Seconds Consumption. Fig.2 shows the total flux swing at the plasma boundary together with the inductive flux and resistive flux on axis
as functions of plasma current for the same 6MA discharge illustrated in Fig.1. The maximum volt-sec available is also shown for maximum premagnetisation for the two cases (a) maximum uniform reverse current in primary stack and (b) maximum reverse current in central pancakes and maximum reverse current in top/bottom pancakes. For the first 5 seconds the resistive consumption is negligible but rises as \( q_0 \) reaches 3 at 4.5MA. In this discharge the current ramp is only 0.25MA/sec and a total of 6 volt-sec is consumed resistively by the start of the flat top. Clearly at this rate there are insufficient volt-sec to reach 7MA at low \( q_0 \). However a current ramp of 0.5MA/sec with simultaneous toroidal field ramp has been demonstrated at low \( q_0 \) up to 6MA. This scenario has sufficient volt-sec left to reach 7MA. For \( q_0 > 3 \) the resistive volt-sec is negligible during the current rise and at 5MA a flat top of 9 sec has been demonstrated.

1.4 Heating During the Current Rise. RF heating has been applied during the current rise before the onset of sawteeth. Very high electron temperatures result, and a saving of about 1 volt-sec resistive loss is found. Although the change in \( I_0 \) is small the onset of sawteeth is delayed by several seconds indicating either a change in the current profile near the magnetic axis or a stabilisation mechanism similar to the monster sawtooth [2].

2. HIGH CURRENT X-POINT PLASMAS. The formation of a magnetic separatrix is possible in JET at plasma currents 3-4MA by the combined effect of leakage fields from the primary and shaping coils, resulting in elongation < 2. Higher currents can be achieved in the single null configuration by displacing the plasma axis away from the mid-plane and unbalancing the main vertical field coils. Fig.3 shows a 3.8MA X-point obtained in this way. This discharge entered the H-mode with neutral beam heating. An X-point has been produced at 4.5MA but in this case the separatrix was close to the lower belt limiter. A 5MA X-point should be possible with larger imbalance in the vertical field coils.

The X-point plasmas are no more vertically unstable than limiter plasmas with the same current and elongation. However, in the asymmetric configuration, the vertical motion is usually towards the X-point and the velocity is larger. Values of \( I_0 \text{AB} = 2.5 \text{MA/(M-Tesla)} \) have been achieved for X-points compared with 2.1 for limiter plasmas.

3. COMPARISON OF LIMITER AND X-POINT CONFINEMENT. Heating experiments are currently underway for high plasma currents in both limiter and X-point configurations. Values of \( n_0(T)T_1(0)I_{T1}=1.8 \text{E20m}^{-7} \text{keV-sec} \) have been achieved in recent 5MA limiter experiments which is similar to 3MA H-mode data. At 3.8MA \( n_0(T)T_1(0)I_{T1}=3 \text{E20} \) has been achieved in the H-mode. In both 5MA limiter and 3.8MA H-mode the plasma energy content reaches ~7MJ. However, whereas the limiter case requires a power input of 15MW, the 3.8MA X-point requires only 8MW to reach this stored energy.

References

Fig. 1
Empirical stability diagram showing:

- Stability boundaries in $L_{\parallel}$-$q_\psi$ plane
- Trajectory in $L_{\parallel}$-$q_\psi$ plane of 6MA pulse 13944
- Amplitude of $B_0$ versus $q_\psi$
- Quasi stationary mode amplitude versus $q_\psi$
- □ 1 sec after breakdown
- ● 1 sec intervals
- △ start of current flat top

Fig. 2
Flux at plasma boundary versus plasma current (MA) for shot 13944

- Total flux change
- Plasma inductive flux change
- Plasma resistive flux change
- Volt-sec limit (a) see text
- Volt-sec limit (b) see text
- □ 1 sec after breakdown
- ● 1 sec intervals
- △ start of current flat top
Fig. 3
Magnetic surfaces for H-mode shot 14663. $I_p = 3.8$MA, $B_T = 2.1$ Tesla.
Note the single x-point at the top.
ION TEMPERATURE PROFILES AND ION ENERGY TRANSPORT IN JET DURING ADDITIONAL HEATING AND H-MODES


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EXPERIMENTAL $T_i$ PROFILES

The ion temperature is deduced in JET using various diagnostic systems in order to cover all possible plasma scenarios. Ion temperature profiles have been in the past deduced from a model dependent analysis of the data obtained with an array of four passive Neutral Particle Analysers (NPA) only during ohmic and ICRH heated discharges [1].

In this paper we report on initial $T_i$ profile measurements during Neutral Beam Injection. These have been obtained from both the multichord line-of-sight system of the visible charge exchange (CX) spectroscopy diagnostic which uses the JET Neutral Beams as diagnostic beams [2,3], and the upgraded version of the analysis of the NPA data.

With the CX spectroscopy system, radial profiles of ion temperature and toroidal velocity [4] are deduced from the Doppler broadened recombination spectra of fully stripped plasma ions such as deuterium, helium, carbon and oxygen. A multichord array in the near equatorial plane intersects the neutral beams at 8 positions between $R=2.3\,m$ and $R=4.1\,m$ with a spatial resolution of $\pm 10\,cm$. An additional single vertical line intersecting the neutral beams in the plasma centre enables to compare the temperature of different ions and to investigate possible anisotropic velocity distribution functions.

In the first period of operation, ion temperature profiles were measured during H-mode phases of X-point discharges [5] and during combined RF and NB heating. Representative peaking factors of parabolic ion temperature profiles are between 1 and 2, the profiles becoming flatter during H-modes. Fig.1 gives an example of the time evolution of the ion temperature profile (based on $0^+ (10 \text{ to } 9)$), while in Fig.2 profiles at different times are compared. Hollow ion temperature profiles can sometimes be observed. A more detailed investigation will be possible when a new system using a finer mesh of lines-of-sight employing 15 chords becomes operational.

In Fig.3 profiles obtained from the NPA data in different NBI scenarios are shown. They are obtained taking advantage of the fact that the JET NPA system can simultaneously measure hydrogen and deuterium particles. While the deuterium spectra are simulated taking into account a Fokker-Planck description of the slowing down of the beam particles [6], the hydrogen spectra remain Maxwellian and help in the reconstruction of the profile. In this way profiles during a hot ion mode discharge
(Fig.3a), during NBI into a limiter discharge (Fig.3b) and during NBI into X-point discharges both in L-mode phase (Fig.3c) and in H-mode phase (Fig.3d) have been obtained.

A broadening of the profile is observed at the transition from L to H mode, consistent with the CX spectroscopy.

TRANSPORT ANALYSIS

Initial results of ion transport analysis based on Ti-profiles measured by NPA were presented in [7]. The availability of an extended data base for ohmic and RF-heated plasmas allows us now to establish those conclusions more firmly. Furthermore, we are in a position to discuss preliminary results on local ion transport during NBI.

The predictive transport code JETTO connected to JET data banks has been used, adopting the same modelling techniques described in [7], unless otherwise specified. The main conclusion in [7], ie. the evidence for anomalous ion transport, has been confirmed by further analysis of ohmic and RF-heated discharges. It remains true even when the LIDAR electron temperature profile is taken as experimental reference instead of the ECE measurement (the former can be as much as 25% lower than the latter in some of the considered cases).

Figure 4 shows a representative example of our results, corresponding to the flat-top phase of a "monster sawtooth" for a deuterium plasma with 7.5 MW on-axis RF heating (H minority) and peak electron temperature Teo = 7 keV (from LIDAR). Electron energy transport has been modelled here using a non-linear heat flux model (\(\dot{\gamma}e \propto \gamma p\)) [8], that yields Te-profiles in agreement with the experimental data, and virtually insensitive to the variations in Ti(p) shown in Fig.4. The three computed Ti-profiles in Fig.4 have been obtained assuming: a) \(\chi_i = \chi_e\); b) \(\chi_i = \chi_i(n_i) + \chi_i,neo\), where \(\chi_i(n_i)\) is the anomalous conductivity due to the excitation of ion temperature gradient driven modes [7], while \(\chi_i,neo\) is the Chang-Hinton neoclassical coefficient [9]; c) \(\chi_i = 15 \chi_i,neo\).

It is apparent that a strong ion thermal conductivity in the outer half of the plasma is necessary to reproduce the measured Ti(p) : \(\chi_i = \chi_e\) appears to be the best available prescription.

The same prescription, when used in the simulation of an X-point discharge with NB heating, proves successful in reproducing Ti(p) from NPA during the L- as well as during the H-phase.

We should stress the preliminary nature of these calculations, and the fact that we are simulating the behaviour of a plasma undergoing transients (Fig.5).

In the transport code simulation, the time evolution of the electron temperature is satisfactorily reproduced by using an empirical electron thermal conductivity model based on the "profile consistency" constraint [10], combined with Bohm-like transport in the region where q<k, to simulate the average effect of sawteeth.

Figure 6 compares the experimental NPA profiles during L and H mode with those computed using \(\chi_i = \chi_e\).

During the early stage of the H-mode, to which the profile in Fig.6 refers, the thermal conductivity throughout the plasma is found to be comparable to its ohmic value. At the earlier L-mode time in Fig.6, \(\chi\) is somewhat larger, but it evolves in time and no discontinuity is observed in transport at the transition to the H-mode.
FIGURE 1

Time behaviour of $T_i$ profiles from CX spectroscopy

FIGURE 2

$T_i$ profiles from CX spectroscopy:

a) L-mode
b) H-mode, end of monster sawtooth
c) after sawtooth crash

FIGURE 3

$T_i$ profiles from NPA analysis in different NBI scenarios. Also shown are $T_i$ profiles from ECE (dashed lines).

a - # 10169
b - # 10981
c and d - # 10755
Experimental $T_i$ profiles compared with those computed with:

- $a - \chi_i = \chi_e$
- $b - \chi_i = \chi_{i}(n_i) + \chi_{i,neo}$
- $c - \chi_i = 15 \chi_{i,neo}$

for a 'monster sawtooth' with $B = 2.7$ T, $I = 2$ MA, $n_e = 3 \times 10^{19}$ m$^{-3}$, $P_e = 7$ MW.

**FIGURE 5** — Time evolution of volume average electron temperature and density, $H$ radiation and auxiliary power. Arrows indicate the times for which $T_i$ profiles are shown in fig. 3c, d. At these times NBI deposition profiles are centrally peaked and broad, with 60-70% of power being absorbed by plasma ions.

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INTRODUCTION

The emission of γ-rays from excited reaction products in JET plasmas has been monitored systematically since the first successful observations in 1987 [1]. Fusion γ-rays, i.e. γ-rays emitted by fusion products from fuel ions, are routinely recorded for monitoring the fusion yield from D-3He plasmas, as reported in a separate contribution to this conference [2]. Here we concentrate on γ-rays emitted by reaction products from interactions of RF driven minority ions with plasma impurities (mainly carbon and oxygen).

The experimental set-up is essentially the same as the one used in 1987 [1]. Two detectors, a 125 mm diameter by 125 mm long NaI (Tl) and a 75 mm by 75 mm BGO scintillator, are located in well shielded positions in the roof laboratory and view the plasma vertically along similar lines of sight. The results from both detectors are in reasonable agreement.

H MINORITY HEATING

A typical spectrum obtained during H minority heating is shown in figure 1. The most readily understood lines are those generated by RF accelerated H ions inelastically scattered from carbon and oxygen impurities in the plasma. For the excitation process to be energetically possible these ions have to exceed threshold energies of 4.8 MeV and 7.5 MeV for carbon and oxygen, respectively.

The most prominent lines, however, are caused by second harmonic driven deuterons reacting with 12C to yield 13C in an excited state. The intensity of the 3 MeV line as monitored with a single channel analyzer follows closely the sawtooth behaviour of the neutron yield (figure 2) and only emerges from the background when RF heating is applied, thus proving the correlation between RF heating and γ emission. The observation of the 3.8 MeV γ-line requires a certain fraction of the deuterons to exceed a threshold energy of 1.3 MeV; further analysis is hampered by the apparently erratic behaviour of the underlying cross section as well as the relative uncertainty in the density of the carbon ions. However, a simple model based on the ratio of the intensities of the 3.6, 3.8 MeV doublet to the 3 MeV line and the knowledge that the cross section leading to the doublet is approximately equal to the cross section leading to the 3 MeV state, yields deuteron tail temperatures between 300 and 500 keV for RF power levels between 9 and 15 MW. The tail temperature has been defined here as the e-folding energy of an assumed exponentially decaying energy distribution function.
The fact that the magnitude of the $^{12}\text{C}(d,n)^{13}\text{N}$ cross section for deuteron energies in the range of interest is comparable to the $^{12}\text{C}(d,p)^{13}\text{C}+\gamma$ cross section, coupled with the observation of 3 MeV $\gamma$ yields in excess of $10^{13}$ $\gamma$/s/sec, leads to the conclusion that, at present power levels (9-15 MW) and impurity concentrations ($\chi$ 10% C), about 10% of the observed neutrons originate from interactions with carbon. Moreover (d,d) neutrons created by the interaction of the deuteron tail with thermal deuterons from the bulk of the plasma contribute significantly to the total neutron yield. Unfortunately the energy spectrum of these neutrons is very broad, extending from $\sim$ 1.5 MeV to 3.5 MeV, which makes observations with our high resolution neutron spectrometers difficult. Nevertheless, comparing the intensity of the 2.5 MeV neutron peak to the total intensity of fast neutrons shows that up to 50% of the observed neutrons are of non-thermal origin under conditions of intensive H minority heating.

**Fig. 1:** Gamma spectrum recorded during H minority ICRF heating ($I_p = 3$ MA, $B_T = 2.8$ T and $P_{IF} = 15$ MW). The energy level scheme for $^{13}\text{C}$ has been taken from reference 4.

**Fig. 2:** The 3 MeV gamma radiation exhibits sawtooth oscillations and appears only when ICRF heating is applied (at 9 sec into the discharge).

**Fig. 3:** Vertical neutron and $\gamma$ profiles obtained during NB heating (4 MW) (figs. a and b) and later in the same discharge during NB plus ICRH (H minority) heating (4 MW and 13 MW) (figs. a' and b').
Finally, on one occasion the JET neutron profile monitor [3], which also monitors γ-rays on a regular basis, had been tuned to the 3 MeV γ-line and the resulting profile (figure 3b) demonstrates the central power deposition of RF energy. Assuming that the carbon impurity profile is flat then a power deposition volume of 8 m^3 can be deduced in reasonable agreement with other observations and calculations [5].

^3^He MINORITY HEATING

Spectra obtained during ^3^He minority heating are much richer in lines as can be seen from figure 4. All prominent γ-lines can be identified as emanating from the ^12^C(^3^He,p)^1^N+γ reaction and transitions from 8 of the 9 first excited states in ^1^N were observed. The highest γ energy observed so far is 7,028 keV; in order to populate the corresponding level the energy of the ^3^He ions has to exceed a threshold of 2.81 MeV. This is the first direct evidence for ^3^He ions reaching such high energies during RF heating. For this line to be seen, the RF power has to exceed 8 MW.

From the relative intensities of the various lines and knowledge of the cross-sections, the full distribution function of the fast ^3^He ions could be reconstructed. This however, is a major task which has only been started recently. In the mean-while, a crude model based upon the comparison of the 6.4 MeV and 5.1 MeV line intensities and energy averaged cross sections has been applied to a 2 MA discharge with 10 MW RF power (#14618). For this particular discharge we deduce a tail temperature of around 1.5 MeV.

The quite common observation of inverted sawteeth on the outer γ channels of the neutron profile monitor (figure 5) yields a further proof that these γ-rays are emitted by the plasma. It also shows that energetic ^3^He ions are expelled from the middle of the plasma at the sawtooth crash. This effect is particularly prominent with monster sawteeth.

CONCLUSION

The reactions of energetic, RF-heated, H, D and ^3^H minority ions with carbon and oxygen impurity ions have been monitored by observing γ-radiation from their excited reaction products. The importance of carbon as the major impurity in JET plasmas is amply confirmed. An indication of the energy of the minority ions is obtained from a study of reaction thresholds; for ICRF powers in excess of 8 MW, minority ions are accelerated beyond 7.5 MeV for protons (H minority, fundamental), 1.3 MeV for deuterons (H-minority, 2nd harmonic) and 2.8 MeV for ^3^He (^3^He-minority, fundamental). Equivalent tail temperatures in the MeV range have been observed for the minority ions.

Sawtooothing of individual gamma-ray lines shows the expulsion of fast ions from the middle of the plasma. Inverted sawteeth in the gamma-emission from the outer regions of the plasma provide conclusive evidence that these γ-rays were emitted from the plasma.

Finally, comparison of the 2.5 MeV neutron peak recorded with a neutron spectrometer with the total neutron yield shows that, under conditions of intense RF heating, a significant fraction of the emitted neutrons (50% or more) are of non-thermal origin.
Fig. 4: Gamma spectrum recorded during \(^9\text{H}\) minority ICRF heating (\(I = 2\) MA, \(B_r = 3.4\) T and \(P_{\text{FF}} = 10\) MW). The energy level scheme for \(^{14}\text{N}\) has been taken from reference 4.

Fig. 5: Gamma-rays recorded with the neutron profile monitor. The inner channels eg. b and d, show normal sawtooth behaviour whereas the outer channels, eg. a and c, display inverted sawteeth.

REFERENCES
The confinement and slowing down of fast tritons has been investigated by measuring the ratio of 14 MeV and 2.5 MeV neutron production rates. Tritons of 1.0 MeV are produced in the $d + d \rightarrow t + p$ reaction at the same rate as the 2.5 MeV neutrons from the $d + d \rightarrow ^3\text{He} + n$ reaction. The majority of these tritons will remain confined in the plasma and slow down to thermal energies through Coulomb collisions with electrons and ions; a fraction will undergo fusion reactions $t + d \rightarrow ^4\text{He} + n$, in which the 14 MeV neutrons are emitted. The fraction of the tritons which burn-up is essentially equal to the ratio of 14 MeV to 2.5 MeV neutron production.

The 2.5 MeV neutron emission is obtained from a set of fission chambers for which the calibration uncertainty is about ±10%. The 14 MeV neutron production is measured by means of the activation of Copper samples placed close to the plasma using a pneumatic transport system which returns them, after each discharge, to a counting room for the measurement of the induced $\gamma$-activity. The reaction $^{63}\text{Cu}(n,2n)^{62}\text{Cu}$ is employed; it has a threshold energy of 10.9 MeV and a fairly high cross section (0.44 b at 14 MeV). The experimental set-up and the first measurements, performed in 1986, are reported in /1/. The absolute calibration of the activation technique has to be calculated in order to relate the total neutron emission to the local neutron flux at the irradiation position. The 1986 triton burn-up measurements, presented in /1/, made use of response coefficients calculated with FURNACE /2/, a code based on ray tracing and the discrete ordinates method. The absolute calibration has been revised recently using the neutron transport Monte Carlo code MCNP /3/ to provide a more detailed description of the vessel wall and the holes through the surrounding structure, through which the irradiation ends are inserted. Neutron activation response coefficients were calculated for three different samples (Indium, Zinc and Copper) for deuterium plasmas producing 2.5 and 14 MeV neutrons. The accuracy of the calibration was tested by comparing the results for the 2.5 MeV neutron yields with those obtained with the fission chambers; an overall agreement at the ±15% level was demonstrated.
For the 14 MeV neutrons, the response coefficient calculated with MCNP for Copper was 1.6 times lower than that calculated with FURNACE. This reduction is in part due to use of more accurate activation cross sections (≈20%) and to the inclusion of effects due to local inhomogeneities at the irradiation ends (10%). A small unexplained discrepancy remains. For the present work, we adopt the more recent MCNP results. The overall accuracy for the Copper response coefficients, including modelling, calculation statistics, and cross section uncertainties, is estimated to be about ±10%.

Further measurements have been obtained in 1987; ohmically-heated plasmas (with plasma currents up to 6 MA) and combinations of NBI and ICRH with up to 21 MW of additional heating power have been studied. All 1986 (corrected) and 1987 data are shown in figure 1 vs. the plasma current. Only measurements with statistical errors below 20% have been retained, these errors being actually less than 10% for most data points. The triton burn-up ratios fall in the range 0.5 to 1.5% depending upon plasma conditions. These conditions varied in the ranges: $1.8 \leq B_{\parallel} \leq 3.4$ T, $1 \leq I_{D} \leq 6$ MA, $3.7 \leq T_{e} \leq 7$ keV, $1 \leq n_{e} \leq 7 \times 10^{19}$ m$^{-3}$, $1.6 \leq Z_{\text{eff}} \leq 8$, the main impurities being Carbon and Oxygen. Finally the neutron yield for these discharges varied between $10^{13}$ and $4 \times 10^{15}$.

The triton burn-up ratio, $\rho$, is proportional to the product of the triton confinement fraction $f_{c}$, the deuterium density $n_{D}$ and the $d+t$ fusion probability (which depends on the triton slowing down rate). The $d+t$ fusion cross section has a maximum around $E_{T} = 170$ keV, well below the triton initial energy, so that the burn-up is sensitive to the slowing down mechanism for fast tritons. The main motivation for studying triton burn-up is to test the validity of the classical slowing down model for MeV ions without invoking anomalous losses or displacements of ions from their classical drift orbits. For this purpose, the burn-up measurements have been compared with model calculations based on classical triton confinement and slowing down in the plasma. Two codes have been used, each tending to stress different aspects of the problem. The first is a time-independent code, SOCRATE /4/, which takes into account the prompt-loss fraction of the tritons and the excursions from the flux surfaces on which they are born. The slowing down rate and the $d+t$ fusion probability is calculated for each confined triton moving on the guiding centre drift orbit corresponding to the initial energy. The second is a time-dependent code, TRAP-T /5/, which assumes that the tritons stay on the flux surface on which they are born. Since the slowing down for tritons can be as long as 1 sec in JET, the time variation of plasma parameters has to be considered, especially for discharges with strong additional heating. Both codes calculate slowing down as due to electron and ion Coulomb drag, other collisional effects such as pitch angle scattering and energy diffusion being neglected. Both use the magnetic equilibrium flux surfaces and the plasma parameters $T_{e}(r)$, $n_{e}(r)$ and $Z_{\text{eff}}$, provided by other diagnostics, as input data. The triton source distribution is deduced on the assumption that $T_{i}(r) = T_{e}(r)T_{i}(o)/T_{e}(o)$. This approximation may be rather crude in the case of NBI heating where many of the tritons are produced by beam-plasma interactions. However, the 2.5 MeV neutron emission profiles are measured on JET /6/ and no important
differences in the profiles are observed between ohmic and NBI heated discharges. The ratio \( n_D/n_e \) is deduced from visible bremsstrahlung measurements of \( Z_{\text{eff}} \) and from the concentration of impurity species from UV spectroscopic measurements. Uncertainties in \( Z_{\text{eff}} \) propagate increasingly larger errors in the ratio \( n_D/n_e \) for higher \( Z_{\text{eff}} \) values; therefore, a restricted set of the 1987 discharges has been selected for which \( 2 \leq Z_{\text{eff}} \leq 3 \) so that \( n_D/n_e \) varied between 0.83 (±7%) and 0.65 (±15%). For the discharges so selected, the total plasma current was 3 or 4 MA and the calculated prompt loss fraction was always less than 3%. Both ohmic and NBI and/or RF heated discharges with up to 15 MW additional power were included. Predictions for the triton burn-up were obtained for these discharges from the two codes; they are in remarkably good agreement (within a few per cent) in all cases, provided that the time independent code is run for several times representative of the discharge temporal evolution. This agreement indicates that the neglect of orbit drifts is a good approximation for JET plasma with \( I_p \geq 3 \) MA.

The comparison between experimental and theoretical burn-up ratios for the 1987 data is shown in figure 2. The error bars for the experimental values retain only the statistical errors; the systematic error is about ±15%. The error bars for the calculated burn-up values take into account the uncertainties in the input data for \( Z_{\text{eff}} \) and \( T_e \). The electron temperature used is that measured by the LIDAR diagnostic which has recently become available /7/. It is based on Thomson scattering and does not suffer, unlike ECE, from problems associated with a magnetic field dependent calibration. The uncertainties in \( T_e \), which propagate almost linearly in the burn-up ratio through the thermalization time of the tritons, is ±10%. Errors arising from approximations in the classical models adopted in the codes are difficult to assess. Both codes neglect losses due to pitch angle scattering but, since the loss cone is small, their effect is expected to be negligible.

It can be seen from Fig. 2 that agreement has been found between the experimental measurements and the theoretical calculations within the overall uncertainty of ±20% for both values. No particular trend has been found for additionally heated as compared to ohmic discharges. The present results, using the 1987 data, correlate rather better with the line for \( P_{\text{exp}} = P_{\text{th}} \) than did the 1986 data reported in ref. [8]. The two sets of data are not inconsistent with each other but nevertheless there is an apparent improvement which can be attributed mainly to the use of LIDAR instead of ECE temperature data, and also to the selection of more reliable burn-up data (ie. \( P_{\text{exp}} \) for \( Z_{\text{eff}} \leq 3 \)). Since the major source of uncertainty for high \( Z_{\text{eff}} \) discharges (ie. \( Z_{\text{eff}} > 3 \)) lies in the ratio \( n_D/n_e \), confirmation of the classical nature of the triton burn-up process now permits us to invert the problem and to consider the burn-up measurements as contributing to the evaluation of \( n_D/n_e \).

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Figure 1:
Illustrating the range of conditions over which measurements of the triton burn-up fraction have been obtained.

Figure 2:
Comparison of experimental with calculated triton burn-up fractions for the subset of measurements with $Z_{\text{eff}} \leq 3$. The solid line for $\rho_{\text{exp}} = \rho_{\text{th}}$ is provided as a guide to the eye.
THE EVOLUTION OF $Z_{\text{eff}}(r)$ PROFILES IN JET PLASMAS

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1. INTRODUCTION

The understanding and control of impurity production in JET plasmas remains a fundamental aim of the experiment. To this end, the global impurity content in the device is routinely monitored throughout each discharge by making measurements of the average ion charge $Z_{\text{eff}}$, using visible continuum emission collected along two discrete lines of sight.

The use of a 15-telescope array aligned in a common poloidal cross-section has permitted the temporal evolution of $Z_{\text{eff}}(r)$ to be studied, under a variety of plasma conditions. In this paper, results are reported pertaining to a study of plasmas fuelled by gas or pellet injection and heated ohmically, by the application of ICRF or by NBI.

2. APPARATUS & ANALYSIS

The apparatus and method of analysis have been described in detail previously [1], albeit for a 13-channel version. The present system covers the same field of view, with two extra channels aligned on the zone where the magnetic separatrix is formed during X-point operation.

Each channel in the array measures the brightness $B = \frac{1}{4\pi} \int e(l) dl$, at 523.5 nm, along a chord through the plasma. Using the technique of Abel inversion, the brightnesses are transformed into the radial profile of continuum emissivity, $e(r)$. Since $e(r) = Z_{\text{eff}}(r) n_e(r)/T_e^{1/2}(r)$, knowledge of $n_e(r)$ and $T_e(r)$ permits the radial profile of $Z_{\text{eff}}$ to be determined and its temporal evolution to be followed.

3. RESULTS AND DISCUSSION

(i) Ohmic Discharges with Gas Fuelling

The precise value of $Z_{\text{eff}}$ determined at any time during a discharge, for any given setting of machine parameters, is sensitive to the history of previous machine operation and cleaning procedures. This is also the case when additional heating is applied or D₂ pellets injected.

During the ohmic phase of a discharge, and following the establishment of a stationary density profile, $Z_{\text{eff}}(r)$ is usually peaked on axis with $Z_{\text{eff}}(a)/Z_{\text{eff}}(0) \leq 1.4$ and $Z_{\text{eff}}(a)/Z_{\text{eff}}(a) \leq 2.5$. Traces (a) of Figures 1, 2, 3 and 4 are representative examples of such profiles. For stationary discharge conditions, $Z_{\text{eff}}$ decreases with increasing $n_e$ and increases with $I_p$. Provided the vessel is well conditioned the data are represented by the relation $Z_{\text{eff}} = 1/(n_e/J)$, where $J$ is the plasma current density.

(ii) Discharges Heated by NBI

a) Limiter Discharges

The injection of a beam of D atoms causes a
reduction of $Z_{\text{eff}}$, the amount depending on the power of the injected beam, i.e. on the flux of injected neutrals. Figure 1 shows two $Z_{\text{eff}}(r)$ profiles: (a) is prior to NBI while (b) was obtained after 5s of injection of a 5.5 MW beam of 80 keV D atoms. The relative shape of the $Z_{\text{eff}}(r)$ profile is little changed by the beam, but there is a reduction in the absolute values, $Z_{\text{eff}}$ decreasing from -2.6 to 2.3.

Despite a beam fuelling rate of $6 \times 10^{20} \text{s}^{-1}$, which increases the neutral source in the plasma centre by two orders of magnitude or more [2], recycling at the edge dominates the global particle balance. Recycling produces an electron influx that is $-\frac{4}{3} \times$ larger than the beam fuelling rate. The injected D atoms dilute the impurities, thereby reducing $Z_{\text{eff}}$.

b) X-point Discharges When the plasma is limited by a magnetic separatrix (X-point discharge) and the H-mode of confinement is achieved following the start of NBI, a different form of evolution of $Z_{\text{eff}}(r)$ is observed from that reported in (ii)a. Figure 2 shows 3 profiles obtained during a discharge in which 7.5 MW of NBI were applied at 12.5s.

Profile (a) pertains to the ohmic phase of a discharge in contact with the belt limiters. Profiles (b) and (c) were obtained during the H-mode phase - (b) was evaluated 0.05s after the L to H transition while (c) occurs 1.1s later. During the H-mode, the profile of $Z_{\text{eff}}(r)$ steepens while the absolute values increase steadily - $Z_{\text{eff}}$ increases from -2.8 at 13.1s to -3.7 at 14.2s. The profile is slightly hollow, as confirmed by soft X-ray [3] and spectroscopic [4] measurements. An improved global particle confinement time by a factor of 3-5 has been deduced from $D_\alpha$ measurements [1] and from impurity transport analysis [4].

(iii) Discharges Heated by ICRF The application of ICRF results in increased hydrogenic and impurity influxes. At power levels $\leq 4$ MW both types of influx increase in proportion to the applied power and $Z_{\text{eff}}$ is unchanged in value. However, at higher power levels $Z_{\text{eff}}$ increases during the RF pulse and $Z_{\text{eff}}(r)$ evolves.

Figure 3 shows two $Z_{\text{eff}}(r)$ profiles: (a) pertains to the ohmic phase while (b) was evaluated after 4s of heating at a power of 11 MW. The increased impurity production at the plasma edge leads to an increase of $Z_{\text{eff}}$ from -3.4 to 4.8. $Z_{\text{eff}}(r)$ increases significantly in the outer region and its overall shape flattens with a slightly-hollow centre. Following switch-off of the RF source $Z_{\text{eff}}$ decreases as the density rapidly decays, both to stabilise close to their ohmic values.

(iv) Discharges Fuelled by Pellet Injection Pellet injection has a dramatic effect on the temporal evolution of $Z_{\text{eff}}(r)$, due to the abrupt deposition of D atoms in numbers comparable to the plasma electron content, on a sub-millisecond time scale.

Multi-pellet fuelling has achieved values of $Z_{\text{eff}}(o)$ of 1.3 - 1.5, with $n_e(o) \sim 9 \times 10^{19} \text{m}^{-3}$, following the injection at 0.5s intervals of 5 pellets during current ramp-up in a 3 MA limiter discharge [5]. Each pellet had an average electron content of $7 \times 10^{20}$ - the initial plasma electron content was $10^{21}$. A single large pellet yields similar results, as reported in [1]; wherein the injection of $4.5 \times 10^{21}$ D atoms into a plasma of electron content $1.3 \times 10^{21}$ reduced $Z_{\text{eff}}(o)$ by > 3.

The effect of a single small pellet is less pronounced, but it well illustrates the plasma response to injection. In Figure 4, a time sequence of $4 Z_{\text{eff}}(r)$ profiles is shown for a plasma of initial electron content
- 1.4 x 10^{21} into which a D_{2} pellet containing ~ 2.2 \times 10^{21} atoms was injected. Profile (a) was obtained before injection. Profile (b) pertains to 0.1s after injection, which reduces Z_{\text{eff}}(o), by ~ 40\%, and also Z_{\text{eff}} at the plasma edge. However, at R ~ 3.6m, corresponding to the location of the q=1 surface, Z_{\text{eff}}(r) has barely changed. The prompt reduction in Z_{\text{eff}} is consistent with a fixed plasma impurity content being further diluted by the injected D ions. Following injection, Z_{\text{eff}}(r) fills in, traces (c) and (d). Z_{\text{eff}} increases until after 3-4s it is similar to that for a gas-fuelled discharge at the new, higher density.

4. CONCLUSIONS

The temporal evolution of Z_{\text{eff}}(r) in JET plasmas has been studied using a 15-channel poloidal array to measure the visible bremsstrahlung emission. This has given an insight into the global impurity behaviour during ohmic, NBI and ICRF heating, and in the case of pellet injection.

The imprecisions in calibrating the array lead to systematic errors in the measured signals. The resulting absolute errors in the Z_{\text{eff}}(r) profiles are estimated to be up to 25\% on axis. At the plasma edge n_{e}(r) is not well determined, leading to further uncertainties in Z_{\text{eff}}(r), which is not usually evaluated for the outermost 0.3m. The relative errors in Z_{\text{eff}}(r) are assessed to be about half the absolute errors, permitting changes in profile shape to be followed with reasonable confidence over most of the minor radius.

Under all conditions, recycling at the plasma edge plays an important role. During NBI in limiter discharges, Z_{\text{eff}} decreases but its profile shape is unchanged due to recycling being dominant over the central neutral source, in the global particle balance. In plasmas in which an H-mode is established Z_{\text{eff}} rises steadily and Z_{\text{eff}}(r) becomes steeper, although with a hollow centre. At powers ≥ 4 MW, ICRH leads to a gradual increase in Z_{\text{eff}} with a flattening of Z_{\text{eff}}(r), through increased impurity production at the plasma edge. Pellet injection, where deep penetration is achieved, causes a large abrupt decrease in Z_{\text{eff}}(o), through dilution of the core impurities. However, on a time scale of 3-4s recycling establishes a higher Z_{\text{eff}} which is consistent with edge fuelling.

None of the scenarios above yield Z_{\text{eff}}(r) profiles which indicate neoclassical accumulation of the light impurity species, such as C and O, in the plasma centre. However, metallic impurities such as Ni could accumulate significantly without being discernible on the Z_{\text{eff}}(r) profiles, since their initial concentrations are low [4].

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FIG. 1: $Z_{\text{eff}}(r)$ profiles for pulse #10499. (a) $t=6s$, during the ohmic phase of the limiter discharge. (b) $t=12s$, during NBI. $P_{\text{NBI}} = 5.5\,\text{MW}$. $I_p = 4\,\text{MA}$.

FIG. 2: Radial profiles of $Z_{\text{eff}}$ for pulse #14822. (a) $t=5.1s$, ohmic phase with plasma on limiters. (b) $t=13.1s$, 0.05s after start of H-mode. (c) H-mode, $t = 14.2s$. $P_{\text{NBI}} = 7.5\,\text{MW}$. $I_p = 3\,\text{MA}$.

FIG. 3: $Z_{\text{eff}}(r)$ profiles for pulse #13543. (a) $t=10s$, ohmic phase of limiter discharge. (b) $t=14s$, after 3s of 11MW of ICRH. 8 antennae in dipole configuration. $I_p = 3.0-3.5\,\text{MA}$.

FIG. 4: Radial profiles of $Z_{\text{eff}}$ for pulse #9522. (a) $t=6s$ (pellet injection at 7s). (b) $t=7.1s$. (c) $t=8s$. (d) $t=12s$. Profile evolves rapidly and $Z_{\text{eff}}$ increases. $I_p = 3\,\text{MA}$.
MULTI-PELLET INJECTION ON JET*


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and

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Introduction

JET and the US Department of Energy have jointly installed and are jointly operating a Multi-Pellet Injector for fuelling experiments on JET under the umbrella of the Bilateral Agreement on Fusion Research. The actual pellet launcher (built by Oak Ridge National Laboratory) can deliver multiple pellets of frozen fuel, from three different barrels in parallel in the sizes 2.7, 4 and 6 mm diameter (length equals diameter approximately) at a maximum frequency of several per second with nominal speed of up to 1500 ms⁻¹, which are injected horizontally into the midplane of the tokamak. More information can be found in [1] which also deals with the fuelling aspects of the pellet experiments described in the following. The JET Multi-Pellet Injector was brought into operation in last summer and systematic investigations with pellet injection have started later in the year with the aim to increase the central plasma in a more controlled way and to purify the plasma.

Since it is known that for large hot plasmas pellet penetration depths are likely to be insufficient to reach the centre of the discharge the overall strategy in the initial experiments has been to set up a high density target plasma in an ohmic discharge with central density peaking and to heat this subsequently in hope that the density profile will be maintained for a reasonable length of time. Mainly 2.7 and 4 mm deuterium pellets have so far been injected into the current flat top and into the current ramp, for plasma currents in the range of 3 to 5 MA in limiter discharges with toroidal fields of 2.1 to 3.4 T.

Review of initial experiments

Following the above mentioned strategy a large number (>100) tokamak discharges have been performed the evaluation of which is now being done but is hampered by the existence of particular problems with diagnostics to cope with rapid, multiple changes induced by the pellet events. It is therefore a tedious and time-consuming procedure to derive from the raw signals the time evolution of radial profiles needed for more detailed and subtle evaluation.

*This work has been performed under a collaboration agreement between the JET Joint Undertaking and the U.S. Department of Energy.
Peaked density profiles with high central density were established in the first example: Fig. 1 shows the time evolution of the electron density versus the major radius of a 2.5 MA/3T discharge (shot 13572) without auxiliary heating (as constructed from FIR laser interferometry and optical Bremsstrahlung measurement). One 4 mm pellet is injected into the early current flat top at 4.5 s and a combination of 2.7 and 4 mm pellets following each other in $10^{-3}$ s distance at 5.5 s create a very peaked profile with a density on axis of more than $1.2 \times 10^{20} m^{-3}$ (average $4.1 \times 10^{19} m^{-3}$) decaying over 2 s to $8.1 \times 10^{19} m^{-3}$ still with a peaking factor $n_e(0)/n_e$ of 3 (as supported by the LIDAR Thomson scattering profiles; the hollow profile at 4.5 s may be a feature due to the limited number of signal traces available). The electron temperature on axis (being about twice the volume average one) drops with the first pellet from 3 to 2 keV and with the subsequent combined pellets to 1 keV but recovers in about 2 s leading at 7.5 s to a central electron pressure peak of 0.25 bar; tentative estimates give values around 2 during this time for $Z_{eff}$ (lower still for the pellet event itself) and around 0.6 sec for the energy confinement time, virtually constant over the density decay time.

Fig. 1: Electron Density Profile over major radius $R$ and time $t$ for shot No. 13572

If the pellets do not penetrate sufficiently into the plasma to peak the central density strongly they nevertheless increase the average density on a faster time scale than gas feeding and to or even above the density limit. An example for this is shot 13544, a 3.3 MA, 2.8T discharge with three 4 mm pellets, Fig. 2, followed up by on-axis RF heating of 16 MW at 42 MHz, H-minority ICRH. This drives the plasma stored kinetic energy up
to about 6 MJ with a radially very flat central density of \(6 \times 10^{19} \text{ m}^{-3}\). This profile holds on in time during the RF pulse - conditions are maintained for about two seconds - at a level which can also in the end be reached by gas puffing; however, the high D-D reaction rate of \(1.4 \times 10^{15} \text{ s}^{-1}\) (no neutral beam heating) suggests a relatively clean plasma with a high deuteron content. The central electron pressure is exceeding .75 bar at a central electron temperature of 8 keV.

When pellets are too closely spaced in time locked modes - MHD activities with formation of almost stationary magnetic islands associated in this case with high-density operation - lead to rapid density pump-out and usually deterioration of peakedness \[2\]. Good profile peaking of density can be obtained by fuelling the tokamak discharge very early in the current ramp. Shot 14550 - the 3-dimensional density evolution can be found in \[1\] - shows a build-up to \(9 \times 10^{19} \text{ m}^{-3}\) by injection of 7 2.7 mm pellets for a 3 MA/3T discharge, fig. 3. The central temperature is kept by the cooling of the pellets to about 1 keV before it takes off to 8 keV with RF (8MW) and NB (5MW) being applied at 3.5s; immediately the central density decay sets in and the peaking disappears with the total particle contents roughly maintained. The interesting part is the high electron pressure of .65 bar max and the high D-D reaction rate approaching \(4 \times 10^{15} \text{ s}^{-1}\). The plasma energy is 4MJ but decaying in line with the temperature decay after 4.5s; with this collapse the central pressure as well as the D-D reaction rate are declining while the 13 MW are still maintained. A similar
discharge - shot 14555 - exhibits very similar features: its D-D rate peaks at \(4.3 \times 10^{15} \text{s}^{-1}\) when the central electron temperature reaches 7.5 keV with combined auxiliary heating of 10MW ICRH and 5MW NB, it does however decay before the first temperature collapse .8s later and seems more to decline in parallel with the central density in this case. The D-D rate is still about \(2.10^{15} \text{s}^{-1}\) 1.5s later when the LIDAR measures the radial electron pressure profile - fig. 4 - with about 0.5 bar peak and the highest pressure gradient, i.e. 2.2 bar \(\text{m}^{-1}\), we have seen so far. Both shots exhibit global pressure values - as inferred from diamagnetic loop measurements - in excess of 35% of the Troyon limit.

Conclusions
The first series of multi-pellet injection experiments on JET have so far shown that with central deposition (usually to within one third of the minor radius) peaked density profiles with central values exceeding \(10^{20} \text{m}^{-3}\) can be reached. These peaked profiles can persist for times in the order of seconds in the pure ohmic case. Auxiliary heating - on-axis radio frequency and combined RF and neutral beam heating (up to levels of 16MW) - seems to accelerate the slow central density decay considerably but generates for short times (some fraction of a second) high central electron pressure values (up to .75 bar) in a core with a high D-D reaction rate (up to \(4.10^{18} \text{s}^{-1}\) at about 6keV electron temperature) presumably because of the high deuterion contents brought about by the injection of clean pellets. The next steps in the programme will be - apart from generally widening the operational range of injection application - to attempt to hold on to the peaked profiles by varying the heating and fuelling profiles and to understand the physics underlying the formation and the decay of this interesting core.


THE JET MULTIPLE PELLET LAUNCHER AND FUELING OF JET PLASMAS BY MULTIPLE PELLET INJECTION*


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Pellet Injector System Description and Performance

A three-barrel repeating pneumatic pellet launcher developed at ORNL is the principal component of a new plasma fueling system for the Joint European Torus (JET) [1,2]. This versatile device consists of three independent machine-gun-like mechanisms equipped with high-speed extruders to provide solid deuterium to each gun assembly. The injector features nominal pellet sizes of 2.7 mm, 4.0 mm, and 6.0 mm, giving ideal volume-average plasma density increments of $0.82 \times 10^{19} \text{m}^{-3}$, $2.66 \times 10^{19} \text{m}^{-3}$, and $8.9 \times 10^{19} \text{m}^{-3}$, respectively, and has been qualified at repetition rates of 5 Hz, 2.5 Hz, and 1 Hz, respectively. Each gun can operate (individually or simultaneously) at the design repetition rate for 15-s pulses. Pellet speeds in the repeating mode average 1300 m/s; in the single-shot mode, the performance is close to 1500 m/s. Additional components of the ORNL pellet launcher system include: (1) an instrumented propellant and fuel gas feed system; (2) injector diagnostics, including a fiber-optic pellet detection system and optical systems for remote monitoring of solid hydrogen extrusions and high-speed flash photography of pellets; and (3) a data acquisition and remote control system consisting of a programmable logic controller (PLC) and a computer/CAMAC-based operator interface and data acquisition system.

The balance of the installation at JET consists of the following JET-supplied subsystems [3]: (1) a launcher-torus vacuum interface for differential pumping of propellant gas and extrudate fuel featuring a 50-m$^3$ vacuum chamber equipped with an $8 \times 10^5$ L/s cryocondensation pump; (2) a liquid helium delivery, storage, and recovery system and fuel and propellant gas distribution systems; (3) a fire control sequencer that provides timed trigger pulses for initiation of the extrusion process and programmable firing of pellets; (4) a microwave cavity-based pellet mass measurement system and an instrumented target array to facilitate aiming of all three guns; and (5) a PLC-based control system and computer operator interface for these subsystems.

*This work has been performed under a collaboration agreement between the JET Joint Undertaking and the U.S. Department of Energy.

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The fueling system became operational on JET in October 1987. To date, plasma fueling experiments have been performed with the 2.7- and 4.0-mm guns operating in the multipellet mode.

**Pellet Penetration and Particle Deposition Profiles**

Penetration of the 2.7- and 4.0-mm pellets for ohmic plasmas without significant populations of suprathermal electrons agrees with the neutral and plasma shielding model [4], as shown in Fig. 1. Measured penetration is determined from vertical soft x-ray data in all cases. The calculated penetration uses $T_e$ profiles from second harmonic ECE data and $n_e$ profiles from six vertical chords of an FIR interferometer. The 2.7-mm pellets have a pellet-by-pellet mass correction in the calculated penetration using the signal from the microwave cavity. The larger scatter in the 4-mm data may be due to pellet mass variation in the experiment; the calculated penetration is based on an average mass. Earlier JET data from 3.6- and 4.6-mm pellet experiments using the IPP Garching single-pellet injector show similar agreement with the model [5].

The gross particle deposition has been determined by volume integration of the density profiles after pellet injection. This gives, on average, $6.6 \times 10^{20}$ and $2.3 \times 10^{21}$ particles for the 2.7- and 4-mm pellets, respectively, which corresponds to $\approx 70$--$75\%$ of the pellet inventory as determined by the ideal pellet dimensions. The discrepancy results from a loss of pellet mass by erosion of the pellet diameter in the gun barrels.

Details of the particle deposition in ohmic plasmas are illustrated in Fig. 2 by LIDAR Thomson scattering profiles taken within 20 ms of pellet injection. Figure 2a shows the plasma density profiles after the first and third pellets in a sequence of three 4-mm pellets injected into a 3-MA ohmic discharge at 1 Hz. The profile is strongly inverted after the first pellet (which penetrates to $R = 3.4$ m) but more centrally peaked after the third pellet, which penetrates to the magnetic axis. Central penetration is accomplished as the central electron temperature decreases from 4 keV to 2.5 keV during the fueling pulse. As shown in Fig. 2b, a deeper deposition profile can be achieved by injecting a 2.7-mm pellet a few milliseconds after the 4-mm pellet. This technique was used to produce the highly peaked density profile shown in Fig. 2c. The subsequent evolution of the density profiles in ohmic and auxiliary heated discharges is discussed by Kupschus et al. [6].

**Plasma Fueling Experiments**

One of the primary objectives of the experimental program on JET is to produce clean, centrally peaked, high-density target plasmas for central and off-axis heating (ICRF and NBI) experiments. To date, the 2.7- and 4-mm injectors have been used for this purpose in X-point and limiter plasmas at up to 5 MA of plasma current. Plasma density buildup has been demonstrated in the startup and flattop phases of ohmic limiter and X-point discharges using 2.7-mm pellets for pulse lengths exceeding 2.5 s. Figure 3 illustrates the plasma density evolution from FIR data in response to injection of 2.7-mm pellets at 2.5 Hz starting 1 s into the discharge (at $I_p = 1.5$ MA) and continuing until the start of the current flattop phase of a 3-MA limiter plasma. In this discharge, a centrally peaked density profile is maintained in time as the second and subsequent pellets penetrate to and somewhat beyond the magnetic axis. This was accomplished by adjusting the fueling rate so that the central electron temperature was maintained in the range of 1.2--1.3 keV throughout the fueling pulse. The volume-average density increases linearly during pellet fueling, and 75% of the particle input rate is accounted for in the rate of rise of the plasma particle inventory. Peak to volume-average density ratios in the range of 2--2.5 are maintained during the fueling pulse, compared with 1.3 for gas puffing cases. The
peaking factor is sensitive to pellet penetration, as demonstrated by comparing Fig. 3 (shot 14550) and Fig. 4 (shot 14545). The conditions of Fig. 4 differ from those of Fig. 3 in that the central electron temperature is higher (2–2.2 keV) during the fueling pulse, resulting in reduced pellet penetration (to $R = 3.25–3.35$ m). The density profiles are consequently broader, giving peaking factors of 1.75 and central densities 50% smaller than the values on shot 14550. While the volume-average density increments during pellet injection for these two cases are similar, the density decay after pellet injection is somewhat more pronounced on shot 14545, giving a 15% smaller volume-average density at the end of the pellet fueling pulse. The response of these profiles to auxiliary heating, which starts at 3.5 s, is discussed by Kupschus et al. [6] at this conference.

Long-pulse plasma fueling has also been performed in the flattop phases of ohmic discharges using primarily the 2.7-mm gun. In limiter plasmas, the volume-average density can be maintained at $3.5 \times 10^{19}$ m$^{-3}$ at an injection frequency of only 1 Hz. Preliminary experiments in ohmic X-point discharges indicate that pellet fueling provides better control over the density level than gas puffing but that the fueling efficiency is lower (i.e., the plasma pumpout is faster) than in the high-recycle limiter plasmas. For 2.7-mm pellets that penetrate to $R = 3.5–3.6$ m (i.e., about halfway to the axis), injection...
frequencies of 2 Hz are required to maintain the volume-average density in the range of $2 \times 10^{19} \text{ m}^{-3}$. A more pronounced effect has been obtained by injecting a 4-mm pellet (which penetrates to $R = 3.15 \text{ m}$) followed by 2.7-mm pellets at 5 Hz. In this case, a central density of $5 \times 10^{19} \text{ m}^{-3}$ was sustained with a profile shape similar to that of Fig. 4. For conditions typical of the flattop phases of JET discharges, deep pellet penetration is not achieved with 2.7-mm pellets alone, and to date we have not observed density profile peaking as strong as that observed on ASDEX [7] with partial penetration.

**Summary**

A new multipellet long-pulse plasma fueling system is in operation on JET. In the initial experimental phase, a variety of plasma density profile shapes have been produced with peak to average values ranging up to 2.5 and peak plasma density up to $1.2 \times 10^{20} \text{ m}^{-3}$.

**References**

1. INTRODUCTION

An important objective in tokamak research is to find suitable descriptions for the measured thermal and particle fluxes, and to identify the underlying mechanism of transport. Many detailed models for transport in tokamaks have been developed, which make specific predictions for correlations between thermal and particle transport coefficients. To exclude some of the competing models, accurate measurements of the correlations are required. Thermal and particle transport in JET have been described by expressions of the form [1,2]

\[ Q = -\chi n V T + Q_p; \quad \Gamma = -D n + \Gamma_p \]  

(1)

\( Q \) and \( \Gamma \) are the heat and particle fluxes, \( \chi \) and \( D \) are the respective diffusivities, \( Q_p \) and \( \Gamma_p \) are convective (or pinch) fluxes. In order to reduce uncertainties arising from shot-to-shot and spatial variations it is necessary to determine the coefficients simultaneously at the same spatial point. In this paper we describe evaluation of thermal and particle transport coefficients by three different methods satisfying the above requirements: (a) From the velocity and damping of heat and density pulses propagating outwards following a sawtooth collapse. (b) Measurement of inward propagation of electron density and temperature perturbations produced when a small pellet is injected into the plasma. (c) The method of 'flux gradient' analysis applied to transient conditions. In the following we present determinations of \( \chi \) and \( D \) for Ohmically heated discharges in deuterium, limited by the outer belt limiters, with \( I_p = 3 \text{ MA}, 2.8 \leq B_\phi(T) \leq 3.4, \) and \( 1.5 \leq \bar{n}_e(10^{19} \text{ m}^{-3}) \leq 2.7. \)

2. DETERMINATION OF \( \chi_e \) AND \( D_e \) FROM SAWTOOTH PROPAGATION

The local electron thermal diffusivity \( \chi_e \) is determined by analysis of propagation of electron temperature perturbations in the region outside the mixing radius after a sawtooth collapse [3]. \( T_e(r,t) \) is determined from measurements of electron cyclotron emission, analyzed with a grating polychromator (KK2). Two parameters are derived, a heat pulse velocity and a damping rate of the pulse amplitude. \( \chi_e \) is determined by comparison with simulations of these using a diffusive model including sources. Fig.1 shows a representative parameters of the heat pulse propagation. Measurements are made in steady state, in a region bounded by the normalized minor radius 0.65r/as0.8, and averaged over a period of \( = 1 \text{s}. \) From analysis of a number of discharges a value of \( \chi_e = 2.9 \pm 0.4 \text{ m}^2/\text{s} \) is deduced.
The electron particle diffusivity $D_e$ is determined in an analogous manner by measuring the propagation of a density pulse following a sawtooth collapse [4]. The density perturbation $\delta n_e(r,t)$ is observed using a multichannel reflectometer [5]. The diffusing perturbation is clearly seen at all positions outside the mixing radius $r_m$. The measured pulse delay time is compared with predictions of a diffusive model in which an equilibrium electron density profile is periodically flattened inside $r_m$. An edge recycling coefficient $R=1$, and a pinch flux $T_p = -D_n r/a^2$ are assumed. Fig. 2 shows a comparison of the modelled and measured density pulse propagation. $D_e$ has been varied to match the measured delay in the region, $r_m < r < 0.85a$. Analysis of several discharges gives $D_e \approx 0.4 \text{ m}^2/\text{s}$ with an estimated error of 30%. The above determinations of $D_e$ and $\chi_e$ have been performed in the same discharges, covering the same spatial and temporal region, and yield $\chi_e/D_e \approx 7$.

3. DETERMINATION OF $\chi_e$ BY PELLET INJECTION

Here, the local thermal diffusivity is determined by analysis of the propagation of a 'cold front' into the region $r < r_p$, where $r_p = 0.6a$ is the pellet penetration radius. $T_e(r,t)$ at $r_p < r < a$ is measured using the KK2 instrument. In a previous investigation we have shown that the assumption of diffusive propagation of the 'cold front' into the region $r > r_p$ may be justified [6]. $\chi_e$ is determined by comparison with simulations using a diffusive model including sources. The initial condition, the temperature and density profiles immediately after pellet injection, is accurately determined from measurements of $T_e(r)$, $n_e(r)$, pellet mass, ablation rate and penetration depth. During evolution of $T_e(r,t)$ in the spatial region of interest, $\Delta t \leq 60 \text{ ms}$, the perturbed density profile is assumed to be unchanged; the evolution of $n_e$ is much slower than that of $T_e$, justifying this assumption. The model $T_e(r,t)$ evolution is most sensitive only to the local value of $\chi_e$, enabling an accurate determination. Fig. 3 shows a comparison of the measured and modelled $T_e(r,t)$ at radii $r < r_p$. A value of $\chi_e = 1.6 \pm 0.3 \text{ m}^2/\text{s}$ is deduced at $0.45a < r < 0.55$. This value is close to that determined, under similar conditions immediately after pellet injection, using the heat pulse propagation method [7]. A heat pinch term $Q$, is not included in the modelling employed here. However, preliminary indications are that a credible global power balance for electrons cannot be constructed without such a flux. An exactly analogous determination of $D_e$ is in preparation, but has not been completed yet.

4. DETERMINATION OF $\chi_e$ AND $D_e$ USING THE 'FLUX GRADIENT' ANALYSIS

From eq.1 we have that a plot of $Q$ vs $-nVT$ has slope $\chi$ and an intercept $Q_p$. Similarly for $D$ and $T_p$. Such plots were analyzed for steady-state heat flux for plasmas under various conditions [1]. This technique can be extended to transient conditions by plotting successive time points, yielding the transport coefficients in a single pulse. Transients in density can be induced by injection of a small (2.7 mm) pellet or application of ICRF power. Analysis of these gives $D_e = 0.5 \pm 0.1 \text{ m}^2/\text{s}$. The corresponding thermal diffusivity is deduced to be $\chi_e = 2-3 \text{ m}^2/\text{s}$. The discharges analyzed have plasma parameters very close to those in section 2. Analysis of a number of discharges using this method yields $3\chi_e/D_e \leq 8$. Moreover, both these coefficients show an inverse dependence on the plasma current $I_p$, fig.4 and ref[1].
Table I summarizes all the measurements described in sections 2-4.

5. DISCUSSION

The consistent behaviour of the measured $x$ and $D$ suggests that these are linked, and that $Q_p$ and $\Gamma_p$ may be manifestations of the same transport mechanism [8]. Employing expressions derived in [8], we deduce that $2T_n \ln n / n T = Q_p$. The measured values of $Q_p$ [1] and $\Gamma_p$ are in rough agreement with this expression. The derivations of [8] imply that $\Gamma_p \propto (-VT/T)$ and that $Q_p \propto (-Vn/n)$. The dependence of $\Gamma_p$ on the $T_e$ scale-length can be inferred from analysis of the $n_e$ profile evolution during transients in which the $T_e$ profile undergoes a significant change. Peaking of the $n_e$ profile after sawtooth collapse and during current ramp-down are two such instances. Fig.5 shows the measured and modelled evolution of the line-integrated density in a sawtoothing discharge; if $\Gamma_p \propto (-VT/T)$, then the 30ms hesitation observed in the density rise following a sawtooth collapse is reproduced. Fig.6 shows a plot of $\Gamma_p/nD$ against $-VT/T$ during contraction of the plasma column in current ramp-down. $\Gamma_p$ is calculated as $\Gamma_p(r) = 2n/r \int [S(r') - dn(r')/dt] r' dr' + D(r)Vn(r)$. Fig.6 demonstrates the close correlation between these two parameters.

<table>
<thead>
<tr>
<th>$x_e$</th>
<th>$D_e$</th>
<th>$-Q_p$</th>
<th>$-\Gamma_p$</th>
<th>Method</th>
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<tbody>
<tr>
<td>m$^2$·s$^{-1}$</td>
<td>m$^2$·s$^{-1}$</td>
<td>kW·m$^{-2}$</td>
<td>$10^{19}$ m$^{-2}$·s$^{-1}$</td>
<td></td>
</tr>
<tr>
<td>2.9 ± 0.4</td>
<td>0.4 ± 0.12</td>
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<td>-</td>
<td>(a) sec.2, r/a -0.7</td>
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<tr>
<td>1.6 ± 0.3</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>(b) sec.3, r/a -0.5</td>
</tr>
<tr>
<td>2 - 3</td>
<td>0.5 ± 0.1</td>
<td>4.3 ± 0.6 [1]</td>
<td>1.5 ± 0.2</td>
<td>(c) sec.4, r/a -0.7</td>
</tr>
</tbody>
</table>

6. REFERENCES


ACKNOWLEDGEMENT

Part of the work discussed in this paper makes use of data obtained during pellet injection experiments performed under a collaboration agreement between the JET Joint Undertaking and the U.S. Department of Energy.
FIG. 1 The peak amplitude of the heat pulse, $\Delta T_e$, and the delay in the arrival time of the peak, $\Delta t$, vs r/a. Fitted lines give average speed and damping of the heat pulse.

FIG. 2 Delay $\Delta t$ between sawtooth collapse and arrival of density pulse. Line is the prediction with $D_e = 0.39 m^2 s^{-1}$.

FIG. 3 Evolution of $T_e$ at different radii r/a after injection of a pellet. The smooth lines are the modelled $T_e$ evolution at the measurement radii.

FIG. 4 Average Diffusion coefficient $D(r/a > 0.5)$ vs plasma current $I_\phi$.

FIG. 5 Fluctuations of line-integrated density, $n_e l$, during sawtooth activity: (A) measured, (B) modelled with a particle pinch proportional to the electron temperature scale-length.

FIG. 6 Evolution of $T_e/nD$ with electron temperature scale-length $\nabla T/T$ during decay of $I_\phi$ in discharge termination.
POLARIMETRIC MEASUREMENTS OF THE q-PROFILE

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ABSTRACT
First polarimetric measurements of the poloidal field distribution on JET are presented. q₀ appears to be significantly below 1 in sawtoothing discharges. ICRF stabilization of sawteeth leads to even lower q₀'s.

1. DESCRIPTION OF THE DIAGNOSTIC

The six vertical channels of the JET far-infrared multichannel interferometer have been modified to also make polarimetric measurements (Faraday rotation). The configuration is essentially as described in reference 1, and consists of a simultaneous measurement of the probing beam intensity along two orthogonal directions of polarization.

In JET, however, the half-wave plates which are used for calibrating the system are located before the input of the beams into the torus (rather than at the exits). Also, the recombination of the probing and reference beams and the separation of the unrotated and rotated components is done in 2 stages, using a polarization-independent beam splitter for the former and a tungsten wire grid for the latter. (See figure 1).

The "interferometer" signal amplitude is proportional to cos θ, where θ is the angle of Faraday rotation, while the "polarimeter" signal is proportional to sin θ. The first of these is squared to produce a DC signal proportional to cos² θ, the second is multiplied by the first in a phase-sensitive detector to produce a DC signal proportional to sin θ cos θ. The two DC signals are digitized, and their ratio used to determine tan θ. This method of measurement requires a calibration, which is performed by rotating quartz half-wave plates (located at the point of entry of the beams into the torus) to produce a rotation of the beam polarization in steps of 0.4°. The measurement of the Faraday rotation angle is made with an error < 5%, and a sensitivity of ~ 1%. The time resolution is 1-10 msec., determined by the integration time of the electronics. The lower limit on the time resolution is set by the 100 kHz modulation frequency of the polari-interferometer.

2. ANALYSIS

The Faraday rotation angle is proportional to \( \int n_e B_{||} \, dx \), where \( n_e \) is the electronic density and \( B_{||} \) is the magnetic field along the probing chord. This implies that polarimetric measurements must be used in conjunction with other diagnostics to infer the poloidal field distribution. In particular, the non-circularity of JET cross-sections means that the poloidal field enters into the Abel inversion of the chord-integrated data both as a parameter (determining the shape of the flux surfaces) and
as the unknown to be determined. The results presented here are derived using two approaches. In the first, the flux geometry is determined from magnetic measurements at the plasma boundary via an equilibrium identification code, IDENTC [2]. The poloidal field distribution is then determined by a generalized Abel inversion in this geometry. The electron density profiles are obtained from the interferometric data, or from the LIDAR diagnostic [3]. The agreement between these is generally very good (< 10%). The Abel Inversion is made in 2 steps. The discharge cross-section is divided into six nested elements, determined by the flux surfaces of tangency of the six polarimeter chords. A first calculation of the poloidal field distribution is made in this geometry, and used to calculate the expected angle of Faraday rotation on six vertical "virtual" chords, tangent to the same flux surfaces, but on the opposite side of the magnetic axis to the corresponding real chords. A Chebyshev polynomial is fitted to the complete set of Faraday rotation angles (real, virtual + zero points at the plasma boundary) and the second Abel inversion is performed on this smoothed data. This method ensures that the smoothed data conforms both to the data and to the transformation properties implied by the flux geometry. (See Fig. 2.)

The major radius of the magnetic axis, \( R_0 \), and the elongation of the flux surfaces near the axis, \( \kappa_0 \), are directly related to the value of \( q_0 \) which is inferred by Abel inversion in a given flux geometry. In fact, for a set of vertical probing chords, since the density profile is flat near the axis, the inferred value of the vertical field, \( B_\parallel \), is inversely proportional to the assumed \( \kappa_0 \). The inferred toroidal field at the axis, \( B_{\perp 0} \), is inversely proportional to the assumed \( R_0 \), and

\[
q_0 = \kappa_0 \frac{r}{R_0} \frac{B_{\perp 0}}{B_\parallel} - \frac{\kappa_0^2}{R_0^2}
\]

The position of the magnetic axis determined by the equilibrium reconstruction code is in good agreement with that determined by soft X-ray tomography [4]. However, the elongation of the central flux surface in the equilibrium reconstruction is systematically greater (5-10%) than the elongation of the SXR iso-emissivity surfaces. If the latter are taken to represent the true elongation of the flux surfaces, a lower value of \( q_0 \) is deduced. (See Section 3).

Uncertainties in the absolute calibration, the electron density profile, and the flux geometry, as well as the intrinsic sparseness of the data, lead to an error in \( q_0 \) of \( \pm 15\% \). Nevertheless, as discussed above, much of this error is systematic and changes in the \( q \)-profile can be measured with considerably better accuracy. Initial experiments using a horizontal displacement of the plasma column across the lines of sight to reduce the sparseness of the data have not yielded significant changes in the value of \( q_0 \) which is determined, suggesting that sparseness does not introduce large errors in this parameter (although this may not hold for the \( q \)-profile as a whole).

The second method of analysis is an integrated one, in which a solution to the Grad-Shafranov equation is found with the source term optimised to best fit both the polarimetric and magnetic measurements. Accordingly an objective function is built as the sum of the squares of the differences between calculated and measured quantities. Its minimum is found by expressing its gradient with respect to the free parameters, in terms of a set of generalized Lagrange multipliers, which are solutions of
the so-called adjoint problem. The techniques used are Finite Elements for the partial differential equations and Conjugate Gradients for the minimisation. This approach is still being developed, in particular with respect to the treatment of the electron density profile. Preliminary results show that the Faraday rotation information improves the knowledge of the innermost region of the plasma, and stabilizes the solution against measurement errors. A case of consistent reconstruction is shown in Fig. 3.

3. RESULTS

Figure 4 shows the evolution of \( q_0 \) during a 4 MA, 2.8 T, ohmic discharge. The elongation of the inner flux surfaces has been taken to be that given by the tomographic reconstruction of the soft X-ray emissivity. \( q_0 \) decreases rapidly during the ramp-up of the plasma current, and reaches a value of 1 at 5.5 seconds, nearly simultaneously with the onset of sawtoothing, as evidenced by the soft X-ray emission. \( q_0 \) then continues to decrease, ultimately saturating at a value of ~0.7. The steady-state value of \( q_0 \) does not appear to depend on the flat-top plasma current or the safety factor at the plasma boundary. During a sawtooth, \( q_0 \) changes little (<5%). In fact, the change in \( q_0 \) is at present confused by uncertainties related to the simultaneous redistribution of the density profile.

Figure 5 shows the evolution of \( q_0 \) during a 2 MA, 2.1 T, radio frequency heated discharge. The interval from 7.5 to 9.2 seconds corresponds to a period of sawtooth stabilization, known as a "monster sawtooth" [5]. At the start of the RF pulse the Faraday rotation angles increase, largely due to the RF-induced density rise and, in the case of chords inboard of the magnetic axis, to the displacement of the magnetic axis with increasing \( B_0 \). During the monster sawtooth, \( q_0 \) decrease by ~15%. A similar decrease is also inferred by IDENTC from magnetic data alone. These results are in good agreement with a field diffusion calculation performed using the TRANSP code [6], assuming neo-classical resistivity. (Note that the initial value of \( q_0 \) in the TRANSP calculation is prescribed by the sawtoothing model that was adopted, which resets \( q_0 \) to 1 after each sawtooth).

At the collapse of the monster sawtooth, the density profile is markedly perturbed, and this makes analysis of the \( q \)-profile more uncertain. Nevertheless it appears that when sawtoothing is re-established, \( q_0 \) returns to approximately its pre-monster value.

Taken together, the above results suggest that ordinarily the steady-state value of \( q_0 \) is determined by the sawtooth mechanism, perhaps indirectly via the temperature profile. When sawteeth are suppressed, \( q_0 \) can evolve to lower values by field-diffusion.

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FIG. 1: Schematic of the Faraday rotation measurement.

FIG. 2: Example of data analysis:

a) Electron density profile. Solid line: Abel inversion; dashed line: LIDAR data.
b) Closed: Measured Faraday angles; solid circles: mirrored angles; solid line: polynomial fit.
c) Poloidal magnetic field.
d) q-profile.

FIG. 3: Solid line: q-profile from a consistent reconstruction; dashed line: q-profile from Abel inversion.

FIG. 4: Evolution of $q_e$ (circles) and SXR intensity in an ohmic discharge.

FIG. 5: Evolution of the Faraday angle on a central chord of the electron temperature and of $q_e$ (measured and modelled) during an ICRH pulse with sawtooth stabilization.
OPERATION REGIME AND CONFINEMENT SCALING OF NEUTRAL BEAM HEATED 
JT-60 DISCHARGES

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1. OPERATION REGIME OF JT-60 DISCHARGE
During the last run period (Jun-Oct 1987), experiments were 
conducted putting emphasis on high plasma current operations to 
achieve better confinement. For this purpose, graphite-made limiters 
and armor tiles were introduced, along with the improvement of power 
supplies. The radiation loss in neutral beam heated limiter discharge 
has been lowered to ~ 30% of total input power (~ 20% in divertor 
discharge). As a result, plasma current up to 3.2 MA was achieved 
in limiter discharge and 2.7 MA in divertor discharge. The maximum 
stored energy \( W = 3.1 \text{ MJ} \) (± 10%) was obtained in a limiter discharge 
with \( I_p = 3.1 \text{ MA} \), \( P_{\text{abs}} = 23 \text{ MW} \), \( n_e = 1 \times 10^{20} \text{ m}^{-3} \) and \( T_i(0) = 3.9 \text{ keV} \). The 
fusion product \( n_e(0) \cdot T_i(0) \) achieved \( 6 \times 10^{19} \text{ m}^{-2} \cdot \text{s} \cdot \text{keV} \) with \( T_i(0) = 3-4 \text{ keV} \) in hydrogen plasmas with perpendicularly \( H^0 \) beam injection.

The vacuum vessel was kept at high temperature (200–300 °C) 
throughout the run period, which moderately regulated the wall 
pumping effect for density control. During current ramp-up phase, 
an \( m=3/n=1 \) locked mode occurs at \( q_s \), slightly lower than three followed 
by a disruption. This can be avoided by injecting sufficient \( N_B \) in 
ramp-up phase. The current ramp-up rate was optimized so as to 
minimize the volt-sec consumption and obtain a stable discharge. 
Lower hybrid wave was also applied to save the volt-sec consumption.

The operation region in \( n_e-I_p \) space is shown in Fig. 1. The 
maximum attainable density is roughly proportional to the plasma 
current. For a fixed current higher than 2 MA, the density is bounded 
within a certain region. The lower bound comes from the inevitable 
particle fuelling from \( N_B \) necessary to suppress the locked mode. 
Although the density can be raised by substantial gas puffing, it 
is limited by occurrence of MARFE. The maximum density of \( 1.2 \times 10^{20} \text{ m}^{-3} \) 
is achieved in a limiter discharge of \( I_p = 3.15 \text{ MA} \). The Murakami 
parameter reaches a value of \( 7.5 \times 10^{19} \text{ m}^{-2} \cdot \text{T}^{-1} \) at \( q_s = 2.2 \). In a divertor 
discharge, the maximum density of \( 6.6 \times 10^{19} \text{ m}^{-3} \) is obtained by 
intensive gas puffing of 30 Pam's⁻¹.

2. EMPIRICAL CONFINEMENT SCALING
Long pulse and high power auxiliary heating is one of the 
outstanding features of JT-60 discharge. Up to 20 MW of \( N_B \) has been 
injected into hydrogen plasmas with duration up to 6 s. This makes
JT-60's confinement nature far from that of Joule plasmas. And the resultant scaling found in this high power regime shows more of an asymptotic feature. The total energy contents are obtained from magnetic measurements which are consistent with those from kinetic data analysis /1/. The overall confinement property has not been much affected by the change of the first wall material.

In low power heating ($P_{abs}$ ≤ 5 MW), global confinement time $\tau_g$ has slight dependence on electron density. In a high power region ($P_{abs}$ ≥ 10 MW), $\tau_g$ becomes almost independent of $n_e$, as previously reported /2/.

Figure 2 shows the total stored energy as a function of the total absorbed power (shine through and reionization losses are subtracted). From this we see that $W$ has an offset linear relationship to $P_{abs}$ and that the incremental confinement time $\tau_{g\text{inc}} = \delta W/\delta P_{abs}$, is roughly independent of plasma current ($54 \pm 4$ ms for $a_p \geq 0.85$ m). Figure 3 shows the intercept $W(0)$ ($W$ extrapolated to $P_{abs}=0$) as a function of the plasma current. As in the figure, the offset increases with the plasma current as $W(0) = 0.19 I_p^{3/2}$ MJ. It is notable that the improvement of confinement in high current discharges mainly comes from the offset part $W(0)$ and that $W(0)$ is predominantly determined by the plasma current. Although the change in electron density may affect the power deposition profile or the transport property, its influence is not explicitly seen in the global confinement. Kinetic measurements show that higher current discharge has higher $T_e$ and broader $T_n$ profile /1/.

Both power-law type and offset-linear type scalings fit equally well to JT-60 data. However, performing current scans for fixed power yields different current dependence for different power input (Fig. 4). That is, although the power law scaling can reproduce a global confinement for a given set of parameters, it may not reflect precise dependence on each parameter. The same thing can be seen in earlier experiments on smaller machines that exhibits stronger current dependence for lower power input /3/.

The global confinement improves with increasing the plasma minor radius $a_p$ as shown in Fig. 5. These data are taken from divertor discharges in which the sum $a_p+R_p$ ($R_p$ is the major radius) is kept approximately constant. Also shown in the figure is a curve predicted by the Goldston scaling /4/. If the confinement scaling is expressed in terms of power laws as $\tau_g = R_p^{\alpha}a_p^{\beta}$, the JT-60 results suggests that $R_p^{\beta} > a_p^{\alpha}$. Hence $\beta$ can never be negative unless $\alpha$ is also negative. In any case, the Goldston scaling cannot reproduce the JT-60 results properly in size scaling. Due to the lack of data with sufficient $a_p$ scans, a definite size scaling has not yet been found. Still, a standard regression analysis using the offset linear scaling tells that the incremental confinement time relates to $a_p$, as $\tau_{g\text{inc}} = 0.062 a_p^{-1/8}$ s, which is consistent with the Shimomura-Odajima scaling. Moreover, comparison with JET results /5/ implies that the incremental confinement time should be proportional to the plasma cross section.

From the above arguments the empirical confinement scaling of JT-60 can be expressed in terms of offset linear scaling as

$$\tau_g = 0.19 I_p^{3/2}/P_{abs} + \tau_{g\text{inc}}$$

(1)
The regression analysis shows $-\frac{1}{6} \ln^2 - 0.0624^1.8 s$. Comparison of the global confinement time from magnetic measurement with the scaling law Eq. (1) is shown in Fig. 6.

3. CONCLUSIONS

It is found that the confinement of high power auxiliary heated plasma is predominantly determined by the plasma current. And the offset linear scaling represents the JT-60 confinement results better than the power law scaling, in that it can reproduce both the current and power dependence more properly. The confinement improves with increasing the plasma minor radius, which is consistent with the Shimomura-Odajima scaling.

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Fig. 1. Operation region in $n_e-I_p$ space. Closed and open symbols show limiter and divertor discharges, respectively.

Fig. 3. Offset part of the total stored energy $W(0)$ as a function of plasma current.
Fig. 2. Total stored energy versus absorbed power. Lines show best linear fits for each plasma current.

Fig. 4. Confinement time as a function of plasma current, exhibiting weaker $I_p$ dependence for higher input power.

Fig. 5. Dependence of confinement time on plasma minor radius for divertor discharges, showing contradiction to the Goldston scaling.

Fig. 6. Confinement time plotted against the offset linear scaling law (Eq. (1)).
ENERGY CONFINEMENT ANALYSIS OF NEUTRAL BEAM HEATED JT-60 DISCHARGES

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1. Introduction
In June to October 1987, neutral beam and/or rf heating experiments were
performed in JT-60 hydrogen plasmas with a graphite inner wall instead of TiC coated Mo
one, and with wall temperature at 200 ~ 300 °C. The operation region was extended to
higher levels of plasma current and density of up to 3.2 MA and 1.2x10^{20} m^{-3}, respectively
[1]. The contamination of metallic impurities (e.g. titanium) decreased in ohmically and
neutral beam (NB) heated discharges compared with that obtained in the discharges with
TiC coated Mo limiter. The heating experiments of near-perpendicular NB injection
(Ho --> H^+) were undertaken with beam energy (E_b) both of about 40 kV and of about 70
kV. According to a difference in beam energy, the variance of electron temperature profile
shape was observed for discharges with almost the same q_{eff}. Energy confinement time of
NB heated plasmas was improved not only in higher I_p discharges but also in higher E_b
ones.

2. Energy Confinement with low and high beam energy
Between NB heated discharges with low E_b (~40 kV) and those with high E_b (~70 kV) central electron temperature T_e in high E_b injection is higher than that in low E_b one,
which is irrespective of divertor or limiter configuration. The central ion temperature
measured by He beam scattering also showed the same tendency. The electron temperature
profile peakedness T_{e0}/T_e in low E_b injection is much lower than that in high E_b one.
Figure 2 shows that the difference of T_e (r) between both injections occurs in r < a/2 region.
The n_e dependence of the effective ionic charge Z_{eff} does not change in low and high E_b
injected discharges as plotted in Fig. 1. Also the sawtooth amplitude did not differ in both
cases.

The power balance and energy confinement in these discharges were analyzed by
using the time-independent 1-D transport code (LOOK / OFMC / SCOOP ) including a
Monte-Carlo calculation for power deposition and particle source of NB injected into plasma,
and beam component of plasma energy content [2]. For the ion transport, \chi_i = 5 \chi_i^C.H. [3]
was assumed. In Fig. 3, the convection loss evaluated at r < a/2 region is shown as a
function of n_e for low and high E_b injected discharges. It is noted that the convection loss is
enhanced in low E_b injection. This enhancement can be explained by the increase in particle
fueling of low E_b injection. In this case, the power input from beam to electron, P_{be}, is small
compared with that in high E_b injection case, but the power input from ion to electron, P_{ie},
increases because of lower \( T_e \). Though the net power input to electron, \( P_{he} + P_{te} + P_{OH} \), estimated in \( r < a/2 \) region in low \( E_b \) injected discharges is almost equal to that in high \( E_b \) ones, the enhanced convection loss may make electron temperature profile shape flattened as shown in Fig. 2. For the ion transport in \( r < a/2 \) region, the net power input to ion, \( P_{hi} - P_{te} \), in high \( E_b \) case is comparable with that in low \( E_b \) one for discharges with similar \( \bar{n}_e \) and \( P_{abs} \), while the measured central ion temperature in high \( E_b \) injection was higher than that in low \( E_b \) one. This may be explained by higher \( T_e \) and less convection loss in high \( E_b \) injected discharges compared with those in low \( E_b \) ones. Because the enhanced convection loss decreases with \( \bar{n}_e \), the global electron confinement time in \( r < a/2 \) region is improved with \( \bar{n}_e \): 50 msec ( \( E_b \sim 40 \text{ kV} \) ) and 70 msec ( \( E_b \sim 70 \text{ kV} \) ) at low \( \bar{n}_e \) of \( 2.5 \times 10^{19} \text{ m}^{-3} \), and 80 msec ( \( E_b \sim 40 \text{ kV} \) ) and 90 msec ( \( E_b \sim 70 \text{ kV} \) ) at high \( \bar{n}_e \) of \( 4.5 \times 10^{19} \text{ m}^{-3} \) for 1.5 MA discharges with the absorbed power, \( P_{abs} \), of about 15 MW.

On the other hand, \( T_e \) profile shape is almost independent of \( \bar{n}_e \) especially in low \( E_b \) injection discharges as shown in Fig. 1, though NB power deposition profile varies largely with \( \bar{n}_e \): extreme edge heating at high \( \bar{n}_e \) and center heating at low \( \bar{n}_e \) (profile consistency). This result suggests that \( \chi_e \) in \( r < a/2 \) region may be improved with \( \bar{n}_e \).

Figure 4 shows the comparison of the stored energy contents and the global energy confinement time between in low \( E_b \) case and in high \( E_b \) one. It should be noted that the stored energy of beam component, \( W_b \), in low \( E_b \) injected discharges is by a factor of two less than that in high \( E_b \) ones. This tendency is also recognized in the stored energy of thermal components, but their differences are small. The reduction of the stored energy results in the degradation of global energy confinement time, \( \tau_{E,G} \), in low \( E_b \) injected discharges as shown in Fig. 4 (bottom). It is interesting to note that \( \tau_{E,G} \) in low \( E_b \) case depends on \( \bar{n}_e \), while \( \tau_{E,G} \) in high \( E_b \) case is independent of \( \bar{n}_e \) as previously reported [5]. A small amount of the stored energy of beam component mainly leads to such a \( \bar{n}_e \) dependence of \( \tau_{E,G} \) in low \( E_b \) injected discharges.

3. Energy confinement of high \( I_p \) plasma

Figure 5 shows the \( \bar{n}_e \) dependence of \( T_{eo}, \bar{T}_e/\bar{T}_e \) and \( Z_{eff} \) for high power NB heated discharges ( \( P_{abs} = 15 \sim 21 \text{ MW} \) ) with high \( E_b \) ( \( 70 \sim 75 \text{ kV} \) ), indicating that with increase in \( I_p \) a similar \( T_{eo} \) is obtained even in higher density region, and electron temperature profile shape becomes flattened. Moreover the density profile shape is very flat at high \( \bar{n}_e \) in JT-60, then the plasma stored energy of thermal component increases with \( I_p \) and \( \bar{n}_e \) for a given \( P_{abs} \), while the total stored energy including beam component, \( W_{TOT} \), increases only with \( I_p \). It should be noted that the thermal stored energy itself is insensitive to \( \bar{n}_e \) in 3 MA NB heated discharges ( \( P_{abs} \sim 20 \text{ MW} \) ) with high \( \bar{n}_e \) above \( 8 \times 10^{19} \text{ m}^{-3} \), because \( T_{eo} \) scales as \( 1/\bar{n}_e \) and the difference of \( T_{io}-T_{eo} \) is small. Also in this case the stored energy of beam component occupies only a few percent of the total stored energy kinetically evaluated. The maximum total stored energy and its global energy confinement time are about 3 MJ and about 150 msec, respectively.

In Fig. 6, the total stored energy kinetically evaluated is plotted against \( P_{abs} \), which shows that \( W_{TOT} \) is improved with \( I_p \), and that the incremental energy confinement time, \( \tau_{E,inc} \), seems to be about 60 msec for the plasma current in the range of \( 1.5 \leq I_p \leq 3 \text{ MA} \). The offset linear dependence of \( W_{TOT} \) on \( P_{abs} \) is consistent with that estimated by the diamagnetic loop measurement [6] in an absorbed power range of \( 9 \leq P_{abs} \leq 22 \text{ MW} \). Though the global energy confinement time shows a typical L-mode degradation against \( P_{abs} \), it is improved with plasma current.
4. Conclusions

From a point of view of energy confinement, the characteristics of recent NB heated plasmas in JT-60 are presented and analyzed. Main results of them are:

1) Improvement of energy confinement was obtained in NB heated discharges both with high \( E_b \) injection and with high \( I_p \) operation. The former improvement can be explained by increase in the stored energy of beam component, and the latter one by increase in the stored energy of thermal component with the broadened \( T_{eo} \) profile shape.

2) In low \( E_b \) injection heating the total plasma stored energy for a given absorbed power, therefore energy confinement time, depends on \( \bar{n}_e \) because of the reduction of beam stored energy, while they are insensitive to \( \bar{n}_e \) in high \( E_b \) one.

3) Electron temperature profile shape varies with beam energy of NB injection: broadened profile shape (\( r < a/2 \)) is obtained in low \( E_b \) injection heating. This may be explained by the enhanced convection loss in low \( E_b \) injected discharges.

4) Electron temperature profile consistency for the variance of NB power deposition profile by \( \bar{n}_e \) scan may be related to possible improvement of \( T_{eo} \) with \( \bar{n}_e \).

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The authors would like to express their appreciation to Drs. S. Mori and K. Tomabechi for their continued encouragement and support.

References


Fig. 1 Comparisons of \( T_{eo}, \frac{T_{eo}}{<T_e>}, \) and \( Z_{eff} \) against \( \bar{n}_e \) between low \( E_b \) injection and high \( E_b \) one.

Fig. 2 Electron temperature profile in low and high \( E_b \) injection.
Fig. 3 Convection loss versus $\bar{n}_e$ in low and high $E_b$ injections.

Fig. 4 Comparison of stored energy contents between in low $E_b$ injection and high $E_b$ one.

Fig. 5 Comparisons of discharge characteristics of plasma current of 1.5 MA to 3 MA.

Fig. 6 Total stored energy versus absorbed power.
HIGH ENERGY ION TAIL FORMATION AND ITS BEHAVIOR
IN ADDITIONALLY HEATED JT-60 PLASMAS

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Introduction

JT-60 plasmas\(^1\) were additionally heated mainly by near-perpendicular neutral beam injection (NBI). In high power neutral beam heating by low and high energy injections, high energy ion tail formation and its behavior have been studied with charge-exchange diagnostics.

Charge-exchange diagnostics

Charge-exchange measurements were carried out with three neutral particle analysers and these sight lines are shown in Fig. 1. Two passive neutral particle analysers, electrostatic type (CX(a1)), 0.1 \(\sim\) 110 keV and E/B type (CX(a2)), 0.1 \(\sim\) 300 keV, were installed in the same plane inclining at an angle of 5° with respect to the poloidal plane and view a plasma at different spatial locations by angles of 38° and 29° to the equatorial plane, respectively. The other analyser, E/B type (CX(v)), 0.1 \(\sim\) 400 keV, sights a plasma vertically and the observation field of the analyser crosses a heating-beam zone near the center of the plasma. Active charge-exchange measurement is performed by the location. These two nearly perpendicular and one vertical analysers detect only particles with \(u_\parallel/u\approx0\sim0.1\.

Experimental results and discussions

1) High energy injection

In the high energy injection of 60 \(\sim\) 70 keV, energy relaxation time of beam-produced fast ions was measured in the parameter range of \(T_e(0)=1.6\sim3.5\) keV, \(n_e=(1.5\sim7.6)\times10^{19}\) m\(^{-3}\) and \(P_{\text{NBI}}=8.6\sim19.3\) MW from the temporal evolution of charge-exchange spectra after neutral beams turned off. The relaxation time \(\tau_{eL}\) was defined as a duration in which the charge-exchange spectrum returns to a Maxwellian distribution after the turn-off of neutral beam injection. Figure 2 shows the dependence of the relaxation time on \(T_e(0)^{3/2}/n_e\). The observed energy relaxation time was approximately proportional to the parameter and the tendency was clear when the electron temperature was nearly constant at 3 \(\sim\) 3.5 keV. A significant difference was not observed in the relaxation time due to plasma configurations (limiter or divertor) and injection directions of neutral beams (co-, counter- or nearly balanced injection) in the observed parameter range. The beam slowing-down times calculated for electron temperatures of 1 \(\sim\) 4 keV based on the classical collision process were also shown by the solid lines in Fig. 2. The measured relaxation time was comparable to the beam slowing-down time.

In order to investigate the slowing-down process precisely, the observed charge-exchange spectrum was compared with a one-dimensional...
bounce-averaged Fokker-Planck calculation. Figure 3 compares the measured charge-exchange spectrum with the calculation in high energy injection at 70 keV into a low-density plasma of a limiter configuration. In the figure, both spectra are normalized at the injection energy. The injected power was up to 22 MW and the beam power fraction $P(E):P(E/2):P(E/3)$ was 80:12:8. The line averaged electron density was $2.5 \times 10^{19} \text{m}^{-3}$ and the central electron temperature was 5 keV. The central ion temperature was determined to be 11 keV by an active beam scattering method. In the calculation, the ion temperature profile was assumed to be similar to that of the electrons. The obtained charge-exchange spectrum shows a good agreement with the Fokker-Planck calculation. From observed energy relaxation time and an agreement between the experiment and the simulation, the slowing-down process of beam-produced fast ions was well explained by the classical collision process.

2) Low energy injection

When neutral beams were injected at low energy of $32 \sim 40 \text{keV}$ into a limited plasma ($r_p = 78 \sim 80 \text{ cm}$) with a large clearance between the plasma boundary and the outer wall and low plasma current of $0.7 \sim 1 \text{ MA}$, injected ions were accelerated up to 90 keV by beam-driven ion cyclotron range of frequency (ICRF) wave. The limiter configuration accompanied with the beam acceleration is shown by the outermost magnetic surface in Fig.1.

Typical charge-exchange spectra with the high energy ion tail for three sight lines are shown in Fig.4. Three spectra in a case without the high energy tail are also shown for a comparison in the figure. Charge-exchange spectra measured by near perpendicular analysers, CX(a1) and CX(a2), show a remarkable acceleration of beam ions, whereas spectrum from a vertical analyser, CX(v), showed no remarkable difference. Since CX(v) sights a heating-beam zone where the neutral density is increased, charge-exchange flux is enhanced by a factor of $\sim 10$ when beams in the observation field are injected. So CX(v) mainly measures charge-exchange flux from the central part of the plasma. Charge-exchange spectra measured from three directions suggest that the beam acceleration occurs in the peripheral region of the plasma. Figure 5 shows the dependence of tail temperatures, $T_{tail}$, measured by CX(a1) on the injected neutral beam power for different electron densities. The tail temperature increased unexpectedly with the injected power and electron density when the injected power was above $\sim 12 \text{MW}$. The tail temperatures measured by CX(v) shown by hatched region increased gradually with the neutral beam power. In the cases of injections with co- or counter-beams alone, the high energy ion tail was not observed because injection power could not exceed the threshold power of $12 \text{MW}$.

Accompanied with the high energy ion tail formation, ICRF wave was observed by electrostatic probes located just outside the first wall (Fig.1). The ICRF wave was observed only with beam injection and measured continuously during a beam pulse. This indicates that the neutral beam injection drove the ICRF waves. As shown in Fig.6 with a temporal evolution of average electron density, the frequency spectra have a prominent peak near $104 \text{ MHz}$, which was corresponding to $2\omega_{ci}$ of outer peripheral region of the plasma. In contrast to previous experiments on JFT-2M/2/1, PDX/3/ and JET/4/, prominent peaks of fundamental and other higher harmonic frequency waves were not observed but one shot in cases of low energy injection on JT-60. Figure 6 also shows that the wave intensity increased with electron density, and
then the central frequency shifted slightly to lower frequency side. This down-shift indicates that the location, where the ICRF wave was excited, shifted to outer region due to the density increase or the increase of beam ion component, \( n_B/n_e \).

When the neutral beams are injected at low energy, beam ion density increased in the outer part of the plasma. And in the case of near-perpendicular injection to the magnetic field lines, it is generally shown that such a way of injection give rise to a considerable amounts of anisotropy in the outer region, where pitch-angle scattering is less dominant than in the central part. Some neutral beam heating experiments mentioned above showed that ICRF waves were driven by a positive slope of velocity space, \( \delta f/\delta v_n > 0 \), and anisotropy of pressure or ion temperature, \( T_i/T_{ib} \). Instead of a temperature anisotropy, the anisotropy was evaluated by beam temperatures defined as \( T_{bi} = \langle E_i \rangle \) and \( T_{b2} = 2 \langle E_i \rangle \), which were the velocity-space average perpendicular and parallel energies of beam ions. From the velocity distribution determined by a bounce-averaged Fokker-Planck code without considering the interaction of the ICRF wave with beam ions, \( T_{bi}/T_{b2} \) was expected to be 6~8 at \( r/a > 0.8 \). It is considered that the high beam ion density and the anisotropy in the peripheral region are some of possible causes of the ICRF wave excitation and the resulted ICRF wave acceleratated beam ions remarkably.

**Summaries and conclusions**

In the case of high energy injection, the observed energy relaxation time was shown to be comparable to the beam slowing-down time predicted from classical collision process. Precise comparison of charge-exchange spectrum with the Fokker-Planck calculation showed that the slowing-down process was well explained classically. On the other hand, when neutral beams were injected at low energy with high power, beam ion density increased at outer parts of the plasma and considerable amounts of anisotropy in velocity space was expected in the peripheral region. It is considered that this anisotropy in the velocity space excited ICRF wave and it accelerated beam ions unexpectedly. Investigation of the excitation mechanism of the ICRF wave, including waves driven by fusion products, and the interaction of the wave with plasma particles should be one of developing subjects in present-day large tokamaks for reactor-grade plasmas.

**Acknowledgements**

The authors would like to express their gratitude to Drs. S. Mori, K. Tomabechi, M. Yoshikawa and S. Tamura for their continued encouragement and support.

**References**

Fig. 1: Sight lines of charge-exchange analysers.

Fig. 2: Dependence of energy loss time $\tau_{EL}$ on $T_e^{3/2}/n_e$.

Fig. 3 (left): Comparison of charge-exchange spectrum with Fokker-Planck calculation.

Fig. 4: Charge-exchange spectra with (open) and without (filled) a high energy tail.

Fig. 5: Dependence of tail temperatures on neutral beam power.

Fig. 6: Frequency spectra of beam-driven ICRF wave.
INITIAL RESULTS FROM THE PBX-M TOKAMAK


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Summary

During the initial period of PBX-M operation, discharges with indentations up to 20 to 24% and plasma currents up to 450-600kA at $B_t = 1.2T$ have been produced. During OH operation, the sawtooth period is seen to increase from 5 msec up to 40 msec as the indentation increases from 3% to 16%. Typical H-mode transitions, with a sudden drop in the $I_{HA}$ emission and increase in electron density, have been observed with $P_{inj} > 2.5$MW.

Introduction

The goal of the PBX-M experiment is to demonstrate access to the 2nd regime of ideal MHD stability against ballooning and internal kinks and to achieve high $\beta$ plasmas relevant to attractive reactors. This goal is being attempted with strongly indented plasma shaping (indentation up to 30%) and by employing both active and passive controls to enhance the stability properties against the MHD modes. A long term goal of the program is to establish physics understanding of 2nd stability regime and to design an optimized reactor that will operate in the 2nd stability regime.

The passive control system in PBX-M is a stainless steel covered aluminum conducting shell which surrounds the plasma surface almost entirely except for a 40 cm gap for the neutral beam injection and diagnostic access. These passive plates are expected to enhance the plasma stability against various surface modes. The feedback control system is used also to control the separatrix position. Active controls for the pressure and current profiles, in order to produce and maintain profiles favorable for operation in the 2nd stability regime, will be available in early CY89. These include 2 to 4MW of Ion Bernstein Wave Heating and 2MW of Lower Hybrid Current Drive.

Plasma formation and shaping control

Plasma formation is done by first creating circular plasmas; these are then shaped into the bean-shaped configuration as the plasma current is increased to 400 to 600kA. The final plasma size is 30cm wide on the horizontal midplane plane, 60cm high, and the indentation can range from 15 to 25% depending upon the relative strength of the indentation field coil current to the plasma current.
For the shaping control feedback system, the evolution of plasma shape is monitored by measuring the flux and B-field at 20 locations facing the plasma on the passive stabilizers. The symmetric components of these signals are used for the shaping control and the asymmetric components are for the vertical position control. An analog computer is used to decompose the symmetric components into three dominant eigenvalues of the shape deformation. These components are the dipole, corresponding to the plasma radial shift, the hexapole, giving the magnitude of indentation, and the third moment controls the width of plasma around the lobe area. The system requires approximately 70 summing circuits and 50 coefficients which can be varied conveniently through a terminal. At present, the system is linear and does not compensate the shielding effects due to the passive stabilizers. Nonetheless, the system successfully maintains the plasma shape within the passive stabilizer geometrical limit. With a modest plasma current ramp the plasma current can reach 600kA. For example, as shown in Fig. 1 the plasma has the indentation of 24% with $q_{\text{edge}}=3.5$, $B_t=12.5$ and $I_p=590$ kA.

The internal plasma shaping is also monitored by the soft X-ray emission measured tangentially with a pin hole camera. The preliminary analysis of this measurement indicates that near the 1/2 plasma radius the elongation is approximately 2, which is more than that measured previously in PBX.

![Fig. 1 The plasma flux surfaces of 590kA: $B_t=12.2$ KG, indentation $=24\%$, $q_{\text{edge}}=3.5$.](image)

**Vertical position control**

The vertical positional instability is a potential danger for the non-circular plasma configuration. However, in PBX-M, the nearly perfectly surrounding conducting shell
provides strong stability against MHD time scale vertical motion. In the typical experiments the external field varies strongly on the horizontal plane and \((\partial B_z/\partial R)/(B_z/R)\) is approximately -3 to -4 on the magnetic axis. This field gradient is about a factor of 2 higher than that in PBX. The feedback system of vertical position requires some compensation to avoid the overstable situation, where the plasma column oscillates vertically with the time period comparable to the passive stabilizer \(L/R\) time constant. For example, if the sensor located outside the passive stabilizer is used for the detecting the vertical position, it is necessary to add a small but finite subtraction of feedback current itself to stabilize the oscillation.

**MHD behavior**

A few interesting features of MHD, particularly on the sawtooth behavior were observed in discharges with higher indentation. During the OH phase, the sawtooth period was initially 5 msec, increasing in time and reaching 40 msec just before the current ramp down phase (Fig. 2). Although it is possible that the factor of 1.5 to 3 increase of the plasma density may prolong the crash period \(/L\), the observed density increase is not sufficient to explain the observed prolonged periods. One possible mechanism is that the gradual increase of indentation strength from 3% to 16% may

![Graphs showing current, density, and X-ray emission over time](image)

*Fig. 2* Time history of an OH plasma showing the current (top), density (middle), near-central soft X-ray emission (bottom), and indentation (%).

change the onset condition of sawtooth crash by adding more triangularity around \(q=1\) surface as was suggested by Bussac and Pellat [2]. Hard X-ray bursts, indicating a loss of runaways, was often observed during the sawtooth crash (Fig. 2 at 530 and 565
msec). The precursor oscillation was \( m=1/n=1 \) is seen only during the first few sawteeth, and it reappears only during neutral beam injection. However, a successor oscillation with frequencies of 1 to 2 kHz is often observed after the crash but it decays within several cycles. The crash appears to originate at a preferential toroidal and poloidal location.

**H-mode**

Clear transitions into the H-mode were observed in separatrix limited discharges. Usual confirmation of the transition was made from the TV pictures where the sudden sharpening the plasma edge and decrease intensity over the entire field of view was seen. The H-mode transitions seen in PBX-M were fairly typical, as seen in Fig. 3. For the discharge shown here, the transition occurred at \( t=475 \) msec, during neutral beam injection of >2.5MW. A threshold of 2.5MW was essential for the transition. The transition typically occurred at the crash of a sawtooth, and was characterized by the sudden decrease in the H-alpha emission and rise in electron density. A sawtooth crash was not necessary to trigger the transition, however, several discharges exhibited transitions (characterized by the drop in the H\( _\alpha \) emission) that occurred 1 to 2 msec after a sawtooth crash or not related to a sawtooth crash at all. In addition, many of the transitions were slow; that is, the drop in the H\( _\alpha \) emission occurred over a period of 0.5 to 2msec. Finally, little change in the MHD activity, as measured by the toroidal and poloidal Mirnov coil arrays, was seen across the transition. Although the analysis is still underway, if anything there was an increase in the broad band activities(30-50KHz) across the transition.

![Fig. 3. Time history of a 2.5 MW NBI discharge showing the H-mode transition at \( t=425 \) msec.](image)

**References**

ION TEMPERATURE AND ENERGY BALANCE IN OHMIC FT DISCHARGES

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INTRODUCTION

A systematic exploration of the FT operating Hugill diagram has been performed in the range $2.4 < q < 4.5$, $1. < nR/B < 3.7 \times 10^{19}$ Wb$^{-1}$ at $B = 6$ T [1]. A set of repetitive plasma discharges has been obtained at 14 selected locations in the Hugill diagram. The ion temperature has been measured with a 7 channel, mass resolving neutral particle analyser (NPA) and by four neutron yield monitors. The NPA looks at the plasma perpendicularly to the magnetic field and radial scans of the charge exchange neutral emission have been systematically completed, shot by shot at the selected locations in the Hugill diagram. Due to the perpendicular direction of observation the NPA detects neutrals coming from ions with a very high perpendicular velocity which are trapped in the local mirrors of the toroidal field ripple. At the high energies observed by the NPA ($3 < E/T < 9$) trapped ions are subjected to the vertical $B \times VB$ drift and it is well known [2] that if the measurement is performed in the drift direction very high apparent temperatures are measured even at the plasma edge. In order to minimise this effect the profile measurements are performed in the plasma half section opposite to the drift direction.

The ion temperature from the neutron yield monitors is available for all the plasma discharges while the peak value from the NPA cannot be obtained in the high density, low temperature cases due to the plasma opacity to neutrals. When both peak values are available a good agreement is found. The dependence of the peak ion temperature $T_i(0)$ from the average plasma density $n$ and the current $I$ can be empirically expressed by $T_i(0) = 2.1 \times 10^7 / n^{0.4}$ (keV, MA, $10^{14}$ cm$^{-3}$). In the next section the measured peak ion temperature are compared with the results of the ion plasma balance using the neoclassical thermal conductivity.

The experimental ion temperature profiles are very wide in all the selected locations on the Hugill diagram and the measured ion temperature in the outer part of the plasma seems to be significantly higher than the electron temperature and typically in the range 0.4, 0.6 KeV. These results can not be reproduced by a neoclassical code and suggest to be due to the fact that even in the plasma half plane opposite to the drift direction the distribution of the high energy perpendicular ions is different from the local bulk distribution. The reason for this can be seen in the ripple convective effect [3] and in the last section the results of a simple model are illustrated.

ION POWER BALANCE

The ion power balance equation $P_{ei} - P_{ree} - P_{ex} - P_{conv} - P_{re} = 0$ has been solved in cylindrical geometry using the experimental value of the electron temperature by ECE and soft-X measurement. Figure 1 shows both the experimental and the computed
values of the peak ion temperature in the different locations on the Hugill diagram. The computed values range refer to an anomalous neoclassical factor between 0.5 and 2.0 which seem to be compatible with the experimental error in all the cases.

**RIPPLE TRAPPED ION MODEL**

The deformation of the high energy distribution of the perpendicular ions in the direction opposite to the vertical drift is due to the ripple convective effect. The ripple trapped ions drifting vertically towards the plasma edge can be detrapped from the ripple magnetic well and in this case they become banana ions which can reach the opposite location on the vertical plane. In this last location, by collisions they can become trapped again in the ripple magnetic well. A simple model has been tentatively developed to quantify this effect. At a given energy the ion distribution function has been divided in three parts passing ions $N_p$, banana ions $N_b$ and ripple trapped $N_t$. The angular limit between the three components are given by $\sqrt{\Delta}$ and $\sqrt{\delta}$ where $\Delta$ is the toroidal magnetic well depth for banana ions and $\delta$ is the ripple well depth. An angular balance equation can be written for trapped and banana fractions at a given vertical location $y$

$$V_{\parallel} \frac{d}{dy} N_t(y) = -\left( \frac{1}{\tau_{tb}} \right) N_t(y) + \frac{1}{\tau_{bt}} N_b(y)$$

$$0 = \frac{1}{\tau_{tb}} N_t(y) - \left( \frac{1}{\tau_{bt}} + \frac{1}{\tau_{bp}} \right) N_b(y) + \frac{1}{\tau_{rp}} N_p(y)$$
where the times $t_{tb}$, $t_{bl}$, $t_{bp}$ provide the rate of exchange between the fractions. The passing ions are supposed to be maxwellian at the local temperature. The banana density at the location $y$ is supposed to coincide with the banana having their turning points in $y$. An similar equation can be written for banana ions in $-y$

$$0 = \frac{1}{t_{tb}} N_t(-y) - \left( \frac{1}{t_{bl}} + \frac{1}{t_{bp}} \right) N_b(-y) + \frac{1}{t_{pb}} N_p(-y)$$

$V_d$ is the ripple trapped drift velocity. The collisional times $t_{tb}$, $t_{bp}$ are known $t_{bl}$, $t_{bl}(\sqrt{\Delta - \sqrt{\delta})^2}$. The inverse collisional times $t_{bl}$, $t_{bp}$ can be found by imposing that in the case $V_d = 0.0$ the angular balance equations provide an isotropic ion distribution. The previous system can be solved and the proper ion distribution function can be used to compute the neutral emission. In Fig. 2 a typical experimental result for 5 radial positions is shown while in Fig. 3 the computed spectra both for a standard maxwellian model and the ripple convective model are plotted.

The comparison shows that the experimental results are intermediate between the two model results. In fact the considered model do not allows for any diffusion loss or thermalization of banana ions so that banana density is overestimate in the outer part of the plasma.

**CONCLUSIONS**

In the explored parameter range of FT the peak ion temperature seems to be compatible with neoclassical predictions inside the experimental uncertainties. The measured neutral particle profiles can not be directly used to derive the bulk ion
temperature profile and the simple previous model reproduces qualitatively but not quantitatively the measured results.

REFERENCES

INTRODUCTION

A study of impurity transport as a function of the major plasma parameters has been performed in the FT Tokamak by the impurity injection method. Chromium was chosen as the test element for this experimental campaign because it has a good number of spectral lines observable in the vacuum-ultraviolet (VUV) spectral region and is not fully ionized even in the central region of an FT discharge.

The atomic beam of Chromium was produced by laser ablation from a target and its diffusion in the plasma could be followed by monitoring the soft-X and VUV emissions. The intensity of the source was kept low enough to perturbate the plasma current not more than 5%. In particular the time decay of the emission was correlated to the other discharge characteristics and a numerical model was used to simulate the complete time evolution of the soft x-ray signals and to evaluate the parameters characterizing the impurity transport in the plasma.

Assuming a reference discharge defined by: \(B_T = 6\,\text{T}, I_p = 300\,\text{kA}, n_e = 1.2 \times 10^{14}\,\text{cm}^{-3}\), we scanned around it the magnetic field: \(4\,\text{T} < B_T < 8\,\text{T}\), the density: \(0.8 < n_e < 2.1 \times 10^{14}\,\text{cm}^{-3}\) and the \(q\) parameter: \(0.14 < 1/q_L < 0.31\). Each discharge was repeated several times to reduce reproducibility errors.

IMPURITY CONFINEMENT TIME

A typical decay time of the x-ray diodes signal from various lines of sight is shown in Fig. 1 where it appears that in most cases \(t(r)\) is constant inside a radius

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**Fig. 1** - Soft X-ray decay time vs line of sight coordinate

**Fig. 2** - Central impurity decay time vs the safety factor
corresponding roughly to the $q = 1$ surface and drops to a smaller value outside this region. In agreement with the results from other tokamaks [1,3] we found that the central decay time $\tau(0)$ does not depend on plasma density. Also a comparison of two discharges in hydrogen and deuterium agrees with a linear dependence on the background gas mass [1,2]. On the other hand the dependence of $\tau(0)$ from the safety factor exhibited in our case a different behaviour depending on the way the $q$ value was varied; from Fig. 2 it appears that the the experimental points align on two different straight lines one of which, obtained at fixed plasma current, shows the same dependence as in [1,3], while the other, corresponding to fixed magnetic field, has the opposite slope. A trial to find a single law describing the previous data was made by plotting $\tau(0)$ against the $B_T \times I_p$ product, this is shown in Fig. 3 where all the experimental points obtained up to now dispose fairly well on a straight line.

**NOT SAWTOOTHING DISCHARGES**

The results discussed in this communication have been obtained in stationary sawtoothing discharges with parameters not too close to the operational limits of the tokamak. It has already been observed that when the tokamak operating conditions are not strictly specified, confinement can be a not unique function of the main discharge parameters [4,5,6]. Evidences of this were obtained occasionally in FT for some discharges, slightly differing from those we looked for, during which sawtooth activity did not start up. Central x-ray emission for two of them, obtained at $n_e = 1.2 \times 10^{14}$ cm$^{-3}$, are compared in Fig. 4 showing that not sawtooothing discharges have much longer impurity confinement times than the corresponding sawtooothing ones.

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**Fig. 3** - Central impurity decay time vs $B_T \times I_p$

**Fig. 4** - Comparison of central soft x-ray emissions for two similar discharges ($B_T = 6T$, $I_p = 300$ kA, $n_e = 1.2 \times 10^{14}$ cm$^{-3}$) with and without sawtooth activity
NUMERICAL SIMULATION

In order to obtain information about the transport mechanism of the impurities in the plasma we have numerically evaluated the time behaviour of the impurity density by solving their diffusion equation. According to the current model of anomalous impurity transport [7], we assumed that the transport flux is described by a diffusion coefficient $D$ and a convective velocity $v$, equal for all the ionized states:

$$\dot{n} = D \frac{n}{r} - v \frac{n}{r/a}$$

By taking $D$, $v$ and the temperature and density profiles of the discharge as input parameters, the numerical code evaluates some VUV line brightnesses and soft x-ray emissions for Chromium including K and L lines, bremsstrahlung and recombination continuum. The best choice for $D$ and $v$ is obtained by simulating both the rise and decay phase of the impurity x-ray emission along the observed lines of view. Nevertheless if we assume that $D$ and $v$ are constant along the plasma cross section we are only able to simulate discharges with a weak sawtooth activity; for example, a satisfactory fit of the signals from x-ray diodes viewing along the central chord and at 7 cm from the centre is obtained choosing $D = 5000 \text{ cm/s}$ and $v = -600 \text{ cm/s}$, for the discharge defined by $B_T = 6 \text{T}$, $I_p = 300 \text{kA}$, $n_e = 2.1 \times 10^{14} \text{ cm}^{-3}$.

Conversely, in strongly sawtoothing discharges we observe that the rise time of the experimental signal is much longer for central chords than for chords passing outside the $q=1$ surface. This led us to assume a different value of the transport parameters inside this surfaces. Figure 5 refers to a discharge of this kind and

![Graphs showing comparison of experimental soft X-ray at r = 0 cm (a) and at r = 7 cm (b) with numerical predictions.](image)

**Fig. 5** - Comparison of experimental soft X-ray at $r = 0 \text{ cm}$ (a) and at $r = 7 \text{ cm}$ (b) with numerical predictions ($D = 7000 \text{ cm/s}$, $v = -3850 \text{ cm/s}$ outside $q = 1$ and $D = 1700 \text{ cm/s}$, $v = -190 \text{ cm/s}$ inside $q = 1$) for a discharge with strong sawtooth activity ($B_T = 6 \text{T}$, $I_p = 300 \text{kA}$, $n_e = 1.6 \times 10^{14} \text{ cm}^{-3}$)
illustrates the result obtained for $D = 7000 \text{ cm/s}$, $v = -3850 \text{ cm/s}$ outside the $q=1$ surface and $D = 1700 \text{ cm/s}$, $v = -190 \text{ cm/s}$ inside that surface. This result appears consistent with the hypothesis that sawtooth activity acts against impurity penetration in the plasma centre strongly reducing the inward velocity in that region [8,9].

CONCLUSIONS

Impurity confinement scaling measured on FT agrees with known results from other tokamaks except for the dependence on plasma current. Numerical simulation requires different coefficients inside and outside the $q = 1$ surface in the general case, nevertheless when sawtooth effect is observed to be weak a uniform treatment of the transport gives satisfactory results.

REFERENCES

EXPERIMENTAL OBSERVATION OF ION-TEMPERATURE-GRADIENT-DRIVEN TURBULENCE IN THE TEXT TOKAMAK


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Previous measurements on the Texas Experimental Tokamak (TEXT) have shown that the global energy confinement time ($\tau_E$) saturates with increasing plasma density. This trend, which has been observed in numerous other ohmically-heated tokamaks, has been attributed to the increased dominance of the ion thermal loss channel over the electron channel at sufficiently high densities. This paper describes the identification of an ion mode (i.e., microturbulence propagating in the ion diamagnetic drift direction) in the spectrum of tokamak density fluctuations which correlates with $\tau_E$ saturation and may be a signature of ion-temperature-gradient-driven turbulence. In addition, a profile-consistent microinstability-based model is used in the BALDUR-transport code (time dependent, 1D) for the purpose of interpreting anomalous confinement properties. Fluctuation measurements are made with a twin-frequency multichannel far-infrared laser scattering system which can measure the $S(k,\omega) \equiv \langle \hat{n}(k,\omega) \rangle^2$ spectra at six separate wave vectors during a single tokamak discharge.

Microturbulence poloidal frequency spectra for low ($n_e \approx 2 \times 10^{13}$ cm$^{-3}$), medium ($n_e \approx 4 \times 10^{13}$ cm$^{-3}$), and high ($n_e \approx 8 \times 10^{13}$ cm$^{-3}$) density gas-fueled TEXT discharges are shown in Fig. 1. The scattering volumes at each $k_\parallel$ are located above the midplane along a vertical chord at the major radius $R = 1m$. The heterodyne receiver system permits resolution of the wave propagation direction with negative (positive) frequency corresponding to the electron (ion) diamagnetic drift direction. For each wave vector examined, there is a distinct large-amplitude broadband ($\Delta \omega/\omega = 1$) peak at negative frequencies. However, as the plasma density is increased, an additional peak in the microturbulence spectra is also observed at positive frequencies corresponding to the ion diamagnetic drift direction. Since the fluctuations are measured in the laboratory frame of reference, it should be noted that plasma rotation effects induced by a negative plasma potential would serve to shift the spectra to more positive frequencies. The appearance of the ion feature in the fluctuation spectra occurs simultaneously with the saturation of $\tau_E$ and is a signature of the ion-temperature-gradient-driven instability. This instability occurs when the parameter $\eta_i (= L_{ni}/L_{ti}$, ratio of ion density to temperature scale lengths) increases beyond a threshold level in the range $\eta_{ic} = 1-2$. 
Figure 1. Microturbulence poloidal frequency spectra for $k_\theta = 7, 9, 12$ cm$^{-1}$, $I_p = 350$ kA, $B_T = 2.8$ T, and (a) $n_e = 2 \times 10^{13}$ cm$^{-3}$, (b) $n_e = 4 \times 10^{13}$ cm$^{-3}$, and (c) $n_e = 8 \times 10^{13}$ cm$^{-3}$. Vertical axis are in arbitrary units.
By injecting pellets into the plasma, it is possible to obtain high-density discharges with sharply peaked $n_e(r)$ when compared to gas-fueling. This in turn reduces the density scale length $L_n$ and can potentially be used to drive $\eta_i$ below the critical level for instability. Figures 2 (a) and (b) show the fluctuation frequency spectra at times before ($\bar{n}_e = 3 \times 10^{13}$ cm$^{-3}$) and after ($\bar{n}_e = 6 \times 10^{13}$ cm$^{-3}$) the pellet injection. For each case, the fluctuations are observed to propagate predominantly in the electron diamagnetic drift direction suggesting the $\eta_i$ instability has not been excited. In contrast, a comparable high-density gas-fueled discharge contains appreciable levels of the turbulence propagating in both directions, as shown in Fig. 2(c). Energy confinement properties of the pellet discharge are under active investigation but presently unknown.

Figure 2. Microturbulence poloidal frequency spectra for $k_\perp = 9$ cm$^{-1}$, $L_p = 250$ kA, $B_p = 2.8$ T, and (a) pre-pellet; $\bar{n}_e = 3 \times 10^{13}$ cm$^{-3}$, (b) post-pellet; $\bar{n}_e = 6 \times 10^{13}$ cm$^{-3}$, and (c) high-density gas-fueled discharge; $\bar{n}_e = 6 \times 10^{13}$ cm$^{-3}$. Vertical axis are in arbitrary units.

In the preceding discussions, it has been pointed out that the presence of an ion feature in the measured fluctuation spectrum is consistent with theoretical predictions for the appearance of ion temperature gradient driven drift instabilities ($\eta_i$ modes). Moreover, the observed confinement properties for gas-fueled plasmas appear to exhibit significant improvement when the ion feature is absent. These trends have motivated a systematic simulation study of the relevant TEXT discharges using microinstability-based models for the anomalous electron and ion thermal diffusivities, $\chi_e$ and $\chi_i$, in the BALDUR transport code. The trapped-electron modes in both collisional and collisionless regimes are represented in $\chi_e$ and, when the parameter $\eta_i$ exceeds $\eta_{ic}$, the toroidal ion temperature gradient modes account for enhancements to $\chi_i$. For $\eta_i < \eta_{ic}$, the ion losses are taken to be only neoclassical. As described in detail in Refs. 5 and 6, this model was calibrated on deuterium plasmas from TFTR and successfully applied to the simulation of numerous discharges from TFTR, ASDEX, and ALCATOR-C. In the present calculations, which deal with hydrogen plasmas, a scaling factor $Z^2/M$ was used in the expression for $\chi_i$ to accommodate the familiar empirical mass-charge dependence observed in most tokamak experiments which fall in the saturated ohmic regime. Preliminary results from simulations of helium discharges on TFTR have also yielded good agreement when this empirical factor was utilized.
Results from the transport code simulations for $\tau_E$, $T_{eo}$, and $T_{io}$ (central electron and ion temperatures) are compared with corresponding experimental results displayed in Table I. The entries are representative cases from the database of Fig. 1 and Ref. 7, where the fluctuation properties in gas-fueled discharges were investigated. In both the medium density case ($n_e = 4 \times 10^{13}$ cm$^{-3}$) where the ion feature was weak, and the high density case ($n_e = 8 \times 10^{13}$ cm$^{-3}$), where the fluctuation measurements indicate active ion modes, the simulation and experimental results are found to be in reasonable agreement. It is important to point out here that for the high density case, where the ion feature in the fluctuation spectra is strongest, agreement was possible only if the anomalous $\chi_1^{7\mu}$ was turned on in the simulation, thereby decreasing $\tau_E$ by a factor of two. For the medium density discharge, where the ion feature is small, the effect of $\chi_1^{7\mu}$ on confinement is significantly reduced.

Table 1. Comparison of BALDUR simulations with TEXT experimental results for two gas-fueled discharges.

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From these results it is clear that there is a strong correlation between fluctuation measurements of the ion feature, $\tau_E$ saturation in TEXT, and predictions from the microinstability-based transport model including anomalous $\chi_1^{7\mu}$ effects. Future work will include using the simulation model for pellet discharges and low density gas-fueled discharges.

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Coupling of Particle and Heat Transport
Measured via Sawtooth Induced Pulse Propagation


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Detailed measurements of sawtooth induced density-pulse propagation have been made on the TEXT tokamak utilizing a high-resolution multichannel interferometer. Particles ejected from the plasma center to the mixing radius \( r_{\text{mix}} \) during the sawtooth crash are observed to be transported to the plasma edge via a diffusive process after the sawtooth crash. This picture is consistent with the heat-pulse analog as measured via soft x-ray (SXR) detector arrays on TEXT. In addition, the density perturbation induced by the sawtooth disruption is of same order as the electron temperature perturbation. Analysis of the “propagation” process, based on a diffusive model, has led to estimates of the electron particle diffusion coefficient, \( D_{e}^{DP} \) (Density Pulse), and the electron thermal diffusivity coefficient, \( \chi_{e}^{HP} \) (Heat Pulse). On TEXT, the magnitude of \( D_{e}^{DP} \) is found to be comparable to \( \chi_{e}^{HP} \), suggesting that heat and particle transport observed after the sawtooth crash might be coupled. In this paper, measurements of \( D_{e}^{DP} \) and \( \chi_{e}^{HP} \) obtained from the same set of ohmic, gas-fueled TEXT discharges will be presented and compared in order to investigate the coupling of the heat and particle transport.

Figure 1 shows a comparison between the raw sawtooth signals from the multichannel interferometer and the SXR diode array. In both cases, signals outside the inversion radius, \( r_{q=1} \), clearly exhibit a heat or density pulse time delay as \( r \) increases. Note the inversion radii of density (\( \sim \pm 9 \text{ cm} \)) and SXR sawteeth (\( \sim 13 \) and \( \sim 7 \text{ cm} \)) are different. This can be partially, but not completely, explained in terms of chord-averaging effects of the interferometer. As can be seen from the traces in Fig. 1, the speed of the “propagation” has some fluctuation from sawtooth to sawtooth, for both measurements. Therefore, a typical sawtooth representing the steady-state region of the discharge is generated through a coherent signal averaging procedure. In addition, an inversion process is applied to the chord-averaged density sawtooth in order to obtain the local density perturbation. Fluctuations in the SXR signal are treated as being proportional to the local electron temperature. However, strictly speaking, it also is a chord-averaged measurement in addition to being a complicated function of density and impurity concentrations as well as electron temperature. In Figure 2, the magnitudes of the local density and temperature perturbations at \( r = r_{\text{mix}} \) are plotted against each other for various discharge conditions. It is important to observe that the temperature perturbation induced by the sawtooth disruption is accompanied by a density perturbation of comparable magnitude, although the rate of increase is larger for the density perturbation. At sufficiently low plasma densities, the sawtooth density perturbation becomes small and comparable to the inherent detection phase noise level while SXR sawteeth are clearly observed. Note, however, that the magnitudes of density perturbation in discharges with a low \( q_{a} (\sim 2.0) \) and/or \( n_{e0} \) close to the density limit deviate significantly from this trend. Typically, \( \Delta n_{e} = 1 \sim 3\% \) and \( \Delta T_{e} = 3 \sim 6\% \) at the mixing radius for discharges discussed in this paper.
A comparison between the typical time-of-flight characteristics of the heat and
the local density pulse is made in Figure 3(a), for the TEXT discharge; $B_i = 20 kG$,
$I_p = 200 kA$, $q_a \sim 3.5$, $\bar{n}_{ce} = 3 \cdot 10^13 cm^{-3}$. The data are quite similar, suggesting that
heat and particles are transported together at nearly the same rate within error bars.
The magnitudes of $D_{e}^{DFP}$ and $\chi_c^{HP}$ deduced from the slopes are 1.5 and 2.0 $m^2/sec$,
respectively. These values are both larger by a factor of 2 to 3 than $D_{e}^{OGP}$ deduced
from the oscillating gas puffing experiments$^4$ and $\chi_c^{PB}$ deduced from the steady-state
power balance calculations$^5$. The OGP results are thought to be indicative of the
steady-state due to their slow time scale ($f^{-1}_{mag} \geq \tau_E$) and edge fueling characteristics.
As $q_a$ decreases, the $t_p \sim n^2$ relationship for density and temperature pulses gradually
deviate. Figure 3(b) compares two discharges where $q_a$ is changed from $\sim 2.5$ to
$\sim 2.0$ by increasing $I_p$ from 200kA to 250kA. The density and temperature data differ
significantly from each other for $q_a \sim 2.0$, whereas the slopes for $q_a \sim 2.5$ are basically
the same. This indicates that the heat and density pulse no longer diffuse out at the
same rate when $q_a$ is as low as $\sim 2.0$.

To further investigate the heat and particle transport after the sawtooth crash, the
scaling behaviour of $D_{e}^{DFP}$ and $\chi_c^{HP}$ with discharge parameters are studied for the same
set of ohmic, gas-fueled TEXT discharges using the same observation time window.
In Figure 4, the measurement values of $D_{e}^{DFP}$ and $\chi_c^{HP}$ are plotted against $q_a^{-1}$
for the discharges surveyed. Both coefficients decrease with increasing $q_a$, however, the
magnitude of $\chi_c^{HP}$ is not always 1.5 times that of $D_{e}^{DFP}$, which is expected for strongly
coupled transport. Note that the $\bar{n}_{ce}$ scaling is not separated out in this plot. To see
the effects of $q_a$ and $\bar{n}_{ce}$ more clearly, a regression analysis was performed yielding the
scalings as

$$D_{e}^{DFP} \sim q_a^{-2.7} \bar{n}_{ce}^{-1};$$
$$\chi_c^{HP} \sim q_a^{-2.1} \bar{n}_{ce}^{-3}.$$  

The trends of each scaling are qualitatively the same, i.e., both coefficients are much
more sensitive to the magnitude of $q_a$ than $\bar{n}_{ce}$. This suggests that particle and heat
diffusion after the sawtooth crash may be coupled. In contrast, previous steady-state
heat and particle transport measurements on TEXT show a different scaling$^{4,5}$: $\chi_c^{PB} \sim
D_{e}^{OGP} \sim (q_a \bar{n}_{ce})^{-1}$. Here the transport coefficients vary just as sensitively with the
density as with $q_a$.

Quantitatively, however, the differences between the exponents in the scalings of
$D_{e}^{DFP}$ and $\chi_c^{HP}$ indicate that the heat and particle diffusion are not always coupled
to each other. The ratio of the transport coefficients scales as $D_{e}^{DFP}/\chi_c^{HP} \sim q_a^{-0.6} \bar{n}_{ce}^{+0.2}$,
indicating that heat and particle transport are gradually decoupled as $q_a$ is lowered
and/or, to a lesser degree, as $\bar{n}_{ce}$ is increased. This deviation is apparent in the
different slopes in Fig.3. Also, for discharges with a very low $q_a$ or $\bar{n}_{ce}$ close to the
high-density limit, the size of the density sawtooth crash differs significantly from the
correlation observed for other discharges (recall Fig.2). In such discharges, indications
of a non-diffusive process or substantial particle source effects have been observed in
the time development of the sawtooth density perturbation profile$^6$.

A possible explanation for the difference between $\chi_c^{HP}$ and $D_{e}^{DFP}$ is that the
local electron temperature perturbation is not accurately reflected in the SXR sawtooth
data (for reasons mentioned earlier). In addition, the diffusive model used for sawtooth heat pulse propagation neglects the density perturbation when solving the heat equation. This may introduce appreciable errors as the density perturbation is comparable in magnitude to the temperature perturbation. In contrast, the diffusive model employed for sawtooth density pulse propagation is completely decoupled from temperature effects. Finally, it remains to be determined through more detailed analysis whether or not a density perturbation can significantly affect estimates of $\chi_e^{HP}$.

In summary, we observe that the temperature perturbation generated from the sawtooth crash is typically accompanied by a density perturbation of comparable magnitude. Both the density and the temperature perturbations are found to be transported out through the confinement zone of the TEXT tokamak in a manner consistent with a diffusive process. The magnitudes and the scalings of $D_e^{DP}$ and $\chi_e^{HP}$ are found to be similar but distinctively different from their steady-state counterparts. However, as $q_a$ is lowered from $\sim 3.0$ to $\sim 2.0$, gradual separation of the heat and density pulse propagation is observed.

This research is supported by the U.S. Department of Energy, through contract No. DE-AC05-78ET55043.


Figure 1. Raw density(a) and SXR(b) sawtooth signals from a discharge with $B_T = 20kG, I_p = 300kA, \bar{n}_e = 3 \cdot 10^{13} cm^{-3}$. Channel spacing of the interferometer and SXR array is 3cm and 1.3cm, respectively. The plasma beyond $r = -15cm$ is inaccessible to the interferometer due to port constraints.
Figure 2. Correlation of density and SXR crash size. Point A is from a discharge at the density limit; B from a discharge with $q_a \approx 2.0$.

Figure 3. The time when the maximum perturbation is reached, $t_p$, is plotted against the minor radius squared, $r^2$, for heat and density pulses. The slope indicates the rate of heat and particle diffusion. (a) $B_T = 20kG$, $I_p = 200kA$, $\bar{n}_{e0} = 3 \cdot 10^{13} cm^{-3}$; Heat and particle diffusion are approximately equal. (b) As $q_a$ decreases, SXR and density pulse data diverge.

Figure 4. $D_{DP}$ and $\chi_{e HP}$ vs. $q_a^{-1}$. Note a linear fit to the datapoints does not go through the origin, indicating a nonlinear dependence on $q_a^{-1}$.
HEAVY ION BEAM PROBE MEASUREMENTS OF SPACE POTENTIAL AND ELECTROSTATIC FLUCTUATIONS IN TEXT WITH A RESONANT MAGNETIC FIELD


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A heavy ion beam probe (HIBP) has been used to measure the plasma equilibrium space potential $\Phi$ (relative to the vacuum vessel ground), and properties of the normalized fluctuating density component $\tilde{n}/n$. The results were obtained during ohmic heating discharges on TEXT[2], a tokamak with major radius $R = 1m$ and minor radius $a = 0.26m$ defined by a poloidal ring limiter. We report the effects observed on both these parameters when externally applied resonant fields are used to produce a mixture of magnetic islands and stochastic regions[3]. Most of the results were obtained with a plasma current $I_p = 200KA$, a toroidal field $B \phi \approx 2T$, and a line of sight average density $\bar{n}_e \approx 2 \times 10^{19}m^{-3}$.

The resonant fields are produced using discrete poloidal saddle coils placed around and outside the vacuum vessel. For the experiments described here the dominant toroidal mode number $n = 2$, with multiple poloidal mode numbers $m \approx 7$. An example of the calculated field structure is shown in Figure 2b), which is also discussed later with reference to Figure 2a). The extent of the individual magnetic island structures is shown at the toroidal location of the HIBP for a saddle coil current $I_h = 4KA$ (flux surface averaged $<b_\perp/b_\phi> \approx 0.05\%$ at the limiter radius). The figure is drawn for a restricted range of poloidal angle $\theta$ and radius $r$, corresponding to the area which is diagnosed by the HIBP.

The $m$ numbers of the islands are given on the right-hand ordinate. Islands with $m = 4$ through 7 are seen; higher $m$ numbers exist but for clarity are not shown. The islands with $m/n = 7/2$ and 6/2 overlap, producing stochastic regions. Increasing $I_h$ produces further island overlap and larger stochastic regions: reversing $I_h$ interchanges island X and O points. Profile measurements show $T_e$ and $|\nabla T_e|$ are reduced in the region affected by the perturbing fields ($r > 0.7$), but $T_i$ and $n_e$ are reduced by $< 20\%$.

The equilibrium plasma potential $\Phi$, obtained without the resonant fields, is shown by the $\diamond$ symbols in Figure 1a) as a function of flux surface coordinate $\rho$ ($\rho = 1$ corresponds to the limiter position, $\rho = 0$ the plasma axis). The results were obtained over a range of poloidal angles $\Theta$. Uncertainties in the absolute values of $\Phi$, a result of errors in the heavy ion beam energy calibration, are removed by normalizing the HIBP results to the results obtained from Langmuir probes with $\rho > 1$, where $\Phi > 0$. For $\rho < 0.95$, $\Phi < 0$, and the radial electric field $E_r \approx -4000Vm^{-1}$. Such a potential well has been previously
observed[4], and shown to be compatible with the ion momentum balance equation:

\[ Z_i \left[ E_r + v_\theta B_\phi - v_B B_\theta \right] - \nabla p_i = 0 \]

where subscripts \( \theta, \phi \) refer to poloidal, toroidal coordinate, \( p_i \) is the ion pressure, \( n_i \) the ion density, and \( Z_i \) the ion charge. In Figure 1a) the expected potential from \( \nabla p_i \) (labeled \( \Phi \nabla p \)) and from \( v_\theta B_\phi \) (labeled \( \Phi v^\theta \)) are shown, assuming that \( v_\theta \) is given by the neoclassical expression [5], \( n_i = n_e \), and \( \Phi v^\theta (\rho = 1) = \Phi v^\theta (\rho = 0) = 0 \). As noted previously[4], the contribution from the theoretical \( \Phi v^\theta \) is small, and no toroidal rotation \( (v_\phi = 0) \) is needed to explain the results (i.e., \( \Phi \approx \Phi v^\phi \) for \( \rho > 0.7 \)).

Figure 1b) shows the spatial variation of the change in potential \( \Delta \Phi \) following application of the resonant fields. Two cases are shown: \( I_h = 2KA \) (broken line) and \( I_h = 4KA \) (solid line). With \( I_h > 1.5KA \) (the value where island overlap and stochastic onsets) \( \Delta \Phi \) increases. The largest change to \( E_r \) is at the edge, \( \rho_{\text{crit}}(I_h) < \rho < 1 \), where \( \rho_{\text{crit}}(I_h) \approx 0.9 \) corresponds approximately to the position within which there is no island overlap (i.e., there is no island overlap for \( \rho < \rho_{\text{crit}} \)). With \( I_h = 4KA \) Figure 1a), symbols, shows the resulting value \( \Phi \approx 150V \) for \( \rho > 0.6 \). Such a profile would be consistent with the ion momentum balance with \( v_\phi \approx 0, v_\theta \approx 0 \) only if the ion pressure gradient is flattened for \( \rho > 0.75 \). This is inconsistent with the limited \( T_i(r) \) and \( n_e(r) \) profile results. The region where \( E_r > 0 \), normally restricted to \( \rho > 0.95 \), is expanded to \( \rho > 0.85 \), i.e. in the stochastic layer and behind the limiter. No evidence of structure associated with the presence of the computed coherent magnetic islands has been found.

The HIBP is also used to measure properties of the density fluctuation levels. In particular \( \tilde{n}/n(r, \theta, \omega) \), with frequency \( \omega/2\pi \) between 0 and 500 KHz, has been found with and without the resonant fields applied. Without resonant fields \( \tilde{n}/n \) (integrated over \( \omega \) and \( k \)) is poloidally uniform to within 10% for \( \theta \) between 0 and 65° (0° corresponds to the outer equator, 90° to the top), and decreases from \( \tilde{n}/n \approx 0.16 \) at \( \rho = 1 \) to \( \approx 0.01 \) at \( \rho = 0 \). Figure 2a) shows contours of the fractional change \( \Delta(\tilde{n}/n)/(\tilde{n}_0/n_0) \) with \( I_h = 4KA \), with subscript \( o \) referring to values with no perturbing field, \( I_h = 0KA \). Figure 2b) shows the computed extent of individual magnetic islands at the location of the measurements of \( \Delta(\tilde{n}/n)/(\tilde{n}_0/n_0) \). Poloidal asymmetries are seen, with increases in \( \tilde{n}/n \) of up to 60% correlating with island X points and decreases of 20% corresponding with island O points. If the electrostatic turbulence driven particle flux is also increased at the X points, this will act to reduce the density gradients in this region. Test particle models of particle transport in the mixed magnetic island and stochastic regions should include this effect.

The phase velocity \( v_{ph} \) between two points poloidally separated by \( \approx 2cm \) has also been found. At \( \rho = 0.7, v_{ph} \approx 3.2x10^3 \text{ms}^{-1} \) in the electron diamagnetic drift direction, and
is unaffected by the resonant fields. However, at $\rho = 0.92$ the resonant fields ($I_h = 4$KA) change $v_{ph}$ from $\approx 3.8 \times 10^3 \text{ms}^{-1}$ in the electron diamagnetic drift direction to $\approx 2.7 \times 10^3 \text{ms}^{-1}$ in the ion drift direction: i.e., in the stochastic edge the propagation direction is reversed. These values are not consistent with either the electron diamagnetic drift velocity $v_{de} < 2 \times 10^3 \text{ms}^{-1}$ (electron drift direction) or the sum of the electron drift and neoclassical poloidal rotation velocity ($v_{de} + v_{\phi \text{neo}}$) $< 2.5 \times 10^3 \text{ms}^{-1}$ (electron drift direction).

In conclusion, the presence of externally applied resonant magnetic fields changes $E_r$ in the stochastic edge region. Here the plasma potential changes from $\Phi < 0$ to $\Phi > 0$. While the results without perturbing fields are consistent with the small neoclassical poloidal velocity and no toroidal rotation, this is not the case with the perturbing fields. The density fluctuation level $\tilde{n}/n$ is increased near magnetic island X points and decreased near island O points. The phase velocity of the fluctuations in the poloidal direction is affected by the presence of a stochastic layer, changing from the normal electron diamagnetic to ion drift direction. The phase velocities are not consistent with the assumption of neoclassical poloidal rotation and no toroidal rotation.

* Work performed under contract for the Department of Energy.

References

Fig 1a) The variation of plasma space potential $\Phi$ with flux surface coordinate $\rho$, * without perturbing field and ■ with perturbing field ($I_h = 4$KA). The potential contributions for the case without perturbing field from $\nabla p_i$ (labeled $\Phi^{\nabla p}$) and from $v_0B_0$ (labeled $\Phi^{v_0B}$) are also shown.

Fig 1b) The change in plasma space potential as a function of flux surface coordinate $\rho$, for $I_h = 2$KA (broken line) and 4KA (solid line).

Fig 2a) Contours of the increase in normalized density fluctuation level, $\Delta(n/n)/<n_0/n_0>$, over a range of radius $r$ and poloidal angle $\theta$. Subscript $0$ refers to values with $I_h = 0$KA.

Fig 2b) The computed extent of the magnetic islands in $(r,\theta)$ space for $I_h = 4$KA. The right hand ordinate gives the $m$ numbers of the islands.
OBSERVATION OF ANOMALOUS ION HEATING IN THE TJ-I TOKAMAK

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INTRODUCTION

The ion temperature is being studied in the TJ-I tokamak by two independent methods, charge exchange (CX) and Doppler spectroscopy (DS), the latter with spatial resolution. The main motivation was initially to support impurity confinement studies in this tokamak (1), where we have found that ion energy and particle confinement times seem to be linked, and deserves a more detailed study. The present work is mainly oriented to study and try to understand the anomalous behaviour of central impurity temperatures as a function of density in the TJ-I tokamak. In addition, renewed interest in experimental ion temperature profiling in tokamaks has been motivated by the finding that experimental and theoretical profiles show significant discrepancies (2), and so invalidating the conventional method used to perform ion power balance in tokamaks.

EXPERIMENTAL

Measurements herein presented have been carried out in the TJ-I tokamak (R = 30 cm, a = 9.5 cm) which has been operated for this experiment with plasma currents around 40 kA, toroidal field between 1 and 1.5 T, and line averaged electron densities between 0.5 and 3 E 13 cm⁻³. The central ion temperature has been measured by a neutral spectrometer which samples the neutral energy distribution along a line of sight forming 30° with the tokamak major radius and contained in the equatorial plane. It is a monochannel apparatus in which parallel plates are swept by a rapid high voltage ramp allowing one to scan the neutral energy range of interest in a few milliseconds. The optical monochromator performs impurity linewidth measurements, using a refractor plate, along any selected plasma chord. Both scans are synchronized at a preprogrammed time of the tokamak discharge using an electronic method described in ref. (3). Doppler measurements were made using two different lines belonging to TJ-I central ions: CV 2271 Å and CVI 5292 Å.
RESULTS AND DISCUSSION

Figure 1 shows Doppler temperatures, obtained with the CV 2271 Å line, and CX temperatures as a function of electron plasma density. A noticeable discrepancy is observed at the lower densities. Whereas the CX temperature increases slightly on increasing density, following basically the Artsimovich law, the Doppler temperature decreases on increasing density.

It should be noted that the anomalous heating effect observed at the lower densities takes place only in impurity ions but not in protons. This result suggests that we are not seeing an anomalous heating mechanism similar to that reported in references (4) and (5). On the other hand, this anomalous effect is observed with the same intensity for lines originating in low and highly excited levels. Consequently, it is difficult to attribute it to anomalous broadening, i.e. a quasi-static Stark effect due to non-thermal electric field fluctuations (6). Part of the broadening could be due to gross mass motion and microturbulence at the centre. This possibility, which could be ruled out if measurements with ions of different charge to mass ratio were available, cannot be completely discarded, but it seems unlikely since it is independent of how the machine is operated.

![Fig. 1. Comparison between central ion temperatures obtained by CX and DS in a TJ-I density scan.](image1)

![Fig. 2. Charge exchange losses that give account of temperature discrepancy.](image2)

An alternative interpretation of the present results can be based on the distinct influence of charge exchange loss processes on protons and high Z impurity ions, in a plasma in which the proton impurity thermalization time is not negligible compared with particle
confinement time. We have studied whether the consequences of this hypothesis are consistent with the other tokamak data. Assuming that the impurity and proton energy confinement times are roughly equal, the ratio of impurity to proton temperature would be given by:

\[
\frac{T_I}{T_P} = \frac{1}{1 - \frac{Q_{\text{CX}}}{Q_{\text{ei}}}}
\]

where \(Q_{\text{CX}}\) and \(Q_{\text{ei}}\) are the power densities, at the plasma centre, for charge exchange and electron-ion heat transfer processes respectively. From this simple formula and using experimental data, the ratio \(Q_{\text{CX}}/Q_{\text{ei}}\) can be estimated. Figure 2 shows the results of this analysis. To judge the correctness of this view, we have determined with a consistent procedure \(Q_{\text{ei}}\) in order to obtain the absolute value of \(Q_{\text{CX}}\), and then to deduce the consistent central neutral densities which would account for this loss term. This simple model gives account of the order of magnitude of the neutral concentration at the centre, from 2 to 5 \(\times\) 10^9 cm^{-3}; but it does not predict its relative electron density dependence. From CX data, the ratio of neutral density at the lowest and highest densities has been estimated to be 10. This discrepancy suggests that in addition to the aforementioned charge exchange contribution, other mechanisms should be considered; e.g., a different convection loss term for impurities and protons.

The most puzzling point to account for is why this effect has not been observed in other plasmas with the same collisionality regime as TJ-I (plateau for protons and Pfirsch-Schlüter for impurities). We believe that the reasons are as follows: a) few comparisons are reported in the literature between central ion temperatures measured by two independent techniques covering a density range like that in this work; b) due to the electron temperature range of TJ-I plasmas (typically 300 eV) and its small radius, we can detect, in the visible and near UV range, lines belonging to central ions which make this type of study possible; c) a practical point that has proved to be important in TJ-I to unambiguously detect this effect, has been to take data along several chords around the expected position of the magnetic axis. In TJ-I, like in many others, we do not have a fine tuning of the vertical position, so in practical tokamak operation the ion temperature peaking position has been the sole sensitive diagnostic for small, but relevant for these measurements, vertical plasma displacements. In Fig. 3, we present two ion temperature profiles obtained by spatially scanning the linewidth
of the 5292 Å line of CVI. The central hotter line can be empirically controlled by adjusting, using a trial and error procedure, the radial field.

**Fig. 3.** Ion temperature profiles (CVI 5292 Å) for two different settings of the radial field, 8 on left and 12 on right (a. u.).

In conclusion, an anomalous central impurity ion temperature behaviour has been observed in the TJ-I tokamak. The different influence of charge exchange and convection losses for protons and impurity ions have been proposed to explain the effect.

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INTRODUCTION

The role of turbulence in tokamak confinement properties is an important topic that concentrates many efforts, both in the theoretical and experimental area, trying to correlate fluctuations in different plasma magnitudes with plasma confinement properties (1, and reference there in).

In this paper an attempt to correlate the particle confinement time value and electrostatic fluctuations in the TJ-I tokamak \((R_0 = .3 m, a = .1 m, B_T \leq 1.5 T, I_p \leq 75 kA)\) is presented. By measuring density and floating plasma potential fluctuations, with Langmuir probes at the plasma edge, and deducing from these measurements the total flux of particles leaving the plasma due to electrostatic fluctuations, we try to infer if these losses are enough to explain the particle confinement properties of the device or other fluxes, probably due to electromagnetic turbulence, must be invoked. A microwave reflectometer has been used to determine the electron density and its fluctuations with radial resolution as additional information and as verification for the probe measurements.

PROBE MEASUREMENTS

Density and floating potential fluctuations have been studied by means of a square array (2 x 2 mm.) of four Langmuir probes. The array can be introduced in the discharge up to 1 cm outside the chamber wall, that in TJ-I acts as toroidal limiter along the equatorial midplane. Position of the array is 2.5 cm above this plane, in the outer part of the torus. Two probes were biased 65 V into the ion saturation regime to determine density fluctuations. The remaining two probes were unbiased in order to measure the floating potential at two different poloidal positions.

The electronic system used in these measurements has a 300 KHz bandwidth. Data acquisition was accomplished by using a four channel camac module, (LC8210), with 8 KB memory per channel. Data were digitized at a 500 KHz sampling rate. Conventional FFT was used for data analysis.

Traces for plasma current intensity, \(I_p\) and Langmuir probes, working in saturation and without bias, for a typical TJ-I discharge are shown at Fig. 1.
Figure 2 (a,b) shows power spectra for density and potential signals obtained during 4 ms of the discharge, 6 ms after discharge initiation. They have been determined neglecting electron temperature fluctuations. Data were analysed in 1 ms time slices and averaged for the whole period.

Both spectra are basically constant up to 30 kHz and decrease for higher frequencies with an \( \omega \) dependence given by:

\[
P(\omega) \approx \omega^{-\alpha} \quad (\alpha = 1.5 \pm 0.5)
\]

value similar to those obtained in other tokamaks.

Coherence between density and potential fluctuations is presented at Fig.2c. There is a moderate coherence for frequencies below 30 KHz, and very low at higher ones. Density in the surroundings of the probes was calculated from the saturation current and it is in the order of 0.2 - 0.5 \( \times 10^{12} \) cm\(^{-3} \).

An estimate of the particle confinement time due to electrostatic fluctuations can be made from the total particle flux \( \Gamma \), given by (2):

\[
d\Gamma(\omega) = \frac{\kappa(\omega)}{B} \cdot |\gamma_n(\omega)| \cdot \sin \theta_n(\omega) \cdot \vec{n} \cdot \vec{\psi} \cdot d\omega
\]
where $\kappa(\omega)$, $\gamma_{n\phi}(\omega)$, $\theta_{n\phi}(\omega)$ are the average wavenumber, coherence and phase angle between density and potential, respectively. $\bar{\eta}$ and $\bar{\Phi}$ are the fluctuation r.m.s. values.

Fig. 3 shows the spectrally resolved particle flux. Integrated for all the spectrum, a total flux value of $\Gamma \approx 6 \times 10^{15} \text{cm}^{-2} \text{s}^{-1}$ is determined from the probe data, assuming that the flux is poloidally and toroidally symmetric.

The particle confinement time is given by $\tau_p = n_e a / 2 \Gamma$ where $n_e$ is the chord-average density ($= 2 \times 10^{12} \text{cm}^{-3}$), deduced from reflectometry measurements, and $a$ is the plasma minor radius ($a = 9.5 \text{cm}$). The obtained particle confinement time ($\tau_p = 3 \text{ ms}$) is, taking into account the uncertainties for this estimation, in agreement with a previous spectroscopic estimate of $\tau_p$, in the order of $1 \text{ ms}$ (3).

**REFLECTOMETRY MEASUREMENTS.**

A swept microwave reflectometer working in the band 30-55 GHz has been used in these experiments to measure the radial profile of the electron density (4).

By keeping the frequency fixed during the discharge, density fluctuations could be detected at the position of the cutoff layer for that particular frequency. Changing the frequency, in a repetitive shot-to-shot basis, the radial profile of those fluctuations was deduced. Mode X propagation was selected due to its higher sensitivity to fluctuations.

Fig. 4 shows the radial profile deduced from the reflectometer data for the typical discharge used in these experiments.

The radial dependence of the density fluctuations has been studied changing the reflectometer frequency in successive repetitive discharges. Results agree with probe findings.
about the decay of spectra with $\omega^{-\alpha}$ ($\alpha \approx 1.5$) and its broadening when going deeper in the plasma.

Also these measurements suggest, that the level of density fluctuations is not only important at the plasma edge but inside the discharge, probably due to the presence of a high MHD level of fluctuations, as it is shown at the figure 5 where the amplitude, in the reflectometer signal spectra, at the frequency of the MHD activity, is plotted. A dominant component with the maximum at half the plasma radius appears for that frequency.

**FIG. 5**

**SUMMARY**

Density fluctuations have been measured in the TJ-I tokamak by Langmuir probes and microwave reflectometry. Consistency in the experimental observations have been obtained: frequency spectra broaden when going inside the discharge and high frequency dependence can be fitted by $P(\omega) = \omega^{-\alpha}$ ($\alpha = 1.5 \pm 0.5$).

Combining these measurements with floating potential fluctuations an estimation of the particle confinement time due to electrostatic turbulence was made. Its value could suggest that particle confinement in TJ-I is mainly determined by this kind of turbulence

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PLASMA DISRUPTIONS IN TOKAMAK TBR-1


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Abstract: Disruptive instabilities were investigated by using soft x-ray and pick-up coils diagnostics. Sawteeth parameters were obtained. Major disruptions were caused by m/n=7/1 and 1/1 modes coupling while an ergodization region between q=2 and q=3 islands would develop a minor disruption.

Internal (sawteeth) and external (minor and major) disruptions were identified in the discharges of the small tokamak TBR-1 (main parameters: R = 0.30 m, a = 0.03 m, Bφ = 4.0 T and I = 10 kA, with discharges duration ranging up to 7 ms). Two basic diagnostics were used: one composed by 20 magnetic pick-up coils (16 equally spaced in the poloidal direction and 4 in the toroidal direction) and another one to detect the soft x-ray emission from the plasma. The soft x-ray system has six surface barrier detectors (ORTEC CR 19-50-100) placed inside an imaging camera, each one of them viewing a chord of the plasma column through a polypropylene covered slot.

Typical observed sawteeth oscillations are shown in fig. 1-b. Their period is τ = 130 µs and no transients are verified in the loop voltage (fig. 1-c) during the sawteeth crashes. They show the expected m=1 oscillations superimposed to them.

The predicted time for the sawteeth crashes, accordingly to the Kadomtsev's model (2), would be τc = 18 µs. The experimental measurements at the TBR-1 agree quite well with this value except for high Z eff discharges, when the crash time almost triple. In this latter case the sawteeth falls are smooth showing up the presence of m=1/n=1 oscillations in all its way down (fig. 2-c), rather than a very sharp decay as it is usually observed (figs. 2-a and 2-b).

The 1/1 mode does not always increase in amplitude, while the sawtooth is rising up. Instead, it is usually observed the maintenance of the amplitude level and in some cases even a decrease is found out (fig. 2).

Comparisons of the sawteeth periods with those predicted by scaling laws were also done. Estimations obtained by using the expressions from refs (3) and (4) give τ = 87 µs and τ = 210 µs, respectively, while the experimental values ranged from 100 µs to 170 µs.

Another experimental feature is related to the observation of a sawtooth relaxation, occurring more to the end of the discharge, given place to a strong 1/1 oscillation. Future investigations shall be carried out to verify if this behaviour is really caused by current profile modifications leading to q(0) > 1.

The major disruptions at the TBR-1 discharges, with 3.05s² < 4.0, are characterized by a very noticeable growing precursor mhd activity, a negative spike in the loop voltage and a small displacement of the plasma column towards the inner part of the torus. The typical duration of the disruptive
event is $T_D = 200 \mu$s. The perturbed poloidal field ratio, right before the occurrence of the negative spike, is $B_0 / B_0 = 4\%$ and the precursor oscillation growth rate ranges from $1.4 \times 10^3 s^{-1}$ to $8.9 \times 10^4 s^{-1}$. The Fourier analysis of the poloidal magnetic pick-up coils signals showed that 2/1 is always the dominant and responsible mode for the analysed disruptive events. A precursor 3/2 mode predicted in some theories (5) has not been observed before the disruptions. Calculations made about the $q=2/1$ resonant magnetic surface localization and the half-width of its corresponding islands (6) showed that there is no possibility of their interaction with the limiter. On the other hand, the 2/1 and 1/1 mode coupling could be a realistic mechanism through which the disruptive instability would initiate. This is so because the frequencies of these modes change and, prior to the negative spike, the 1/1 oscillations in the soft x-ray signals are often observed to have the same frequency of the 2/1 mode, as it is shown in fig. 3.

The TBR-1 device has not yet a feed-back plasma position controlling system to prevent the column to wander around the toroidal geometric center. Whenever the best value of the vertical field is not achieved, it is very usual to obtain short discharges like the one presented at fig. 4, where a sequence of minor disruptions is observed. If the right field conditions are settled, longer discharges (4–7 ms) are obtained and only two or three minor disruptions are observed during or right after the plasma current rise. Fourier analysis indicates that the dominant 2/1 or 3/1 modes, separately, are the main precursor modes associated with this kind of disruption. Their growth rate is $5 \times 10^3 s^{-1}$ with $B_0 / B_0 = 2\%$, typically.

Sometimes a precursor 1/1 oscillation can be detected by soft x-ray diagnostic system but it is much more frequent to observe the complete absence of such mode prior to a minor disruption. Therefore, the 2/1 and 1/1 mode coupling mechanism can not be satisfactorily accounted to be responsible for these disruptions.

Again, calculating the location and half-width of the $q=2$ and $q=3$ islands it was shown that a direct interaction with the limiter should be discarded. In this way, the most probable mechanism associated with the minor disruptions concerns an interaction between 2/1 and 3/1 modes through an ergodization of the magnetic field lines, as proposed elsewhere (7, 8).

In TBR-1 the 1/1 mode amplitude does not always increase during a sawtooth period and, for high $Z_{eff}$ discharges, this mode is observed during the sawtooth crashes. Major disruptions were preceded by 1/1 and 2/1 modes coupling while minor disruptions seems to be related to the magnetic surfaces break-up. *Partially supported by CNPq.

References:
Fig. 1: a) Plasma current evolution, b) typical sawteeth oscillations (central channel) and c) loop voltage.

Fig. 2: Examples of sawtooth activity in three different discharges. The 1/1 oscillation amplitude does not always increase during a sawtooth period. The crashes are smoother for discharges with high $Z_{\text{eff}}(c)$. 
Fig. 3: 1/1 soft x-ray (a) and a dominant 2/1 mhd oscillations prior to a major disruption indicated by a negative spike in the loop voltage (c).

Fig. 4: a) Plasma current evolution, b) soft x-ray emission and c) loop voltage during a discharge in which successive minor disruptions occurred.
THE POWER DEPENDENCE OF $\tau_E$ IN THE H-MODE OF ASDEX

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INTRODUCTION: The global energy confinement time $\tau_E$ of H-mode plasmas is found in ASDEX to be power independent for hydrogen injection into deuterium plasmas /1/. DIII-D reports the same behaviour for deuterium injection into deuterium plasmas /2/. These findings are contrary to those from JET /3/ and JFT-2M /4/, where degradation of $\tau_E$ in the H-mode is observed as in the L-mode. This causes great uncertainty in predicting H-mode confinement of next-generation tokamak experiments. In this paper we try to analyze the power dependence of $\tau_E$ in the quiescent H-mode, $H^*$, in comparison with regular H-mode results. The $H^*$-mode may display the intrinsic H-mode confinement properties because ELMs do not additionally contribute to the energy losses. The power scaling studies of JET were done without ELMs, but the published results of ASDEX and DIII-D were obtained with ELMs.

POWER SCALING OF $\tau_E$ IN THE $H^*$-MODE: Under the conditions $I_p=0.32 \text{MA}$, $B_t=2.2 \text{T}$ power scans in the quiescent H-mode were possible in ASDEX both for $H^0$ and $D^0$ injection into deuterium plasmas in the restricted power range of $1.75 \leq P_{NI} \leq 3.5 \text{ MW}$. Without ELMs an additional difficulty in the $\tau_E$ analysis arises because the $H^*$-phase does not reach the steady state. During the beam pulse the $H^*$-phase is terminated by a thermal quench caused by large impurity radiation which finally matches the power input (globally as well as locally in the plasma centre) and which initiates an intermediate L-phase /5/. The duration of the $H^*$-phase depends on the heating power and increases from 95ms at 1.75MW to 125ms at 3.5MW. During the $H^*$-phase the particle content increases linearly at a rate corresponding to 2-3 times the beam fuelling. The impurity radiation increases nearly exponentially. After the $H^*$-transition, $\beta_p$ increases once more owing to the improved confinement; when the impurity radiation starts to affect the energy balance, $\beta_p$ rolls over and decreases already during the $H^*$-phase.

Because of the non-steady-state conditions and the central radiation issue, the $\tau_E$ analysis was done in three steps: $\tau_E$ is evaluated according to $\tau_E^{(1)} = \frac{E}{(P_{\text{tot}} - \text{d}E/\text{d}t)}$ to correct for the lack of stationarity. These values allow the comparison with those from other experiments. In a second step the following relation is used:

$$\tau_E^{(2)} = \tau_E^{(1)} \cdot \left[ 1 - \frac{\int r \int p_{\text{rad}}(\rho) \rho \text{d}\rho \text{d}r}{\int r \int p_{\text{heat}}(\rho) \rho \text{d}\rho \text{d}r} \right]^{-1} \quad (1)$$
Equ. (1) corrects for the radiation power emitted in the radial zone where the radiation - and heating power densities overlap. Mere edge radiation does not lead to a correction. The $\tau_E^{(2)}$ values should represent the actual transport properties of the quiescent H*-mode. There is a certain arbitrariness in the choice of $\tau_E^{(2)}$ as it is strongly varying with time. In a final step, full transport analysis (using TRANSPI) is done for two H*-discharges, the one with the lowest and the one with the highest heating power, to have a further check on the data and the analysis.

Figure 1 shows the power dependence of the global confinement times $\tau_E^{(1)}$ and $\tau_E^{(2)}$ in the H*-mode, evaluated as described above in comparison with $\tau_E^{(3)}$ values of the regular H-mode with ELMs. As H-phases with ELMs reach steady state, $\tau_E$ is calculated in this case simply from $\tau_E = E / P_{tot}$. While the H-mode does not show any distinct power dependence ($\tau_E^{(4)}$ (ms) = 52.2 $\cdot$ P (MW)$^{-0.05}$) with $H^0$ into $D^+$, the quiescent H-mode does. The following power dependence is obtained for the H*-phase (based, however, on 4 data points only): $\tau_E^{(1)}$ = 86 $\cdot$ P$^{-0.45}$ for $H^0$ into $D^+$ and $\tau_E^{(1)}$ = 132 $\cdot$ P$^{-0.75}$ for $D^0$ into $D^+$. The curve in Fig. 1 is the power fit to the $D^0$ into $D^+$ cases; the horizontal line guides the eye to the results obtained with ELMs.

Figure 2 shows the time dependence of $\tau_E$ as obtained from the full transport analysis and of $\tau_E^{(1)}$. The confinement times are plotted from the OH-phase into the L- and finally into the H*-phase. The results are shown for the extreme power cases. The $\tau_E$ values in Fig. 2 are obtained taking into account the radiation losses and correspond to $\tau_E^{(2)}$ as obtained from equ. (1) / 5 /.

In summary we also find in ASDEX a power dependence of $\tau_E$ in the quiescent H*-mode as in JET. As such a dependence is not observed in the regular H-mode, it can be speculated that $\tau_E$ in this case is predominantly determined by the energy losses caused by ELMs superimposed on the heat transport losses so that the overall confinement time is power-independent. We shall try to analyze this possibility in the following.

THE EFFECT OF ELMs ON THE GLOBAL CONFINEMENT: As ELMs are an external mode, the global confinement is affected. Figure 3 shows a discharge where quiescent phases follow those with ELMs. As soon as ELMs set in, the particle and energy contents decrease. It is difficult to assess the energy losses per ELM. From the changes in slope of the $\Delta p$ trace in Fig. 3 we conclude that $\tau_E^{\text{ELM}}$ (representing the ELM losses) is about 110 ms.

The energy lost at an ELM can be determined in three ways: $\Delta p$ can be measured directly by using the equilibrium coils, which are placed within the vacuum vessel and which have sufficient time response. Typically 5% of the energy content is lost by an ELM. Another possibility is to determine the energy loss via the effect of an ELM on the plasma profiles. For $n_e$ and $T_e$ continuous measurement is possible. An ELM affects both $n_e$ and $T_e$ from the plasma edge to about $r = 15$ cm. The relative amplitude increases with radius for $n_e$ and $T_e$ in about the same way with $\Delta n_e / n_e \approx \Delta T_e / T_e = 10$% at $r = 30$ cm. On the assumption that the $T_i$ profile is affected in the same way by an ELM, we can integrate the energy loss and conclude that about 8% of the energy is lost. A further possibility to evaluate the energy loss of an ELM is by measuring the power flux into the divertor chamber onto the target plate. This measurement is rather inaccurate and can only serve as a consistency check. Typically an ELM leads to power deposition at the target plate with a peak power density of about 1 kW/cm² and a width of about 4 cm. The duration of increased energy loss during an ELM is approximately 0.4 ms; the energy loss is assessed to
about 10 kJ in rough agreement with the other estimates. The particle loss per ELM is typically $1 \cdot 10^{19}$ corresponding to $\Delta N_e / N_e = 5\%$

The repetition time of ELMs depends somewhat on the heating power. At low power ELMs appear erratically and the ELM period fluctuates. At high power, ELMs appear in a more regular form and the repetition time becomes constant at about 6 ms. It was not possible to determine the amplitude distribution of the ELMs. The low-power, high-frequency ELMs are clearly smaller in amplitude. Assuming a constant ELM amplitude of $\Delta E / E$ of 6\% and a repetition time $t_{ELM} = 6$ ms and superimposing these losses onto the transport losses (see Fig.1) we can calculate the power dependence of the global confinement time: $\tau_E^{-1} = \tau_E^{(2)} - 1 + \Delta E / E \cdot t_{ELM}^{-1}$. The results calculated in this way are also plotted in Fig. 1. They roughly agree with the measured $\tau_E$ data obtained with ELMs.

SUMMARY: First the global energy confinement time is shown to decrease with power when the H-mode is operated in the quiescent H-phase as in JET. This conclusion is drawn, however, from a rather weak experimental basis because of our limitations in operating the H-mode without ELMs. It is shown that the lack of power dependence in the regular H-mode with ELMs could be due to the power-dependent microscopic transport losses being superimposed on those caused by ELMs.

The ratio between $\tau_E$ in the quiescent H-mode of ASDEX and JET (both with $D^0 \rightarrow D^+$) is about 20, which is a factor of 2 more than the ratio of the currents, indicating an additional size scaling given approximately by the major radius.

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FIGURE CAPTIONS:

FIG.1 Global energy confinement time in the quiescent ($H^*$) and the regular ($H$) H-mode with ELMs versus heating power. The curve is a power fit to the $D^0 \rightarrow D^+$ case, the horizontal line guides the eye to the $H$-cases with $H^0 \rightarrow D^+$. The circles are calculated from the $H^0 \rightarrow D^+$ $\tau_E$ values with the radiation correction.

FIG.2 Energy confinement time ($\tau_E$) and replacement time ($\tau_p^{(2)}$) of two $H^*$ discharges with different heating powers, as obtained from TRANSP.

FIG.3 Line - density $n_e$, divertor $H_\alpha$ radiation and $\beta_{pol}$ for a discharge with intermittent quiescent and ELM - active phases.
Fig. 1

\[ \tau_E (\text{ms}) \]

\[ P_{\text{TOT}} (\text{MW}) \]

- \( \times \) \( H^0 \rightarrow D^+ (H) \)
- \( \square \) \( D^0 \rightarrow D^+ (H^+) \)
- \( \diamond \) \( H^0 \rightarrow D^+ (H^+) \)
- \( \blacksquare \) \( H^0 \rightarrow D^+ (H^+) \) radiation corrected
- \( \bigcirc \) calculated from (■) with ELMs simulated

Fig. 2

\[ \tau_E (s) \]

\[ P_{\text{NI}} = 1.77 \text{ MW} \]

Fig. 3

\[ H_\alpha (\text{a.u.}) \cdot \bar{n}_e (10^{13} \text{ cm}^{-3}) \]

\[ \beta_p \]
TRANSPORT ANALYSIS OF THE L-TO-H TRANSITION IN ASDEX BY COMPUTER SIMULATION

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ABSTRACT: The transport properties and ideal ballooning stability during the L-phase and a burst-free period at the beginning of the H-phase are explored by computer modelling. It is found that the diffusivities $\chi_e$ and $D$ are reduced in the steep gradient zone by a typical factor of six a few ms after the H-transition. Local transport in the inner plasma improves at an early stage by a factor of about two. Prior to and after the H-transition both electrons and ions are in the banana regime. Ideal ballooning modes are shown to be stable everywhere, including the edge zone.

INTRODUCTION: The first H-regime studies /1, 2, 3/ already revealed that the L-to-H transition starts at the plasma edge and that processes near the separatrix are of crucial importance. At present, the roles of the divertor chamber, scrape-off layer and steep gradient zone in the H-mode transition are still poorly understood. In order to get more insight, simulations of ASDEX discharges are carried out by modified versions of the BALDUR predictive transport code /4, 5/. The code is appropriate for analysing edge processes since it includes a scrape-off layer model and runs with a non-equidistant radial grid capable of resolving the very steep density and temperature gradients close to the separatrix ($r_s = 40$ cm). Anomalous energy and particle fluxes are modelled by local (flux-surface-averaged), empirical electron heat diffusivities $\chi_e$, diffusion coefficients $D$ and inward drift velocities $v_{in}$. The conditions and processes prior to and at the H-transition are studied, with special attention being paid to the plasma periphery. Computed time developments of electron density and temperature profiles and beta poloidal are compared with measured results. The search for critical parameters which might trigger the H-mode /6/ is being continued. Self-consistent modelling of the processes in the divertor and scrape-off plasmas /7/ is not attempted, since sufficient reliable information is lacking.

TRANSPORT BEHAVIOUR OF THE INTERIOR AND EDGE PLASMA: An H-discharge with line-averaged density $n_e = 3.3 \times 10^{13}$ cm$^{-3}$, plasma current $I_p = 320$ kA, toroidal magnetic field $B_t = 1.85$ T and neutral injection power $P_{NI} = 3.45$ MW (H$^0$→D$^+$) is analysed (see Fig. 1). The discharge shows a relatively late L-to-H transition at $t^* = 1.229$ s succeeded by a burst-free period of 35 ms. Both $n_e$ and $\beta_p$ ($\beta_p$) almost reach a saturated L-state and quickly rise during the quiescent H-phase. The time developments of electron temperature and density profiles measured by periodic multichannel Thomson scattering (circles) are presented in Figs. 2 and 3. The solid curves are modelled with the $\chi_e(r)$ given in Fig. 4, $D(r) = 0.5 \chi_e(r)$ and $v_{in} = 233$ cm s$^{-1}$.
The results presented in Fig. 4 were obtained with the Spitzer resistivity, which yields better agreement with the measured loop voltage and the expected q value on axis than the neoclassical resistivity. In the L-phase and the quiescent H-phase, the plasma is everywhere stable to ideal ballooning modes even in the steep gradient zone. At 1.26 s, i.e. a few ms prior to the first burst, the stability limit is not reached at the edge if the 25% higher \((\partial p/\partial r)_c\) due to small aspect ratio is taken into account.

Earlier analyses /13, 14/ are confirmed by the present result that the pressure gradient in the pre-transition L-phase is everywhere smaller than the critical value for ideal ballooning modes. This finding is incompatible with the model assumption made for the L-phase in Ref. /15/.

The profile parameter \(\eta_\text{e} = d \ln T_\text{e}/d \ln n_\text{e}\), which is important for drift instabilities, does not change much in the L- and H-phases (see Fig. 4). The modest variation even at the edge results from raising both the density and temperature gradients in the zone A. In the initial phase of the H-transition particle flow blocking at the separatrix should yield transiently flat density profiles. The corresponding higher \(\eta_\text{e}\) or \(\eta_\text{f}\) values possibly trigger the reduction of the diffusivities in the steep gradient zone.

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with wall radius \( r_W = 49 \text{ cm} \). The density and temperature gradients near the separatrix were determined on ASDEX on a shot-to-shot basis by an edge Thomson scattering device /8/. The steep gradient zone \( \Lambda \) extends over about 2.5 cm and persists throughout the H-phase. A few ms after \( t^* \) the electron temperature measured by the ECE diagnostics begins to rise, indicating reduced diffusivities in the edge zone at this very early stage. In the simulations reduced values of \( x_e \) and \( D \) in the zone \( \Lambda \) are thus applied for \( t \geq t^* \). It is found that the electron heat diffusivity has to be diminished by a factor of about six (see Fig. 4) in order to obtain the measured temperature pedestal.

The scrape-off region is modelled by classical \( x_e^0 \), subsonic flow

Mach number \( M = v_m/V_s = 0.11 \) with \( V_s = [2(T_e + T_i)/m_i]^{1/2} \) and cross-field diffusivities \( x_e \text{ SOL} = 1.2 \times 10^{-2} \text{ cm}^2 \text{s}^{-1} \) and \( D \text{ SOL} = 2 \times 10^3 \text{ cm}^2 \text{s}^{-1} \), yielding the measured decay lengths \( \lambda_T = 1 \text{ cm} \) and \( \lambda_N = 1.5 \text{ cm} \).

Figure 4 shows that \( x_e^0 \) in the inner plasma is reduced by a typical factor of two in relation to \( x_e^0 \), which agrees with earlier transport analyses /9, 10/. The results of Fig. 2 are obtained with \( x_e = x_e^0 \) for \( t \geq t^* \). Obviously, the modelling fits the measured \( T_e \) profiles. In addition, good fits to the measured \( D \) and \( x_e \) temperature profiles are obtained with ion heat diffusivities three times the neoclassical values /11/. The conclusion that both \( x_e \) and \( D \) in the zone \( \Lambda \) drop immediately after the H-transition agrees with results on DIII-D, where the density and temperature gradients inside the separatrix were found to rise steeply within a few ms after \( t^* \).

It is obvious from Fig. 3, however, that the transport model fails to simulate the measured high edge density and the density shoulder in the initial burst-free H-phase. Extensive studies have shown that a further reduction of \( D \) in the zone \( \Lambda \) raises the density gradient. A higher edge density is only obtained by strongly reducing the particle flow across the separatrix. This process seems to be crucial for the L-to-H transition.

The electron collisionality factor \( v_{ei} \) (see Fig. 4) and the ion collisionality factor \( (v_{ij} < 0.6) \) do not change a lot in the L- and H-phases. The plasma at the periphery does not proceed to a different collisionality regime. At the edge, both electrons and ions are in the banana regime \( (v_{ei}, j 1) \) already in the L-phase. It is concluded that a different collisionality regime is not reached, which contradicts the assumption in Ref. /12/ concerning the ion energy transport mechanism across the separatrix.

Ideal ballooning stability is studied by the transport code, which evaluates the local criterion for large aspect ratio

\[
-\frac{2R_0 q^2}{B_1^2} \frac{\partial p}{\partial \eta} = f(s)
\]

where \( (\partial p/\partial \eta) \) is the critical pressure gradient and \( f(s) \) is a known function of the dimensionless shear \( s = (r/q) d q/dr \). At finite aspect ratio (\( A = 4.2 \)) critical pressure gradients are obtained which are typically 25% higher. The profiles of \( j_t, q \) and \( s \) result from solving the diffusion equation for \( B_0 \). Both the thermal pressure and the anisotropic beam pressure are included.
**Fig. 1:**
Time evolution of $\bar{n}_e$, $H_\alpha$ intensity and $\beta_p$ in an H-discharge.

**Fig. 2:**
Comparison of measured (circles) and computed (solid curves) electron temperature profiles.

**Fig. 3:**
As in Fig. 2, but electron density profiles.

**Fig. 4:**
Profiles of $\chi_e$, electron collisionality factor $\chi_e$, ideal ballooning stability parameter $\delta p/\delta r/(\delta p/\delta r)_c$ and $n_e/d$ in $T_e/d$ in $n_e$ in L- and H-phases.
Several experiments were performed in TFTR to attempt to achieve an H-mode using the toroidal inner limiter. TFTR was operated at low $q_{\psi}(a)$, with the plasma current being ramped slowly during neutral beam injection. The discharges exhibited a phenomenon which we call the "$q$-mode" when $q_{\psi}(a)$ passed close to low-order rational numbers such as 3. The electron density profile broadened, while the $D_\alpha$ emission detected at some of the views decreased abruptly. These constitute the signature of the $q$-mode.

**Phenomenology** - Figure 1 shows plasma conditions for a $q$-mode discharge where 9 MW of $D^0$ beam power were co-injected into $D^+$ plasma from 2 to 4 s. The chord-integral electron density increases at 3.1 s when the transition occurs. Concurrently, the $D_\alpha$ emission decreases and $q_{\psi}(a)$ decreases below 3.15.

![Figure 1](image_url)

**Fig. 1** Overview of a discharge with a major radius of 233 cm, minor radius of 67 cm, and a transition to the $q$-mode occurring at 3.10 s. (a): chord-integral electron density, $D_\alpha$ emission at 30° above the midplane, and plasma current; (b): safety factor, $\beta_p^{\text{pia}}$, and the measured neutron production rate.

The $D_\alpha$ emission from various chordal views, shown in Fig. 2, exhibit periodic modulation with different phases in the different views. The sum of the $D_\alpha$ signals drops 30%, and then oscillates. The soft X-ray emission and the fluctuations of the poloidal magnetic field, $B_\phi$, measured by some of the Mirnov coils have similar modulations. These observations indicate a helical structure...
rotating toroidally in the co-direction before 3.15 s, then slowing down, and finally stopping or locking at 3.25 s.

The electron temperature in the region ~ 1-8 cm from the inner limiter decreases and flattens when the transition occurs and has modulations similar to those of the Dα emissions. The ion temperature profile was measured by Doppler broadening of photons from the Kα lines in Ni and from the charge-exchange recombination to C+5 (n = 7 → n = 6). The central ion temperature decreases after the transition. The electron density profile measured through vertical chords, broadens toward the inner limiter, then rises after the transition, but the Thompson scattering measurement along the midplane does not show a broadening near the limiter.

The charge-exchange efflux of energetic (~ 10-100 keV) beam ions drops precipitously after the transition. Co-injection was used for most of the experiments aimed at obtaining the q-mode. For these discharges, the central toroidal rotation velocity, \( v_\theta(0,t) \) decreases by ~ 50% simultaneously with the decrease in the frequency of the oscillations of the Dα emissions. Several attempts with approximately balanced injection and with slightly more counter-injection than co-injection also resulted in the q-mode, but \( v_\theta(0,t) \) was small and did not change significantly at the transition. Also \( \beta_P^{\text{dia}} \) decreased or remained unchanged, unlike the co-injection cases.

The visible bremsstrahlung signal increased at the transition, often ~ 100%. The radiated power profile measured by a bolometer array had a hollow profile, with the radiating ring moving radially inward and becoming slightly narrower after the transition. The total radiated power was ~ 1.5 MW and decreased transiently ~ 10%, then increased after the transition.

The q-mode has been observed in a variety of plasma conditions. The neutral beam power threshold was ~ 6 MW. Most of the observations have \( q_w(a) = 3.0-3.2 \), and for these, \( I_p \) is ~ 1-1.5 MA. Transitions were also seen with \( q_w(a) = 2.5-2.6 \), but the signatures were poor (with brief Dα transitions and small increases in \( n_e \)). Hints of transitions with \( q_w(a) \sim 4, 5, \) and 9 have been seen with very small \( n_e \) increases.

The range in \( n_e \) prior to the transition is limited, so neither upper nor lower \( n_e \) thresholds can be deduced from the data. Usually no gas-puffing (except for a pre-fill) was used, though several shots with gas puffs during neutral beam injection had transitions. The discharge in Figures 1-2 had D3 gas-puffing of 40 Tl/s between 2.5 and 2.95 s. The lowest value of \( n_e \) prior to
the transition was $2.4 \times 10^{19}/m^3$.

**Analysis** - The Mirnov coil signals were analysed to find resonant modes. During some $q$-mode discharges, slowly rotating modes were identified with the ratio of the poloidal and toroidal periodicities, $m/n$ being less than, but close to $q_{\psi}(a)$. For instance, a 3/1 mode was identified for the discharge in Figs. 1-2. The helicity of the mode agrees with the toroidal and poloidal time delays for the various measurements of the modulations in the $D_{\alpha}$ emissions. The intensity of the mode, indicated by the magnitude of the $B_\theta$ modulations, grew from the start of the transition until the mode appeared to stop or lock at $\sim 3.25$ s. Small amplitude locked modes and modes with $m > 5$ are difficult to identify accurately in TFTR, so possibly all of the transitions had locked modes near the last closed flux (LCF) surface with $m/n < q_{\psi}(a)$.

Detailed transport analysis was done with the TRANSP code. Values for the energy transport coefficients, $\chi_e$ and $\chi_i$ were calculated from the observed $T_e$ and $T_i$ profiles, and the heat convection multiplier was assumed to be 3/2. The wall recycling ionization rate was calculated with the neutrals code DEGAS using the $D_{\alpha}$ measurements. This rate was used in TRANSP to calculate the deuterium confinement time $\tau_D$ and the wall recycling coefficient, $R$.

The calculated $q_{\psi}(r)$ profiles and the toroidal current density profiles $j_\phi(r)$ were used to calculate the radial eigenfunctions of the resistive linear tearing mode. The measured amplitude of the Mirnov signals and the radial mode structure were used to calculate the helical flux contours. Figure 3 shows the resulting islands for the discharge in Figs 1-2. The peak in the $D_{\alpha}$ emission, in $B_\theta$, and in $T_e$ correspond to the X-point between the islands. The island location intersects the limiter, whose presence, however, is not included in the model. The width is in rough agreement with the range where $T_e$ decreases and oscillates coherently with the $D_{\alpha}$ emission.

The effective ion charge, $Z_{\text{eff}}(t)$, calculated from the visible bremsstrahlung emissions increase from about 2.5 to 3.3. Carbon and oxygen line emission increases as well. Modeling of these data with the MIST code indicates that the carbon concentration doubles at the transition. The oxygen concentration was $\sim 1/10$ the carbon concentration and metallic impurities had negligible relative concentrations ($\sim 10^{-3}$ that of carbon).

The increase in $Z_{\text{eff}}$ is consistent with most of the increase in the volume-averaged electron density coming from carbon. The increase in density reduces the slowing down time for beam ions, thus the calculated beam power deposited in electrons and impurities increases at the transition. The kinetic energy stored in the fast, co-moving ions from the neutral

![Fig. 3 Plot of constant helical flux at 3.12 s showing 3/1 islands and the location of the inner limiter. Chords where the $D_{\alpha}$ emissions and the Mirnov signal in Fig. 2 were measured are indicated.](image-url)
beams decreases, in qualitative agreement with the decrease in charge-exchange efflux. In the TRANSP calculations, the beam ion and plasma rotation energy were converted to thermal energy, causing a transient increase in heating power when the transition occurred.

The calculated global thermal energy confinement time before the transition is 15% larger than the prediction from Goldston's L-mode scaling for $H^0 \rightarrow D^+$ discharges, and it increases by ~10% at the transition due to the increased thermal energy. Most of the other q-mode discharges analyzed had $\tau_e$ closer to the L-mode scaling value before the transition, and had smaller increases during the transition. The magnitudes of the deuterium and electron confinement times, $\tau_D$ and $\tau_e$, are comparable to $\tau_E$, and rise at the transition, but remain higher after the transition. The calculated $R$ increases from ~0.9 to ~1 at the transition, then decreases to ~0.8 about ~2 s after the transition.

Discussion - Apparently, a necessary condition for the q-mode is the occurrence of resonant surface modes which rotate slowly or lock. Observations in other tokamaks of islands and of locked modes indicate that both are deleterious to confinement, unlike the case of the q-mode.

The increase in the carbon concentration is believed to result from increases in the carbon source and in the carbon confinement time, $\tau_C$, which presumably is comparable to that of $\tau_D$. A likely source is that the LCF surface becomes distorted by the islands, modifying the region of contact of the plasma with the limiter. An amorphous layer of carbon deposits onto low flux areas of the limiter during or after each discharge. The distorted plasma boundary increases the plasma flux onto these areas, and could increase the removal of carbon from these areas. Since these areas are above and below the midplane, the increased electron source would be in regions measured by the infrared laser interferometry along vertical chords near the inner limiter, but not along the midplane measured by Thomson scattering, consistent with the observations.

The reduction of $D^0_\alpha$ emission indicates a decrease in the rate of $D^0$ recycled from the limiter. This rate is $N_D R/\tau_D$, where $N_D$ is the number of $D^+$ within the LCF surface. The simulations indicate that $N_D$ is relatively constant, and there is an increase of $\tau_D$ and a reduction of $R$. The latter could be caused either indirectly by increased co-deposition of deuterium onto the new carbon layers caused by the increased efflux of carbon, or directly by the distortion of the LCF surface. During steady-state portions of discharges, the region of contact appears to become saturated with deuterium, but if the contact changes, unsaturated regions could become available to absorb more deuterium. During growing plasmas, the plasma boundary comes in contact with unsaturated limiter areas, so $R$ and $D^0_\alpha$ should decrease similarly.

Comparisons - The q-mode has many of the features of the H-mode. Besides the D$\alpha$ transition and the power threshold, both exhibit increases in $\tau_E$, $\tau_D$, carbon concentration, and $Z_{eff}$, broadening of $n_e(r)$, and decreases in the amplitudes of the charge-exchange efflux and the ion saturation currents in the scrape-off layer. Differences are the requirement for the q-mode of special $q_W(a)$, the drop in $T_e(r)$ near the edge, and the absence of edge relaxation modes. Also in the H-mode, the increase in $\tau_E$ generally lasts longer.

We wish to thank D. Heifetz for help. This work was supported by the U.S. Department of Energy Contract number DE-AC02-76-CHO-3073.

CONFINEMENT STUDIES OF H-MODE IN DIVERTOR/LIMITER DISCHARGES ON JFT-2M


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ABSTRACT. Detailed studies of the H-mode in a single-null divertor configuration were performed and the H-H like transition with ELM were observed. The effects of the divertor length on the H-mode were studied.

1. INTRODUCTION AND EXPERIMENTAL SETUP

JFT-2M[1] has capability of elongated limiter discharges and open divertor discharges, and in both configurations the high confinement mode (H-mode) has been observed[2]. In this paper, four cases with different divertor length(id=9.6, 6.2 and 2.0 cm) which means distance between the null point and divertor plates, are studied. Co- and counter-injection of hydrogen neutral beam with each maximum power of 0.8MW are used. Estimated net input powers are 90% of torus input for co-injection and 70% for counter-injection. The peripheral electron temperature is measured by the heterodyne radiometer [3] (90GHz) which is calibrated by the laser scattering and the soft X-ray energy analysis. The time response of the radiometer is 0.01ms and the effect of the wall reflection coefficient (80-90%) raises the optical depth of the gray radiation from the peripheral plasma to near the black body level.

2. OBSERVATION OF ELM AND H-H TRANSITION

Time evolutions of plasma parameters in a single-null divertor discharge (id=6cm, Ip=240kA, Bt=12.6kG, k=1.4, deuterium-gas) with NBI co-injection (0.79MW) are shown in Fig.1. Do, ECE and MHD signals in Fig.2-(b) suggest four-type changes (T1-T4). Fig.4-(a)-(c) show the changes of peripheral temperature distribution by ECE at T1, T2(T3) and T4. T1 is the usual L-H transition in which the peripheral temperature profile shows the stepwise changes, synchronizing with saw teeth crashes and also m=2 magnetic fluctuations[4] are stabilized. Fig.2 shows the electron temperature profile after L-H transition by laser scattering which has temperature in
consistency with ECE measurement. Such a profile corresponds to the tearing mode stable profile[5] in contrast with joule or L-mode profile and this suggests the possibility that m=2 magnetic fluctuation is tearing mode. Electron density at temperature pedestal by laser scattering is about 2-3x10^{13}/cc which is sufficiently higher value for temperature measurement by ECE. In T2 fluctuations (ELM:several kHz) in ECE signal at steep gradient region are observed and as seen in Fig.4(b) H-mode profile changes from H1 to H2 with more reduction of D_α intensity, magnetic fluctuations (Fig.1(b)) and peripheral SBD signal (Fig.2(d)). A increase of FeXV signal in H2 state indicates the higher particle confinement than H1 and corresponds to increase of radiation from center part(Fig.3). Transition in T2 means a short H2->H1-H2 transition and causes the D_α spike. This mode very frequently appears for shorter divertor length case and sometimes causes the energy degradation. T4 is usual H-L transition and due to radiation cooling etc. a steep gradient (>200ev/cm) relaxes to a gentle gradient(<100eV/cm) with increase of D_α, MHD fluctuation and sawteeth activity. For longer divertor length cases H-H transition is hard to appear and sometimes ELM causes the D_α burst and energy degradation.

3. The effects of divertor length on H-mode
The divertor length were changed from 9cm to -2cm (limiter). In joule phase m=2 mhd fluctuation and low-Z impurity increase with decrease of divertor length I_d. The dependence of the threshold heating power needed for L-H transition, remote cooling from divertor plasma, and stored energy on divertor length for NB1 co-injection are illustrated in Fig.5. The maximum divertor density for co and counter-injection is also shown in Fig.5 and time evolution of divertor plasma density profile for I_d=9cm divertor configuration is shown in Fig.6 and maximum divertor plasma density is about 3.5x10^{213}/cc. This region corresponds to the intermediate recycling region.

ACKNOWLEDGEMENTS
The authors are grateful to Drs. Y. Tanaka, M. Tanaka, K. Tomabechi and S. Mori for continuous encouragement.

REFERENCES
Fig. 1: Time evolutions of plasma parameters in a single-null divertor discharge ($l_d=6\text{ cm}$, $I_d=240\text{ kA}$, $B_t=12.6\text{ kG}$, $k=1.4, \text{deuterium-gas}$). MHD signal is envelope of magnetic fluctuation.

Fig. 2: The electron temperature profile after L–H transition by laser scattering.

Fig. 3: Time evolution of radiation profile measured by bolometer array.
Fig. 5: Threshold power, remote cooling, divertor density and stored energy vs. divertor length.

Fig. 6: Time evolution of divertor plasma density profile for 1d=9cm divertor discharges.

Fig. 4: The changes of peripheral temperature profile measured by ECE at T1, T2(T3) and T4 in Fig. 2.
Neutral Beam Current Driven Operation of the DIII–D Tokamak


General Atomics, San Diego, California, U.S.A.

Neutral beam current drive experiments in the DIII–D tokamak with a single null poloidal divertor are described. A plasma current of 0.34 MA has been sustained entirely by neutral beams with H-mode quality energy confinement. Poloidal beta values reach 3.5 without disruption or coherent magnetic activity, suggesting that these plasmas may be entering the second stability regime.

This paper presents results from the DIII–D tokamak in which the plasma current was sustained entirely by neutral beams for 1.5 sec without assistance from the ohmic heating transformer. After ohmic startup, the ohmic heating coil current was held constant (no loop voltage at the plasma surface) so that the plasma current could freely adjust to the neutral beam drive. The DIII–D tokamak was operated with a single-null divertor configuration having a 1.69 m major radius, 0.6 m minor radius, and 1.7 vertical elongation. The on-axis toroidal field was 2.1 T. Experiments were carried out with a helium plasma having a line-averaged density $n_e = 2 \times 10^{19}$ m$^{-3}$ and a central temperature of approximately 2 keV. Eight 75 kV hydrogen beams$^1$ injected in the same direction as the plasma current. Four beams intersected the vacuum system axis at 47$^\circ$ and four beams intersected at 63$^\circ$.

Plasma parameters are shown in Fig. 1 as a function of time. Initially a 0.22 MA ohmic discharge was established without sawteeth, indicating an on-axis safety factor $q_0 > 1$. At 1.1 sec the ohmic heating primary coil current was held constant, so without beam injection the plasma current decayed, as shown by the dashed line of Fig. 1a. With 10 MW of absorbed neutral beam injection (Fig. 1b), the plasma current increased to 0.34 MA. The rapid plasma current increase at 1.1 sec was largely due to an ohmic boost associated with the increasing flux from the vertical equilibrium coils which adjust to the increasing poloidal beta. During the period when the current was sustained, the loop voltage (Fig. 1c) was zero, excepting periodic voltage spikes associated with ELM-like relaxation phenomena observed on outer radial chords of soft x-ray emission and on an external magnetic probe. Neutral beam injection increased the total plasma energy determined from magnetic and diamagnetic measurements, as shown in Fig. 1e. The poloidal beta, shown in Fig. 1f, reached 3.5.

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Two transport codes that include current drive physics models have been used to model these experiments. These codes are the General Atornics ONETWO code and the ACCOME code from the Japan Atomic Energy Research Institute. Results from these codes, given in Table 1, indicate that the current is predominately fast ion driven rather than bootstrap.

An important aspect of these results is that the energy confinement was of H-mode quality in a noninductive current drive tokamak. The 24 ms energy confinement time of this low-current high-power discharge (72 ms/MA) slightly exceeds the 67 ms/MA DIII–D H-mode scaling obtained with 8.4 MW hydrogen beam injection into a deuterium plasma. The 24 ms energy confinement is 2.4 times longer than Kaye-Goldston L-mode scaling. ELM phenomena characteristic of H-mode appeared although the density rise common to H-mode did not occur.

Seven discharges with MHD $\beta_p > 3.4$ and diamagnetic $\beta_p > 3.0$ have been produced without disruptions or coherent $n \neq 0$ modes. The $\beta_p = 3.5$ discharge shown in Figs. 1 and 2 has an inverse aspect ratio $\epsilon = 0.31$, so that $\epsilon \beta_p = 1.1$. While these values of $\beta_p$ and $\epsilon \beta_p$ are among the highest achieved in a tokamak, they are not remarkable in themselves. What is remarkable is the absence of the large amplitude MHD modes, observed in ISX–B$^3$ and Doublet III$^4$ at high $\beta_p$. The only coherent MHD activity in the present discharges is associated with ELMs which are similar to those observed at lower $\beta_p$. In these discharges the toroidal beta was 0.5%.

Figure 2 shows EPITD magnetic analysis$^5$ of measurements from 41 flux loops and 25 magnetic probes distributed around the walls of DIII–D. Shown is the flux surface equilibrium, radial current density, and $q$ profiles at two times: (a) before beam injection, and (b) in steady state when $\beta_p = 3.5$. This magnetic analysis is based on a third-degree polynomial in magnetic flux. The radial current profile exhibits (see Fig. 2b) a 0.2 m outward Shafranov shift and an extreme outward peaking which, theoretically, are necessary for second regime stability.$^7$–$^{12}$

The issue of whether these plasmas entered the second stable regime revolves around the exact value of the axial safety factor $q_0$. MHD equilibrium analysis of external magnetic measurements can accurately determine the edge $q$ (or $q$ at the 95% flux surface, $q_{95}$), the magnetic axis shift, the poloidal beta, $\beta_p$, and the plasma internal inductance, $\ell_i$. Information on $q_0$ is only weakly obtained and is dependent on the functional forms assumed in fitting the current profile. Our best estimates are that the axial $q_0$ during NBI current drive is $3 \pm 1$ (determined by minimizing chi-squared within one standard deviation). The absence of soft x-ray sawteeth supports $q_0 > 1$ in the ohmically-heated plasma and the drop in $\ell_i$ seen in Fig. 1g shows substantial current profile broadening during neutral beam injection that would further increase $q_0$. Ballooning mode$^9$ analysis indicates that the plasma is entering the second stability region$^7$–$^{12}$ if $q_0 > 2.5$.

The authors would like to thank our DIII–D colleagues for assistance in carrying out and interpreting these experiments. In addition, we wish to acknowledge valuable discussions with J. Kesner, M.E. Mauel, and A.M.M. Todd.

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References

2. Poloidal Beta is defined as $\beta_p = 2 \mu_0 \langle p \rangle / B_0^2$ where $\mu_0 = 4 \pi \times 10^{-7}$ H/m, $\langle p \rangle$ is the total plasma pressure integrated over the plasma volume, and $B_0 = \mu_0 I_p / (\text{poloidal circumference})$.

Table 1

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$^{(a)}$ Does not include bootstrap current effect from hot ion deposition.
Fig. 1. Time dependence of (a) plasma current, (b) neutral beam injection power, (c) loop voltage, (d) line-averaged density, (e) plasma energy, (f) poloidal beta, and (g) internal inductance.

Fig. 2. A comparison between ohmic (1.08 s, dashed line) and beam driven (1.4 s, solid line) EFIT magnetic analysis of (a) flux surfaces, (b) midplane plasma current density, and (c) safety factor $q$. 
H–Mode Study in DIII–D+


General Atomics, San Diego, California, Lawrence Livermore National Laboratory, Japan Atomic Energy Research Institute, Hitachi, Ltd.

Introduction

H–mode has been observed in DIII–D single-null divertor discharges (R=1.67 m, a=0.63 m, b/a ~ 1.9, $B_T \leq 2.1$ T, $I_p \leq 2.0$ MA, $P_{\text{total}} \leq 10$ MW) [1,2]. Many observed features of the H–mode are similar to those of other devices. A major goal of DIII–D research is to understand the underlying H–mode physics mechanisms. This understanding may be critical for further improvement of the H–mode discharge performance. In this paper, we discuss new and detailed experimental findings, closely related to the H–mode mechanisms.

H–Mode Accessibility

To achieve the H–mode transition, a minimum input power and a minimum density are required. These minimum values depend on various plasma parameters. The detailed parametric dependence has not yet been established. However, the H–mode threshold power is lowest for low $B_T$, deuterium-plasma, and $\nabla B$ drift toward the X–point. At $B_T=1.1$T, the H–transition has been achieved with low power ($0.8$ MW) ECH alone [3]. Furthermore, a temporary H–mode phase lasting ~30 msec is triggered by a sawtooth crash in purely ohmically-heated ($1.4$MW, $B_T=0.9$T), low q discharges.

Profiles of the Electron Temperature and Density

Figure 1 shows the temporal evolution of the electron temperature and density profiles in a typical H–mode discharge ($B_T=2.1$ T, $I_p=1.26$ MA, $P_{\text{total}}=7$ MW) [4]. Within several msec after the transition, the density profile becomes flat with a very steep edge gradient and further develops into a hollow profile. The averaged density continuously increases during the entire H–phase until an edge localized mode (ELM) occurs. The edge electron temperature rises rapidly from 200 eV to 350 eV within ~50 msec after the transition and then saturates. The edge density gradient is very high (~$3.5 \times 10^{13}$ cm$^{-4}$), while the edge temperature gradient is modest (~100 eV/cm), compared with those reported by ASDEX ($\nabla n = 0.5 \times 10^{13}$ cm$^{-4}$, $\nabla T=200$ eV/cm) [5]. This difference in the edge profiles, particularly the edge $\nabla T$ may be an important factor in explaining the difference in the normalized energy confinement ($\tau_E/I_p$). $\tau_E/I_p$ for ASDEX is at least 50% higher than that for DIII–D[5].

Edge Localized Mode

Two kinds of ELMs, giant and mini, have been seen during the H–mode phase. A giant ELM is an edge relaxation mechanism that prevents the density and the stored energy from rising continuously and dumps typically 10% to 15% of the total energy and particles stored in the plasma onto the divertor plate within 10 msec [6]. The effect of the giant ELM extends well beyond the edge region and influences the global energy confinement by ~30% (lower $\tau_E$ for frequent ELMs). This effect becomes severe at low q and low $B_T$ where unfavorable (high

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frequency and large amplitude) ELM behavior, perhaps interacting with an expanded sawtooth zone, reduces $r_E$. The pressure gradient prior to a giant ELM is found to scale approximately as $I_p^2$ and to be close to the first ideal ballooning stability limit for a wide range of the plasma current, as shown in Fig. 2 [7]. Thus it is reasonable to believe that the ballooning mode triggers an ELM and destroys the steep edge gradients. Fig. 3 shows the time sequence of the H-ELM transition which occurs at $t=2.1$ sec in the discharge of Fig. 1. Occasionally clear and coherent precursor oscillations appear as the first event in an ELM. The oscillation amplitude is high on the outer midplane, The toroidal mode number ($n$) is $\sim 6$ and the frequency is $\sim 40$ kHz. This coherent mode may be the ballooning mode (further investigation is needed to be conclusive). This precursor appears to lead to an L-phase, characterized by a high amplitude of incoherent magnetic fluctuations (described later). A sudden change in the edge soft X-ray signal and in the divertor $H\alpha$ signal occur after the rise of the incoherent fluctuations. The ELM phase typically lasts 10 msec, but for some cases, particularly at high $I_p$ ($\sim 2$ MA), the ELM phase continues and becomes a clear L-phase, which lasts more than 100 msec.

In addition to the giant ELMs described above, repetitive mini-ELMs with a period of $\sim 1$ msec appear every few msec when the input power is slightly above the power threshold or when the distance between the X-point and the divertor plate/ or between the separatrix and the limiter is small. The amount of particles and energy expelled during a mini-ELM is small, less than 1% of the total particles and energy stored in the plasma. However, continuous mini-ELMs cause the density rise to saturate. $r_E$ with such ELMs is comparable to values with infrequent giant ELMs.

**Energy Confinement Scaling [2,8]**

Most of the experiments have been done with hydrogen beams into deuterium plasma ($H^0$-$D^+$) which causes hydrogen dilution. With increasing hydrogen dilution in the deuterium plasma, the energy confinement decreases. Part of the reason for the isotope dependence of $r_E$ may be that ELM behavior becomes more unfavorable as the plasma changes from pure D to pure H (see Fig. 4). For deuterium cases, $r_E$ is independent of $P_{total}$, $n_{av}$, and $B_T$. $r_E$ is 120 msec at $I_p = 1$ MA. With significant dilution (between 30% and 50%), $r_E$ decreases with increasing $P_{total}$. $r_E$ increases linearly with $I_p$ and $r_E/I_p = 85$ msec/MA (67 msec/MA) at $P_{total} = 5.7$ MW (8.4 MW). For pure hydrogen discharges ($H^0$-$H^+$), $r_E$ drops from ohmic value to 60 msec at $I_p = 1$ MA and $P_{total} = 7$ MW. $r_E$ in pure helium discharges ($He^0$-$He^{++}$) is found to be similar to that of pure hydrogen discharges.

**Edge Incoherent Magnetic Fluctuations**

Edge incoherent magnetic fluctuations correlate well with the edge transport, particularly at the transition and may be responsible for rapid transport at the edge during the L-mode discharges. At the L-H or H-ELM transition, changes of all other parameters observable with our present diagnostics do not precede the rapid change in the magnetic fluctuation amplitude, as shown in Fig. 3. The poloidal amplitude distribution of the incoherent magnetic fluctuations for well diverted discharges is found to be very peculiar. The amplitude is high near the divertor hit spot where the separatrix intersects the divertor plate and low at other locations. An interpretation of this is that an image current induced by the edge magnetic fluctuation just inside the separatrix flows in the scrape-off layer, shielding the magnetic fluctuation from probes outside the plasma. The image current then flows into the divertor plate, resulting in a high fluctuation level at the divertor hit spots. The abrupt change in the fluctuation amplitude at the transition is simultaneous (within 20 μsec) at all poloidal locations, suggesting that the magnetic turbulence has a poloidally extended structure. Because the magnetic fluctuations
precede the change in the \( H_\alpha \) emission from the divertor region, we feel that the magnetic fluctuations do not originate in the divertor region. The microtearing mode, amplified by bootstrap current, is a good candidate for the observed magnetic turbulence. It not only explains a strong \( \beta_p \) dependence of the fluctuation amplitude, observed in the earlier Doublet III limiter (L-mode) discharges [9], but also predicts existence of the low amplitude regime (H-mode) at edge temperatures above a threshold.

References

Fig. 1. (a) Density and \( H_\alpha \) behavior at the H-transition. (b) time evolution of the electron temperature and density profiles (measured by Thomson scattering), (c) time evolution of the edge electron temperature and density gradients in an H-mode discharge.
Fig. 2, Normalized electron pressure gradient (just inside the separatrix) prior to ELM is plotted as a function of \( S/q_{95}^2 \) where \( S \) is the shear and \( q_{95} \) is the \( q \) value at the flux surface with 95% of the separatrix flux value. The open points are the theoretical (first) ballooning stability limit \( (T_e(r) = T_i(r) \) is assumed).

Fig. 3, Time sequence of the H-ELM transition at \( t=2.1 \) sec in the discharge of Fig. 1. \( b(\text{out mid}) \) is the signal of the Mirnov coil on the outer midplane, which exhibits coherent precursor oscillations. \( b(\text{div}) \) is the signal of the Mirnov coil near the divertor hit spot, which exhibits incoherent fluctuations.

Fig. 4, ELM characteristics change with ion species. Solid line: divertor \( H_\alpha \) emission (A.U.), dotted line: density (\( \times 10^{13} \text{ cm}^{-3} \)) \( (B_T = 2.1 \text{ T, } I_p = 1 \text{ MA}) \).
STUDIES OF ENERGY TRANSPORT IN JET H-MODES

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1. INTRODUCTION
The H-mode discharges in JET separatrix configurations with a single null X-point have the best confinement properties [1,2]. However, even these discharges exhibit a decrease of the global confinement time with increasing neutral beam injection power. Part of this degradation with power can be attributed to increased impurity radiation and poor beam penetration at the higher densities concomitant with improved particle confinement and the higher power levels. In addition, these discharges are not in steady-state, since the density increases steadily with time until impurity radiation precipitates a transition back to an L-mode.

In the present paper, the local heat transport properties in the interior of ohmic, L- and H-phases of 2MA discharges are determined (Section 2) using both the time-dependent energy balance code, TRANSP [3], and the timeslice code, QFLUX [2]. In addition, the global confinement properties of higher current discharges (≤3.8MA) are analysed in Section 3.

2. LOCAL TRANSPORT ANALYSES
The input data to TRANSP and QFLUX comprises the magnetic flux surface geometry, temporal variation of the electron density profile (FIR interferometer), electron temperature profile (ECE), central ion temperature (X-ray crystal spectrometer measurement of Ni^{27+}, neutron yield and NPA data), radiated power profile (bolometer), the plasma current, loop voltage, Z_{eff} (visible bremsstrahlung) and edge particle confinement time (H_{n} monitors). Neutral beam heating is calculated by Monte Carlo methods in TRANSP and using the PENCIL code in QFLUX.

The ion temperature profile is not available routinely. It is calculated in TRANSP by assuming an ion thermal diffusivity either (a) of the neo-classical form enhanced by a factor adjusted continuously in time so as to regulate the central ion temperature at the measured value, or (b) proportional to the electron thermal diffusivity. Both prescriptions give good agreement with the measured diamagnetic stored total energy and reasonable agreement with the neutron yield; the ion temperature profiles using (b) are nearer to the NPA profile data. The ion temperature in QFLUX is assumed to have the same profile as the electron temperature, but normalised to the measured central ion temperature. Electron and ion contributions to the heat flow (which depend on the ion temperature profile through energy equipartition) are not separated and results are presented for the total heat flux, Q, across a flux surface of area, S, in terms of n_{e}V_{ke}T_{e}S. An effective total heat diffusivity, \chi = Q/n_{e}V_{ke}T_{e}S, is defined.
TRANSP Results: For JET pulse #10755 (2MA, 2.1T, P_{NBI} ≤ 9MW) the temporal variation of Q with n_e T_e S at three radial positions in the plasma interior (normalised radius, ρ = 0.5, 0.75, 0.9) is shown in Figure 1 (at times avoiding the collapse of a sawtooth) through the ohmic, L- and H-phases. The temporal variation of χ (at all times) at the three radial positions is shown in Figure 2(a). At the start of the L-mode χ is well above its ohmic value but decreases subsequently to a level close to its ohmic value prior to the transition to the H-mode. χ is maintained close to this level during the H-phase. The spatial profiles of χ at selected times throughout the discharge are shown in Fig 2(b).

QFLUX Results: For the same JET pulse, QFLUX results at a normalised radius, ρ = 0.5, and a limited number of time points show reasonable agreement with the results of TRANSP (Figure 1). Figure 3 shows an accumulation of data from a number of discharges at 2MA and 2.1T at the normalised radius, ρ = 0.5. In Figure 3(a), the data is distinguished according to the phase of the discharge, with χ decreasing continuously from the L- to H-phase. In Figure 3(b) the results are distinguished with respect to density and indicate that the transition from the L- to H-phase is concomitant with an increase in density that is observed. At constant density, χ increases with power, while at constant power, χ is reduced with increasing density. This result substantiates earlier indications [2].
Fig 3: QFLUX results for data accumulated at several times during a number of pulses distinguished according to (a) discharge phase and (b) density.

3. GLOBAL ENERGY CONFINEMENT

The local transport analyses discussed in Section 2 were carried out for 2MA H-mode discharges since most of the H-modes obtained in 1986 were at this current. Since then a larger number of H-mode discharges with currents in the range 3 to 4MA have been obtained.

Figure 4 shows the global energy confinement time $\tau_E = W/(P_{tot} - W)$ as a function of the total net input power ($P_{tot} - W$) for H-mode discharges with currents of 3MA and above. $W$ is the diamagnetic measurement of the stored energy. Most of the data arise from one specific day of operation in February 1988 on which the neutral beam power, $P_{NBI} < 7.5$ MW. These data are supplemented by the 3MA data obtained in November 1986 ($P_{NBI} < 10$ MW) [2]. Since in all these discharges the plasma density was increasing steadily with time the data set has been restricted to $W < 0.2$ $P_{tot}$, relatively close to steady-state conditions. Two density ranges (3-4 and 4.5 x $10^{19}$ m$^{-3}$) are distinguished. At low powers the data seem to split into two branches with data at higher density corresponding to higher confinement times. At higher powers such a density dependence is less apparent.

Figure 4 also shows the first H-mode data obtained at a plasma current of 3.8MA. At a total heating power $P_{tot} = 8$ MW this discharge resulted in record values of plasma energy (7MJ), energy confinement time with appreciable additional heating (0.9s) and fusion product $(n_1(0)\tau_E T_1(0) = 2.4 \times 10^{29} \text{m}^{-3} \cdot \text{s} \cdot \text{keV})$. 
4. CONCLUSIONS

Local transport analyses with both the TRANSP and QFLUX codes indicate that during the L-phase of JET single null X-point discharges the total heat transport coefficient in the plasma interior decreases from its initially high value to a level close to the ohmic value. This level does not change significantly during the transition to the H-phase, implying that the higher stored energy associated with the H-phase results from improved confinement at the very edge of the plasma. On the basis of a larger data set obtained at different times during different discharges, the underlying trend towards reduced confinement in the plasma interior with increasing input power is ameliorated, in part, by improved confinement at higher density.

The global analysis indicates that confinement during the H-phase continues to improve with current (up to 3.8MA), but still degrades with increasing neutral beam power.

![Characteristics of a JET discharge in which ELM-phases separate a train of shorter H-phases.](image)

It is worth noting that conditions closer to steady-state have been achieved in discharges that exhibit ELMs separating a train of shorter H-phases. The duration of these ELM-phases can be affected by varying the radial distance between the separatrix and the belt limiters. As shown in Figure 5, for a certain optimum distance between separatrix and limiters, this procedure allows the plasma density and radiation to be kept closer to steady-state without too large an adverse effect on the global energy confinement time.

ACKNOWLEDGEMENTS

The authors would like to acknowledge R J Goldston and D C McCune of the Princeton Plasma Physics Laboratory for the use of the TRANSP code and C H Best for supervising the installation of TRANSP at JET.

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THE JET H MODE

A Tanga, D Bartlett, M Bures, A Gibson, N Gottardi, C Gowers,
P J Harbour, A Hubbard, M Keilhacker, E Lazzaro, P D Morgan, P Noll,
P H Rebut, N Salmon, D Stork, D D R Summers, A J Tagle, P R Thomas,
K Thomsen, M von Hellermann, M L Watkins

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Magnetic configuration

H mode has been obtained in JET with single null configuration and neutral beam heating. The single null magnetic configuration was obtained by using the multipolar field produced by shaping windings and primary coil stray fields. The plasma equilibrium flux plot of a discharge with plasma current of 4 MA is shown in Fig. 1. Developments of higher current scenarios with single null configuration are based on differential currents in the main equilibrium coils and have been described. Ranges of parameters of single null discharges are summarised in Table I. The elongation of single null discharges is larger than limiter discharges because of the need to keep the plasma detached from the bottom belt limiter. The magnetic axis of the plasma is displaced by 0.1 ± 0.2 m below the midplane; consequently plasma-antenna distance for single null configuration at high plasma current is larger than 0.1 m, causing values coupling resistance of the order of 10 Ω or less. With single null configuration the operational space was confined by a low q limit of q* = 5.2p R (1+R2); for values of q* 2 to 3 there was still a low density disruptive limit of n_eff (m^-3) = 3.5x10^19 Ip/MA. In the single null configuration JET has been operated with the ion (VB) drift direction towards the X-point.

Features of the H-mode

Scans of neutral beam power, plasma current, toroidal field and X-point position have shown that threshold power for the H-mode transition depends strongly on toroidal field, with an empirical fit of the neutral beam power P_th (MW) = 0.6(B_T(T))^2 ± 0.5. The minimum power threshold was 3 MW with a toroidal field of 1.8 T. There was no clear dependence on plasma current, within the power resolution of one neutral beam source of ±1 MW. The transition to the H-mode was achieved for X-point to the dump plates distances <5 cm. Within the limited neutral beam power available of 7 MW it was not possible to achieve H-mode with He^+ target plasma, suggesting a higher threshold for helium plasma, similar to D-111-D(3). In deuterium, the H-mode transition was achieved with a minimum target plasma density of n_eff (m^-3) = 0.5x10^19 Ip/MA.

The short (0.1-0.45 s) L phase is followed by a quiescent H phase, lasting up to 2.5 s. The collapse of the H-phase is caused by the build-up
of radiation losses, in the bulk plasma, which increase with plasma density until the total input power less radiation losses is equal to the power threshold. The duration of the quiescent H phase increases with (a) neutral beam power, (b) clean machine conditions and (c) He conditioned vessel surfaces leading to lower plasma recycling. Termination of the H phase is followed by either a full plasma disruption or reverting to an L phase. For plasma-limiter distances less than 5 cm a series of short (0.5 s) H and L phases is observed, with distances of the order of 10 cm within a Neutral Beam pulse of 5 s usually there are two or three H and L phases lasting $1 + 2$ s; finally for distances larger than 10 cm there is just one long H phase ending with plasma disruption.

At the lowest toroidal field of 1.8 T, the L phase lasts for just the slowing down time of the injected particles, while at 2.8 it lasts 0.6 s, (NB power 7+8 MW). The L phase is characterized by edge relaxations which increase in amplitude and decrease in frequency until at a sawtooth crash, the L to H transition occurs (as detailed elsewhere(4)). The edge relaxations are localized at $r/a \approx 0.90$, according to the soft X-ray array. Each relaxation produces the loss of approximately 0.4% of the total plasma particles and 0.2% of total energy. The L to H transition is marked by a 20% increase of the electron temperature, confined to the separatrix, taking place in a time of the order of 10 ms.

The time evolution of electron temperature and plasma density are shown in Fig. 2. This figure is made up of a series of LIDAR profiles(5) taken on a sequence of subsequent reproducible plasma pulses, resulting a profile every 0.1 s. NB heating starts at 12.5 seconds, during the flat top of a discharge of 3 MA plasma current and with a toroidal field of 2.1 T. At the L - H transition, at 12.8 s the plasma density profile, already flat in the ohmic phase, exhibits the build-up of a steep edge gradient.

The density increase, throughout the H phase, occurs with a profile little altered in shape. It should be noted that flat density profiles are a common feature of both limiter and x-point JET discharges at low values of $q_{\text{cy}}$. The sequence of electron temperature profiles shows the increased edge temperature and the perturbations caused by sawteeth activity. Electron temperature profiles in H-mode discharges are broad. For example, the values of the ratio of peak-to-average temperature is 1.7, virtually independent of $q_{\text{cy}}$; in contrast to limiter discharges where, during the flat top, the ratio of peak to average electron temperature varies as $T_e/\langle T_e \rangle \approx \frac{3}{4} q_{\text{cy}} + 1(6)$. H-modes usually have sawteeth activity. However temporary stabilization has been observed in a similar way than monster sawteeth observed in limiter discharges, with neutral beam power of 7 MW. Temporary stabilization produced an increase in the peaking of the temperature profile.

### Confinement time

The global confinement time ($\tau_e$) increases in the H-mode, and is roughly a factor of two larger than in limiter discharges at the same plasma current and 50% larger than during the L phase(1). In the current range 1.0 to 3.8 MA confinement times scale proportionally to plasma current and are not sensitive to toroidal field. The gain in confinement of the H mode, compared to the L mode, is mainly due to higher average
densities and partially to higher electron temperatures. The ion central
temperature, measured by x-ray crystal and c-x spectroscopy, exceeded the
electron temperature in the early part of the H-mode, at lower densities,
while approaching the central electron temperature later at higher
densities. In the early phase of the JET H-mode, after the L to H
transition, the confinement time includes a strong contribution due to the
large value of the derivative of total stored energy, of the order of 50%
of the total input power. In this phase the power deposition profile of
beams is centrally peaked. Gradually, the rate of increase of stored
energy decreases and, at the same time, the beam power deposition profile
becomes quite peripheral. A possible density dependence of energy
confinement time is hidden in the non stationary nature of the discharge.
Cleaner discharges tend to have lower radiation losses, with longer
H-phases, and larger values of global confinement times. Local transport
studies indicate that the total energy diffusivity in the plasma
interior returns to ohmic level already during the evolution of the L
phase. Abel inverted bolometer radiation profiles typically show a steady
increase of radiation during the H phase at the periphery of the plasma.

During long pulses, where a succession of H and L phases is observed, a
gradual increase in the central radiation is also observed, suggesting a
gradual drift toward the centre of metal impurities. When a monster
sawtooth during H-mode is present, the radiation profiles show a slight
reduction of radiation losses across the radius.

Conclusions

In JET H mode the improvement in confinement is of the order of a
factor of two compared to limiter discharges at the same level of plasma
current. The threshold power scales greater then $B_i$. Global confinement
time scales with input power and plasma current, roughly twice the Goldston
scaling, but this should be offset against the increasingly off-axis
heating due to bad beam deposition profiles at high neutral beam power.
The values of $q_{cyt}$ for H-mode discharges range between 1.8 and 4.5, compared
with a range of 2 and 10 for limiter discharges. Sawtooth reconnection
radii are in the range of $r/a = 0.5+0.7$. It is possible that sawtooth
activity has an effect on plasma confinement and on density and temperature
profiles. So far ICRH have been unsuccessful in producing or sustaining
an H-mode because of a) low coupling resistance due to large plasma-
antennae separation and b) deleterious effect of increased impurity influx
caused by the antennae. Future experiments will be carried out at larger
plasma current and higher neutral beam power. The range of H mode
discharges to will also be extended higher values of toroidal field.

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  Diego, USA (Nov 1987).
(4) A Hubbard et al. This conference.
(5) C Gowers et al. This conference.
Table I

<table>
<thead>
<tr>
<th>Description</th>
<th>Value</th>
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<td>Plasma current range in single null configuration</td>
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<tr>
<td>Plasma current range with H mode</td>
<td>1-3.8 MA</td>
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<tr>
<td>Toroidal field range</td>
<td>1.7-2.8 T</td>
</tr>
<tr>
<td>Neutral Beam power range for H-mode</td>
<td>3-9 MW</td>
</tr>
<tr>
<td>Neutral Beam energy range</td>
<td>70-80 keV</td>
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<td>Target discharge gas</td>
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<tr>
<td>Neutral Beams</td>
<td>D+</td>
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<tr>
<td>Maximum line average density with H mode</td>
<td>6x10^{19} m^{-3}</td>
</tr>
</tbody>
</table>

Fig. 1
Flux plot for pulse No: 14328
Plasma current 4 MA
Toroidal field 2.4 T
Plasma density 2.10^{19} m^{-3}

Fig. 2
a) Time evolution of the plasma electron density profile
b) Time evolution of the plasma electron temperature profiles
LIDAR profiles obtained from a series of reproducible discharges. Plasma current = 3.0 MA, Toroidal field = 2.1 T, NB power 7 MW, NB pulse starts at 12.5 s, L-H transition at 12.8 s.

References (continued) (7) M Keilhacker et al. This conference.
(8) K H Behringer et al. This conference.
INTRODUCTION

H-mode operation of the JET tokamak was established in 1986[1]. Since then the phenomenon has been studied for a range of powers and particle energies of the Neutral Beam heating system. The operational regime has also been extended to higher plasma current[2]. The transition from L to H-mode after the initiation of NB injection, and the subsequent evolution of the profiles of the plasma parameters, have been investigated using a range of the existing diagnostics including bolometry, interferometry, soft x-ray tomography, charge-exchange and $D_α$ visible spectroscopy, ECE and neutral particle analysis, and the new JET LIDAR-Thomson scattering[3] and FIR polarimetry systems[4]. The results of this study and of a preliminary analysis of the ballooning mode stability of the measured pressure profiles are reported and discussed below.

PROFILE EVOLUTION MEASUREMENTS

Figure 1 shows the typical evolution of the main plasma parameters and H-mode signatures during the development of the discharge from an ohmic X-point to L-mode, to H-mode (with Sawtooth oscillations) and back to an ohmic X-point discharge. In this series of discharges the transition from L to H-mode is characterised by a significant fall in the $D_α$ signal (vertical chord) after a series of "spikes". The spikes have been correlated with mhd modes near the plasma edge (Edge Localised Modes or ELMs). In about half the cases in the series the L to H transition and the termination of ELMs appear to coincide with a sawtooth crash.

The central region of the electron temperature profile, which can be slightly hollow during the ohmic X-point phase, figure 2a, quickly fills in and rises to close to 4 kev just after the transition phase. It remains around this value until the NB heating is terminated when it falls back again to the 2 kev range in about 100ms, a time comparable to the slowing down time of the fast ions. It is clear from the outer region of the $T_e$ profiles that during this series of discharges a pedestal of about 400ev forms at R-4.0m during the L to H transition and this remains throughout the H-mode period.
Fig. 1 The different phases of development of a normal 3MA, 2.1T H-mode: a) total beam power (MW) b) line density ($x10^{20} \text{ m}^{-2}$) c) central ECE signal showing sawteeth (keV) d) $D_\alpha$ signal (arb. units)

Fig. 2 LIDAR electron a) temperature b) density and c) pressure profiles during ohmic X-point (12.3s) and H-mode (14.3s) periods of #14830. (bars denote measurement errors)
A more dramatic change is observed in the density profiles, figure 2b. During the transition the central peak flattens as the density rises and a steep edge gradient is formed. The overall level rises throughout the H-Mode period due to the fuelling effect of the beams giving gradients in the edge region up to $2 \times 10^{20} \text{ m}^{-4}$. During the sawtooth rise phase the profile becomes slightly hollow and then flattens again immediately after the crash.

The LIDAR density profiles have been used in the analysis of the polarimetry results and these indicate that the q-profile broadens during the H-mode period.

$T_i(R)$ from multi-channel CXRS during the H-mode is similar in shape to $T_e(R)$ with central values in the range 4-6kev. When the ion component is included with the electron pressure profiles derived from the LIDAR $T_e$ and $n_e$ data, figure 2c, near the end of the H-mode, the central and global values for $\beta$ are -2.8% and 2% respectively, close to the highest values yet recorded on JET. The pressure profile has been found to be ballooning mode stable when examined using the analysis of reference [7], with the central region close to the stability limit.

The signals from the bolometer camera during the X-point phase are generally very asymmetric with much higher intensity on the lines of sight nearest the inner torus wall. In general this prevents Abel inversion for the X-point period. However, once the transition to H-mode has taken place, symmetry is restored (ie the total radiation is constant on a flux surface) and the generalised Abel inversion[5] can be carried out. The evolution of the total radiation profile is shown in figure 3. Again the profile exhibits a central depression throughout the H-mode, although it does tend to fill in somewhat as it evolves. The profile shape and the modest increase in the overall radiated power level indicate that significant impurity build up in the plasma centre is not occurring. This is supported for the light impurities by $Z_{\text{eff}}$ profile results.

The hollowness during sawtoothing is also a feature of the soft X-ray profiles figure 4, and again in contrast to results on other machines, the increase in signal intensity is consistent with the density rise due to fuelling by the beams and not indicative of any significant increase in impurity radiation.

**IMPURITY ACCUMULATION** In the majority of the observed H-modes, which last for ~2s, no significant impurity accumulation in the centre of the discharge has been seen. This may be related to the presence of a large $q=1$ region in these sawteething H-modes and also to the long timescales that can be expected for the accumulation process on a large tokamak. This result is in contrast to ASDEX where in sawtooth and ELM-free H-modes, accumulation of iron was observed when the H-mode duration was sufficiently long (>150ms)[6].
SUMMARY The development of hollow $n_e$, soft x-ray emission and total radiation profiles is characteristic of JET H-modes. Impurity accumulation is generally not observed in the normal sawtoothing, ELM free discharges. A high central $\beta$ is achieved and a preliminary analysis of the stability against ballooning modes indicates that, near the centre, the pressure profile is close to the limit but within the normal stability region.

References
ANALYSIS OF CURRENT AND PRESSURE PROFILES IN JET H-MODE DISCHARGES

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INTRODUCTION.

High confinement discharges have been obtained in JET with a magnetic separatrix configuration with a single stagnation point [1] at currents up to 3 MA and intense neutral beam (NBI) heating. When the NBI power exceeds about 2.5 times the ohmic power a transition occurs from the low confinement (L) regime to the high energy confinement (H) mode. The characteristic signatures of H-modes are, as observed previously in ASDEX [2] the sharp increase in density, the appearance of pedestals in the measured temperature profiles and a rather quiescent MHD behaviour. A significant difference between JET H-modes and those of ASDEX is the persistence in JET of sawtoothing conditions with q(0)<1 across the L-H transition. In ASDEX the disappearance of sawteeth was related to a flattening of the current indicated by reduction of the internal inductance.

An analysis of the X-ray emission signals, of the interferometric and LIDAR measurements of the density, of the ECE and LIDAR measurements of the electron temperature and pressure profiles consistent with the magnetic equilibrium, shows that JET H-modes may be characterised by a long non-steady state condition with unconventional density and pressure profiles which are flat or hollow within the q=1 surface, which expands after the L to H transition. Fig.1 shows the LIDAR, pressure and q profile for the 3 MA H-mode shot 14815, in the L and H phases.

The pressure and current profiles obtained in this class of H-modes allow sawtoothing, with q(0)<1 while keeping stability against localised or ballooning instabilities in the low shear plasma core according to Mercier criterion.

In the confinement region, outside the q=1 surface, there is evidence that the current and pressure profiles are related in a way similar to that of the theoretical model by B. Kadomtsev [4], in which the H-mode appears to be a branch of possible equilibria with minimal total plasma energy for a given total plasma current. The implications of the model are that only incoherent perturbations which leave the total current constant, are responsible for the relaxation to a minimum energy state, while low number modes can still be unstable. However this is not consistent with JET H-mode data as JET H-modes are sawtoothing because the core equilibrium still corresponds to an L-mode with degraded confinement which should hence be obtained by helicity constraint in the mixing region. Close to the Troyon limit finite aspect ratio and finite beta effects can be identified in the current density profile (Pfirsch-Schlueter current) not accounted for by the large aspect ratio model.
H-MODE EQUILIBRIUM AND MINIMUM ENERGY PRINCIPLE.

It has been shown that general variational principles lead to universal tokamak current and pressure profiles [3-5]. To describe JET H-modes we use a principle of minimisation for the plasma total energy:

\[ W = \int \text{d}V \left[ \frac{1}{2} \left( \nabla \psi \right)^2 + P \right]_{\gamma-1} \]

(1)

with the following prescription chosen in accord with observations. In (1) \( P(\psi) \) is the plasma pressure and \( F/R = B_\tau \) the toroidal field. Specifically we consider that: a) sawtoothing inside the mixing radius \( r = r_s \) is related to a unique helicity \( K = J \cdot A \cdot B \text{d}V \) for \( 0 < r < r_s \). The plasma energy \( W \) should be minimised with constant \( K \) (L-mode). b) in the confinement region \( r > r_s \) we use Kadomtsev minimisation of the total plasma energy \( W \) keeping the current in this confinement region constant (H-mode). In fact the constraints a) and b) are equivalent to that of keeping constant the total plasma current:

\[ I = \int \text{d}V \left[ \text{R} dP/d\psi + (F/R) dF/d\psi \right] \]

(2)

c) the equilibrium must be matched across \( r = r_s \). The equilibrium for \( r < r_s \) obtained from a) using a Lagrange multiplier \( \lambda \), and choosing \( \Lambda_z(r_s) = 0 \), implies:

\[ J_z = R B_z = \dot{J}_o; \quad P = P_o \]

(3)

both constant. In the confinement region \( r_s < r < a \), for large aspect ratio, and \( \beta_p - 1 \) the minimisation scheme b) leads to a Grad-Schlüter-Shafranov equation of the Liouville type:

\[ \text{div} \left( \nabla \psi / \lambda \right) = -Q \exp {\left( \psi / \lambda \right)} \]

(4)

In azimuthal symmetry the solution of (4) with \( \lambda = \beta_p \cdot \psi / B \tau \) leads to the profiles [3,4] which have the remarkable consistency property:

\[ q = q_o + H(r - r_s)(r^2 - r_s^2)/a^*; \quad J = J_o q_o^2 / q^2; \quad P = P_o q_o^2 / q^2 \]

(5)

with \( a^2 q_o = I_p / \pi J_o r_s^2; \quad (a/a^*)^2 = q_s - q_o \). The equilibrium profiles of pressure or current are characterised by the scale length \( a^* \) or the peakage parameter which is, for \( r_s / a^* \ll 1 \):

\[ <P>_P / P_o = q_o / q_a \]

(6)

For a given value of the parameter \( Q = (B^2 / \beta_p (a) / 2\mu_o) / P_o \) there are two possible solutions since:

\[ \left( <P>_P / P_o \right)^2 - <P>_P / P_o + Q / B = 0 \]

(7)

and the solution with larger scale length (larger \( <P>_P / P_o \)) could be interpreted as an ASDEX type H-mode which is sawtooth free with \( q_o \geq 1 \). In JET discharges however, \( q_a / q_o \) has remained constant \( \approx 4 \) through the L-H transition selecting a branch of equilibria within the mixing region relaxes to profiles with constant helicity.
EXPERIMENTAL DATA.

Fig. 2 and Fig. 4 show the actual current and pressure profiles for the L-H transition in shot 10766, plotted against the theoretical model of Eqs. (5). It appears that over most of the plasma cross section the experimental H mode profiles fulfill the consistency relations (5) even if the space dependence is different from that of the theoretical models. This behaviour is quite general in the H phase but it can also occur during the L phase. Fig. 4 shows the experimental temperature peaking $\langle T_e \rangle / T_0$ plotted versus $q_o / q_A$ for a large number of JET-H mode discharges. The experimental points appear to be bracketed between the limits (a) of the ohmic profile consistency principle $\langle T_e \rangle / T_0 - 3 q_o / 2 q_A$ and (b) the limit (6) for $\langle P \rangle / P_0$. This might be an indication of the equivalence of the two principles. Finally, we consider whether data agree with the model [6] of the H-mode as a transition of the outer plasma layers into the second stability region for ballooning modes.

A ballooning stability diagram in the plane $[\varphi / B_n, -2(R^2 / B_n) q^2 (dp/dr)]$ is shown in Fig. 5 for shot 10767 for a surface close to the separatrix. Fig. 6 shows the J profile consistency of shot 10767.

The L phase of the trajectory is all in the first stability region, and the L-H transition is only marginally approaching the second stability.

CONCLUSIONS.

The "diagnostic" integrals of motion I and K have been used to describe JET H-mode as states of constrained minimum energy. They appear to be one of two possible branches of stable finite $\beta_p$ equilibria, the other being sawtooth free as in ASDEX.

REFERENCES


Fig. 1 Pressure and $q$ profiles for shot number 14815 at $t - 14$ sec.
SPONTANEOUS TRANSITIONS IN THE TEMPERATURE OF A TOKAMAK PLASMA WITH SEPARATRIX

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1. INTRODUCTION

Tokamak experiments with a magnetic separatrix often exhibit spontaneous transitions in the edge electron temperature at the onset of an H-mode [1-3]. Such a transition might result, for example, from reduced edge losses due to the reduced recycling associated with the H-mode [4] or as a result of the edge plasma becoming stable again to ballooning modes in the second stable regime [5-6]. In the present paper, a two species plasma model is developed which exhibits a spontaneous transition in the edge electron temperature. The basis for the model is the critical electron temperature model [7] in which the anomalous transport disappears in the region of very high shear near a separatrix. The model has been shown to provide an acceptable interpretation of JET experimental data in terms of the global scaling of the stored electron energy and detailed profile analysis using transport code simulations [8,9]. In particular, a single formulation accounts for both the anomalous behaviour in ohmic plasmas and the degradation of confinement with additional heating (L-regime). The present work extends this formulation to the simulation of the electron temperature profiles and the improved energy confinement observed in H-mode plasmas in JET separatrix configurations.

2. TWO SPECIES PLASMA MODEL

Ion power balance

The ion power balance inside the separatrix may be written:

\[ P_i + \int n \nu_{ei} k (T_e - T_i) dV = Q_i \]  (1)

where \( P_i \) is the power directly heating the ions, \( n \) is the plasma density, \( T_e \) and \( T_i \) are the electron and ion temperatures and \( \nu_{ei} \) is the energy equipartition frequency. The ion heat flux, \( Q_i \), across the separatrix of surface area, \( S \) is:

\[ Q_i = - \int n_i \nu_k T_i dS = n_i \frac{k(T_{ix} - T_{ia}) S}{\lambda} \]

where \( \nu_k \) is the ion heat diffusivity, and \( T_{ix} - T_{ia} \) is the ion temperature difference over a distance \( \lambda \) in the vicinity of the separatrix.
To determine fully the energy equipartition term in equation (1) requires a full transport calculation of the temperature profiles. In the present analysis, this term is approximated by \( n v e i k(T_{e x} - T_{ix}) V \), where it is assumed that the integral is dominated by a contribution over a volume \( V \) near the separatrix. The ion power balance can then be written in terms of \( T_{ex} \) and \( T_{ix} \):

\[
P_i + n v e i k(T_{e x} - T_{ix}) V = n \chi_i k(T_{ix} - T_{ia}) S / \lambda
\]

and the ion heat flux is given by:

\[
Q_i = P_i + n v e i k T_{e x} V + n \chi_i k T_{ia} S / \lambda
\]

Electron power balance

The electron power balance at the separatrix may be written:

\[
Q_e = -\int n v e i k T_{e x} dS = n \chi_e k(T_{e x} - T_{ea}) S / \lambda
\]

where \( \chi_e \) is the electron heat diffusivity and \( T_{ex} - T_{ea} \) is the electron temperature difference over the distance \( \lambda \).

Total power balance

The total power balance at the separatrix can be expressed in terms of \( Q_e \) alone by eliminating \( T_{ex} \) from equations (2) and (3):

\[
P = Q_e + Q_i = Q_e + P_i + \frac{n v e i e V(T_{ea} + Q_e / S n x e) + n \chi_i e T_{ia} S / \lambda}{1 + v e i \ V \lambda / \chi_i S} - \frac{n \chi_i e T_{ia} S / \lambda}{\lambda}
\]

With the simplifying assumptions that the edge temperatures \( (T_{ea} \) and \( T_{ia} \)) and power to the ions \( (P_i) \) are zero and with \( v e i = n v_0 / T_{ex}^{3/2} \), equation (4) may be written in the form:

\[
P = Q_e + \frac{Q_e}{Q_e^{3/2} + q}
\]

where \( P = P/p \), \( Q_e = Q_e/p \), with the power normalisation, \( p \), and the parameter, \( q \), defined as:

\[
p = (n^2 v e V(n e S x e / \lambda)^{3/2})^{2/3}
q = \chi_e / \chi_i
\]

3. INTERPRETATION OF THE RESULTS

Given the electron and ion heat diffusivities the model determines the electron power flow, \( Q_e \), and hence the edge electron temperature (equation (3)) for a given power flow, \( P \), across the separatrix. The general solution of equation (5) is shown in Fig.1 for \( q \) in the range 0.0001 \( \leq q \leq 1.0 \). For \( \chi_e > \chi_i \), \( Q_e \) and the edge electron temperature increase monotonically with \( P \) (q = 1.0 curve in Fig.1). With decreasing \( \chi_e / \chi_i \),
the rate of increase of $Q_e$ with $P$ is reduced at low $Q_e$ and the ions transport a larger fraction of the total input power ($q = 0.1$ curve). For $\chi_i >> \chi_e$ the ions transport a substantial fraction of the total input power ($q = 0.0001$ curve).

Of course the ion power flow is itself dependent on the electron temperature (through energy equipartition). As a result, for certain input powers, three solutions can exist (eg, $q = 0.01$ curve), two of which are stable ($S_1$ & $S_2$) and one unstable ($U$) to perturbations in the edge temperature. Under these conditions a transition between the two stable solutions will manifest itself in a spontaneous change in the electron temperature and the electron power flow across the separatrix becomes significant. These solutions occur above a certain power threshold.

4. RELEVANCE TO THE JET H-MODE

The transport models of [8,9] lead to temperature profiles in fair agreement with experiment. Inside the separatrix the electron temperature profile is determined largely by the critical temperature gradient and $\chi_i - \chi_e$. Under these conditions $q$ $-1$ and $p$ $-9$MW, typically. For an input power of $5$MW, $Q_e$ $2.7$MW and the edge temperature $40$eV.

Near the separatrix, however, the heat diffusivities will be reduced due to the local magnetic shear, ultimately being limited by neoclassical transport. $\chi_i$ will change less than $\chi_e$. For typical JET conditions ($q$ $0.03$ in Fig. 1 and $p$ $3$MW) a power in excess of $5.4$MW is needed to make the transition from a low edge temperature ($-250$eV) with ion transport dominating the power flow to a high edge temperature ($-1$keV) with electron transport becoming more important. Such a spontaneous change in the electron temperature is identified with the pedestal observed to develop within $-0.05-0.1$m of the JET separatrix when a transition is made from the L-H regime at an input power, $P$ $5$MW.

The simplified 1-D electron energy transport code [8,9] has been used to simulate the electron temperature profiles of JET Pulse #10755. The electron transport model presented at the Varenna Theory Conference (but with the anomalous diffusivity multiplied by the shape parameter $(R/r)(\sqrt{V_T_T_e/r_e})(T_1/T_e)^2$ [9]) has been used in conjunction with the boundary condition given above for the edge electron temperature. The density, auxiliary heating and radiated power profiles are represented by analytical functions based on experimental data. Neoclassical corrections to the Spitzer resistivity are accounted for by a radial dependence in $Z_{eff}$. The electron temperature, the resistivity and the current density are assumed to be constant within the $q=1$ surface. The computed and experimental profiles show satisfactory agreement (Fig.2) in the ohmic, L- and H-phases. It should be noted that the same transport equations are used to simulate all phases of the discharge.

5. SUMMARY

The two species plasma model developed exhibits a spontaneous transition in the electron temperature in the high shear region in the edge of a tokamak plasma with separatrix where the transport is ultimately neoclassical. The transition occurs above a certain power threshold,
predicted to be about 5MW for JET, in accord with the power threshold for transition to the H-mode. Furthermore, the electron temperature profiles and energy confinement in ohmic, L- and H-phases of discharges in JET separatrix configurations are well simulated by the critical electron temperature gradient model, which contains intrinsically a degradation of confinement with additional heating.

References

Fig 1: Solution to equation (5) for various values of the parameter q.

Fig 2: Experimental (—) and computed (—) electron temperature profiles at times 10.5s (OH), 11.7s (L-phase, P=4MW) and 12.7s (H-phase, P=8MW) during JET pulse #10755.
**Introduction**. Spatial profiles of soft X-ray intensity \((E \geq 1.75 \text{ keV})\) are measured on JET using one vertical and one horizontal X-ray diode array (100 detectors total). Using tomographic reconstruction techniques spatial distributions of local X-ray emissivity are then derived. Remarkable differences in the X-ray emissivity distributions have been observed depending on the type of plasma discharge. X-ray distributions measured during limiter discharges usually have a gaussian shape. In deep contrast, extremely peaked profiles of X-ray emissivity, with most of the emission coming from a narrow central region, have been observed after pellet injection and flat or even hollow emissivity profiles have been observed during the H-mode of X-point discharges. In part 1 of this paper simulations which reproduce well both the shape and absolute emissivity of the measured profiles are presented. They show that peaked or hollow emissivity profiles reflect similar changes in the density distributions of the different species.

The vertical X-ray camera which directly views the trajectory of pellets injected into the plasma shows very intense bremsstrahlung emission from the interactions of plasma electrons with pellet particles. In part 2 of this paper a model which accounts for the main aspects of the observations will be discussed.

1. **Simulation of emissivity profile**.

**Peaked profile** – Fig. 1 shows the X-ray distribution before \((t=4.44 \text{ s})\) and after \((t=7.52 \text{ s})\) the injection of two D pellets (4 mm pellet at 4.5 s and 2.7 mm pellet at 5.5 s) into an ohmically heated plasma \((\# \ 13572, I=2.5 \text{ MA}, B=2.8 \text{T})\). The first pellet, cooling the plasma, allows the second pellet to penetrate deep into the plasma, leading to a dramatic increase of the central emissivity and consequently to a strong peaking of the X-ray distribution. Very similar observations made on other experiments have already been reported\(^{1,2}\).

![Fig. 1 - X-ray emissivity distribution before and after pellet injection.](image-url)
Fig. 2 shows the radial emissivity profile of the peaked distribution \((t = 7.52 \text{ s})\) together with two different simulations using the radiation code IONEQ. Taking into account the transmission of the filter the code calculates for each radiating species a radial profile of emissivity and these are then summed up to yield a total profile. The radial charge state distribution for each species is calculated assuming coronal equilibrium and using the impurity concentrations derived from PHA spectra and the profiles of \(T_e\) and \(n_e\) measured with the LIDAR Thomson scattering diagnostic. The measured profile of \(n_e\) is also very peaked. In the first simulation (curve 1) which reproduces well the shape and absolute emissivity of the measured profile the density distributions of the different plasma species are assumed to be as \(n_e\). The calculation shows that the observed emission is mainly due to free-free radiation from fully ionized C and O. In the second simulation (curve 2) the radial density distribution of D is taken as \(n_e\) but a much broader pre-pellet distribution is used for the C, O, Cl and Ni impurities. A peaked D density profile is to be expected as the second pellet deposits, in this case, a large part of its particles in the plasma center. The peaking of the impurities is more surprising and indicates an increase of the inward convection velocity. This result can be understood according to the neoclassical theory which predicts that the large density gradient of the D ions should drive the impurities inward. However since the neoclassical theory has not been adequate to describe the majority of experiments this interpretation should be looked at with caution and more work is needed to confirm it.

**Hollow profile** – Hollow X-ray profiles are observed in certain X-point discharges correlated with hollow electron density profiles. Fig. 3 shows such X-ray emissivity profile measured during an H-mode together with two simulations. The simulation which reproduces the measured profile well (curve 1) was obtained using the hollow electron density profile measured by interferometry and assuming for the other plasma species the same shape of density profile. A second simulation (curve 2), assuming flat density profiles over the central region, was run to verify that the hollowness in the X-ray profile was not due to line radiation from Ni or Cl but rather a direct consequence of the hollow densities.

![Fig. 2 - Profile of X-ray emissivity after pellet injection and two simulations assuming peaked or broad density profiles for the impurities.](image)

![Fig. 3 - Measurement and simulation of hollow X-ray emissivity profile during an H-mode.](image)
2. X-ray ablation emission. Detailed observations on D pellets injected into JET have been made with the vertical soft X-ray camera which directly views the pellet trajectory. Strong emission (Fig.4) is seen from the pellet-plasma interaction region due to bremsstrahlung radiation from collisions between the plasma electrons and the ablated pellet particles. The observations have been used extensively to determine the pellet velocity and depth of penetration. The time dependence of a single channel has shown that the emission originates from within a region with a diameter $2 r_c = 7 \text{ cm}$ in the major radial direction and $r_c$ is taken as the critical radius at which the ablatant flows along the magnetic field lines. The length of the hose of ablatant along the field lines was determined from toroidally spaced X-ray detectors as much less than 2 m. Measurements with different Be filters showed that the effective temperature of the electrons in the region of the pellet was close to that of the plasma before pellet injection.

![Fig. 4 - The ablation of a pellet as seen by the vertical X-ray camera.](image)

The pellet ablation model of Parks and Turnbull\textsuperscript{4} has been used to calculate the absolute intensity of the X-rays. For the shots considered here the model also gave an accurate prediction of the pellet range although a more complicated model\textsuperscript{5} is generally required to predict the pellet range in JET. In this model the incident plasma electrons interact with the ablated particles which expand spherically with high density and low temperature near the pellet surface. In a later refinement to the theory\textsuperscript{6}, the ablated particles flow along the magnetic field lines at $r_e = 2.5 r_p$, but we prefer to keep $r_e$ as a free parameter as the measurements show a much larger (7×) value. Inside $r_c$ we use the analytic asymptotic solution of PT but outside we have found the corresponding solution for cylindrical geometry. The temperature of the incident electrons differs significantly from $T_e$ close to the pellet surface only. As the X-ray emission comes from much larger radii it is reasonable to use the plasma value of $T_e$ everywhere. The ablatant density is

$$
\rho = \rho_e \frac{1}{a^{1/3}} \left( \frac{r_e}{r} \right)^{7/3}, \quad r < r_c; \quad \rho = \rho_e \frac{1}{a^{1/3}} \frac{r_e^{7/3}}{r_c^2 z^{1/3}}, \quad r > r_c;
$$

where $a^{1/3}$ is a constant near unity, $r_e$ and $\rho_e$ are the radius and the density at the sonic surface and $z$ a variable along the field lines. As the velocity of the ablated particles is much less than the plasma electrons the X-ray emission may be calculated from the hydrogenic bremsstrahlung formula

$$
P = 1.69 \times 10^{-32} \frac{P}{m} n_e \sqrt{T_e} e^{-E_c/T_e} \quad (\text{Wcm}^{-3})
$$

where $E_c$ is the energy cutoff. Integrating over the field of view (Fig.5) gives
\[ P_x = 7.9 \times 10^{-38} \frac{D}{m} r_e^{2/3} \left\{ \frac{(D/2)^{2/3} - r_e^{2/3}}{2} \right\} n_e \sqrt{T_e} e^{-E_e/T_e} \quad (W) \]

This expression (accurate to \( \sim 20\% \)) depends only weakly on \( r_p \), confirming that there is little emission from the region near to the pellet. \( P_x \) has been compared with experiment (Fig. 6) for the case of a 3.6 mm D pellet injected into a D plasma with \( I = 3 \, \text{MA}, \, B = 2.8 \, \text{T}, \, n_e = 1.35 \times 10^{19} \, \text{m}^{-3} \) and \( T_e = 4.8 \, \text{keV} \). The calculated emission generally exceeds the observed value in the outer part of the plasma, but decreases to well below the observed value at the end of the pellet trajectory. The observation that \( P_x \) always increases sharply at the end of the trajectory is not adequately explained by this model.

![Fig. 5 - Schematic of the ablating plasma within the field of view of an X-ray detector (D x l). For \( r > r_e \) the particles flow along the field lines (z direction).](image)

![Fig. 6 - Comparison between measured (each data point corresponds to a particular detector) and calculated soft X-ray power versus radial position.](image)

In conclusion: (1) the X-ray measurements support the concept of the formation of a plasma hose, but its radius is larger than predicted and it is quite short (< 2m). (2) the PT model provides an approximate estimate of the observed emission except towards the end of the pellet's range where the calculated value is much too small, (3) the X-ray energy spectrum is approximately as expected.

**References**

ANOMALOUS TRANSPORTS AND RELAXATION OF TOROIDAL ROTATION IN PLASMA

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In spite of the intensive studies of anomalous transports in tokamak plasmas, these phenomena are still far from complete comprehension. The main objective of a given paper is to classify anomalous transport on a basis of the known theoretical and experimental results, to single out characteristic spatial and temporal scales and, proceeding from dimensional considerations, to derive unknown transport coefficients.

One can find two different mechanisms of anomalous transport. The first one is related with a turbulent stochastization of field lines on a collisionless skin-layer scale, $C/\omega_{pe}$, due to non-zero collisionless Landau dampint, $[1] - [3]$. Turbulences on such transversal scales can be represented as a set of dual vortices strongly-extended along the field lines and moving along them with a thermal electron velocities, $\nu_e$. Their amplitudes should satisfy the condition $e\gamma/e = 1$ $[4]$. The characteristic scale of electron displacements across the field, $\Delta r \sim C/\omega_{pe}$ as $\nu_e \ll \gamma/\omega_{pe}$. The frequency of their collisions with the vortices is $\omega_{ef} \sim eK_n\nu_e \sim e\nu_e/(\gamma R)$. The factor $\varepsilon$ allows for an interaction of electron not with a wave resonant at a given surface, for which $K_{II} = 0$, but with a satellite, for which $K_{II} \sim (qR)^{-1}$ and the amplitude is $\varepsilon$-times lower $[2]$. From considerations of dimensionality one can obtain the electron heat conduction and electron diffusion coefficient

$$D_{\alpha n} = \chi_{\alpha n} = \Delta R^2 \omega_{ef} = \frac{\varepsilon^2}{\nu_{pe}^2} \frac{\nu_e}{\gamma R} \varepsilon$$

(1)
The motion of ions as such low-scale vortices \( \left( \frac{\gamma}{c} > \frac{1}{\omega_p e} \right) \) will be slightly perturbed, except of the collisions with vortex top which can reject an ion. A turbulence model for ions can be represented as a set of scattering centres with a distance between them equal to \( qR \) in a longitudinal direction, and to \( c/\omega_p e \) in a transversal direction. A characteristics ion displacement, \( \beta_s \), characteristic frequency of collisions with vortices is the same as that for electrons. Diffusion, heat conduction and viscosity coefficients for ions are estimated as:

\[
D_{an}^i \sim \nu_{an}^i \sim \beta_{an}^2 \sim \frac{\nu_\parallel}{qR} \epsilon, \quad \beta_{an}^2 = \frac{T_e + T_i}{m_i \omega_{Bi}}
\]

Note that ions are in the Pfirsch-Schlüter regime, \( \omega_{\parallel i} > \frac{\nu_\parallel}{q^2} \epsilon^{1/2} \). Therefore, one should probably add a factor \( (1+q^2) \) to the expression (2), and the effects of interaction between the trapped particles and the drift waves with \( K_\perp \sim \beta_s^{-1} \), described below, are possible for electrons only.

As the plasma diffusion is determined by a lower of two diffusion coefficients, \( D_i \) and \( D_e \), one has \( D_{an} \approx D_{an}^e \).

Similar qualitative estimates have a sense only when they are confirmed by more strict theoretical and experimental results. The electron heat conduction (1) was more accurately obtained in \([1] - [3]\) and fits well the experimental Merezhkin-Mukhovatov scaling law or the neo-ALCATOR one. An anomalous viscosity calculated in \([5]\), \([6]\) coincides with the viscosity (2) and describes well not only the experiments, comparison with which is given in \([5]\), but the later ones \([7]\). Therefore one can hope that a new result for ion heat conduction (2) be also true. The problem of anomalous resistance is a more complicated one. Formally, from the dimensionality consideration one can write

\[
\sigma_{\parallel n} \sim \frac{\omega_p^2}{\omega_c^2} \sim \omega_p^2 \frac{qR}{v_e' c}
\]

In a one-dimensional case, a passing particle flying over a potential hill loses, first, its velocity and then picks it
up to a previous value. Such a particle does not lose its momentum, and the estimate (3) is not true. In case of interaction with vortices, an electron, approaching the hill top, is not only decelerated but shifted in a transversal direction and can be removed beyond the limits of a given vortex to a distance and no more be accelerated up to the previous velocity. In this case, the estimate (3) turns out to be true. Such an anomalous conduction in some regimes in tokamaks turns out to be lower than a classical one. Their ratio has the form $G_{el} / G_{an} \sim c \lambda Z_n Z_y / (e^{\gamma Z n y}) (\lambda$ is the Coulomb logarithm, $Z$ is the effective charge). The estimate (3) needs an accurate theoretical and experimental verification.

The second transport mechanism is an interaction of trapped electrons with larger scale drift waves $k_1 \sim \theta_i^{-1}$. The spectral maximum in drift oscillations observed in a tokamak refers to this region. Passing electrons in such waves are shifted almost reversible [2]. Trapped electrons interact with the waves in a resonance way at a frequency of toroidal precession, $\omega_\perp = \omega_\perp \frac{\theta_i^2}{2} \omega_\perp$ (where $\omega_\perp$ is the drift frequency). Their displacement turns out to be irreversible. By the order of magnitude one can write $D_{an}^{el} \sim \theta_i^{1/2} \omega_\perp / \omega_\perp \approx \theta_i^{1/2} / \omega_\perp$. A characteristic velocity can be estimated as $\omega_\perp \theta_i^{1/2}$; a frequency, as $\omega_\perp$ at $\nu_\perp \sim \nu_\perp$. As a result, one has

$$D_{an}^{el} \sim \gamma^{el} \sim \frac{\rho_s \nu_\perp}{c_l^2 \nu_\perp^2} \nu_\perp^{1/2}$$

Such a coefficient has been obtained from a quasi-linear approximation in [8]. Pay attention to a temperature pinching effect related with a strong dependence of $\omega_\perp$ on $\nu_\perp$ [9]. A flux of heat through an electron channel and that of particles can be directed towards an increase in $T_\perp$. Anomalous transport by trapped electrons does not depend on density and can play a decisive role at high densities, resulting in a saturation of the $T_\perp$-growth with a rise in $n$. One dimensional set of transport equations has the form

$$\frac{\partial n}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} r \frac{\partial}{\partial r} \frac{1}{\nu_i} n / \nu_i = (D_{an}^{el} - D_{an}^{el}) \frac{\nu_i}{c_l^2 \nu_\perp} \nu_i \left( \frac{D_{an}^{el}}{c_l^2 \nu_\perp} \right) \frac{\nu_i}{c_l^2 \nu_\perp}$$

(5)
\[ \frac{3}{2} \frac{\partial}{\partial t} (n T_e) = \frac{1}{r} \frac{\partial}{\partial r} r \left[ (\chi_n^{e} + \chi_n^{i}) \frac{\partial n T_e}{\partial r} + \frac{n}{T_e} \left( \frac{3}{2} \chi_n^{e} - P_r \chi_n^{i} \right) \frac{\partial n}{\partial r} \right] + \Delta e_i + Q_e \]  

\[ \frac{2}{\partial t} (n T_i) = \frac{4}{r} \frac{\partial}{\partial r} \chi_n^{i} \frac{\partial}{\partial r} (n T_i) \]  

\[ \frac{\partial u_i}{\partial t} = \frac{1}{r} \frac{\partial}{\partial r} r \left( \frac{u_i}{n} + \nu_n^{i} \right) \frac{\partial u}{\partial r} + \frac{\Gamma}{m_i n} \]  

Here, \( U \) is the toroidal rotation speed, \( P_N \) and \( P_T \) are the pinching parameters for particles and heat, \( 1 \leq P_N \leq 3/2 \), \( 0 \leq P_T \leq 3/2 \), \( \Delta e_i \) is the heat exchange between electrons and ions, \( N, Q_e, Q_i \) and \( F \) are the sources. Convective removal of ion heat is neglected here.

REFERENCES

CURRENT TRANSPORT IN A CHAOTIC MAGNETIC FIELD AND SELF-SUSTAINMENT OF ISLANDS

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Introduction

In previous work [1], it has been proposed that the anomalous transport observed in Tokamaks could be a consequence of the magnetic topology prevailing in such devices (i.e., the existence of small magnetic islands surrounded by regions where the field lines have an ergodic behaviour). The consequences for current transport and self-sustainment of the topology are considered in this paper. In particular, it is found that the diffusion of the current liberated at each sawtooth crash can sustain the topology between the inversion radius and the plasma boundary, thus extending the conclusion of Ref. [2] to most of the plasma volume.

Trajectories of Magnetic Field Lines in Cylindrical Coordinates (r, \( \theta \), R\( \phi \))

We suppose that the magnetic field can be split into two terms: a dominant term \( B_0 \) with cylindrical symmetry (\( B_{0r} = 0 \)) and a smaller term \( B_i \), such that \( B_i^2 \ll B_0^2 \). \( B_i \) is responsible for chaotic behaviour of the magnetic field lines.

The following equations could be extended to any geometry by a suitable choice of coordinates.

We define

\[
\psi = \int_0^r r' B_{0\phi} \, dr' = \int_0^r r' B_{\phi} \, dr
\]

(1)

\( \psi = \text{const} \) is the magnetic surface of the unperturbed field and \( L \) is the length along a magnetic field line: \( d\phi = (B_\phi / BR) \, dL \).

The trajectories of a field line are given by

\[
dr = (RB_e / B_\phi) d\phi, \quad d\theta = (RB_e / B_\phi) d\phi
\]

(2)

which defines a transformation \( M = (M_\psi, \phi) \) (for \( \phi = 0; \ M_\psi = M_\phi \)). In particular: \( \psi = \psi (\psi_0, \Theta_0, \phi) \) or inverting the previous expression \( \Theta_0 - \Theta_{0i} (\psi_0, \psi, \phi) \); \( \Theta \) could be a multivalue function (index \( i \)). In the same way, inverting initial and current positions along a magnetic line: \( \Theta = \Theta_i (\psi, \psi_0, - \phi) \).

Flux conservation imposes: \( d\psi \, d\Theta = d\psi_0 \, d\Theta_0 \), which can be written:

\[
\frac{\partial \Theta_i}{\partial \psi} = \frac{\partial \Theta_{0i}}{\partial \psi} (\geq 0)
\]

(3)

and

\[
\int_0^\infty d\psi_0 \left( \Sigma_i \frac{\partial \Theta_i}{\partial \psi_0} \right) = \int_0^\infty d\psi \left( \Sigma_i \frac{\partial \Theta_{0i}}{\partial \psi} \right) = 2\pi
\]

(4)
Let us define: 
\[ F(\psi, \psi_0, \phi) = \frac{1}{2\pi} \sum_i \frac{\partial \theta_{oi}}{\partial \psi} \]  

**Invariance along a magnetic line:**
A function \( f \) constant along a magnetic line satisfies the equation:
\[ \mathbf{B} \cdot \nabla f = 0 \]
which gives
\[ \frac{\partial f}{\partial \phi} + r \mathbf{R} \frac{\partial f}{\partial \psi} + \frac{1}{r \mathbf{R}_\psi} \frac{\partial f}{\partial \theta} = 0 \]  

\( \theta_{oi} \) is invariant along a magnetic line: (initial condition)
\[ \sum_i \frac{\partial}{\partial \psi} \left[ \frac{\partial \theta_{oi}}{\partial \phi} + r \mathbf{R} \frac{\partial \theta_{oi}}{\partial \psi} \right] = 0 \quad \text{or} \quad \frac{\partial F}{\partial \phi} + \frac{1}{r} \left( \sum_i r \mathbf{R} \frac{\partial \theta_{oi}}{\partial \psi} \right) = 0 \]  

In a chaotic region, the number of value 1 increases equally over the whole region as \( \phi \) increases and as a first term
\[ \sum_i r \mathbf{R} \frac{\partial \theta_{oi}}{\partial \psi} \rightarrow -r K(\psi) \frac{\partial F}{\partial \phi} \]  
and
\[ \frac{\partial F}{\partial \phi} = \frac{1}{r} \left( r K(\psi) \frac{\partial F}{\partial \psi} \right) = 0 \]  

\( F \) is the probability of finding a magnetic line starting at \( \psi_0 \) in an interval \( d\psi \) around \( \psi \) after a number of turns \( = (\phi/2\pi) \).

Eq. (9) is a diffusion equation of magnetic lines to be compared with the similar equation in [1]
\[ r K(\psi) = \frac{A^2 r}{2 L_0} \left( \frac{d\psi}{dr} \right)^2 \]  
where \( A \) is the distance between two chains of islands and \( L_0 \) the mean length of magnetic line to cross one chain of islands. If \( m \) is the poloidal mode number and \( \mu \) the bandwidth
\[ \Delta = (3/2 \mu m^2) \left( \frac{d\psi}{dr} \right)^{-1}, \quad L_0 = 2\pi R (2\mu m/3) g(\gamma) \]  
with \( \gamma \) the overlapping coefficient: \( \gamma = (\epsilon_1 + \epsilon_2)/2A \) and \( g(\gamma) = 1 \). For \( \gamma = 1.2 \) \( (g' < 0) \) with these definitions, where \( r = 1/q \) is the rotational transform
\[ r K(\psi) = \left[ 0.27/\mu^2 m^2 g(\gamma) \right] \left( \frac{d\psi}{dr} \right)^2 \]  

**Application to the current**

The current flowing along the field line obeys in the first approximation
\[ \mathbf{B} \cdot \nabla J_n = 0 \]  

We shall suppose that this is true along \( N \) turns of the torus.

The diffusion of the magnetic field lines across the magnetic surfaces leads to a transport of current. We define a mean resistivity \( \bar{n} = V/2\pi RJ_\phi \)
so that
\[ \bar{n} B_\phi \frac{V}{2\pi R} = \frac{V}{2\pi R} \int_{-\pi N}^{+\pi N} n(\psi) B_\phi \frac{d\phi}{2\pi} d\phi \]
where \( V \) is the voltage per turn. Eq. (13) leads to \((J_n/B) - (J_\phi/B_\phi)\) is constant along a field line and follows Eq. (9):

\[
\frac{\partial}{\partial \psi} \left( \frac{B_\phi}{J_\phi} \right) + \frac{\partial}{\partial \psi} \left( \frac{B_\phi}{J_\phi} \right) \frac{B_\phi}{2\pi R} \frac{B_\phi}{J_\phi} = 0
\]

which taking into account Eq. (14) gives

\[
2\pi N \frac{\partial}{\partial \psi} \left( \frac{B_\phi}{J_\phi} \right) \frac{B_\phi}{2\pi R} \left( \frac{B_\phi}{J_\phi} \right) = \left( \frac{B_\phi}{J_\phi} \right)
\]

where \((B_\phi \eta/R)\) is a mean value over a magnetic surface \( \psi \).

The conditions at the limits are that the flow of current is null:

\[
\frac{\partial}{\partial \psi} \left( \frac{B_\phi}{J_\phi} \right) = 0 \quad \text{for} \quad \psi = 0 \quad \text{and} \quad \psi = \psi_a
\]

where \( a \) is the radius of the plasma.

By integration, Eq. (16) gives

\[
\int_a f_a \frac{B_\phi}{J_\phi} \eta(\psi) B_\phi 2\pi R d\phi = \int_a f_a \frac{B_\phi}{J_\phi}
\]

Conclusions

Several conclusions can be deduced from these equations:

(a) When \((\pi N R^2/L_o a^2)\) is large compared to 1 (small \( m \) numbers and larger islands), Eq. (16) becomes:

\[
\frac{B_\phi}{J_\phi} = \text{const} = \frac{1}{\psi_a} \int_a f_a \psi(\psi) B_\phi 2\pi R d\phi
\]

which give the Taylor condition for reverse field pinch [3]. The next order term can be deduced from Eq. (16). Eq. (19) could also describe the state of the plasma after a disruption in a tokamak; nevertheless, the transport of current will take some time to allow the eddy current in the islands to die away. The effect on the heat transport should be much faster.

(b) When \((\pi N R^2/L_o a^2)\) is small compared to 1 (the \( m \) number is large and the size of overlapping islands is small), we can define

\[
\delta J_\phi = J_\phi - \frac{V}{2\pi R N} = 2 \frac{\pi N}{B_\phi} \frac{\partial}{\partial \psi} \left( \frac{B_\phi}{J_\phi} \right)
\]

which produces an extra current in the chaotic region if

\[
\frac{\partial}{\partial \psi} \left( \frac{B_\phi}{J_\phi} \right) > 0
\]

The critical profile in this case for large \( m \) numbers could be written with the hypothesis that \( B_\phi \) and \( RK \) are constant:
\[ \frac{\partial}{\partial r^2} (J_\phi) = -\alpha J_\phi^2 \]

which could be integrated to give:

\[ J_\phi = J_0/(ar^2+1) \]  \hspace{1cm} (22)

The self-consistency of the islands if \( m \) is lower but not too small (if not the tearing condition has to be taken into account \([2,4]\)) is given in Ref.\([1]\) (Ampère's Law).

\[ L_\phi \mu_0 \delta J = B_\phi \]  \hspace{1cm} (23)

which linked to Eq.\((23)\) leads to the self-consistency condition; if the islands are sustained only by the transport of current, this imposes a curvature of the current greater than that given in Eq.\((22)\). In Eq.\((20)\) defining \( \delta J \), only current diffusion has been taken into consideration but other effects can play a dominant role such as a contribution of fast particles. Thermodynamic effects linked to a flow of plasma along the chaotic field lines may also sustain the islands. Local electric fields and non-ambipolar diffusion may be important. In this presentation, the estimation of \( N \) remains a problem.

\( \delta \) the current created between the two internal disruptions could be redistributed with this mechanism along the chaotic field line: at the edge of the \( q = 1 \) surface, an extra production of current \( \delta J_8 \) could be calculated:

\[ 2\pi \psi(\delta J_8/B_\phi) = \delta I_8 = J_0^2 \int \left[ (\eta(\psi) - \eta \psi) 4\pi^2 R \right] / Vd\psi \]  \hspace{1cm} (24)

This current could assist in maintaining the chaotic region.

References


DENSITY FLUCTUATIONS IN FT TOKAMAK

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Measurements of density fluctuations in FT, by means of infrared (IR) scattering, are presented. The data are taken along a vertical diameter of the plasma poloidal cross section and compared with results previously obtained on a chord close to the edge. Frequency and K-spectra show similar characteristics for poloidal and radial propagation. The scattering signal is recorded during the whole shot interval and usually follows the density evolution. Additional heating power with lower hybrid radiofrequency (LHR) (8 GHz) does not produce any evident effect on the signal. A modification of the fluctuations spectral density is observed when marfes propagate into the diagnostic view line.

1. EXPERIMENTAL SET UP AND PARAMETERS RANGE.

The scattering system is essentially the same as that described in Ref. [1]. A 3 W, 10 μm, gaussian CO₂ laser beam is weakly focused by a concave mirror to a beam waist of 5 mm in the plasma center; a portion of it (150 mW), after an external path, is

![Graph](image)

Fig. 1 - Fluctuation level during flat top (a) and ramping (b) density discharges at 60 kG. LF channel: 10-30 kHz, HF channel: 260-360 kHz, both in independent au. The upper traces show the plasma current ($I_p$) and the average electron density ($n_e$).
sent onto a GeCu photoconductive detector for homodyne detection of the scattered power. The signal is analysed by bandpass RMS detectors, covering the frequency range 2-1000 kHz, and then recorded during the whole discharge. Measurements have been performed using the conventional scattering geometry, in the K-vector range 12-50 cm⁻¹, together with the far-forward technique in order to extend this range down to 5 cm⁻¹ (2). Even if both techniques have been used, a relative calibration of the two is not available yet. The scattering volume is a cylinder, with a 5 mm radius, covering a whole vertical plasma diameter: only poloidally propagating fluctuations are probed. At present only a shot by shot scan of the K-spectrum is possible.

2. RESULTS.

In Fig. 1a we report typical signals taken during a sawtoothing steady state discharge: after a fast increase, following the density build up, the fluctuations level, both in the high and low frequency channel, keeps practically constant during the flat top. A ramping density discharge is shown in Fig. 1b: here again the scattering signals follow the trend of the electron density apart from a tendency to saturate towards the end of the discharge. At 200 ms the effect of a strong gas puff is evident in particular in the high frequency channel.

Frequency spectra for two different K-vectors are reported in Fig. 2: there is no significant decrease of the spectral density at low frequency side with the resolution indicated by the error bars, while the spectral width is larger at higher K. The dashed lines indicate, for comparison, the roll off slope already found for fluctuations close to the edge (r/a= 0.8) mainly propagating in the radial direction (1). Both characteristics are reasonably unchanged when passing from one to the other

![Frequency spectra during flat top. n_e = 1.6x10^{11} cm^{-3}; I_p = 300 kA; B_T = 60 kG.](image-url)
configuration. In Fig. 3 frequency integrated power spectrum is shown vs K. Again the roll-off slope found in Ref. [1] seems to approximately fit the data.

No significant change in the fluctuations level was found during the injection of LHR power for electron heating up to 230 kW at 8 GHz.

![Figure 3](image-url)

**Fig. 3** - Frequency integrated spectral density vs K-vector. \(n_e = 1.6 \times 10^{14} \text{ cm}^{-3} \), \(I_p = 300 \text{ kA} \), \(B_T = 60 \text{ kG} \).

![Figure 4](image-url)

**Fig. 4** - Time evolution of scattering signals in presence of a mafse (lower traces in independent au). LF = 2–10 kHz; HF = 680–960 kHz; K = 5 cm\(^{-1}\). The He signal is in au.
Figure 4 shows the behaviour of density fluctuations in a discharge disrupting for density limit in which a marfe is present: when the cold dense region associated with the marfe propagates towards larger major radius the interferometer, viewing a vertical chord passing through the center, looses fringes and the fluctuations suffer a modification of their frequency spectrum as indicated by the different behaviour of the two different channels.

3. SUMMARY

The scattering signal, especially at low frequency, essentially follows the electron density evolution both in the flat top and ramping phase. Frequency and K-spectra of density fluctuations with poloidal K, integrated along a vertical chord passing through the plasma center, show similar characteristics to those with mainly radial K measured closer to the edge in a previous work. This confirms the K isotropization already found on other machines [3,4]. A modification of the frequency spectrum associated with marfe has been observed and is now under investigation.

REFERENCES

The magnetic diffusion in a plasma machine (e.g., TOKAMAK) plays an important role in understanding the characteristics of the initial phase current penetration (which appears to proceed anomalously fast) and the phase distribution of the plasma [1,2]. It is important to understand the nature and the scaling of the instabilities which might be responsible for these characteristics, since on the one hand, fast current penetration becomes more crucial in the large devices of the future generation and in reactors, while on the other hand distributions are accompanied by a local flattening of the electron temperature profile or rapid plasma losses.

In the present work, we consider a cylindrical inhomogeneous plasma and the magnetic field $B(r,t)$ in the plasma is considered to be an unknown function of space and time, while the temperature profile of the plasma is governed by the linear relation:

$$ T = T_0 (1 + \delta t) f(r) $$

where $T_0$ is the background temperature of the plasma, $\delta$ is a scaling factor, and $t$ is the time.

Starting from Maxwell's equations:

$$ \nabla \times \vec{E} = -\frac{1}{c} \frac{\partial \vec{B}}{\partial t}, \quad \nabla \times \vec{B} = \frac{\epsilon}{\mu} \vec{J}, \quad \vec{J} = \sigma \vec{E} $$

We can find the magnetic diffusion equation of the plasma as:

$$ \frac{\partial \vec{B}}{\partial t} = \frac{\partial}{\partial t} \left[ \frac{c^2}{\mu B} \frac{1}{r} \frac{\partial}{\partial r} \left( r \vec{B} \right) \right] $$

To solve equation (2) we considered that

$$ \vec{B}(r,t) = \vec{B}_v(r) J(t) $$

$$ \sigma = \sigma_{v0} \left( \frac{1}{T_0} \right)^{3/2} \sigma_{v10} \left[ (1 + \delta t) f(r) \right] $$
where $\vec{B}$ is the magnetic field of the plasma and $\sigma$ is the conductivity.

Substituting from (3) into (2), we get the following equations

$$\left(1 + \gamma t\right)^{3/2} \frac{1}{\dot{\gamma} \frac{\partial}{\partial \gamma}} = \frac{1}{\tau}$$  \hspace{1cm} (4)

$$\frac{C^2}{\gamma \kappa \sigma^2} \frac{1}{\dot{\gamma} \frac{\partial}{\partial \gamma}} \left[ \frac{1}{F^{3/2}(\gamma)} \right] \left( \frac{1}{\dot{\gamma}} \frac{\partial}{\partial \gamma} \left( r \vec{B} \right) \right) = \frac{1}{\tau}$$  \hspace{1cm} (5)

where $\tau$ is a constant.

It is easy to obtain the time dependent part of the magnetic field from (4) as

$$\mathcal{B}(t) = A_0 \exp \left[ -\frac{1}{2} \gamma \tau (1 + \gamma t)^{3/2} \right]$$  \hspace{1cm} (6)

We solve the differential equation (5) in two cases:

i) when $F(\gamma) = F_0$ is constant, we have

$$Z^2 \frac{d^2 \tilde{B}_0}{dZ^2} + Z \frac{d \tilde{B}_0}{dZ} - \tilde{B}_0 \left( 1 + Z^2 \right) = 0$$  \hspace{1cm} (7)

where $Z = \mathcal{F} \left( \frac{r}{\lambda} \right)^{1/2}$, $\mathcal{F} = \frac{1}{\lambda}$, $A^2 = \frac{2 \kappa \alpha^2 \sigma^2 \kappa}{\tau}$, and $\alpha$ is constant.

The solution of equation (7) is the modified Bessel function and have an approximate solution of the form:

$$\tilde{B} \propto I_{1+1} \left( \sqrt{A} \mathcal{F} \right)$$  \hspace{1cm} (8)

if $\tau \rightarrow \infty$ and $A \rightarrow 0$, we have $\tilde{B}_0 \propto \mathcal{F}$, which means that the current density is constant.

ii) When $F(\gamma) = \frac{1}{1 + D}$, $D = \mathcal{F}^2 / 2$ we obtain from equation (5) that:

$$Z^2 \left( 1 - Z^2 \right) \frac{d^2 \Phi}{dZ^2} + \frac{Z}{Z^2} \frac{d \Phi}{dZ} + A \Phi = 0$$  \hspace{1cm} (9)

where $Z^* = \mathcal{F} + 1$, $\Phi = \mathcal{F} \tilde{B}_0$

Equation (9) represent a hypergeometric function, and has the following solution [3].

\[ \mathcal{F} = M \left( 1 - A^2 z^* + A^* \right) \]  \hspace{1cm} (10)

It seems from the model assumption that \( T = T_0 \left( 1 + 8t \right) \) implies the conductivity \( \sigma \propto T^{3/2} \) is grow so the magnetic field \( B \) changes more slowly with time \( t \) leads to a longer time constant (seperation constant) therefore magnetic flux \( \mathcal{W} = \iiint \mathbf{A} \cdot \mathbf{B} \) will be decreased (or perhaps increase less rapidly) this apparently means that the magnetic flux or field lines move radially outward (or perhaps move inward more slowly) as the conductivity \( \sigma \) is grow the plasma becomes a superconductor (for \( t \to \infty \)) (the result of the model assumption of \( T \propto \left( 1 + 8t \right) \)).

References:


THEORETICAL INTERPRETATION OF TURBULENT FLUCTUATIONS AND TRANSPORT IN TEXT
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Abstract. We show that a magnetic fluctuation spectrum with physically reasonable properties can be constructed and used to account for observed potential and density fluctuations, ambipolar particle and heat fluxes obtained experimentally in TEXT.

Introduction. A wealth of experimental data has recently been obtained in the TEXT tokamak relating to fluctuations and anomalous transport (see Wootton at this conference and other references below). The availability of this data - especially in the confinement zone - leads to the possibility of a theoretical interpretation for TEXT. In an earlier paper [1] (see also [2] for a general discussion) we outlined an approach in which we were able to interpret the TFR data. The present paper, which gives an updated version of our theory, satisfactorily interprets many of the observed features of the TEXT data, as well as suggesting new effects which could be sought for.

We adopt a periodic cylinder geometry in which a uniform $B_z$ represents the toroidal field (2T for TEXT). The q-profile, which is taken to be $q(r/a) = 0.8 + 2.2(r/a)^2$, closely fits experiment as described by Wootton [3]. The electron and ion temperature profiles are given by

$$T_{oe,i}(r/a) = T_{oe,i}(o) \exp[-2.25(r/a)^2]$$

where $T_{oe,i}(o) = 0.63$ keV and $T_{oi}(o) = 0.65$ keV. It should be noted that the experiments being interpreted were carried out in hydrogen, with $Z_{eff} = 1.73$. The electron density is taken to be

$$n_{oe}(r/a) = n_{oe}(o) \exp[-1.3(r/a)^2],$$

where $n_{oe}(o) = 4.6 \times 10^{13}$ cm$^{-3}$. The assumed q-profile corresponds to a total plasma current of 208 kA and poloidal field at the limiter of 1.67 T.

In our model we assume that the fluctuating electric field $\vec{E}'$ can be written as

$$\vec{E}(r,t) = \nabla \phi - \frac{1}{c} \frac{\partial}{\partial t} A \left| \begin{array}{c} r \theta \xi \\ \end{array} \right|$$

The electron and ion distribution functions $f_{e,i}(r,v_{\parallel},t)$ are functions of $v_{\parallel}$, which is the particle velocity (random and mean) along the instantaneous magnetic field $\vec{b} = (B_0 + \vec{B})/B$. The functions $f_{e,i}$ satisfy Fokker-Planck equations similar to those derived in our earlier paper [1]. Given $\phi$ and $A_{\parallel}$, and assuming the smallness of $\vec{B}$, these equations are solved as before. Adding the quasi-neutrality condition to the system it is possible to find a relation between $\vec{E}'$ and $\vec{Q}'$. Thus the specification of $\vec{E}'$ for example enables the complete computation of quantities such as $\vec{E}'$, $\vec{Q}'$, and the fluxes $\Gamma_{e,i}(r)$, $\Omega_{e,i}(r)$. For example, $\Gamma_{e}(r)$ is given by

$$\Gamma_{e}(r) = <\int_{-\infty}^{\infty} v_{\parallel} f_{e} dv_{\parallel} > + <\int_{-\infty}^{\infty} \frac{cT_{oe} e \vec{E}'}{eB_z T_{0e}}\>$$

where $<$ denotes averaging over time-scales long compared with drift and collision times, as well as spatial averaging over $\theta$ and $z$. The solution once obtained is independently checked by substituting in the moment
equations derived from the kinetic ones. This self-consistency has been found to be essential in ensuring the ambipolarity of the particle fluxes calculated by our method, irrespective of the spectrum assumed for $S_r$ (or equivalently $\Psi$). The details of the method will be published elsewhere.

Results For calculational convenience we have chosen to prescribe the magnetic fluctuations, $S_r$, which we represent in the form of a Fourier series. The free parameters which occur are determined by comparing the calculated $\Psi(r,t)$ with experiment. The assumed form for $S_r$ includes a continuum of m,n modes. The modes with $\frac{m}{n} = 2$ and $\frac{m}{n} = \frac{5}{2}$, which we believe to be the dominant low order resonances, are treated separately from the continuum. Thus, we write

$$S_r = b_{rc} + b_{r,5/2} + b_{r,2/1},$$

where

$$b_{rc} = \sum_{m=1}^{60} \sum_{n=1}^{60} m^{-1/4} \exp \left[-\frac{1}{4} (m-nq)^2 \right] \cos (m\theta + nq + \omega_0 t),$$

with $\omega_0 = 2000$ rad s$^{-1}$ and $\epsilon = 1.35 \times 10^{-4}$; the summations run over m & n such that $\frac{m}{n} \neq \frac{5}{2}$ and $\frac{m}{n} \neq \frac{2}{1}$. The wavelengths and frequencies of this prescription are consistent with $k_\perp \rho_i < 0.3$ and $\omega_0 \sim \omega$. The quantity $b_{r,5/2}$ is given by the expansion

$$b_{r,5/2} = 10 \sum_{m=1}^{60} \sum_{n=1}^{60} \delta m^{-1/4} \exp \left[-\frac{1}{4} (m-nq)^2 \right] \cos (m\theta + nq + \omega_0 t),$$

where the $\delta$ represents the Kronecker symbol, and $\omega_0$ is

$$\omega_0 = \frac{eB_0 r_1 \rho_i m_j}{eB_j \rho_i r_2,1},$$

amplitude factor. The results arising from this spectrum will now be discussed.

The assumed q-profile indicates that the 5/2 resonance occurs close to $r/a = 0.9$ and the 2,1 resonance at $r/a = 0.75$. The 1,1 resonance, which is not treated explicitly, occurs at $r/a = 0.3$. Our calculation shows that $\langle (E_r)^2 \rangle_{1/2} = 1.4 \times 10^{-3}$ at $r/a = 0.9$, $1.0 \times 10^{-2}$ at $r/a = 0.75$, $6.0 \times 10^{-4}$ at $r/a = 0.3$, respectively. Fig.1 shows the calculated values of $\langle \frac{e\Phi}{r} \rangle_{1/2}$ and $\langle \frac{n}{n_0} \rangle_{1/2}$ plotted as functions of $r/a$ and compared with experiment. It should be noted that the value of $\frac{e\Phi}{r}$ at $r/a = 0.9$ has been used to fix the size of $\epsilon$.

The choices of parameters for the 5/2 and 2/1 resonances are made with reference to the data on the particle and the ion heat fluxes; the treatment of the 5/2 resonance is to ensure that the computed flux matches the radial variation of the measured flux. The treatment of the 2,1 resonance is to account for the observed size of the ion heat flux near this resonance. It must be noted that the experimentalists have apparently filtered the low m,n, low frequency signals from the data in Fig.1.

The code predicts the observed non-adiabaticity $\left( \frac{n}{n_0} < \frac{e\Phi}{r} \right)$ in
amplitude and phase; while this appears to be a generic feature of the kinetic equations, the investigation shows that the magnitudes and relative phases are sensitive to \( n, T, \) and \( q \)-profiles. Note that the calculated fluctuations include the continuum as well as the discrete resonances and are therefore somewhat larger than experiment. In Fig. 2 we compare the calculated particle flux with the measured one. There is qualitative agreement in the radial variation. We note that our theory cannot be used for \( r/a > 1 \) since it does not cover edge-specific effects. It is difficult to say whether the \( 5/2 \) resonance at the assumed amplitude can account for all of the observed particle flux. We can say definitely, that if the \( 5/2 \) resonance is not included, although the non-adiabaticity can be accounted for in terms of \( \Sigma_{cc} \) alone, the particle flux from these modes would not match experiment. According to Schoch et al [4] the low \( m,n \) modes have been filtered out of their \( \Sigma \) spectrum - our results suggest that these modes could be important for particle transport.

Fig. 3 gives the comparison between the calculated and measured total electron heat fluxes. Considering the size of the error bars the agreement is acceptable both in radial variation and magnitude. Although the figures don't show it, theory predicts large fluctuation and heat fluxes at the \( 1,1 \) resonance, and it would be interesting if those could be determined experimentally and compared with theory. We have also checked that the calculated electron heat fluxes are consistent with the experimentally measured ion temperature at \( r/a = 0.5 \), as follows: substracting the anomalous electron loss from the estimated ohmic power at \( r/a = 0.5 \), and setting the difference equal to the classical electron ion exchange, we can calculate the ion temperature using the measured \( T_e \) and \( n \). It comes out to be within 11% of the experimentally measured \( T_i \) (0.3 ± 0.1 keV).

In Fig. 4 we compare the calculated turbulent ion heat flux with the experimentally estimated one. The experimental estimate is difficult and involves large error bars. However, we have found that it quite impossible to account for the ion heat flux as estimated using neoclassical theory. In this respect, TEXT results are in constrast with ohmic heating results from earlier tokamaks apparently operating in the same \( n \) and \( T \) regimes. Our calculations appear to show that particle flux convection can only account for the estimated ion heat flux partially and only fairly close to the boundary. We find that the inclusion of low order resonances in our \( \Sigma \) spectrum to be essential in accounting for the observed ion heat transport. Thus a further round of experiments is clearly called for to determine the dominant components of the low order resonances and their relative amplitudes.

Conclusions We remark that it is by no means obvious a priori that a magnetic fluctuation spectrum can be constructed with physically reasonable amplitudes and frequencies, and estimate self-consistently the measured fluctuations \( \Sigma \) and \( \psi \) together with various particle and energy fluxes, all in reasonable agreement with experiment. Our calculation shows that at least one such spectrum exists, and further theoretical and experimental work is needed to establish it.

In conclusion our model appears capable of providing a unified framework for interpreting simultaneously the turbulent fluctuation and anomalous transport data in TEXT. The only free parameters in the description relate
to two or three constants in the spectrum which are fixed by reference to the maximum experimental value of the potential fluctuations and the radial variation of the particle flux. The calculation is then able to account for several qualitative and quantitative features of the data including non-adiabaticity and the electron and ion heat fluxes. Future experiments involving trace impurity transport under the same conditions and externally induced fluctuations together with more accurate particle and energy flux data could help to test and refine the theory considerably.

References


Fig. 1 Comparison of experimental and theoretical fluctuations.

Fig. 2 Particle flux from source & probe measurements, as compared with theory. (\(\Phi/kT_e\))

Fig. 3 Total electron heat flux - comparison between experiment and theory. (\(q_e\))

Fig. 4 Turbulent total ion heat flux - comparison between theory and experiment. (\(q_{tot}\))
RESISTIVE BIFURCATING STATES RELATED TO AUXILIARY POWER IN A TOKAMAK

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1. The $I/P^{0.5}$ Scaling of the Confinement Time and the Entropy Principle

The functional $S = \int \left( \nabla \times \mathbf{A} \right)^2 dV + \left( \nabla \times \mathbf{A} \right)^2 dV$ (where $\nabla \times \mathbf{A} = (4\pi/\mu)^2$ and $\mu$ is a parameter) is stationary with respect to arbitrary variations of $\mathbf{A}$ (with suitable boundary conditions) when $\mathbf{J}$ satisfies the equation $\nabla \times \mathbf{J} + \mu \mathbf{J} = 0$. The functional $S$ can be related to the probability $P$ of the current density distribution in a coarse-grained configuration space in the frame of a statistical system of cells containing many particles (for a recent thorough discussion of this model see (2)). Under appropriate constraints characterizing a collective equilibrium smeared out over an infinite sea of particle fluctuations, one finds a relation of the form $S = -\int P \ln P dP$, so that $S$ has the meaning of the entropy.

In the case of a cylindrical tokamak the most probable distribution of the axial current $j$ satisfies the zero order Bessel equation $\nabla \times \mathbf{J} + \mu \mathbf{J} = 0$.

In a situation of ohmic relaxation one has $j \sim T^{3/2}$ ($T_{eff}$ is taken as uniform) and $T$ must satisfy the equation

$$\frac{1}{T} \frac{d}{dT} \left( \frac{1}{2} \frac{\gamma}{\mu} + \frac{1}{2} \mu T^{3/2} \right) = 0$$

On the other hand $T$ must also satisfy the equation for the energy balance of the electrons

$$\frac{1}{T} \frac{d}{dT} \left( \frac{1}{2} \frac{\gamma}{\mu} + \frac{1}{2} \mu T^{3/2} \right) + E \frac{d}{dT} + P_A = 0$$

In the purely ohmic case, when the auxiliary power density $P_A$ is zero, one can write this equation in the form

$$\frac{1}{T} \frac{d}{dT} \left( \frac{1}{2} \frac{\gamma}{\mu} + \frac{1}{2} \mu T^{3/2} \right) + \frac{3}{2} \frac{E}{T} \frac{d}{dT} T^{3/2} = 0$$

where $\gamma$ is a constant related to the Spitzer resistivity, $\gamma = A T^{3/2}$.

Equation (3) is compatible with (1) if $\mu = 3 \epsilon / T^{3/2} A m \hat{\chi}$ and $\chi(\epsilon) \sim \hat{\chi}(T_{eff}) T^{3/2}$, where $\hat{\chi}$ and $\hat{\chi}$ are the values at $q(r)=1$. The dependence $\chi(\epsilon) \sim T^{3/2}$ is then the only one which is compatible with a state of maximum probability of the tokamak. However in principle $\hat{\chi}$ can still depend on $\hat{T}$ and we shall put in general $\hat{\chi} \sim \hat{T}^m$, where $m$ is arbitrary at present.

Let us apply the auxiliary power to a relaxed state of maximum entropy. Using the form above for $\chi(\epsilon)$, one obtains from (2) the following equation
If \( p_a \) is uniform, the solution is \( j(r) = \frac{1}{2} c \left( e^{K r} + e^{-K r} \right) - \frac{p_a}{E} \)
where \( H \) and \( K \) are fixed by the conditions \( j(0) = j_a \) and \( j(a) = 0 \).
In the region \( r < \lambda a \), where the plasma is assumed to be sawtooth-unstable, \( j \) is taken to be flat and (4) cannot hold because the plasma is not in a stationary state with maximum entropy. The parameter \( \mu \) can be expressed in terms of \( \Omega / q \), \( \lambda \), \( j \) and \( p_a \) and the scaling of the physical quantities with \( p_a \) can be expressed in terms of \( \mu \). Let us call \( \mu_0 \) the value of \( \mu \) for \( p_a = 0 \). One can show the following scalings

\[
E = E_0 (\mu / \mu_0)^{1/3}, \quad \Omega = \Omega_0 (\mu / \mu_0)^{1/3}, \quad \Omega = \Omega_0 (\mu / \mu_0)^{1/3} m^2/\beta m^2 \]

Experimentally \( \tau \propto 1/\chi \) scales as \( P^{-0.5} \) (Kaye and Goldston /3/). When one tries to represent (5) in the form \( \chi \sim P^x \) with \( x = 0.5 \), one finds that \( m \) should be near to 2. Putting \( m=2 \), the power law is an optimal representation of (5) for \( 2 \leq P / P_a \leq 6 \) (in the auxiliary and ohmic powers). In this range one finds \( x = 0.44 \). The interesting point to note is that \( \tau \sim P^\theta \), for high \( P \), is sensibly proportional to \( q / q_a \). This implies that when \( q \) is constrained to 1, \( \tau \) depends linearly on the total current, consistently with the observations. Fig. 1 shows how \( \tau \sim P^\theta /\chi \) scales with \( P / V^2 \) (in the plasma volume) for different values of \( q / q_a \) and for \( \gamma = 0.20 \), \( j = 0 \). The reason of the \( T^2 \) dependence of \( \tau \) is not clear at present, although it agrees with the temperature scaling of the thermal electron diffusion derived in a thermodynamic model of the saturation of the drift modes in a reactive medium /4/. The other point to be noted is that the form of the current profile, which is trapezoidal (with a tendency to be concave for high \( q \)) is practically independent of the value of \( P_a \), as illustrated in Fig. 2 (solid line, \( P / P_a = 0.19 \), dashed line, \( P / P_a = 12.56 \)) and is also insensitive to the nonuniformity of \( p(r) \), which is smeared out in the integrals involved in the solution of (4). The external injection of power does not move the system from its state of maximum entropy. This seems to be the essential physical aspect of the so-called "profile consistency" /5/.

2. Resistive Bifurcating States Related to Auxiliary Power

When \( P / P_a \) becomes sufficiently large the plasma cannot remain in the resistively relaxed state and becomes unstable. The instability has a resistive nature and in fact it can exist, under different circumstances, also in a purely ohmic plasma. We assume that a critical value \( q_a \) exists at \( r = a \) above which (for \( r > a \)) the thermal diffusivity becomes very large (e.g. due to stochasticity). The radius \( a(t) \) of the surface with the critical \( q \)-value will move slightly in time as the
result of small resistive disturbances of the current profile (related to small local changes of temperature, for instance a sawtooth pulse). Since, in view of the high diffusivity for \( r > a(t) \), the critical q-surface is also a surface of constant temperature, a link is established between the boundary values of \( T \) and \( q \). For reason of continuity, this also holds for the neighbouring internal q-surfaces. The link implies a dependence of the resistivity on the current and a consequent instability, under certain circumstances. Moreover, as far as the dependence of the current can be considered linear in time, the \( q(r,t) = \text{const} \) surfaces are also surfaces of constant poloidal flux /6/. This allows to write a single equation describing the process in terms of the poloidal flux change \( A_1 \):

\[
\frac{d}{dt} A_1 - \frac{1}{c} \left( \frac{\partial A_1}{\partial \lambda} \right) \frac{1}{\pi} \int \frac{d \lambda}{\lambda} \left( \frac{d A_1}{d \lambda} - \frac{1}{\lambda} \frac{d A_1}{d \Lambda} \right) = \begin{cases} \frac{j_0(A_1, \Lambda)}{l} \frac{d A_1}{d \lambda} & \text{for } q \leq a(t) \\ 0 & \text{for } a(t) < q < b \end{cases}
\]

(6)

where

\[
A_1(q, t) + A_1'(q', t) = A_0(a(q), 0)
\]

The derivation of this equation and its bifurcating states where discussed previously /5/. Here we apply the results to the case when \( j_0 \) and \( A_0 \) describe the state of maximum entropy. Fig. 3 gives the instability thresholds in terms of \( P / P_0 \) and \( b/a \), where \( b \) is the location of the material limiters, and for \( \lambda = 0.20, \frac{\hat{q}}{q} = 0.80, a \) \( j = 0 \) and \( b/j = 0.02 \). One sees that the threshold depends critically on the current pedestal. The edge heating disturbances (such that the q-surfaces are moving outwards) do not saturate for the values of \( b/a \) and \( P / P_0 \) near the dashed lines, while nonlinear saturation and a stable bifurcating state occurs in the neighborhood of the solid lines. Fig. 4 illustrates the broader current profile (dashed) after the transition from the relaxed state (solid) in the case \( \hat{q}/q = 0.56, P / P_0 = 7.50 \). The volume current increases slightly and a negative surface current is formed at this stage. The point-dashed profile illustrates the case when all the current is a volume current. One sees that a steep gradient is formed after the transition, in agreement with the observations, but the consistency with the heat transport needs further consideration.

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A COMPARISON BETWEEN PREDICTED AND OBSERVED PELLET PENETRATION DEPTH IN JET OHMIC-HEATED DISCHARGES

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I. Theoretical Model

The theoretically predicted pellet penetration depth is based on the assumption that the shielding effect is primarily due to the slowing down of the plasma electrons in a dense ablated cloud surrounding the pellet. Two versions of the ablation model are considered; (a) the ablant consists of a single species of D₂ molecules only [1] (frozen flow), (b) the ablant is a 4-species mixture of D₂, D, D⁺ and e⁻ (their relative concentrations are determined by the local temperature and pressure in the ablated cloud - equilibrium flow, [2]). When most energy of the incoming electrons is absorbed in a spherical shell of the cloud around the pellet, the time variation of the pellet radius according to these models can be written as

\[
\frac{dr_p}{dt} = -C r_p^{-2\delta} n_e^{1/2} \rho_e^\delta
\]

(1)

The constant C and the power index \( \delta \) depend on the composition of the ablant and the temperature range of \( T_e \) respectively. When \( r_p \) is measured in cm, \( n_e \) in cm\(^{-3} \), \( T_e \) in eV, their values for a deuterium pellet are given in the accompanying table.

<table>
<thead>
<tr>
<th>( T_e ) (eV)</th>
<th>( \delta )</th>
<th>C 1-species ablant</th>
<th>C 4-species ablant</th>
</tr>
</thead>
<tbody>
<tr>
<td>((10^2, 6 \times 10^2))</td>
<td>1.439</td>
<td>(3.639 \times 10^{-8})</td>
<td>(2.763 \times 10^{-8})</td>
</tr>
<tr>
<td>((6 \times 10^2, 10^3))</td>
<td>1.673</td>
<td>(8.064 \times 10^{-9})</td>
<td>(6.343 \times 10^{-9})</td>
</tr>
</tbody>
</table>

*The constant C listed here has been adjusted according to the previous study that only 80% of the incoming electron energy flux is absorbed in a spherical shell of radius \( r_s = 2 r_p \) enveloping the pellet [3]. To study the pellet penetration in JET it is convenient to refer the ablation of the pellet to a frame of reference moving
with the pellet. Thus referring to a parameter
\[ \xi = \frac{r}{r_p} \]

Eq. (1) can be written as
\[ \frac{d\xi}{dx} = \left( \frac{5C}{3U_0} \right) \frac{\rho_0}{v_0} \frac{\xi}{\xi + \rho_0} \frac{\Delta V_\Phi}{\Delta V_e} \]

where \( r_0 \) is the initial pellet radius in cm, \( U_0 \) is the pellet velocity in cm/s. To evaluate \( \xi \) or the variation of the pellet radius \( r_p \) during its crossing from the magnetic surface \( x(\phi + 1) \) to \( x(\phi) \), we take values of \( n_e \) and \( T_e \) defined in the volume element \( \Delta V_\Phi \) bounded by these two surfaces.

The ablated electrons \( \delta n_p \) added to the shells during the crossing of the pellet from the magnetic surface \( x(\phi + 1) \) to \( x(\phi) \) is given by
\[ \delta n_p (\Delta V_{\phi-1}) = \frac{1}{2} N_p (\Delta V_\phi) / \Delta V_\phi \]

where \( N_p \) is the total number of molecules ablated per second in the shell volume \( \Delta V_\phi \). Thus, when \( f = 0 \), all the ablated material is added to the shell volume \( \Delta V_{\phi-1} \) behind the pellet, the pellet experiences the unperturbed plasma state. When \( f = 1 \), all the ablated material is added to the shell volume, \( \Delta V_{\phi-1} \) ahead of the pellet, causing an adiabatic cooling of the plasma temperature for the advancing pellet. Thus, if \( \delta n_p \) is the cold ablated electrons added to the volume element \( \Delta V_{\phi-1} \) and \( E_i \) is the ionization energy, the corresponding temperature perturbation, \( \delta T_e \) and the updated plasma temperature \( T'_e \) in \( \Delta V_{\phi-1} \) are given respectively by
\[ \delta T_e (\Delta V_{\phi-1}) = - \left( T_e (\Delta V_{\phi-1}) + 2E_i / 3 \right) \delta n_p / n_e (\Delta V_{\phi-1}) \]
\[ T'_e (\Delta V_{\phi-1}) = T_e (\Delta V_{\phi-1}) + \delta T_e (\Delta V_{\phi-1}) \]

II. Comparison with Experimental Observations

To compare the theoretically predicted pellet penetration depth with those observed experimentally, we shall confine our studies to those shots of OH discharges only, taking during the Autumn 1986. The plasma state in each shell volume is taken shortly before the pellet injection, thus \( T_e \) processed from ECE data and \( n_e \) processed from the Abel-inverted FIR interferometry chord measurements.
The experimentally determined pellet penetration depth is deduced from arrays of soft x-ray detectors [4]; the detectors are placed 5 cm apart, an uncertainty of observed penetration depth of ±2 cm is to be expected. With the additional shift of the magnetic axis during the pellet injection, the actual uncertainty of the pellet penetration depth could be even greater. According to the convention used in the JET computational code, up to the position of the limiter, the plasma volume is divided into 50 shells, bounded by successive magnetic surfaces of 2 cm apart; the calculated penetration depth is therefore also subjected to an uncertainty of ±2 cm at least. Besides, the pellets used are actually cylinders, the pellet radius used in the computations is based on an equivalent spherical pellet containing the same number of particles, \( N_p \), adjusted according to the total number of particles actually deposited in the plasma volume, and is somewhat lower than the nominal total number of particles contained in the initial pellet [4]. Since we may alternatively define an equivalent spherical radius of the pellet by taking the same volume to surface area ratio, \( V/S \) as that of a sphere, thus there is an uncertainty regarding the proper radius of the pellet in the computation. In spite of these uncertainties, the results shown in Fig. 2 clearly indicates that the ablation model based on the single-species ablatant tends to underestimate the depth actually penetrated by the pellet. (For the pellet used, \( r_p = 2 \text{ mm} \) for the same \( N_p \), and \( r_p = 1.8 \text{ mm} \) for the same \( V/S \), the uncertainty in its equivalent spherical radius can, at most, give an uncertainty in the penetration depth of 1 cm). As shown in the same figure, when the ablatant is taken to be a 4-species mixture, the presence of additional shielding through the dissociation and ionization effect gives a better agreement between the predicted and the observed penetration depth. On the other hand, as shown in Fig.3, a better agreement between the observed and the predicted penetration depth can be also obtained, if the computed penetration depth is based on the 1-species ablatant model but allowing 10% of the ablated cold material to be added to the plasma volume ahead of the advancing pellet.

III. Conclusion

A comparison between the observed and the predicted pellet penetration depth taken in JET Ohmic discharges indicated that the pellet actually penetrates deeper than that predicted by the 1-species ablatant neutral-shielding model. In view of the experimental and theoretical uncertainties mentioned above, the possible cause of the additional shielding effect is still uncertain.

Acknowledgement

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References:
Fig. 2 Comparison of the predicted and the observed pellet penetration depth, $R_d$, measured from the magnetic axis. ▼ Ablatant, 1-species, ○ ablant, 4-species, $f=0$. $T_e$(o), 3.0-4.8 keV, $n_e$, 3.7-7.4 $\times 10^{19}$ cm$^{-2}$.

Fig. 3 Comparison of the predicted and the observed pellet penetration depth, $R_d$, based on the 1-species ablant neutral-shielding model with and without the adiabatic cooling effect. • $f=0$, □ $f=0.1$. 

INTERACTION OF COLD HIGH-DENSITY PARTICLE CLOUDS WITH MAGNETICALLY CONFINED PLASMAS

L.L. Lengyel


The sudden release of a large number of neutral particles subjected to subsequent ionization has found application in fusion research as well as in other plasma-physical areas, e.g. magnetospheric experiments. Although the physical conditions are vastly different in these experiments, there are some fundamental similarities between the respective cloud expansion processes. In all these cases the released neutral particles expand in a spherically symmetric manner until they become ionized. As a result of ionization and subsequent interaction with the magnetic field, the transverse expansion of the cloud is slowed down and brought to a full stop. The radially confined plasma is 'funneled' into magnetic flux tubes: its expansion along the magnetic field lines is practically a vacuum expansion. The expanding plasma distorts the magnetic field: a transient magnetic cavity may form inside the cloud. Since the cloud is heated by the incident plasma particles, the incident energy flux is a function of the species present in the plasma and, in the case of energy carriers with gyro-radii less than the cloud radius, of the diamagnetic state of the particle cloud.

A variable-mass single-fluid single-cell Lagrangian model was developed to describe and analyze the above processes, with implementation primarily to pellet-plasma interaction studies. The rate of the cloud mass variation, i.e. the strength of the mass source attached, is specified by means of a pellet ablation routine coupled to the cloud expansion model: the ablation rate is given at any time instant as a function of the temperature and density of the cloud particles surrounding the mass source.

The model was tested by means of data of earlier magnetospheric barium cloud experiments where the characteristic time and length scales range from 100 s to 1000 s and from 10 km to 500 km, respectively. Sudden particle release (instantaneous ablation) was assumed in this case. Detailed calculations were performed for hydrogen clouds associated with the injection of pellets into tokamak plasmas. The characteristic time and length scales in this case are of the order of 1 μs to 100 μs and 1 mm to 50 mm, respectively (see [1] and [2]).

The time histories of all relevant cloud parameters (which, besides determining the strength of the mass source, represent massive local perturbations for the recipient plasma), of such as its radius, length, temperature, pressure, ionization degree, beta value, lifetime, internal magnetic field strength, etc, are computed for various (total) cloud masses, ablation rates, magnetic field strengths, plasma temperatures, and are analyzed.
Some representative results corresponding to pellet injection scenarios in ASDEX and JET are shown in Figs. 1 and 2. Here the time development of various pellet cloud parameters and local plasma disturbances are shown for a given number of pellet particles deposited between two flux surfaces. If the pellet injection velocity is given, the mass source strength can be calculated on the basis of the residence time of the pellet in the flux tube considered. The number of particles locally deposited is taken either from experimental measurements (ASDEX: Shot No. 18716/1.624 s, see Fig. 1) or ablation calculations (JET: see Ref. [3]). The local ablation rate was computed in this case by means of the ablation model proposed by Houlberg et al. [4]. Constant ablation rate (e.g. linearly increasing ablatant cloud mass) was assumed for the residence time of the pellet in the flux tube considered.

In Fig. 1 the time development of the cloud size (radius, length) and of the transverse velocity of the cloud boundary are shown. The transverse deceleration of the cloud mass involves Alfvén time scales. The maximum cloud radius is defined by the relative magnitudes of the pellet mass locally deposited, the heat flux affecting it (ionization time), and the magnetic field strength applied. The expansion along the magnetic field lines and the return to equilibrium conditions involves hydrodynamic time scales which are by three orders of magnitude larger than the Alfvén time scale. The maximum (calculated) beta reached in the plasma cloud of Fig. 1 was 0.05, the maximum electron density was about $1.2 \times 10^{23} m^{-3}$.

In Fig. 2 the time histories of the magnetic field strength, electron density, and beta are shown for the JET scenario considered. In this case, the plasmoid radius reached after a few overdamped oscillations is about 1.4 cm.

Acknowledgement

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References

Fig. 1: Time development of a pellet cloud in ASDEX (Shot No. 18716/1.624 s): N = 2.53 × 10^18 particles deposited between r = 28 and 29 cm, \( v_p = 650 \text{ m/s} \).
Magnetic field strength in the plasmoid (tesla)

Electron density
(10^22/m^3)

Plasmoid beta
(10^-2)

Fig. 2: Time development of plasmoid characteristics in JET at r = 0.9 m, Δr ≈ 2.0 cm; N = 2.2 x 10^{19} particles are deposited within 20 μs.
PELLET PARTICLE DEPOSITION PROFILES WITH ALLOWANCE FOR NEUTRAL GAS EXPANSION EFFECTS

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**IPP Garching, EURATOM Association, Fed. Rep. of Germany

A numerical model has been developed recently in which the expansion of the pellet cloud along the magnetic field lines as well as in the transverse direction is calculated by means of a self-consistent model that takes into account such effects as finite rate ionization, \( \vec{j} \times \vec{B} \) deceleration, magnetic field convection, and magnetic field diffusion [1]. This model is coupled to a numerical routine in which the local ablation rate is determined as a function of the electron temperature and electron density specified [2]. The total number of particles ablated at a particle flux surface is given by the product of the ablation rate and the residence time of the pellet in the flux tube considered. The expansion and ionization of these particles are followed up numerically until they are stopped at some distance from the pellet. The results of calculations show that the expansion velocity of the neutrals may considerably exceed the traversing velocity of pellets injected into hot plasmas. Hence the neutrals may become ionized and confined to magnetic flux surfaces at some distance from the location of the pellet. The particle deposition profile is thus determined by summing up, at each flux surface, the contributions of particles released by the pellet at the subsequent flux surfaces while traversing them. A numerical algorithm has been developed for this purpose that is described elsewhere [3].

Some representative results corresponding to the injection of single pellets into ASDEX, JET, and NET are given in Figs. 1 to 6. In all these cases the ablation rate was calculated by means of the semi-analytical formula of Ref. [2]:

\[
\frac{d\hat{r}}{d\alpha} \text{Const} \times T_e^{1.64} \times n_e^{0.33}/\hat{r}_p^{0.67}.
\]

The constant appearing in this expression has been chosen in such a manner as to reproduce the pellet penetration depth experimentally observed in ASDEX (Shot no. 18716). Note that unlimited energy reservoir was assumed in the present calculations, i.e. the finiteness of the flux tube volumes in tokamaks was not taken into account.

In all these figures, the number of particles deposited per flux tube is plotted as a function of the radius. The width of a flux tube is defined (in an arbitrary manner) as \( \Delta r = a/40 \), 41 being the number of mesh points used in the calculations. The solid lines shown correspond to local particle deposition whereas broken lines to particle
ASDEX | JET | NET
--- | --- | ---
R (m) | 1.65 | 2.98 | 5.20
\(a\) (m) | 0.40 | 1.20 | 1.62
B (tesla) | 2.2 | 3.0 | 5.5
\(T_{e0}\) (kev) | 30.0 | 4.58 | 0.58
\(T_{e1}\) (keV) | 0.02 | 0.20 | 0.20
\(n_{e0}\) (\(m^{-3}\)) | \(1.5 \times 10^{20}\) | \(1.3 \times 10^{19}\) | \(1.5 \times 10^{20}\)
\(n_{e1}\) (\(m^{-3}\)) | \(0.8 \times 10^{19}\) | \(0.5 \times 10^{19}\) | \(0.6 \times 10^{19}\)

In all these cases, the electron temperature and density profiles were prescribed by means of the expression \(f(r) = f_1 + (f_0 - f_1) \left[1 - \left(r/a\right)^{EX1}\right]^{EX2}\), where for \(T_e\) \(EX1 = EX2 = 2\) and for \(n_e\) \(EX1 = 2\, EX2 = 0.5\), respectively.

Figures 1 and 2 correspond to pellet injection into ASDEX with \(v_p = 800\) m/s. In the reference case (Shot no. 18716, Fig. 1) a pellet of \(N_p = 3.2 \times 10^{19}\) \((r_p = 0.5\) mm\) has been used. For obtaining deep penetration, a pellet size with \(N_p = 5 \times 10^{20}\) has been assumed for Fig. 2. Figure 3 corresponds to pellet injection into JET: here \(N_p = 2 \times 10^{21}\) \((r_p = 2\) mm\) and \(v_p = 1100\) m/s. Figures 4 to 6 correspond to NET scenarios. In the first two cases a rather low pellet velocity, \(v_p = 1100\) m/s, has been used. The first pellet size, \(N_p = 1.6 \times 10^{22}\) \((r_p = 4\) mm\), see Fig. 4) corresponds to a NET refuelling pellet \([4]\). For obtaining intermediate and deep penetrations, the pellet size was increased to \(N_p = 7 \times 10^{23}\) and \(2 \times 10^{24}\) (see Figs. 5 and 6, respectively). In the last case, the pellet velocity was increased to \(2500\) m/s.

The results of these calculations can be summarized as follows:

(a) Within the framework of the analytical (smooth) plasma parameter distributions and the analytical ablation rate function assumed, no notable difference exists between local and non-local particle deposition in the bulk of the plasma. Here non-local deposition has merely a smearing effect.

(b) Non-local deposition may produce notable differences in the wings of the deposition curves particularly at high ablation rates, coupled with limited energy supply.

(c) The existence of significant non-local deposition effects depends upon the amount of the cold mass locally released and the magnitude of the energy flux affecting the neutral cloud. If the resulting ionization process is comparatively slow, significant non-local effects may be produced in the wings of the deposition curves: deeper particle penetration results (see Figs. 2, 5, and 6).

References

Particle deposition in ASDEX:
\( N_P = 3.2 \times 10^{19} \)
\( v_P = 800 \text{ m/s} \)

Fig. 1

Particle deposition in ASDEX:
\( N_P = 5 \times 10^{20} \)
\( v_P = 800 \text{ m/s} \)

Fig. 2

Particle deposition in JET:
\( N_P = 2 \times 10^{21} \)
\( v_P = 1100 \text{ m/s} \)

Fig. 3
Particle deposition in NET:

\[ N_p = 1.6 \times 10^{22} \]
\[ v_p = 1100 \text{ m/s} \]

Fig. 4

Particle deposition in NET:

\[ N_p = 7 \times 10^{23} \]
\[ v_p = 1100 \text{ m/s} \]

Fig. 5

Particle deposition in NET:

\[ N_p = 2 \times 10^{24} \]
\[ v_p = 2500 \text{ m/s} \]

Fig. 6
1. Introduction

Analysis of a weakly dissipative (but inviscid) axisymmetric nearly static MHD equilibrium allows to determine the flux profiles left free in the ideal MHD theory (see [1] and [2]). In the case where no mass injection and auxiliary heating are present and the resistivity \( \eta_{||}, \eta_{\perp} \) as well as the temperature diffusion coefficient \( k_{\perp} \) are assumed to be flux functions, we derive equilibrium relations in the limit of large aspect ratio \( A = R_{o}/a \). In the simplest possible case considering constant transport coefficients ("loop model") we solve the equilibrium equations and obtain simple scaling relations for the central plasma \( \beta \) and the safety factor \( q \). We test the numerical validity of this model for aspect ratios of order unity. Additionally we consider the "Spitzer case" \( i.e. \eta_{||} = \eta_{o}T_{o}^{-3/2} \) and improve the scaling of the loop model according to this case. Finally we compare the predictions of the \( \beta \) scaling laws with measurements.

2. Scaling Relations at Large Aspect Ratio

We consider the usual MHD equations for a weakly resistive axisymmetric plasma involving heat transport \( k_{\perp} \) that is sustained in a nearly static equilibrium by the loop voltage \( (V_{c}) \). Treating this as a perturbation of the ideal static case we can derive a set of 3 equations for the unknown flux profiles \( T_{o}, p_{o} \) and \( I_{o} \) which have to be solved simultaneously with the Grad-Schlüter-Shafranov (GSS) equation for the poloidal flux \( \Psi \) (see [1]-[3]).

Let us now consider the limit of large aspect ratio. First we prescribe the geometry of the \( \Psi = \text{const.} \) surfaces as concentric ellipses.
We set: \( R = A - r \cos \theta; \ Z = K \rho \sin \theta; \ \Psi(R, Z) = \Psi(r); \) \( r \) is the normalized small radius and \( K \) is the vertical elongation of the cross section. Knowing the flux geometry we replace the GSS equation by its flux average to determine \( \Psi(r) \). Additionally we assume low \( \beta \) tokamak scaling (\( q \approx 1, \beta \sim A^{-2} \)) and set \( \eta_{\perp} = 2 \eta_{\parallel} = 2 \eta_o T_o^{-3/2}; \ k_{\perp} = \text{const} \). Evaluating the set of equilibrium relations with respect to large aspect ratio then leads to:

\[
\frac{\lambda_k^2}{2} \frac{1}{r} \frac{d}{dr} \left( r \frac{d\Psi}{dr} \right) = Q(T_o); \quad \Psi(r = 1) = 0; \quad \frac{d\Psi}{dr}(r = 0) = 0 \quad (1)
\]

\[
4\pi \frac{dp_o}{d\Psi} = -A^{-2} Q(T_o) \frac{\lambda_k^2}{4}\left( \frac{d\Psi}{dr} \right)^2; \quad p_o(\Psi = 0) = 0 \quad (2)
\]

\[
\frac{dT_o}{d\Psi} = -\frac{A^{-2} T_o}{4\pi p_o} \int_0^{V_o} dV_o \left( Q^2(T_o) 4 k_{\eta}(T_o) \right); \quad T_o(\Psi_M) = p_o(\Psi_M) \quad (3)
\]

\[Q(T_o) = \frac{2 \nu_e T_o^{3/2}}{\nu_o}; \ k_{\eta}(T_o) = \frac{\eta_o T_o^{-3/2}}{4\pi}; \ \lambda_k^2 = (1 + K^{-2}); \ \Psi_M = \Psi(0) \quad (4)
\]

Here \( 3 k_{\perp} = \chi_e + \chi_i \), where \( \chi_e, \chi_i \) are the usually defined temperature diffusion coefficients, and we omit the \( I_o \) equation because it is decoupled from the rest and of no physical interest here. Integration of the toroidal direction of Ohm's law leads to a condition which determines \( V_{\ell} / \eta_o \) in terms of the total toroidal normalized plasma current \( I \), namely:

\[I = g \pi KA^{-1} Q(T_M); \quad g = \frac{\int T_o^{3/2} d\Psi}{\int T_M^{3/2} d\Psi}; \quad I = \frac{4\pi I_p(MA)}{10 a(m)B(T)} \quad (5)
\]

The equations can be solved iteratively doing the following: Start with a temperature profile \( T_o(\Psi) \); put it into equation (5) to fix \( Q(\Psi) = Q(T_o(\Psi)) \). Then solve equation (1) with the improved \( Q \); use the result to integrate first the pressure equation (2) and then the temperature equation (3). After that restart the procedure.

The solution of the "loop case" is simply obtained by setting \( T_o(\Psi) = T_M = \text{const} \). and performing one complete iteration step. For the central values of \( q \) and \( \beta \) we obtain:

\[q_L = \frac{(1 + K^2) \pi}{IA}; \quad \beta_L = \frac{A^{-2} \bar{I}^4}{4 + 2 \bar{I}^2} \approx \frac{\bar{I}^4}{4A^2}; \quad \bar{I}^2 = \frac{A^2 \bar{I}^2}{\pi^2 K^2 \lambda_k^2} \quad (6)
\]
Concerning the approximation of the $\beta$ formula we used the experimental result that $\tilde{I}^2 \approx 0.13 \ll 2$ for big tokamaks. We calculated a numerical solution for a typical JET discharge (see [4]) and compared it with the scaling relations. We found good agreement, although the results refer to an aspect ratio of only 2.4. In fig.1 we compare the $\beta_L$'s of the scaling law with experimental values and see that they are an order of magnitude too small.

An improvement is the solution of the next iteration step. We start with the temperature profile of the loop model which is of the form: $T_0 \sim \Psi^{\gamma}; \gamma$ depends on $k_{\perp}, I$ and the geometrical coefficients (see [2] and [3]). Now we calculate the factor $g$ of eq. (5). It is $g = 2/(3\gamma + 2)$. The corrected safety factor takes the form $q_M = gq_L$. Using again $\tilde{I}^2 \ll 2$ we obtain for $\beta : \beta_M = G(\gamma)\beta_L; G(\gamma)$ (fig.2) collects the effects of the variation of the resistivity. Because $G$ is very sensitive to slight changes in $\gamma$ we fix it by the demand that the $\beta$ scaling formula gives exactly the measured value for the JET discharge ($\gamma = 0.97$), and we want to look wether the scaling law corrected in this way is usable for other machines. That means we make a kind of profile consistency ansatz for the temperature profiles. The results are plotted in fig.1 and show a significant improvement of the loop model.

3. Conclusions

Investigation of a diffusive equilibrium model involving classical resistivity and constant thermal conductivity led to simple scaling formulae for the basic tokamak parameters. Comparison of the $\beta$ scaling relation with experiments shows satisfactory agreement and can be taken as a hint that the Spitzer case already contains the basic features of long ohmic pulses.

References

Fig. 1 Comparison of the scaling laws with some experimental values:

(1) = ASDEX,
(2) = T10,
(3) = TFTR,
(4) = PLT,
(5) = JET;

\( \Phi = \beta_L, \)
\( \bigcirc = \beta_M, \)
\( \times = \text{experiment} \)

Fig. 2 Correction factor \( G(\gamma) \)
Introduction

Two scenarios for optimising the thermal alpha particle yield in the tokamak are considered. In addition to these thermal alpha particle yields it should be noted that substantial non-thermal contributions to the alpha power, such as beam-beam and beam-plasma interactions [1], can also arise. In general in this paper we will assume that \( Z_{\text{eff}} = 1 \) and that \( T_{\alpha} = T_e \). The significance of these assumptions can be seen from the relation

\[
P_{\alpha} \sim 0.23 B^2 \beta_o^2 \tau_e P_{\text{in}}
\]

Optimisation by Peaking Pressure

The alpha particle heating power may be written as

\[
P_{\alpha} \sim 0.23 B^2 \beta_o^2 \tau_e P_{\text{in}}
\]

Thus for a given toroidal field \( (B_T, \text{ Tesla}) \) and power input \( (P_{\text{in}}, \text{ MW}) \), \( P_{\alpha} \) is optimised by maximising the product of peak beta \( (\beta_o^2) \) and energy confinement time \( (\tau_e \text{, seconds}) \). If \( \tau_e \) is assumed to be independent of the pressure peaking (ratio of peak to volume average pressure, \( P_o / \langle P \rangle \)) then for given stored energy \( (\text{ie } \tau_e \beta_o^2) \) and \( P_{\alpha} \) are increased by peaking the pressure. An operation scenario involving peaked pressure profiles to optimise \( P_{\alpha} \) has been suggested for JET [2]. There is however a decreasing stability limit for \( \beta_o \) as the pressure is peaked [3]. This arises because for a given q-profile, there is a critical pressure gradient \( (P'_\text{crit}) \), on every flux surface, for stability to the ideal ballooning mode. Thus peaking the pressure profile, which involves having \( P' \ll P'_\text{crit} \) on the outer flux surfaces, will lower \( \beta_o \) (and \( \langle \beta \rangle \)). Figure 1 shows how the alpha power determined by this \( \beta_o \) stability limit decreases as the pressure is peaked,
for 7 MA, $B_T=3.4$ T operation in a machine such as JET. Also shown in Fig 1 is the confinement limit on the alpha power determined from the Kaye-Goldston L-mode scaling law [4] with $P_{\text{in}}=40$ MW. It can be seen that in this case (cf [3] for q-profile details), $P_{\alpha}$ is confinement limited for $P_o/<P> < 5$ and stability limited for $P_o/<P> > 5$. The optimal $P_{\alpha}$ occurs when the confinement and stability limits are equal. It should be noted in this case that if the confinement limit is improved by a factor $\sim 2$, by high current X-point H-mode operation (such as is planned in JT-60), then the optimal pressure profiles are less peaked and much higher $P_{\alpha}$ ($\sim 13$ MW) can be achieved.

Fig 1
Confinement and stability limit on $P_{\alpha}$ as pressure is peaked.

Optimising by Reducing The Discharge Aperture

Assuming the Kaye-Goldston L-mode scaling [4] and that $P_{\alpha}$ is confinement limited, we find that $P_{\alpha} \propto R^{2.3} I^{2.5} a^{-3} k^{-0.6} B_T^{-0.2}$ (or $P_{\alpha} \propto a^{1.25} q^{-2.5} k^{3.3} B_T^{2.3} R^{-0.7}$ in terms of the limiter $q_\psi$, taken to be $a^{1.25} k^{1.5} B_T / (R^{1.2} I)$). Thus if machine size is the only constraint, $P_{\alpha}$ is optimised by operating at the highest current consistent with $q_\psi > 2$. If however the current is limited then $P_{\alpha}$ is optimised by moving the plasma to smaller $R$, and reducing 'a' at the largest permissible ellipticity (k), until $q_\psi = 2$ is reached. Figure 2 shows the improvement in $P_{\alpha}$ as the aperture is reduced, for various currents, in a JET size device with $P_{\text{in}}=40$ MW and an assumed elongation limit of 2. The minimum minor radius for each curve arises from the $q_\psi = 2$ limits (using the $q_\psi$ parameterisation due to Todd in Ref [5]). A problem with this reduced aperture inner wall operation is the large evanescent region which makes low field side ICRH and lower hybrid
heating unsuitable. However in these cases ECRH can be particularly effective, as the coupling of the power to the plasma is not then dependent on the location of the boundary. Ray tracing calculations show that the technically simple, low field side launch of the O-mode can be used at 120 GHz. With a launch angle of 70° to the toroidal field, Doppler shifted resonance allows absorption near the magnetic axis of a reduced aperture plasma \( R_o = 2.6 \text{m} \) with an efficiency of 100% \( (5^\circ \text{ cone half angle}, T_e(0) = 10 \text{keV}, n_e(0) = 7.5 \times 10^{19} \text{ m}^{-3}) \). The location of absorption is insensitive to \( T_e(0) \) but can be varied by changing the launch angle.

Operating with a reduced aperture also gives the possibility of using a major radius compression to heat the plasma. Such a compression increases the fusion yield dramatically, since transiently \( P_\alpha \propto C^{6\cdot66} \) (where the compression ratio, \( C = R_{old}/R_{new} \)). Figure 3, shows from a 1-D transport simulation, the predicted increase in alpha power due to a 0.2 sec, \( C=1.2 \) major radius compression for a JET sized device with \( I_p = 5 \text{ MA} \). Two sets of curves are shown in Fig 3, one corresponding to \( P_{in} = 40 \text{ MW before and after the compression, and the other set to } P_{in} = 20 \text{ MW after compression. Additionally for each power input results are shown with and without the alpha power deposited in to heat the plasma; this shows (particularly for } P_{in} = 40 \text{ MW} \) the significant impact of the alpha power and indicates for such scenarios that it should be possible to realistically study alpha particle physics.

Fig 2 Improvement to \( P_\alpha \) from reducing aperture at various currents.
Fig 3
Effect of major radius compression at $t=1.5s$ on $P_\alpha$ with 20 MW and 40 MW input after compression.

Conclusions

Two scenarios for optimising the thermal alpha particle yield in the tokamak have been studied:

1. It has been shown that $P_\alpha$ is optimised by peaking the pressure until the confinement and stability limits on $\beta_p$ are equal.

2. For a given current $P_\alpha$ is optimised by reducing the aperture at the largest permissible ellipticity until the $q=2$ limit is reached. For such reduced aperture inner wall plasmas low field side launch ICRH and lower hybrid heating are unsuitable. A combination of low field side launch ECRH and major radius compression (+ NBI) may however be used to heat these reduced aperture plasmas.

Both scenarios would produce significant alpha particle yields ($Q \chi$), and indicate that realistic studies of thermonuclear alpha-particle physics are possible in the present generation of large tokamaks.

References

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PRELIMINARY RESULTS ON THE ALPHA CONFINEMENT IN NET

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INTRODUCTION

Even for the simple physical assumption of a single alpha particle slowed down by classical coulomb collision in a field of particles of thermonuclear plasma, in presence of magnetic ripple, the guiding center orbit numerical calculations published by various authors show substantial disagreement. For example, the fast alpha particle loss during slowdown for an INTOR like device has been evaluated from about 20% [1,2] to 3% [3,4]. Other substantial discrepancies are also reported for the energy spectra of the lost particles and for the energy density distribution of these losses on the first wall. In a reactor the most deleterious effect of these losses is the additional heat load on the first wall and the subsequent possible, unwanted impurity influx into the plasma.

The reasons for these discrepancies have so far not yet been clarified. However the feeling at the Workshop on alpha-particle effects in ETR, held at DOE, Germantown (USA) in June 1987 [5] was that what were believed to be small differences in the physical models and the numerical codes, may be responsible for large variations in the results. It was thus proposed to perform a numerical Bench-Mark Experiment according to the following physical model.

BENCH-MARK EXPERIMENT

\[
\begin{align*}
R_0 \ (m) & = 5.3 \\
\alpha \ (m) & = 1.2 \\
B_0 \ (T) & = 5.5 \\
T(r) \ (keV) & = 20 \ (1 - \rho^2) \\
n(r) \ (m^3) & = 4 \times 10^{20} \ (1 - \rho^2) \\
J(r) \ (A/m^2) & = 1.32 \times 10^6 \ (1 - \rho^2) \\
Z_{\text{eff}} & = 1 \\
N & = 12 \ \text{toroidal field coils}
\end{align*}
\]

The toroidal field coils are simulated by the segments connecting the following points:

\[
\begin{align*}
R \ (m) &= 2.105 & Z \ (m) &= \pm 3.98 \\
R \ (m) &= 6 & Z \ (m) &= \pm 5.565 \\
R \ (m) &= 9.725 & Z \ (m) &= \pm 2.724
\end{align*}
\]

- The resulting maximum ripple at plasma edge will be about 1%.
- Collisional pitch angle scattering of the \alpha\ particles is assumed.
- The plasma has the shape of a circular torus
- The poloidal projection magnetic flux surfaces are circular and concentric.
- The \alpha\ particles are considered lost, when, for the first time, their guiding centre crosses the plasma edge = wall.
- no other fields as those mentioned above are considered

Calculations will be considered equivalent if:
1) the prompt losses
2) the delayed losses
3) the energy spectrum of the delayed losses
4) the energy loss distribution on the wall

are equivalent

FIRST ORBIT LOSSES

These losses have been calculated, for the Bench-Mark Experiment, by means of our 3D Montecarlo guiding center orbit following code. The full guiding center orbit and particles energy equations were integrated by the fifth order Adam method, the toroidal magnetic field was computed at each step using Biot-Savart law for the 12 filament coils.

Averaged along the toroidal direction we get:

<table>
<thead>
<tr>
<th>Edge ripple</th>
<th>First orbit loss</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>2.15x10^{-2}</td>
</tr>
<tr>
<td>0.94%</td>
<td>2.25x10^{-2}</td>
</tr>
<tr>
<td>1.60%</td>
<td>2.35x10^{-2}</td>
</tr>
</tbody>
</table>

These results are compared with those of [1] in Fig. 1. There is a good agreement up to a ripple of about 1%. For higher ripple value the discrepancy may be due to the difference of the collisionless ripple trapping-detrapping caused by: a) a slightly larger extend of the ripple well zone (cfr. Fig. 2) for [1] which also b) approximate the ripple shape, retaining only the first term of the Fourier expansion.

DELAYED LOSSES

The full guiding center trajectory calculation for a single alpha slowing down up to the critical energy requires about 20 minutes on the CRAY-XMP computer. It implies the calculation of the order of 10^6 steps.

Thus, whatsoever is the care given to the tests, the controls and the verifications we find it difficult to exclude the possibility of the propagation of errors over a such large number of successive steps. Therefore, in order to have a further check of the orbit following method and to get a deeper physical insight of the involved processes, we divide the alpha trajectories in different groups, derive [1] a diffusion coefficient from our orbit following code for each group and compare this experimental coefficient with an analytically derived one.

These groups, represented schematically in the (radius/pitch angle) plan in Fig. 3, subdivide as following the confined particles:

- Circulating costreaming
- Circulating counterstreaming
- Non ergodic banana
- Ergodic banana

\[
\text{Total initially confined alpha} \quad 100\%
\]
The best confined, are the circulating orbits for which the neoclassical diffusion coefficient is valid. The less confined, are the ergodic bananas. In Fig. 4 we show the limits of the ergodic banana region calculated for the Bench-Mark experiment according to the analytical expression derived by [6]. This banana ergodicity, which occurs as the ripple size exceeds a certain critical value is essentially a collisionless...
The corresponding so called ripple-plateau diffusion coefficient [7] is in our case about 4 order of magnitude higher than the neoclassical coefficient. An intermediate diffusion coefficient is obtained for non ergodic bananas. Thus, as one moves, in Fig. 3, from the deeply confined circulating orbit towards the ripple loss cone, the diffusion coefficient increase step by step.

As the slowing down time of the fast alpha is much longer than the ergodic banana confinement time, this group of orbit will be lost before thermalization.

The non ergodic banana confinement time being about of the same order of magnitude as the slowing down time, a premature loss of some of these orbits may be expected. In order to save calculation time we will focus our attention on this critical group.

Banana orbits undergo precession along the toroidal direction. In presence of ripple (even in the absence of ripple well) the tips of the bananas may enter in resonance with the corrugation of the magnetic field and may stop temporarily the precession. An example of such a resonance, which substantially alters the diffusion [8,9], is shown in the poloidal map of Fig. 5.

**FOOTNOTE AND REFERENCES**

DIRECT LOSSES OF \( \alpha \) - PARTICLES IN SPIN POLARISED PLASMAS

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1. Introduction

The dependence of light ion fusion reactions on nuclear spins was recognised early in the development of nuclear physics [1]. More recently, consideration has been given to lowering thresholds for break-even and ignition in tokamaks by exploiting this dependence to enhance reactivities [2]. However, the usefulness of such a scheme is contingent on there being an overall improvement in plasma self-heating power allowing for simultaneous adjustments to product ion losses. We examine approximately the different balances for \(^2\)H - \(^3\)H plasmas confined in a simplified magnetic configuration of concentric circular flux surfaces.

2. Lawson criteria for spin polarised plasmas

Fusion of \(^2\)H + \(^3\)H ions proceeds almost entirely through a single \( J^p = (\frac{3}{2})^+ \) excited state of the intermediate \(^6\)He compound nucleus. Consequently at thermal plasma energies, just the two \( J = \frac{1}{2} \) or \( J = \frac{3}{2} \) orientations of their nuclear spins for zero relative orbital angular momentum contribute to the cross-section, in proportions of 2 : 1 as determined by appropriate ratios of Clebsch - Gordon coefficients [2]. It follows that by aligning, or "polarising", both plasma ion species with their spins co-parallel to a confining magnetic field, the more favourable \( J = \frac{3}{2} \) channel is wholly selected, producing an energy independent increase of roughly 50% in the cross-section and so Maxwellian rate coefficient [2].

Increased reactivity in a fully co-parallel polarised \(^2\)H - \(^3\)H plasma lowers its threshold for ignition, as depicted for optically thin, impurity free cases in Fig.1. Here the Lawson criterion is expressed in terms of a volumetric ratio of plasma thermal energy content to the excess of nuclear fusion power over its radiative cooling rate due to relativistically corrected dipole bremsstrahlung. Reactions \(^2\)H(\(^3\)H,n)\(^4\)He; \(^2\)H(\(^4\)H,n)\(^3\)He; \(^2\)H(\(^2\)H,\(^1\)H)\(^3\)H; \(^3\)H(\(^3\)H,2n)\(^4\)He are included, the first alternately with a conventional rate coefficient [3], corresponding to an unpolarised plasma, and one enhanced by 50% for co-parallel polarisation. It should be noted that spin dependencies of pure deuteron fusions are less well resolved at present, and may be sufficiently energy dependent in the co-parallel mode as to leave their rate coefficients essentially unmodified from unpolarised values [2,4]. Thus they have not been altered in Fig.1. The resulting amelioration of conditions for ignition may appear comparatively modest, but
could be particularly attractive where it transfers ignition to the accessible side of some critical boundary, such as a tokamak $\beta$ limit.

Since $^3$H; $^3$He and intermediate compound nuclei $^6$He; $^6$Li are respective mirror nuclei, $^2$H($^3$He,$^1$H)$^3$He fusion is a very similar reaction, again exhibiting an energy independent increase of almost 50% in the cross-section for full co-parallel polarisation of thermal plasmas [2]. Corresponding ignition curves are also shown in Fig.1, still including unchanged $^2$H + $^2$H reactions. In fact, the additional nuclear charge causes $^2$H + $^3$He and $^2$H + $^2$H cross-sections to cross over at lower kinetic temperatures [3], so that pure deuteron fusions can play a significant role and induce some compositional sensitivity in the optimum characteristic. The optimum unpolarised mixture actually occurs roughly around 70% $^2$H with 30% $^3$He. A 50% to 50% fuel mix has been assumed in Fig.1 in order to maximise gain by co-parallel polarisation in the $^2$H + $^3$He channel. Clearly, for both $^2$H - $^3$H and $^2$H - $^3$He fuels, co-parallel polarisation shifts could similarly be estimated for energy break-even lines, given some set of input demands.

3. Tokamak fusion product prompt loss mechanisms

In the preceding Lawson plots, all charged fusion product energy was assumed to be devoted to plasma self-heating. However, these ions are actually subject to finite loss mechanisms in physical tokamak devices, restricting the true efficiency of energy absorption. Prompt losses of fast product ions before equilibration with the background plasma are primarily associated with departures from axisymmetry in the confining magnetic field [5,6,7].

Product ions trapped in the magnetic field major radial modulation also experience ripple variations arising from the discrete number of toroidal coils. Some particles may become locally trapped in ripple wells themselves, during which field gradient drifts can lead to magnified cross-field diffusion or escape to the vessel wall. This process is

![Lawson ignition curves for 50% $^2$H - $^3$H and $^2$H - $^3$He plasmas, equal temperature ions and electrons.](image1)

![Probability distributions for fusion product emission angle relative to magnetic field direction.](image2)
typically small in practice \([5,6]\), and a more general effect of field ripple is to disturb banana orbits of toroidally trapped particles by causing radial excursions at their anticipated or delayed turning points \([5]\). Such orbits become broadened, or beyond some critical ripple depth even stochastic \([5,7]\), thereby exaggerating banana diffusion and again, especially in the latter situation, rapid interception of material boundaries. A similar but independent effect seems probable if magnetic fluctuations of comparable magnitude coincide with banana end-points. Note also that product ions born on trapped orbits within a banana width of a wall must anyway be immediately lost. All these effects together provide a collisionless direct loss channel which for high energy charged fusion products therefore proceeds faster than either thermalisation or detrapping \([5,6]\), so depriving the plasma of at least some generated fusion power.

In addition to amplifying reactivities, co-parallel spin polarisation (denoted by \(C\)) tends to focus product ejectiles in a plane normal to the magnetic field. A second \(^2\text{H} - ^3\text{H}\) polarisation mode \((P)\) can be formed, having deuterons spins perpendicular to \(B\) and tritons unpolarised, which reproduces wholly unpolarised fusion rates but conversely focusses ejectiles preferentially along \(B\). Taking an unpolarised plasma \((U)\) to yield isotropic ejection, probability distributions of product emission angle \(\phi\) relative to \(B\) in each \(^2\text{H} - ^3\text{H}\) polarisation regime are \([2]\) (see Fig. 2):

\[
P_U(\phi) = \frac{2}{\pi}; \quad P_\perp(\phi) = \frac{4}{\pi} \sin^2(\phi); \quad P_\parallel(\phi) = \frac{4}{5\pi} \left[1 + 3 \cos^2(\phi)\right]; \quad (0 \leq \phi \leq \frac{\pi}{2}). \tag{1}
\]

An immediate consequence concerns the mean Larmor radii of produced \(\alpha\) - particles. Estimating these as proportional to root mean energies of origin perpendicular to \(B\) reveals:

\[
\rho^2 = \frac{m\sqrt{2}E}{Z_0} P(\phi) \sin \phi d\phi \int_0^{\pi/2} P(\phi) \sin \phi d\phi \Rightarrow \rho_U : \rho_C : \rho_P = 1 : \sqrt{1.2} : \sqrt{0.8}. \tag{2}
\]

Hence under co-parallel polarisation each thermalising \(\alpha\) - particle deposits its energy on average over a slightly broader region, which might itself be marginally disadvantageous for ignition. More significantly, though, plasma polarisation must clearly have an accompanying impact on prompt \(\alpha\) losses.

4. Relative self-heating powers of spin polarised plasmas

Without attempting explicitly to quantify collisionless losses for fast \(\alpha\) - particles, we simply assume a fixed portion \(0 < \zeta \leq 1\) of the toroidally trapped fraction always escapes before delivering an invariant and non-trivial part of their fusion energy to the plasma. Relative self-heating powers between \(^2\text{H} - ^3\text{H}\) cases exactly identical except for spin polarisation regime then scale as the differing initial trapped fractions of charged ejectiles.

At a point having strength of magnetic induction \(B\), let ions be born with some distribution \(P(t)\), where \(t = \sqrt{\xi} = \sin^2(\phi)\) is an effective pitch angle variable. Then the fraction trapped may be expressed:

\[
T(B) = \int_{b_\text{MAX}}^{b_{\text{MIN}}} P(t) \frac{1}{\sqrt{1-\zeta}} dt / \int_0^1 P(t) \frac{1}{\sqrt{1-\zeta}} dt, \tag{3}
\]

using adiabatic invariance of single particle energy and magnetic moment. Now assuming concentric circular flux surfaces implies \(B \propto (1 - \epsilon \cos \Theta)\) to leading order in inverse aspect ratio \(\epsilon\), where \(\Theta\) is poloidal angular...
co-ordinate. To this consistent level of simplification, inserting (1) for each polarisation mode thus gives:

\[ T_\theta(\theta) \approx \frac{1}{2} \sqrt{\varepsilon(1+\cos\theta)} \]
\[ T_e(\theta) \approx \frac{3}{2} \sqrt{\varepsilon(1+\cos\theta)} \]
\[ T_p(\theta) \approx \frac{1}{2} \sqrt{\varepsilon(1+\cos\theta)} \]

and averaging finally around surfaces:

\[ \frac{1}{\sqrt{2\pi}} \int_0^{2\pi} T_\theta(\theta) d\theta = \frac{2}{\pi} \sqrt{2\varepsilon} \]
\[ \frac{3}{\sqrt{2\pi}} \int_0^{2\pi} T_e(\theta) d\theta = \frac{3}{2} \frac{\sqrt{2\varepsilon}}{\pi} \]
\[ \frac{1}{\sqrt{2\pi}} \int_0^{2\pi} T_p(\theta) d\theta = \frac{1}{2} \frac{\sqrt{2\varepsilon}}{\pi} \]

In other words, limiting prompt losses for fusion \( \alpha \)-particles relative to an unpolarised plasma are expected to decrease with P-mode polarisation, but increase with co-parallel polarisation, just in proportion moreover to its improved reactivity.

Plasma self-heating power may be estimated as the remaining \( \alpha \) energy rate after collisionless degradation through trapping. Under the foregoing circumstances, co-parallel polarisation only produces a net rise approximately through-out a region satisfying:

\[ \frac{3}{2} \left( 1 - \frac{2}{\sqrt{2\pi}} \right) > \left( 1 - \frac{2}{\sqrt{2\pi}} \right) \implies \varepsilon < \frac{1}{\sqrt{50}} \frac{1}{\pi^2} \approx \frac{0.2}{\pi^2} \]

Outside the critical flux surface at \( \varepsilon_{\text{CRIT}} = 1/\sqrt{50} \), augmentation of direct losses would outweigh that of reaction rates, actually reducing local self-heating compared with an otherwise identical unpolarised case. At least in a limit of large flux surface aspect ratios, overall usefulness of co-parallel polarisation therefore depends on the relative distributions of fusion power inside and outside \( \varepsilon_{\text{CRIT}} \), i.e. the degree of peaking of its profile. This constraint grows stricter as collisionless losses for fast \( \alpha \)-particles become more severe \((f \to 1)\). Note that in contrast P-mode plasma polarisation has no critical flux surface, but rather by inhibiting trapping produces a universal improvement in self-heating over unpolarised performance, even though reaction rates themselves are unchanged.

5. Conclusion

Co-parallel nuclear spin polarisation of a \(^2\text{H} - ^3\text{H}\) tokamak plasma raises fusion reactivity by \(-5\%\). However, less isotropic emission of \( \alpha \)-particles also exacerbates trapping and associated prompt losses. Plasma self-heating relative to an unpolarised case consequently is increased non-uniformly against flux co-ordinate, actually being decreased outside a certain boundary. Alternative P-mode polarisation does not alter reactivity, but by consistently reducing fast \( \alpha \) trapping leads everywhere to a relative improvement in self-heating.

References

EFFECT OF THE LOWER HYBRID WAVE ON THE TOKAMAK DRIFT MODES

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It is known /1-3/ that lower hybrid frequency pump has an essential effect on the drift instabilities. So far this effect was considered only in a local approximation. In fact, however, the drift waves in a tokamak are strongly affected by the magnetic field shear and the toroidal modulation and so the local theory is not applicable. This paper deals with the stability problem for toroidal drift modes in a presence of the HF pump.

In the presence of the lower hybrid pump the toroidal drift eigenmode equation has the form \((-\infty < \eta < \infty)\):

\[
\left[ \frac{d^2}{d\eta^2} + (\zeta_S \Omega)^2 Q(\Omega, \eta) \right] \hat{\phi}(\eta, \zeta) = 0
\]

Here \(\zeta_S = q \sqrt{\beta_0} / \varepsilon_n\), \(q\) - safety factor, \(\varepsilon_n = \eta_n / R\)

\(\tilde{\Omega}^2 = \mu^2 \frac{d\Omega}{d\eta}\), \(R\) is the large radius of the torus,

\(\beta_0 = (k_0 p_0)^2\), \(\Omega = \omega / \omega_0\), \(q_S = \tau_e / m_e \omega_0\),

\(\omega_0 = k_0 T_e / e B \tau_n\),

\(Q = \beta_0 (1 + \tilde{\Omega}^2 \hat{S}) + 1 - \frac{1}{\Omega} + 2 \frac{T_0}{\Omega} (\cos \eta + \frac{1}{2} \tau_0 \sin \eta) -

- \tau^2 \mu (1 + \tilde{\Omega}^2 \hat{S}) \left[ (1 + \tilde{\Omega}^2 \hat{S})^2 - (2 k_0 / k_0^2) \tilde{\Omega}^2 \hat{S}^2 \right]^{-1}

\hat{S} = \tau a^2 \eta / 4\), \(\mu = (U_0^2 \omega_0^2 / c_S^2 \omega_0^2)\), \(U_0 = c E_0 / B_0\),

\(E_0\), \(\omega_0\) and \(k_0\) - electric field strength, frequency and the wave vector of the pump, respectively. Equation (1) is derived in the "ballooning" representation /4/.

In the absence of the pump \( \mu = 0 \) equation (1) reduces to that one obtained earlier in /4-6/.

There exist two drift wave branches in the toroidal geometry: toroidal modification of the Pearlstein-Berk modes /7/ and the modes induced by the toroidicity itself. Effect of the HF pump on the first branch was considered analytically by the matching of asymptotic solutions of the equation (1) modified for the Pearlstein-Berk eigenmodes. Corresponding dispersion relation has the form

\[
\frac{\Gamma'(\alpha)}{\Gamma'(\alpha + \nu + \frac{1}{2})} = -\left(\frac{e^{\pm \infty}}{2^{3/2}}\right)^{1+2\nu} \frac{\Gamma'(\nu)\Gamma'(\nu - \frac{1}{2})\Gamma'(\frac{1}{2} - \nu)}{\Gamma'(-\nu)\Gamma'(\nu + \frac{1}{2})\Gamma'(\nu + \frac{3}{2})} \cdot (1 + \Pi A \cot \Pi \nu)
\]

Here \( \alpha = (x - \gamma)/2 \), \( \nu = -(\alpha + \frac{1}{2}) \), \( \lambda = -\frac{\sqrt{1+4\Lambda}}{2} \)
\( \Lambda = -\mu / k_0 \), \( \Omega = (\frac{1}{2}k_0 \Omega)^2 \), \( \zeta^2 = -i (k_0 / 2 \Omega \lambda) \)
\( \chi = i \lambda^2 (\frac{1}{2} \Omega)^2 \left( -1 + \frac{2\epsilon_0 + 2\epsilon_1}{\Omega - i\delta} \right) \)

and

\[
A = \begin{cases} \frac{2}{\Pi} \tan \frac{\Pi}{2}(\nu + 1) & \text{for even eigenmodes} \\ \frac{2}{\Pi} \cot \frac{\Pi}{2}(\nu + 1) & \text{for odd eigenmodes} \end{cases}
\]

For the matching procedure to be valid, we require \(|\epsilon| < 1\). For definiteness, assume that the destabilizing factor \( \delta \) is associated with trapped electrons-ions collisions. Then for

\[
\omega < \omega_{eff} < \omega_{tr}
\]
we have

\[
\delta \approx \frac{\omega_{tr} - \omega_{eff}}{\omega_{tr}} \left( 2\epsilon \right)^{1/2}
\]

where \( \omega_{eff} = \omega_{tr} \epsilon \), \( \epsilon = \frac{Z_{eff}}{Z} \), \( (2\epsilon)^{1/2} \) is the trapped
electron fraction, $\omega_{Be}$ is their bounce frequency.

Approximate solutions of the eigenvalue condition (2) can be obtained in certain interesting limits:

a) $|\Lambda^2| < \mathcal{C}$. For the lowest $n = 0$ even eigenmode we have

$$\frac{\omega}{\omega_{Be}} \approx \frac{1 - 2\epsilon_n}{1 + b_\theta - \mu_0^2} + i \left\{ \frac{\omega^2 (2\epsilon)^{1/2}}{\omega_{Be} \sqrt{\epsilon e + \epsilon}} (b_\theta + 2\epsilon_n - \mu_0^2) - \frac{\epsilon_n (1 - 2\epsilon_n)}{q (1 + b_\theta - \mu_0^2)} \left[ \frac{\omega^2}{\omega_{Be}} (2\hat{S} - 1) \right]^{1/2} \right\}$$

where $\mu_0^2 = \left( \frac{a}{2} \right)^{3/2} \left( \frac{U_0 \omega_{pi}}{c_s \omega_0} \right)^{4/3} \hat{S}^{2/3} \left( \frac{k_\theta}{2k_\omega} \right)^{1/3}$.

$$\left[ b_\theta \hat{S}^2 + \epsilon_n \left( 2\hat{S} - 1 \right) \omega_{Be} / \omega \right]^{1/4}, \mu_* = \mu \left( \frac{k_\theta}{2k_\omega} \right)^2$$

and we assume $\hat{S}_m \omega << \Re \omega$. It follows from the expression (3) that HF pump suppresses the instability by the ponderomotive frequency shift as well as by the increase of shear damping.

b) $|\Lambda| > \mathcal{C}$ (strong pump). For $\hat{S} > 0,5$ the solution of equation (2) has the form

$$\frac{\omega}{\omega_{Be}} = \frac{1 - 2\epsilon_n}{1 + b_\theta - \mu_i} + i \left\{ \frac{\omega^2 (2\epsilon)^{1/2}}{\omega_{Be} \sqrt{\epsilon e + \epsilon}} (b_\theta + \epsilon_n - \mu_i) - \frac{\epsilon_n (1 - 2\epsilon_n)}{q (1 + b_\theta - \mu_i)} \left[ \frac{\omega^2}{\omega_{Be}} (2\hat{S} - 1) \right]^{1/2} \right\}$$

where $\mu_i = \frac{U_0 \omega_{pi}}{c_s \omega_0} \frac{k_\theta}{2k_\omega} \left[ b_\theta \hat{S}^2 + \epsilon_n \omega_{Be} (2\hat{S} - 1) / \omega \right]^{1/2}$.

For $\hat{S} < 0.5$, $\epsilon_n > b_\theta \hat{S}^2 / \left[ (1 + b_\theta)(1 - 2\hat{S}) + 2b_\theta \hat{S}^2 \right]$ when in the absence of the pump the mode becomes marginally
stable (\( \Delta m \omega = 0 \) for \( \Delta z = 0 \)) /6/, effect of the pump leads to destabilization of the mode:

\[
\frac{\omega}{\omega_{*e}} = \frac{1 - 2 \epsilon_n}{1 + b_e} \int \frac{k_e}{\omega_k} \frac{\epsilon_n(1 - 2 \delta_n \omega_k - b^2)}{\omega_k - \omega_e} \left[ \epsilon_n(1 - 2 \delta_n \omega_k - b^2) \right]^{1/2} (5)
\]

The influence of HF pump on the toroidicity induced modes was considered by numerical integration of (1) for real \( \omega \) and \( \delta_n \), using a shooting method with a sixth-order Numerov scheme. For real \( \omega \) the eigenvalue \( \delta_n \) represents a destabilizing factor which is sufficient for the mode to become marginally stable, thus it is the measure of the shear damping. Figure 1 shows these results. One can see that the toroidicity threshold for these modes excitation (which correspond to \( \delta = 0 \)) is raised in the presence of the pump. Moreover, as in the first case considered above, the ponderomotive frequency shift reduces the instability increment.

References.
STABILITY OF TOROIDICITY INDUCED DRIFT WAVES IN DIVERTOR TOKAMAKS

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The experimental observation of steep pressure gradients at the plasma edge in H-mode discharges, obtained in tokamak with magnetic separatrix, indicates a strong local reduction of the particle and thermal diffusion in comparison with the L-mode situation. It is then natural to investigate how the presence of a magnetic separatrix can affect the stability of the semicollisional electrostatic drift waves, which seem to play a major role in determining anomalous transport. As strong driving mechanisms cannot be invoked at the plasma edge, the stability of these modes is very sensitive to the possibility of inhibiting shear stabilization. Such a possibility can be achieved in toroidal geometry via the coupling between modes centered around different rational surfaces. In a low β tokamak equilibrium with circular magnetic surfaces this mechanism yields a new, toroidicity-induced, branch of the eigenvalue. In the present paper we focus our investigation on the modification of the toroidal coupling and of the stability of this branch, in the collisionless cold-ion limit, due to the presence of a separatrix.

In this limit, by introducing the ballooning transformation [1], the eigenmode equation for the fluctuating electrostatic potential reduces, at the lowest order in the ratio of the perpendicular wavelength and the equilibrium scale length, to the following second order equation in the extended poloidal variable \( \theta \):

\[
\frac{d^2}{d\theta^2} \phi + Q(\theta) \phi = 0
\]

with boundary conditions corresponding to outward energy convection. Here

\[
Q(\theta) = \frac{1}{4} \left( \frac{C^2}{2C} - \frac{C''}{2} + \frac{\Omega b'_0 q^2}{C^2} \frac{1}{\ell_n} \right) \left( 1 + \frac{b_s}{\Omega} \right) \left( 1 + \frac{\omega_{p,\perp}}{\omega_{\text{c},\text{e}}} \right)
\]

with boundary conditions corresponding to outward energy convection. Here

\[
C = b_p / h, \Omega = \omega / \omega_{\text{c},\text{e}} \text{ is the eigenvalue}, \ell (r, \theta) \text{ is the infinitesimal poloidal arc length along the magnetic surface, } B_p = B_{\theta} | b_p (0), r_\theta \text{ and } B_p \text{ are the minor radius and the poloidal field opposite to x-point, } \omega_{p,\perp} = (2C / \rho_B) k \times B - V (P + 1/2 R^2), \text{ and the remaining is standard notation. Following Ref. [2], we model tokamak equilibrium with separatrix prescribing } b_p (0) \text{ on the surface as } b_p (0) = \left[ 1 + g^2 (\sqrt{1 + k - 1})^2 - 2g (\sqrt{1 + k - 1}) \cos (\theta - \gamma) \right] / [1 + k]^{1/2} \text{ with the shape } g \theta (0) / r_i \text{ given by } g^2 (\sqrt{1 + k - 1})^2 [4 + g^2 (\sqrt{1 + k - 1})^2 - 4g (\sqrt{1 + k - 1}) \cos (\theta - \gamma)] = k^2. \text{ Here } \gamma \text{ is the poloidal coordinate of the x-point and } k \text{ labels}
the flux surface ranging from 0 (circular flux surface) to 1 (the separatrix). Moreover we can express \( h(0) = |g^2 + g^2|^{1/2} \),

\[
\beta^0 = \beta_0 \left| \frac{1}{b_p^2} + \frac{b_p^2 p^2}{p^p} \right|
\]

\[
\frac{\omega_{\phi}}{\omega_{*e}} = 2 e_p \left| \frac{\sin u}{b_p} + \cos u b_p \right|
\]

where

\[
p(\theta) = \int_0^1 d\theta' h \left( \frac{\Delta}{b^3_p} + 2 \frac{b_p^p}{p} - 2 \frac{p}{A} \right)
\]

\[
f(\theta) = \frac{p(\theta)}{r_0} = \frac{|g^2 + g^2|^2}{g^2 + 2g^2 - gg''}
\]

\[
A = \frac{k}{2} \sqrt{1+k} \frac{1}{\sqrt{1+k - 1}} , \quad D = \frac{k^2}{n(\sqrt{1+k - 1})} \quad K(k^2), \]

\[
F = \frac{1}{n} \left| (1+k) K(k^2) + \frac{E(k^2)}{1-k} \right|
\]

\( \rho(0) \) being the gaussian radius of curvature, \( u \) the angle between the tangent to the magnetic surface and the major radius direction and \( K \) and \( E \) complete elliptic integrals.

Here \( \theta_0 \) is a complex ballooning parameter to be determined at the next order imposing the stationarity condition [3]

\[
\frac{d\Omega}{d\theta_0} = 0 \quad (3)
\]

The numerical solution of Eq. (1) in the limit \( k \to 0 \) agrees with previous findings [4] and shows the existence of a quasi marginally stable branch corresponding, for \( \xi > 1/2 \), to a double well structure in the potential \( Q(\theta) \). For rational surfaces near the
separatrix, locating the x-point in the equatorial plane, both on the outside (\(y = 0\)) and on the inside (\(y = \eta\)) of the plasma, causes a deepening of the well structure, so enforcing the inhibition of shear damping and the marginal stability result. On the contrary, the location of the x-point on the top of the plasma (\(y = \eta/2\)) produces the flattening of the well and restores the shear damping yielding stabilization of the mode. In Fig. 1 the real (solid line) and imaginary (dashed lines) part of the eigenvalue \(\Omega\) are plotted vs \(k\) for the above mentioned three cases (with \(b_0 = 0.1, \nu_1 = 0.1, s = 1, q = 1\)). In Fig. 2 \(\Omega\) is plotted for the case \(y = \eta/2\). In the other two cases, symmetry considerations allow to conclude that condition (3) is satisfied by \(0_0 = 0\) or \(0_0 = n\). Here \(0_0 = 0\) has been assumed, as it corresponds to the less stable case.

In the limit \(k \rightarrow 1\), Eq. (1) can also be treated analytically as far as modes localized around the x-point are concerned. Assuming the ordering \(x = (0 \cdot y)^{1/2} \sim b_0 = (b_0/\delta) \sim 1, \varepsilon_n = (\nu_1/\delta^{1/2}) \sim 1\) and \(1 \cdot 1/\Omega \sim b_0 \sim \varepsilon_n\), with \(\delta = 1 - k\), and looking for solutions localized at \(1/\delta \ll 1\), Eq (1) reduces to

\[
\frac{d^2 \Phi}{dx^2} + \left(\beta_0 + \beta_1 x + \beta_2 x^2\right) \Phi = 0
\]

with

\[
\beta_0 = -\frac{1}{2} + \frac{\Omega^2 b_0 q^2}{\varepsilon_n} \frac{2(\sqrt{2} + 1)}{2(\sqrt{2} + 1) - 1 - \frac{1}{\Omega} + 2b_0(1 + C_0^2 \cos\gamma)} + \frac{2\sqrt{2}}{\Omega} \nu_1 (\cos\gamma + C_0 \sin\gamma)
\]
Fig. 2: Complex ballooning parameter $\theta_o$ vs the parameter $k$

\[ \beta_1 = \frac{\Omega^2 b_0 q^2}{\bar{\varepsilon}_n^2} 2(\sqrt{2} + 1) \left\{ 4 b_0 c_0 \left(2 - \frac{4}{\pi}\right) - \frac{2 \sqrt{2} \bar{\varepsilon}_n^2}{\Omega} \left[ \left( \frac{4}{\pi} - 1 \right) \sin \gamma - c_0 \cos \gamma \right] \right\} \]

\[ \beta_2 = \frac{5}{4} - \frac{\Omega^2 b_0 q^2}{\bar{\varepsilon}_n^2} 2(\sqrt{2} + 1) \left\{ 1 - \frac{1}{\Omega} + 2 b_0 \left(3 - c_0^2\right) + \frac{2 \sqrt{2} \bar{\varepsilon}_n^2}{\Omega} \left[ \left( \frac{1}{\pi} + \frac{4}{\pi} \right) \cos \gamma + \frac{c_0}{2} \sin \gamma \right] \right\} \]

and $C_0 = C_0(0,\nu)$ to be determined imposing $d\Omega/dC_0 = 0$. In particular, for $\nu = 0$, $C_0 = 0$, and the eigenvalue is $\Omega = 1 - 2b_0 - 2 \sqrt{2} \bar{\varepsilon}_n$.

If, on the other hand, the ordering $|\omega| \sim \delta t/2$, $b_0 \sim \varepsilon_n$, is assumed, Eq. (4) yields $\Omega = (1 - 2 \sqrt{2} \bar{\varepsilon}_n)/2b_0$, which agrees with the result shown in Fig. 1.

**FOOTNOTE AND REFERENCES**

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TOROIDAL ION TEMPERATURE GRADIENT DRIVEN DRIFT MODES WITH DISSIPATIVE TRAPPED ELECTRON EFFECTS

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Introduction

In connection with i.e. auxiliary heating in Tokamaks and other configurations drift modes driven by field curvature combined with temperature gradients are supposed to be essential for the stability and transport properties of the system. The so called $n_1$ mode ($\eta_1 = L_n/L_T$, $L_n$=density characteristic scale length, $L_T$=temperature characteristic scale length) has among other things been suggested to cause confinement degradation in high temperature regions, the break-down of the neo-Alcator scaling at higher densities\(^1\), profile consistency effects and improved confinement phenomena observed at pellet injection experiments. Also in burning fusion plasmas these modes are supposed to play an important role. The so called dissipative trapped electron mode which includes the effects of electron trapping and an electron temperature gradient is used to explain the neo-Alcator scaling at lower densities\(^1\).

The description of the $n_1$ modes in Tokamak geometries has except for some numerical evaluations of kinetic integrals required an expansion in inverse aspect ratio and it has been difficult to derive results valid for a realistic Tokamak aspect ratio. Recently, however, the $n_1$ mode was studied in a two-fluid description taking full toroidal effects into account and using the ballooning mode formalism\(^2\). Simple analytical stability criteria and growthrates were derived.

In this work we study dissipative trapped electron effects on the toroidal $n_1$ mode. We use a fully toroidal two-fluid description where the trapped electron effect is taken into account by a modification of the electron density response\(^3,4\). In a similar way as in (2) we arrive at an eigenvalue equation which in the electrostatic limit is simplified to a quadratic dispersion relation.

Eigenvalue equation

Following the procedure of ref. (2) the basic fluid equations (continuity, motion, and energy) are used to derive the density responses $\delta n_i/n_i$ and $\delta n_e/n_e$ for ions and electrons respectively assuming first a perturbation of the general form $\exp(ikx+jkx-j\omega t)$ where the parallel direction is along the background magnetic field and $k \gg k$. The electron density response is of
the form \( \delta \eta_e / \eta_e = (1-i\omega) \psi / T_e + (\omega \eta_e / \omega) \Phi / T_e \) where \( \psi \) and \( \Phi \) are potentials and \( \alpha = e^3 / 2 \omega_e \eta_e / \nu_e \) accounts for the trapping effect (\( \omega_e \) = electron diamagnetic frequency, \( \nu_e \) = electron-ion collision frequency). The parallel ion dynamics and energy transport is assumed to be negligible whereas the parallel electron heat conductivity is large. By means of the quasineutrality condition we obtain a dispersion relation which is specialized to a tokamak geometry by transforming \( k_r \) and \( k_\theta \) to the relevant operators in the poloidal coordinate (\( \theta \)). We obtain the eigenvalue equation

\[
\phi^{''} + \phi^{'} = (D/k_c^2 \nu_A^2) A \phi / D_e S
\]

where \( \cdot \) denotes differentiation with respect to \( \theta \), \( r=1+s^2 \), \( s \) = the shear parameter and \( A(\alpha, \omega, \theta) \) is of order 1. The equation is fully toroidal - it has been derived without any expansion in quantities which are of the order of the inverse aspect ratio. It is solved by means of the ballooning mode formalism - \( \theta \) is extended to infinity.

**Dispersion relation and marginal stability**

In the electrostatic range \( D/(k_c \nu_A)^2 << 1 \) and the solution of the eigenvalue equation coincides with the solution of \( D_e S(\omega, \theta=0) = 0 \). From the full unexpanded form of this relation we obtain

\[
(A_r+iA_r)\omega^2+(B_r+iB_r)\omega+C_r+iC_r=0
\]

where \( \omega \) has been normalized to \( \omega_{ce} \) and the imaginary parts of the coefficients introduced by the electron trapping are \( A_r = -\alpha \), \( B_r = -(10/3)(\epsilon' / \tau)^2 \alpha \), \( C_r = -(5/3)(\epsilon' / \tau)^2 \alpha \) and \( \epsilon' = \epsilon_n(1-\epsilon) \), \( \epsilon_n = 2L_e / R \), \( \tau = T_e / T_i \). At marginal stability we find \( \omega_r = (1/3)(\epsilon' / \tau)(5^{+} \text{10}) \) and the conditions \((k^2 \rho_r^2 / \tau)(\eta_1 + 1 - 5\epsilon' / 3) + \epsilon' - 1 > 0 \) and

\[
C_r^2 + 70\epsilon''^2 A_r C_r / 9 + 25\epsilon''^4 A_r^2 / 9 - 10\epsilon'' B_r C_r / 3 - 50\epsilon''^3 A_r B_r / 9 + 5\epsilon''^2 B_r^2 / 3 = 0
\]

where \( \epsilon'' = \epsilon' / \tau \) and \( A_r \), \( B_r \) and \( C_r \) are the comparatively simple coefficients obtained without trapping \( \epsilon' \). If these are substituted we get an equation which is quadratic in the parameters \( \eta_1 \), \( k^2 \rho_r^2 \), \( \epsilon' \) and \( \tau \). We note that the stability threshold does not depend on the magnitude of the trapping parameter \( \alpha \). Neglecting Larmor radius effects the condition \( \epsilon' > 1 \) has to be fullfilled in order to have a stable situation. The corresponding stability thresholds of \( \eta_1 \) will be

\[
\eta_1 = 2 + 5 / 3 \cdot 2(1-\epsilon')^2 + 5(\epsilon' / \tau)^2 \quad , \quad \epsilon' > 1
\]

The effect of trapping on the eigenfrequency is illustrated in figs. 1, 2 and 3 (compare fig. 17 of ref (5)).
Fig 1. Eigenfrequency (normalized to $\omega_{re}$) versus density scale length parameter $\epsilon_n$.

$\kappa=0.3, \eta=4, \tau=1, \epsilon=0.2, \eta_c=4, v/\omega_{re}=1.$

Fig 2. Same as fig 1. $v/\omega_{re}=10.$
Main results

- a simple fully toroidal quadratic dispersion relation for electrostatic $\eta_i$ modes including dissipative trapped electron effects.

- an analytical and numerical study of the trapped electron effect on the stability boundaries of $\eta_i$, $\varepsilon$, $k_p$ (Larmor radius parameter), $\epsilon_n = 2L_n/R$ and $\tau = T_e/T_i$. The boundaries are shifted as compared to the case when trapping is not considered. The magnitude of the trapping, however, turns out not to influence the boundaries which suggests that the effect of trapping, though sometimes small, is always important for the $\eta_i$ modes. Generally it sets harder stability conditions on parameters such as density gradient and inverse aspect ratio.

- a study of the trapped electron effects on the growthrate scaling.

References:
(1) Register A., Hasselberg C., Weiland J., CTH-IETF/PP-1987-10
TOKAMAK EQUILIBRIUM DETERMINATION THROUGH FUNCTION PARAMETRIZATION

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A reliable determination of the magnetic flux surfaces in a tokamak plasma lies at the basis of a successful interpretation of most of the diagnostic measurements made on a tokamak. For this purpose, a Grad-Shafranov equilibrium solver is under development for the Rijnhuizen Tokamak experiment RTP. It will identify the flux surfaces from the magnetic measurements made with 12 radial and 12 poloidal field coils and 12 flux loops. To improve the accuracy with which the flux surfaces are determined, the code will make use of additional measurements, if and when they become available. These measurements might include measurements of the sawtooth inversion radius (by means of soft X-ray observations), of the poloidal field profile (by means of Faraday rotation measurements), and of the local current density (by means of tangential Thomson laser scattering). In addition to this, the code will be relatively fast, so that a "fingerprint" will be available soon after termination of a discharge at any required timeslice, and it will be possible to obtain time-resolved information on the equilibrium.

Clearly, these wishes are in discord with the traditional Grad-Shafranov codes, which require a lot of computing time, while not being very flexible with respect to additional measurements. Recent work by Braams, Jilge and Lackner [1,2] has shown, however, that a statistical method, known as function parametrization, provides a means to achieve this goal. This method involves describing a plasma state by a small set of parameters, which are varied over certain well-chosen intervals, such that all or most realistic experimental situations are covered. Then a tokamak equilibrium simulation code, that takes these parameters as input, is run, in order to obtain the corresponding equilibrium. From this, the magnetic (and additional)

<table>
<thead>
<tr>
<th>Table 1: parameters defining a plasma state</th>
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</thead>
<tbody>
<tr>
<td>General expression</td>
</tr>
<tr>
<td>Current density profile</td>
</tr>
<tr>
<td>Pressure profile</td>
</tr>
<tr>
<td>Plasma boundary</td>
</tr>
<tr>
<td>Poloidal flux at boundary</td>
</tr>
<tr>
<td>Plasma position</td>
</tr>
<tr>
<td>Shift of magnetic axis</td>
</tr>
</tbody>
</table>

Here \( \psi \) is the flux, normalized by means of \( \Phi_1 \), and \( \theta \) the poloidal angle.
measurements, such as they would occur in the experiment, are computed. Both the plasma parameters and the simulated measurements are stored in a database, which is then subjected to a statistical analysis that serves to find relationships between the measurements and the plasma parameters. Once these relationships are available, they can be used to interpret real measurements. All relevant plasma parameters may be obtained from a fast evaluation of these relationships, using the measurements as input.

Equilibrium
We make use of the equilibrium code HBT, developed by Goedbloed [4]. Table 1 gives an overview of the parameters this code requires as input to compute the corresponding equilibrium. The quantities $I_0$, $R_0$, and $a_0$ are scaled out of the equilibrium calculations, i.e. they are irrelevant to the equilibrium determination by HBT in the sense that every choice for $I_0$, $R_0$, and $a_0$ yields the same normalized flux surfaces. HBT assumes equatorial symmetry, and we also restrict ourselves to equatorially symmetric plasmas. We do intend, however, to modify HBT such that it applies to non-equatorially symmetric plasmas as well.

Each parameter is varied over a certain interval in a number of steps. Generally, this number can be very low, say 3, in order to obtain information on the plasma parameters with the maximum accuracy allowed by the measurements. The minimal number of simulations that HBT must perform is then $3^{N-M}$, where $N$ is the number of parameters that define a plasma state, and $M$ is the number of parameters that may be scaled out. In the example given in Table 1, this would mean $3^6 = 729$ simulations, which is not excessive.

From the equilibrium computations by HBT we know the flux distribution within the plasma. To obtain the flux function outside the plasma, we use the multipolar moment expansion method, due to Alladio and Chrisanti [5]. First we compute the internal multipolar moments from the current distribution in the plasma. Then we compute the external magnetic fields and the flux function from these multipolar moments, which is now quite easy. Apart from providing an easy method by which to determine the external fields, the multipolar moments have the property that they are mutually independent, which is very useful for function parametrization purposes. Lastly, they also provide insight into the maximum amount of detail that can ultimately be determined from magnetic measurements: Alladio and Chrisanti [5] have shown that for a tokamak with an aspect ratio of 3.16 (cf. RTP: 3.06) and a measurement error level of 1%, it is not possible to determine more than the first 4 internal moments.

For every simulation the plasma parameters, the multipolar moments and the simulated measurements are stored in a database. This database is now subjected to a statistical analysis.

Function parametrization
We represent the plasma parameters, the corresponding multipolar moments, and some other parameters that are of interest, such as $\beta_p$ and $l_p$, that can be computed from the simulation results, collectively by a vector $P$. The measurements are likewise represented by a vector $Q$. We now apply function parametrization to the many occurrences of $P$ alongside $Q$ that we find in the database in order to retrieve the relationships $P = F(Q) + \epsilon$, where $\epsilon$ are small numbers. In general, $F$ will consist of polynomials up to second or third order. A program that performs function parametrization has been written and tested. A flow diagram is shown in Fig. 1.

Subroutine PRESCA performs the so-called "prescaling", i.e. it exploits the a priori physical knowledge of the model in order to make the parameters dimensionless and, if possible, independent of one another (e.g. by computing the multipolar moments from the magnetic measurements, which has the advantage that they may be easily related to the multipolar moments that are generated by the plasma). Also, if possible, it scales the parameters in such a way that expected dependencies between plasma parameters and
measurements become linear, because this will facilitate the fitting procedure.

Subroutine DISPER computes the dispersion matrices of the prescaled plasma parameter and measurement vectors, as well as their eigenvectors and eigenvalues. The dispersion matrices are brought onto their diagonal form by changing to a basis consisting of these eigenvectors. On this new basis, the vectors PREP and PREQ are denoted by TRAP and TRAQ. The collinearities between components of these vectors should now be largely removed. The variances of the components of TRAP and TRAQ within the database are equal to the eigenvalues that have been computed in DISPER. If the variance of a component is small with respect to the other components, it is apparently not important in the description of the measurements, and may be discarded. This procedure is called Principal Component Analysis.

Finally, a standard multidimensional fitting routine fits the vector TRAP to TRAQ. The resulting relationship may be translated into a direct relationship between P and Q. Subsequently, any real measurement Q may be mapped onto the corresponding plasma state P by simply evaluating this relationship. The evaluation also yields a $\chi^2$ which, if deviating too much from the expected value, indicates either a failing diagnostic or an equilibrium that is not covered by the data in the database.

![Flow diagram of function parametrization program](image)

Fig.1 Flow diagram of function parametrization program
Testing function parametrization
The function parametrization program described above has been tested with a wire model. This model computes the vacuum fields of a set of wires in toroidal geometry using Green's function as defined in [3]. We have tested function parametrization with several wire configurations in approximate RTP geometry ($R=0.72 \text{m}$, $a=0.235 \text{m}$). We kept the wire positions fixed (with a number of wires of up to 4) while varying the currents and were able to retrieve the currents from the simulated magnetic measurements with very high precision. This precision is due to the fact that the problem is linear, which makes it ideally suited for function parametrization. Then we varied both the currents and the positions, keeping, however, the wires grouped (i.e., fixed with respect to each other). This is a more realistic model. Now it was not possible to determine the individual currents through the wires, but only suitably defined current moments. The reason for this is that the problem is no longer linear, and, more importantly, no longer has a unique solution: several wire configurations may account for the measured fields to within a given accuracy. The current moments, however, are uniquely defined and mutually independent. The results of this test, along with the definition of the moments, are shown in Table 2, for $\varepsilon = 0$ and $\varepsilon = 0.1$ ($\varepsilon$ is the measurement error, that has been simulated by adding random noise to the simulated measurements. Note that $\varepsilon = 0.1$ is larger than the expected real measurement error, which will be $\leq 0.03$).

These test results look promising and we expect to have a Grad-Shafranov solver, based on the assumption of equatoril symmetry and using the method of function parametrization, ready fairly soon.

<table>
<thead>
<tr>
<th>Moment</th>
<th>Notation</th>
<th>Accuracy ($\varepsilon = 0$)</th>
<th>Accuracy ($\varepsilon = 0.1$)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>$\sum I_i$</td>
<td>0.003%</td>
<td>10%</td>
</tr>
<tr>
<td>1</td>
<td>$\sum I_i r_i$</td>
<td>0.04%</td>
<td>10%</td>
</tr>
<tr>
<td>2</td>
<td>$\sum I_i r_i r_i$</td>
<td>6%</td>
<td>45%</td>
</tr>
<tr>
<td>3</td>
<td>$\sum I_i r_i r_i r_i$</td>
<td>44%</td>
<td>100%</td>
</tr>
</tbody>
</table>

The percentages mentioned here only give a rough indication.

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References

POLOIDAL FIELD SYSTEM AND MHD EQUILIBRIA FOR THE IGNITOR-U EXPERIMENT

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1. Introduction

We present the results of a study of the poloidal field (PF) system for the IGNITOR-U (Upgrade) device /1/. The design of this system represents a crucial problem for a compact, high magnetic field and high plasma current machine.

The analyses performed include evaluation of static magneto-hydrodynamic(MHD) equilibria at different times of the discharge, including start-up.

The PF coil currents were estimated, by starting from a set of optimal plasma confinement configurations and taking into account the relevant inductive flux needs.

Different situations were studied, by verifying that the relevant electromagnetic forces acting on the coils can be withstood and the most promising configuration is described here.

A considerable paramagnetic effect was found due to the high plasma current.

The safety factor at the plasma edge is found to be well above the limit established by the most recent experimental results ($q_{\psi} \gtrsim 3$).

2. MHD Equilibria

The peak design parameters of IGNITOR-U are:

- major radius ($R_o$) = 116.5 cm; $a = 43.3$ cm; $b = 78$ cm; triangularity = 0.29;
- plasma current = 12 MA; toroidal field = 13.2 tesla on $R_o$.

The current pulse of the design includes also a flat-top phase with a 9.4 MA plasma having the same geometrical size and toroidal field 10.3 tesla /1/.

Two different modes of operation are foreseen:

1) plasma bounded by a carbon limiter at the inner edge of the vacuum chamber ($R = 73$ cm).
2) plasma bounded by a magnetic separatrix with X-points at top and bottom of the torus.

We expect that IGNITOR-U will operate at 12 MA with a limiter-bounded plasma and at 9.4 MA with a limiter-bounded plasma or a separatrix defined boundary.

For the limiter-bounded configurations the static MHD equilibria
were computed by a free-boundary code/2/ which solves the Grad-Schlüter-Shafranov equation with a given toroidal current density normalized to yield a specified total plasma current. The toroidal current distribution adopted is (see Fig.2):

$$j_T(\psi, R) = \frac{d}{d \psi_0} \left( R + c/R \right) \psi^K \cdot 4 \cdot R \delta$$

with $c = R^2_0 \left( 1 - \frac{1}{\beta_p} \right) / \beta_p$

In these cases we chose $k = 1.8$ and $\beta_p = .25$. Smaller beta values were adopted for the initial plasma equilibria.

The plasma poloidal current $I_p$ and the corresponding paramagnetic effect were evaluated and found to be important. For the 12MA plasma (see Fig.4) the poloidal current is about $I_p \approx \text{8MA}$ and the paramagnetic effect results to be as high as 11% at $R_0$ (see Fig.3). Thus the maximum toroidal field can be as high as $B_T = 1.32\times 1.11 = 147$ kG.

The safety factor at the plasma edge is $q_e = 3.13$ for the 12MA plasma and $q_e = 3.07$ for the 9.4MA plasma. These values are well above the limit found to guarantee stability against disruptions.

Separatrix configurations can be easily obtained by slightly changing the PF currents (see Fig.5).

The profiles given by these computations can be used as input to an equilibrium code in the flux variable (ESSCO)/3/, which represents the equilibrium package of a 1-1/2 D transport code (JETTO)/4/.

3. Poloidal Field System.

In the design of the poloidal field system the functions of the transformer and of the equilibrium coils were combined, in order to optimize their effects and to reduce the current requirements /5/.

So the transformer coil Tl (see Fig.1) in its negative swing phase produces a shaping effect on the plasma.

The coil position, close to the plasma surface as in previous versions of the design /6/, satisfies the condition required to guarantee the existence of elongated equilibria /7/.

The Tl coil configuration was optimized by increasing the inner radius in order to yield a high flux value with a current density which is consistent with accepted engineering design criteria ($j_{\text{max}} = 11kA/cm^2$).

The considerable value of the flux needs is a strong constraint in the design of the PF system.

The magnetic configuration at the start-up was also investigated (see Fig.1), and a low value of the stray field at the plasma centre (which also is an important constraint) was obtained ($|B_p| (R_0) = 6$ Gauss).

The flux required by the plasma in order to attain a given current value and the flux generated by the PF system in the various phases of the discharge are given in Table I. A margin of 2.5V-s is available for the resistive losses during the first (maximum field) flat-top phase.

Note that the 12MA plasma inductance is 1.85$\mu$H, so that the purely inductive flux can be estimated to be 1.8$x12 = 22.2V-s$.

The system of forces acting on the coils was evaluated at different stages of the discharge. For the end of the first flat-top (12MA plasma) the most pessimistic case, i.e. the one without plasma, was investigated.
The total forces acting on T1 and T2 at the start-up are:
\[ F_r(T1) = 186 \text{MN/m}; \ F_z(T1) = 8 \text{MN/m}; \ F_r(T2) = 139 \text{MN/m}; \ F_z(T2) = -46.7 \text{MN/m}. \]

A complete stress analysis of the structure, including the PF coils, was carried out.

The stress distributions evaluated at different pulse phases are found to be acceptable according to current magnet design criteria.

In conclusion the results obtained show that the expected plasma magnetic configurations can be reached with a good safety margin while respecting realistic design conditions for the adopted magnet systems.

We thank F.J. Helton for her invaluable insight and contribution.

References

/5/ D. J. Strickler, Private Communication.

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Fig.1- Start-up configuration
Fig. 2—Toroidal current density in the equatorial plane

Fig. 3—Magnetic field profiles in the equatorial plane

Fig. 4
Limiter configuration

Fig. 5
Separatrix configuration

I = 12 MA
R₀ = 119 cm.
a = 43.3 cm.
b = 78 cm.
$q_a = 0.91$
$q_\infty = 3.13$

I = 9.4 MA
R₀ = 118 cm.
a = 40.3 cm.
b = 78.4 cm.
$q_\infty = 0.77$
$q(95\%) = 3.66$
1. INTRODUCTION

The stability of toroidal plasmas is often described by approximate MHD equations for a small number of scalar functions ("reduced MHD equations"). We propose a general representation of toroidal MHD in terms of stream-functions and potentials, which facilitates the physical interpretation of various terms and therefore is well suited for deriving different types of reduced equations. We qualitatively discuss stationary plasma states bifurcating from a given equilibrium, and we treat the particular case of a MHD-unstable cylindrical equilibrium.

2. A REPRESENTATION OF TOROIDAL MHD

We consider one-fluid MHD equations including viscosity $F_{\text{vis}}(V)$, resistivity $\eta$, heat flux $q$, and sources $S$. We use cylindrical coordinates $(r, z, \phi)$ and introduce poloidal differential operators acting only on $r, z$; e.g.

$$\nabla_p = \nabla / \partial r + \nabla / \partial z.$$

We choose a suitable time-independent reference magnetic field $B_0$ (e.g., toroidal vacuum field, or an equilibrium field, or a nonlinear stationary solution of MHD equations), and define

$$4\pi j = \nabla \times B_0, \quad b = B_0 / |B_0|, \quad \kappa_0 = b \cdot \nabla b.$$

We then show that quite generally, the magnetic vector potential may be written as [1,2]

$$A = - \nabla \phi_p - R^2 \nabla U \times \nabla, \quad - \Phi \nabla \phi$$

with a gauge condition fixing $\phi_p$. The electromagnetic fields are
\[
B = \begin{pmatrix}
B_x \\
B_y
\end{pmatrix} V_x - V_y \left( \frac{\partial B_z}{\partial y} - \frac{\partial B_y}{\partial z} \right) - V_z \times B, \quad E = - \frac{\partial A}{\partial t} - V \times B + E_{\text{loop}}
\]

The velocity field will be represented in a different way [2]:

\[
V = V_b + V, \quad V \times B = \nu_{v} + \left( R^2 \Lambda, u \right) V - V (\partial u / \partial t) - V \times V
\]

Thus instead of \( V \), we have represented \( V \times B \) by scalar functions. Putting \( \vec{J} = \vec{A} \times \vec{B} \), \( \vec{J}_v = \vec{A} \times \vec{V} \), \( \vec{J}_u = \vec{A} \times \vec{U} \), we derive from the MHD equations a set of six potential equations for \( U, \Phi, J, u, \dot{u}, \Phi \), and seven evolution equations for \( p, \rho, T, J, V, U, S \). In particular (with Coulomb gauge),

\[
\nabla \cdot \vec{J} = \frac{\rho}{\mu_0} \nabla \cdot \nu_{v} + \nabla \cdot (V \times \vec{B} - \vec{j} \times \nabla \Phi), \quad \vec{B} = \vec{B} - \vec{B}_0 \tag{1}
\]

\[
\rho \frac{\partial \nu_{v}}{\partial t} = \vec{B}_0 \cdot \left[ \kappa + \frac{\nu_{v}}{\rho_0} \nu \left( \frac{1}{2} \nu_{v}^2 \right) \right] \nabla \Delta \vec{B}_0 + \frac{\nu_{v}}{\rho_0} \vec{B}_0 \cdot \nabla \left[ B_0^{-2} \left( \frac{\vec{B} \cdot \vec{B}_0}{B_0^2} \right) \right] + \frac{4 \pi \nu_{v}}{\rho_0} \left( B_0 \cdot \vec{j} \right) \left( B_0 \cdot \nabla \Phi \right) - \left( j \cdot \vec{B}_0 \right) \vec{B}_0 \cdot \nabla \left[ B_0^{-2} \left( \frac{\vec{B} \cdot \vec{B}_0}{B_0^2} \right) \right] + \left[ \left( B \cdot \vec{B}_0 \right) \vec{B}_0 + \frac{B_0^2}{
abla \times \vec{B}_0} \right] \cdot \nabla \left[ B_0^{-2} \left( j \times \vec{B} \right) \right] + \nabla \cdot \left( \vec{B}_0 \times \left[ V \cdot \nabla V - \rho^{-1} \left( F \right) \nabla \right] \right) - \vec{B}_0 \cdot \left( \rho^{-1} \vec{V}_0 - B_0^{-2} \vec{V}_0^2 \times \left( j \times \vec{B}_0 - \nabla \Phi \right) \right) \tag{2}
\]

The advantages of our representation are the freedom of the choice of the reference field \( B_0 \) and the possibility of easy interpretation of various terms. In particular the divergence of the curly bracket in eq. (1) is often negligible, hence \( \nabla \cdot \vec{J}_u \), so that the stream function \( \Phi \) represents essentially the \( (E \cdot B) \)-drift. Also the evolution equation (2) for \( \nu \) which is essentially the vorticity density of the perpendicular motion) allows to recognize the driving terms of MHD instabilities (e.g. the first term is the ballooning term). This facilitates the approximations to make for deriving reduced equations.

3. BIFURCATION OF NONLINEAR SOLUTIONS

Dissipative nonlinear systems with given boundary conditions possess in general time-asymptotic solutions (attractors). The bifurcation of such solutions from an equilibrium, which becomes unstable when a given parameter varies, has been treated earlier for the interchange and the tearing instabilities [3]. In the first case, a nonlinear state of stationary convection is reached, in the second a stationary magnetic
island is formed. In the next paragraph we present the case of a stationary kink-type state. In all three cases it is convenient to use a representation in terms of stream functions.

4. BIFURCATION OF KINK-TYPE STATES

The formalism of paragraph 2 can be generalized, under some conditions, to arbitrary (non-orthogonal) coordinate systems, in particular to helical systems. We have considered the case of a cylindrical plasma of radius \(a\) and length \(L\), using an helical coordinate system \((r, \chi = \theta - kr, z)\), with helical symmetry (functions depending only on \(r\) and \(\chi\)) and using \(F_{vis} = \mu \Delta W\) and \(\nabla \cdot q_n = \kappa \Delta p\). Assuming small ratio \(a/L\), constant density, small parallel velocity and neglecting compressibility we derive a reduced system for \(p, \psi\) and \(\Phi\). Neglecting compressibility is justified if \(\eta/\kappa = 8\pi/3\), we therefore assume this relation for simplicity. Considering a static equilibrium with constant axial current \(B_1/Z\pi a\) we can write the following equations for the dimensionless deviations from equilibrium (pressure and magnetic field are normalized to characteristic equilibrium magnitudes, time to Alfvén time and length to \(a\)).

\[
\begin{align*}
\frac{\partial p}{\partial t} &= S^{-1} \Delta p - \left(\frac{4B^2}{\Omega_0^2}\right) \frac{\partial^2 \Phi}{\partial \chi^2} + \frac{1}{r} \frac{\partial}{\partial r} \left( r \frac{\partial \Phi}{\partial r} \right) + \frac{1}{r} \frac{\partial}{\partial r} \left( r \frac{\partial \Phi}{\partial r} \right) + \frac{1}{r} \frac{\partial}{\partial r} \left( r \frac{\partial \Phi}{\partial r} \right) \\
\frac{\partial \psi}{\partial t} &= S^{-1} \Delta \psi + \left( \frac{\delta^* B_1}{\Omega_0} \right) \frac{\partial^2 \Phi}{\partial \chi^2} + \frac{1}{r} \frac{\partial}{\partial r} \left( r \frac{\partial \Phi}{\partial r} \right) + \frac{1}{r} \frac{\partial}{\partial r} \left( r \frac{\partial \Phi}{\partial r} \right) + \frac{1}{r} \frac{\partial}{\partial r} \left( r \frac{\partial \Phi}{\partial r} \right) \\
\frac{\partial \Delta \Phi}{\partial \chi} &= P_R S^{-1} \Delta \Phi - k^2 \beta \frac{\partial p}{\partial \chi} + \left( \frac{\delta^* B_1}{\Omega_0} \right) \frac{\partial \Delta \psi}{\partial \chi} + \frac{1}{r} \frac{\partial}{\partial r} \left( r \frac{\partial \Phi}{\partial r} \right)
\end{align*}
\]

where \(\beta\) is the plasma beta, \(S\) the Lundquist number and \(P_R = \mu/\eta\), with the boundary conditions \(p=\psi=\Delta \psi=\Delta \Phi=0\) at \(r=1\).

In a linear approach we assume normal modes proportional to \(\exp(i m \chi + \gamma t)\). The general solution and the dispersion relation are given by \((Z_m\) being the first root of the Bessel function \(J_m\))

\[
p = -\left(4k/\beta \delta^*\right)/(1-\delta^*) \psi = A J_m (Z_m r) \cos m \chi \ e^{\gamma t}
\]

\[
\gamma^2 + \gamma S^{-1} (P_R + 1) Z_m^2 + S^{-2} P_R Z_m^4 + k^2 m^2 \left( \delta^* - 4k^2 / Z_m^2 \right)/(1-\delta^*) = 0
\]

The analysis of the dispersion relation shows that there are two marginal points \(\delta_1^*\) and \(\delta_2^*\), and that in the vicinity of a marginal point
there is a single real eigenvalue crossing the imaginary axis when $\delta^*$ is varied (fig.1). The growth rate is real between the two marginal points (fig.2). These statements allow to conclude that a nonlinear stationary solution branch bifurcates at each marginal point with exchange of stability between the equilibrium and the new branch.

The arbitrary amplitude A of the linear solution can be determined by the nonlinear terms using an integral equations method and a quasilinear approximation. In the immediate vicinity of the bifurcation points we obtain

$$\frac{1}{8} \sum_{m} \delta^* A^2 = \left[ \delta^* - \delta_1^* \right] \left[ \delta^* - \delta_2^* \right], \quad c_m = \int_0^1 r^{-2} dr \left\{ J_m^2(Z_m r) \right\}$$

This is the signature of a pitchfork bifurcation (fig.3). In order to investigate the behaviour of the branches in the vicinity of $\delta^* = 0$ the effects of compressibility must be taken into account.

Near the bifurcation points, the m=1 stationary state has helically shifted magnetic surfaces, and two convection cells with velocity proportional to $\eta$ (hence small if $\eta$ is small).

REFERENCES
THE FINAL PHASE OF JET DISRUPTIONS

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Abstract. It is widely believed that the loss of plasma energy in a disruption is due to the generation of a turbulent magnetic field. In studying fast disruptions in JET, it is found that the irreversible nature of the disruption is better understood in terms of radiation cooling.

Introduction. JET disruptions often end with a very fast decay of the plasma current. Fig.1 shows the plasma current for a 5MA discharge which terminated in a disruption. The current decay rate reaches 850MA s⁻¹, compared with 0.5MA s⁻¹ for a non-disruptive discharge. The dramatic difference in decay rate shows that a large change in plasma properties occurred in the disruption.

![Fig.1 Plasma current for disrupting discharge.](image1)

![Fig.2 Expanded trace of plasma current and loop voltage.](image2)

Temperature Inferred from Resistance. The plasma resistance during the current decay is calculated by equating the ohmic dissipation to the known input power

\[ I^2R_p = \frac{d}{dt} \left( \frac{1}{2} L_1 I^2 \right) + V_L \]  

(1)

where \( R_p \) is the plasma resistance, \( L_1 \) the internal inductance and \( V_L \) the toroidal loop voltage at the plasma surface. Fig.2 shows an expanded trace of the plasma current and voltage from one loop on the vacuum vessel for the same discharge as Fig.1. \( L_1 \) was calculated to be 2\( \mu \)H before the disruption and the average loop voltage reaches 300V during disruption.
The calculated input power during the current decay reaches ~10GW compared with ~10MW in a non-disruptive discharge. The resulting plasma resistance is = 400µ. If the resistance is due to Spitzer resistivity then, taking \( Z_{\text{eff}} = 3 \), the implied electron temperature is

\[ T_e = 5 \text{eV} \]

Temperature Inferred from Runaway Electrons. The current plateau during the current decay shown in Fig. 2 is a standard feature of fast disruptions on JET, and sometimes persists for 100's of ms. The plateau is due to a current carried by runaway electrons and the number of runaways generated can be used as an indication of the electron temperature.

Calculations of the rate of generation of runaway electrons \([1,2,3]\) show that the important parameter is the ratio of the electric field to a critical value, the Dreicer field, \( E_D \). If the electron temperature is suddenly reduced in the disruption, large electric fields are induced. For a current decay with time constant, \( \tau \), the fractional density of runaways generated is given approximately by \([3]\)

\[ \frac{n_R}{n_e} = \left( \frac{E_D}{E} \right)^{s/3} \exp \left( - \frac{E_D}{4E} \right) \]  

where \( \tau_e \) is the electron collision time. For the observed runaway current of 2MA, approximately 10^-8 of the electrons are runaways and for \( \tau = 10 \text{ms} \) this implies

\[ E_D = \frac{2.6 \times 10^{-9} n_e \pi \Lambda (2\pi R)}{T_e \varpi} = 40 \]  

with \( T_e \) in eV. \( \varpi \) is the ohmic dissipation and \( E_D \) is taken from \([2]\). With \( n_e = 6 \times 10^{19} \text{m}^{-3} \), Eq. (3) gives

\[ T_e = 4 \text{eV} \]

Temperature Achieved by Anomalous Conduction. The equilibrium temperature of a plasma confined in an ergodic magnetic field is estimated by writing the effective conductivity as

\[ \chi_L = \left( \frac{\tilde{B}}{B} \right)^2 v_{te}^2 \tau_e \]  

where \( \tilde{B}/B \) is the level of magnetic field fluctuations, \( v_{te} \) the thermal velocity of electrons. Assuming a parabolic temperature profile and Spitzer resistivity, and balancing the Ohmic heating against the conducted power gives an expression for the central temperature

\[ T_{eo} = 20 \left( \frac{B_0 Z_{\text{eff}}^{2/5}}{\tilde{B}/B} \right) \text{eV} \]

where \( \ln \Lambda = 10 \) has been assumed. Fig. 3 shows a plot of this function for \( Z_{\text{eff}} = 3 \) and \( B = 3 \text{T} \). Even with the perturbed field strength comparable to the poloidal field, the temperature is an order of magnitude larger than the experimental value inferred above. Measured values of \( \tilde{B}/B \) are typically <10% \([4]\).
It is seen, therefore, that the drop in electron temperature at the disruption is not due solely to enhanced transport across the minor radius. The alternative to thermal conduction is radiative cooling. For instance, an influx of impurity atoms at the disruption would cool the plasma while being ionised, then hold the temperature low by radiation.

**Fig. 3** Predicted central temperature versus field fluctuation level. **Fig. 4** Soft X-ray emission for 3ms before disruption.

**Soft X-ray Emission.** The disruption observed in the soft X-ray emission proceeds in two stages, illustrated in Fig. 4. First there is a broadening of the profile, taking ~2ms, followed by a spike in emission, lasting for ~200\(\mu\)s. Following the spike the plasma current increases momentarily and then decays away. Fig. 5 shows this behaviour for one central soft X-ray chord with the plasma current also shown for reference.

The final spike in soft X-ray emission can be understood in terms of a rapid increase in impurity content. For low Z ions and for the X-ray filters used, the emission has the dependence

\[
P_X = n_e n_z T_e^4.
\]  

Incoming neutral atoms increase \(n_z\) but decrease \(T_e\). Writing the impurity content as \(n_z = n_0 (1+\alpha t)\) gives

\[
P_X = P_0 (1+\alpha t) (1 - \frac{n_0 e_i}{W_0} \alpha t)
\]

where \(P_0\) is the initial X-ray intensity, \(E_i\) is the ionisation energy of the incoming neutrals and \(W_0\) is the initial electron energy density.

For carbon or oxygen impurities at typically 4% concentration, Eq. (7) predicts an increase in soft X-ray emission by a factor of 5 to 10. This simple model shows how a large increase in soft X-ray emission can occur, but does not give a mechanism for the impurity influx.

**Carbon III and Visible Bremsstrahlung.** Fig. 6 shows the intensity of a carbon III line (3p-3s transition at 465nm) and of visible bremsstrahlung (at 523.5nm) along lines of sight through the plasma centre. The disruption occurs at 10.522 seconds and there is an increase by a factor of 1,000 in the carbon III intensity, and a factor 200 in the visible
bremsstrahlung. Examination of several chords shows the carbon radiation comes from the plasma bulk rather than an edge shell.

Under these conditions, carbon III emission is a strong function of temperature, peaking around 7eV. The increase in intensity can be understood merely in terms of a drop in electron temperature from \(\approx 500\text{eV}\) to \(\approx 5\text{eV}\). The visible bremsstrahlung emission requires a change in \(Z_{\text{eff}}(n_{\text{e}}d_{\text{l}})^2\) by a factor of 10. For instance if \(Z_{\text{eff}}\) increases by \(-50\%\), then \(n_{\text{e}}d_{\text{l}}\) must increase by a factor \(= 3\). The observed behaviour is consistent with an influx of neutral carbon atoms, such that the carbon density increases to a value comparable with the deuterium ion density, that is \(n_{\text{c}} \approx n_{\text{D}}\).

![Fig. 5](image1.png) Plasma current and emission along one soft X-ray chord during disruption.

![Fig. 6](image2.png) Carbon III line and visible bremsstrahlung emission during disruption.

**Summary**

1. The inferred electron temperature in JET during the current decay of a fast disruption is \(\approx 5\text{eV}\).
2. The lowest temperature which could be achieved by anomalous conduction is \(>50\text{eV}\).
3. Radiation must play a role in the energy loss. A sudden increase in the plasma impurity content is required.
4. The soft X-ray spike at the disruption is consistent with a rapid influx of impurity atoms.
5. Carbon III and visible bremsstrahlung measurements are consistent with a temperature of \(\approx 5\text{eV}\) and a factor \(= 3\) increase in electron density. This is again consistent with an influx of impurity atoms.

**References**

The observed photoneutron yields due to major plasma disruptions in JET fall in the range from $10^{10}$ to $10^{15}$ neutrons, with instantaneous intensities reaching $5 \times 10^{17}$ neutrons/sec. The runaway electron energies can exceed 60 MeV and the currents of high energy electrons can reach 1 MA. On average, the fraction of the available poloidal magnetic field energy that is transferred to high energy runaway electrons is $f = 3 \times 10^{-2} I_p^6$, with plasma current $I_p$ in MA.

1. INTRODUCTION

Major plasma disruptions in JET are frequently accompanied by the acceleration of a strong current of high-energy runaway electrons, which produce photoneutrons when they eventually strike the vacuum vessel wall protection tiles or one of the plasma limiters. Measurement of the photoneutron yield constitutes a simple and non-invasive method of estimating the amount of energy carried by the runaway electrons.

The detailed theory of electron runaway in tokamaks has been given by Knoepfel and Spong [1], who concentrate on the birth process under non-disruptive conditions of relatively low loop voltage. Our present interest is with strong disruptions for which the magnitude of the measured loop voltage may exceed 500 V for periods of 5 ms or more. This implies that the maximum energy to which the runaway electrons may be accelerated easily exceeds 60 MeV. As we have no means of measuring the energy distribution of the runaway electrons, we assume an average energy of 60 MeV. Those electrons which produce photoneutrons must have energies well in excess of 7 MeV, the photoneutron reaction threshold for the elements contained in inconel (the material from which the vacuum vessel is constructed). Also, studies of induced activation in the vessel reveals the existence of radionuclides, such as $^7$Be in the $^{12}$C limiters, for which the reaction thresholds were over 26 MeV. As we have seen, energies calculated from the measured loop voltages are very much higher.

2. OBSERVATIONS

The time-resolved neutron emission was measured with three pairs of $^{235}$U and $^{238}$U fission chambers disposed around the tokamak. The source of
photoneutrons could be highly localized, eg. a limiter. For such a source, the response from a single fission chamber could vary by a factor of 20, depending on the relative positions of source and chamber. In practice, the chambers usually give similar responses (to within a factor of 2) and only rarely show a factor of 10 variation from chamber to chamber. A single strike is therefore unusual. In any case, averaging the results of all three sets of chambers reduces the positional dependence to a factor of 1.5. The calibration to be associated with the chambers is that measured for a deuterium plasma discharge [2] and is known to ± 10% accuracy. The fission chambers are operated in both pulse-counting and current sampling modes and operate over an intensity range from $10^3$ to $10^{22}$ neutrons/sec.

A typical example of the strong neutron emission which arises when a high plasma current discharge is terminated by a sudden disruption is shown in figure 1. The plot shows how the plasma current falls from $-2.5$ MA to $-0.2$ MA in a time interval of only 40 ms. There is a distinct hesitation midway down the plasma current collapse. The plot also shows the strength of the neutron production on a logarithmic scale covering six decades. The duration of the neutron pulse coincides with the hesitation in the plasma current decay. The obvious implication is that a runaway electron current of 1.5 MA has been generated. This hesitation in current decay is less pronounced for disruptions producing fewer neutrons, and imperceptible for those with weak emissions.

Figure 1:
Illustrating the relationship between the rapid decay of the plasma current during a disruption and the strength of the photoneutron emission. The total neutron emission is $5 \times 10^{13}$ neutrons.
A total of 201 disruptions have been studied, covering the period between the startup of the tokamak and December 1986. These disruptions included all those producing at least $10^{11}$ photoneutrons; during this period there were nearly 10,000 discharges.

The variation in photoneutron yield for a particular set of plasma conditions is very considerable and it is difficult to discern how the photoneutron yield is to be predicted. The clearest correlation is obtained by plotting the average yield against plasma current at the instant of disruption, as shown in figure 2.

The yield is found to be fitted reasonably well with the variation $N = 7 \times 10^{12} I^{2.6}$ observed neutrons/disruption ($I$ in MA). The photoneutrons are mostly produced several centimeters deep in the structure surrounding the plasma; it is estimated that only 1/3 of these enter the plasma region and so become available for detection. The true yield of produced neutrons

\[ N = 7 \times 10^{12} I^{2.6} \]

**Figure 2:**
A histogram showing the average number of photoneutrons produced per disruption as a function of plasma current at the instant of the disruption. For each plasma current interval (0.5 MA), the number of disruptions in the class is indicated.
per disruption is therefore \( N = 2 \times 10^{13} I^{1.6} \) (I in MA). The yield per runaway electron can be calculated [3]; for example, for electrons of energy \( E_0 \) (MeV) and a carbon tiled torus with inconel surrounding structures, this calculated yield is

\[
Y = 4.8 \times 10^{-5} E_0 \text{ neutrons/electron}
\]

Thus, the number of electrons per disruption is

\[
N/Y = 4.2 \times 10^{17} I^{1.6} E_0^{-1} \text{ electrons/disruption}
\]

These are relativistic electrons, making \( 1.6 \times 10^7 \) revolutions of the torus per second. They give rise to a current of \( 1.1 \times 10^3 I^{1.6} E_0^{-1} \) kamps. Clearly, runaway electron currents of up to 1 MA are possible.

The total kinetic energy carried by these electrons is

\[
NE_0/Y = 0.07 I^{1.6} \text{ (MJ)}
\]

This is very high, 5 MJ for 5 MA disruptions. The total energy available for accelerating electrons is that associated with the poloidal magnetic field

\[
W_M = \frac{1}{2} LI^2 = 2.4 I^2 \text{ (MJ) with } L = 4.8 \mu\text{H}
\]

On average, the fraction converted to runaway electrons is

\[
f = 2.9 \times 10^{-2} I^{0.6}
\]

This fraction is relatively insensitive to the energy assumed for the runaway electrons, provided it is higher than 25 MeV. This estimate should not be relied upon to better than a factor of 3.

3. CONCLUSIONS

The essential conclusions from this study are

(i) Runaway electron energies are high (over 60 MeV) due to the measured loop voltage in a disruption exceeding 500 V for periods of 5 ms or more.

(ii) The neutron emission scales with plasma current as \( I^{1.6} \), so that a fraction \( 3 \times 10^{-2} I^{0.6} \) (I in MA) of the poloidal magnetic energy is converted to runaway electron production. (It should be noted that the neutron yield from individual disruptions fluctuate considerably about the average yield for a particular plasma current).

(iii) The energy carried by the runaway electrons is considerable (5 MJ for a 5 MA plasma). This energy is normally dissipated at a small number of locations around the vacuum vessel.

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INTRODUCTION

The main impurities in JET plasmas are carbon, oxygen and nickel, with some chlorine occasionally. The nickel concentration $C_{\text{Ni}} = n_{\text{Ni}}/n_e$ depends critically on previous history of ion cyclotron heating (ICRH) and is usually $\leq 10^{-5}$ in ohmic heating campaigns. During ICRH, nickel is released from the antenna screens and subsequently deposited on limiters and protection tiles. After high power ICRH pulses, $C_{\text{Ni}} = 10^{-3}$ has been measured during ohmic phases and several $10^{-2}$ during ICRH. Oxygen concentrations ($C_0 = n_0/n_e$) have recently been 0.5-1% in deuterium (D) plasmas. In helium plasmas, $C_0$ was about a factor 10 lower due to smaller oxygen influxes from walls and limiters. This lower oxygen production by helium (He) clearly indicates a chemical production mechanism. Carbon appears to be produced by physical sputtering and was several % in both D and He discharges with a tendency to increase during additional heating.

Impurity transport in JET has been studied from accidental metal injections and from the radial impurity ion profiles. Confinement times and profiles could usually be described by an anomalous diffusion coefficient $D = 1 \text{ m}^2/\text{s}$ and by an inward drift ($\propto r$) with $v_D(a) = 2 \text{ m/s}$. Under some recent plasma conditions, impurity accumulation was expected, as reported in other experiments in sawtooth-free plasmas, H-mode plasmas and discharges with peaked electron density profiles after pellet injection. Long sawtooth-free periods have been observed in JET during ICRH ("monster sawteeth"). Respective changes of the total ion profiles of nickel, which is most likely to accumulate on axis, have been studied by means of line radiation from different ionization stages originating from various radial locations. These have been chosen according to the electron temperature, $T_e(0)$ ranging from 1 keV after pellet injection to 10 keV during monster sawteeth. The light impurity behaviour has been investigated by means of $Z_{\text{eff}}$ profiles, charge-exchange recombination spectroscopy and from soft X-ray radiation which is often dominated by light impurity bremsstrahlung.

MONSTER SAWTEETH

Monster sawteeth were observed in both D and He plasmas with RF powers of 2-15 MW. Respective nickel concentrations in these plasmas ranged from a few $10^{-5}$ to several $10^{-3}$. Because of deposition and subsequent erosion of nickel, $C_{\text{Ni}}$ usually continued to increase during ICRH in discharges with and without sawteeth. Compared to normal sawtoothing discharges, some 20% increase in $C_{\text{Ni}}$ was observed during monster phases, probably due to a slight increase in global impurity confinement. After the monster crash, impurity confinement often dropped significantly leading to substantial
loss of nickel even with ICRH continuing at the same power level.

Fig.1 shows \( \sigma_{Ni} \) derived from Ni XXVII in the plasma centre and Ni XXV close to the edge during ICRH (\( T_e(0) = 6 - 10 \text{ keV} \)). Both ICRH phases of the pulse lead to higher nickel concentration. During the sawtooth-free period (6.2 - 8.7s), \( \sigma_{Ni} \) continues to increase but in the same way at different radial locations. This means that the nickel ion profile does not change, i.e., no nickel accumulation in the centre. Detailed investigations of sawtooth-free periods up to 3s have not shown any change of impurity transport. It is concluded that the absence of sawteeth, at least during ICRH, does not lead to accumulation of medium-Z impurities in the JET plasma.

**PELLET INJECTION**

Single pellets and sequences of D pellets of different size have been injected into JET discharges (the multi-pellet injector was provided by ORNL under a JET-USDoE Agreement). ICRH and neutral beam heating (NBI) have been applied to the resulting plasmas. Pellet injection usually led to electron density profiles which, after some initial rearrangement, were not significantly different from those produced by gas fuelling. In some cases when pellets penetrated to the plasma centre, strongly peaked \( n_e \) profiles were observed lasting up to 3s. Fig.2 shows an example of profiles peaked for about 3s after the second pellet. \( T_e(0) \) drops to about 0.8 keV and slowly recovers to its pre-pellet value. Two nickel line intensities demonstrate the behaviour of the nickel ions: Ni XXVI, most abundant in the plasma centre below 2.5 keV, and Ni XI from the periphery under all conditions mentioned. While the Ni XI radiation reflects changes in edge density and some reduction in edge nickel content, a substantial increase in line radiation from the plasma centre is observed. Comparison with earlier phases of the pulse, taking into account changes in \( n_e(0) \), shows that the axial nickel density increases by a factor >10. This behaviour has been observed on other tokamaks before [1] and has been interpreted as neoclassical-type accumulation. In the present example, accumulation disappears after about 3s when the \( n_e \) profiles revert to normal. Accumulation has been observed to survive low power NBI, but disappeared immediately after the onset of ICRF heating.

Even under these conditions, the nickel concentration remains very low. \( Z_{eff} \) and soft X-ray profiles are still dominated by light impurities and suggest that light impurity profiles are peaked, too [2]. For a short period after the second pellet in Fig.2, light impurity concentrations are quite low (\( Z_{eff} = 1.5 \)). Therefore, the nickel behaviour expected from neoclassical theory [3] can be predicted without knowledge of the actual concentrations. Calculations demonstrate that the nickel density increase on axis and the observed time scale can be explained by a rearrangement of the nickel content due to the neoclassical inward drift, provided that the anomalous diffusion coefficient is reduced to \( \approx 0.01 \text{ m}^2/\text{s} \).

**H-MODE PLASMAS**

H-mode phases lasting up to about 2s had already been observed in 1986 [4]. These were characterised by a continuous increase in electron density and radiation losses. Radiation profiles were hollow in the centre but showed unusually wide edge shells. Spectroscopic analysis showed radiation losses mainly due to oxygen and, to a lesser extent, to carbon. Nickel may have contributed about 10% in some cases. Detailed investigations showed that the radiation profiles, the unusual behaviour of line intensities from
different nickel ionisation stages and longer particle confinement times could be explained by a high inward drift velocity in a narrow region near the plasma edge and hollow profiles in the centre, a consequence of the non-stationary nature of these discharges. For the latter explanation, the interior diffusion coefficient had to be $0.2 \text{ m}^2/\text{s}$.

Already in 1986, a few H-mode pulses had short periods of edge localised activity (ELMs) most obvious from spikes in $H_{\alpha}$ radiation [4]. With NBI heating extended to 6s in 1987, up to 5 H-mode phases were observed with ELM periods in between. During the first quiet H-mode, lasting again 2s, the 1987 and 1986 results were virtually identical. $\text{Z}_{\text{eff}}$ profiles, now available [5], were hollow as expected. The nickel behaviour during the subsequent phases is described in Fig.3.

During each ELM phase, particles are lost from the plasma periphery stabilising the density increase, on average. Global radiation losses (not shown) also stabilise, but are modulated according to the electron density behaviour. $T_e(0)$ decreases slowly after the initial H-mode phase. After some initial increase, nickel line radiation from the plasma edge (Ni XVII, Ni XVIII) drops during each ELM phase and generally decays to a small value. In contrast, Ni XXVI from the plasma centre continues to increase and a peaked nickel ion profile develops. Simultaneously, radiation profiles fill in (reaching $40 \text{ kW/m}^2$ on axis) and soft X-ray radiation increases. All these observations are an indication of impurity accumulation. However, $n_e$ profiles are flat or even slightly hollow. Global impurity confinement is clearly higher than usual by about a factor 5.

A closer investigation of $\text{Z}_{\text{eff}}$ profiles and soft X-ray profiles shows that nickel does not really develop a profile peaked on axis. The nickel density increases within the inner third of the minor radius, but the profile is actually hollow in the centre. Light impurity profiles remain hollow throughout all these phases and their concentration hardly changes. The plasma edge radiation is mainly due to oxygen, while the radiation increase in the interior is due to nickel. During the final H-mode phase in Fig.3, $\text{CN}_1 = 0.1\%$ results from a spectroscopic analysis in agreement with the local soft X-ray emission. The total nickel line radiation accounts for about half the local radiated power as measured by the bolometer. There is some contribution from chlorine, which tends to increase during the later phases ($c_{\text{Cl}} = 10^{-3}$). The remaining radiation fraction is due to light impurity radiation, i.e. mainly bremsstrahlung.

The observed rearrangement of impurity ions is not readily interpreted in terms of neoclassical transport. Certainly, the non-stationary nature of the edge densities must be considered and several impurity species taken into account.

CONCLUSIONS

Long sawtooth-free periods, observed during ICRH in JET, did not lead to any obvious change in impurity transport. In some cases after pellet injection, when strongly peaked electron density profiles were observed, medium (and low) $-Z$ impurities showed behaviour expected for neoclassical accumulation. ICRH appeared to destroy the accumulation process. During long periods of alternating H-mode and ELM phases, medium-Z impurities developed a generally peaked profile, but with a dip in the centre. Light impurity profiles were generally hollow. This rearrangement of impurity content was much slower than in pellet cases and there is no obvious explanation for this change.
FIG. 1: Nickel concentrations from Ni XXVII in the plasma centre and Ni XXV near the boundary. ICRH power and $T_e(0)$ are also shown. The nickel ion density profile does not change during the monster sawtooth.

FIG. 2: Densities, $T_e(0)$ and nickel line radiation after pellet injection. The increase in Ni XXVI radiation demonstrates a change in particle transport and neo-classical-type accumulation.

FIG. 3: Radiation from different nickel ionisation stages during alternating H-mode and ELM phases. A peaked nickel ion profile develops resembling neo-classical-type accumulation.

REFERENCES
MEASUREMENTS OF 'SNAKES' FOLLOWING MULTIPLE PELLET FUELLING OF JET

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INTRODUCTION

The extremely long lived density perturbation, or snake, first seen by the soft X-ray cameras following single pellet injection into JET, has now been observed following multiple pellet fuelling of JET discharges. The snake is caused by a local density perturbation rotating at a rational q-surface, normally the q=1 surface. The snake can persist for longer than 2 seconds, suggesting that a magnetic island is formed at the rational q-surface, with ablated pellet particles being deposited inside this island. The long snake lifetime implies a change to a new non-axisymmetric equilibrium.

New diagnostics have contributed important new results on the study of snakes. Particularly an array of toroidal soft X-ray cameras and the LIDAR Thomson Scattering system, which makes simultaneous measurements of the electron temperature and density profiles possible. These will be presented in this paper, with particular emphasis on the creation of snakes and the subsequent snake profiles.

CERTAIN NECESSARY CONDITIONS FOR THE CREATION OF SNAKES

The soft X-ray signals for two time periods of 1 millisecond each recorded during the injection of successive D₂ pellets, the second of which created a snake, are shown in Fig.1. The upper signals are from the vertically mounted camera which views the incoming pellet trajectory. The bottom signal is that of a central detector in the horizontally mounted camera which views the injected pellet from behind and gives a measure of the pellet ablation rate. The two soft X-ray cameras are located at the same toroidal position as the multiple D₂ pellet injector. Both pellet trajectories are clearly visible in Fig.1. In Fig.1a the maximum rate of pellet ablation is seen to occur around 60cms, which is outside the sawtooth inversion radius of 45cms, as determined from a tomographic analysis of the soft X-ray data. Throughout this paper the sawtooth inversion radius is taken as being equivalent to the q=1 position. For the second pellet, shown in Fig.1b, the sawtooth inversion radius is determined to be at 42cms. This pellet is observed to penetrate well beyond this radius, with considerable pellet ablation occurring in the region of the q=1 surface. Around such a rational surface, the particles would not be expected to spread rapidly. A local drop in temperature could therefore occur, as only the electrons inside a narrow flux tube which intersects the pellet trajectory would interact with the pellet.
FIG. 1: X-ray signals during the injection of two successive D₂ pellets, the second of which created a snake.

In Fig. 2, LIDAR profiles of the electron temperature and density across the horizontal mid-plane of JET are shown, 19 milliseconds after the injection of a pellet which created a snake. At the sawtooth inversion radius, a clear local density enhancement and temperature depression are seen and it is this density perturbation which is labelled a snake. Local cooling is therefore observed at the q=1 surface, which can lead to a helical current perturbation and the formation of a magnetic island.

Certain necessary, although perhaps not sufficient, conditions for the creation of a snake are, therefore, firstly that a q=1 surface must exist within the plasma, secondly that the pellet must reach this q=1 surface and thirdly that sufficient pellet particles are ablated in the region of the q=1 surface so as to lead to the formation of a magnetic island.

FIG. 2:
Radial electron temperature and density profiles, 19 milliseconds after the injection of a pellet which created a snake.
SNAKE PROFILES AND THE SNAKE LIFETIME

From the soft X-ray cameras and toroidal detectors it is found that the snake has an $m=1, n=1$ structure, with typical dimensions (FWHM) of 25 cm poloidally and 17 cm radially, which is in good agreement with the radial dimension of the perturbation determined from the LIDAR profiles of Fig. 2. From this knowledge of the mode structure, it can be determined from the soft X-ray data that the peak of the snake lies above the horizontal mid-plane at the toroidal location where the LIDAR profiles of Fig. 2 are measured. The maximum density and temperature perturbation at this time is, therefore, considerably larger. Indeed, for this particular snake, the ECE grating polychromator experienced a periodic 2nd harmonic X-mode cutoff at the snake radius. From the maximum frequency at which the cutoff was observed, the density of the perturbation was determined to be $1.5 \times 10^{20} \text{ m}^{-3}$ $\pm 10\%$ compared to the density measured by the LIDAR of $1.15 \times 10^{20} \text{ m}^{-3}$. Further, this is not seen to change over a period of 350 ms, at which time the plasma current was ramped down, implying an effective confinement time for the snake of the order of several seconds.

In Fig. 3, the radial soft X-ray profile measured by the vertically mounted camera is shown. A pellet is injected at 8.5 seconds which creates a snake. This is seen to rotate for a period of 900 ms before slowing down and stopping. This snake is observed to survive through many sawteeth, as is the case with all other snakes. At 10 seconds, when LIDAR profiles were measured, the centre of the snake is determined from tomographic analysis of the soft X-ray data to be just above the horizontal mid-plane at the LIDAR port. These LIDAR electron density and temperature profiles are shown in Fig. 4. The snake is clearly seen in the density profile, centred at 3.35 m. The temperature profile over the corresponding region of this snake shows the temperature falling from about 1700 eV at 3.15 m to 1450 eV at 3.5 m. This does not show any equalization on either side of the snake, as might be expected if the density is confined within a magnetic island.

FIG. 3: Soft X-ray flux around the time of pellet injection showing the 'snake' oscillation.
These LIDAR profiles are, however, taken at a time when the snake is 'locked' in one poloidal location. As the snake is an island deep within the plasma, if it is to 'lock', it must couple to a mode further out in the plasma which can interact with the vessel wall. The magnetic signals show the presence of a large locked mode at this time. It might not, therefore, be expected that the temperature would be equalized at the extremes of the island at this time, as the snake has coupled to another mode. From Fig. 3 it can be seen that, subsequent to these LIDAR profiles, the density perturbation decays away on a timescale of the order of 1 second. In other cases where the snake has been observed to couple to a 'locked' mode, then the density is seen to decay on a similar timescale. The application of on-axis ion cyclotron radio frequency heating (ICRH) is also observed to lead to the decay of the snake. The density is seen to decay on a timescale of the order of 100 milliseconds about 200 milliseconds after the ICRH is applied to the plasma.

SUMMARY

In conclusion, snakes have been observed during multiple pellet fuelling of JET. LIDAR Thomson scattering profiles of electron temperature and density provide clear measurements of the conditions necessary for the formation of a magnetic island at a rational q-surface in which a fraction of the ablated pellet particles can be trapped. ECE measurements in which a persistent periodic cutoff is seen at the snake radius implies a very long effective confinement time for the snake. The presence of 'locked' modes in the plasma, to which the snake can couple, or the application of on-axis ICRH leads to a gradual decay of the snake.

REFERENCES

EFFECTS OF LARGE AMPLITUDE MHD ACTIVITY ON CONFINEMENT IN JET

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Introduction Under many operating conditions in JET, particularly during long periods of additional heating, low m, n MHD activity (usually m=2, n=1 or m=3, n=2) persists with large amplitude (\(B_r/B_0\) (wall) \(\approx\) 0.12) for many seconds. This allows the effects of such large MHD activity on confinement parameters to be assessed. While other loss mechanisms are also important in determining general confinement scalings, these results indicate that large MHD activity can lead to substantial losses of particles, momentum, and energy from the plasma. Other effects such as changes in wall pumping, the plasma boundary, radiative losses, and input power must be carefully excluded from the analysis as they can also alter plasma confinement. For most of this analysis, discharges were chosen that exhibited large amplitude MHD activity arising after a large sawtooth collapse following a period of sawtooth stabilization. The confinement parameters were determined before the sawtooth collapse and again after more than one confinement time following the collapse, when they had reached a new equilibrium value with the same input power in the presence of an MHD mode.

In addition to this analysis, correlations have been made between slowly rotating (\(\sim 3 - 5\) Hz) low m, n (usually m=2, n=1) magnetic activity, called quasi-stationary modes (QSM's), and central m=1, n=1 activity observed on the electron temperature profile. QSM oscillations are observed on the \(H_\alpha/D_\alpha\) signals, the infrared limiter viewing camera, and on ECE signals from the edge right up to the center of the plasma. This strong mode coupling [1] from the center to the edge indicates how such large amplitude MHD activity can affect global confinement properties.

MHD Effects on Particle Confinement Edge influxes as measured by \(H_\alpha/D_\alpha\) and CIII emission signals are modulated up to 50% by slowly rotating QSM's. Similar modulation is also observed on an infrared limiter viewing camera, with good phase correlation with the mode indicating that the bright emission corresponds to the X point of the MHD mode.

Particle pumpout rates tend to increase with the amplitude of the QSM reaching rates two to three times as fast as without the mode under some conditions. Wall pumping and boundary effects, however, can alter the pumpout rates and mask the MHD effects [2]. Figure 1 shows a comparison of two similar X point configuration shots, the first with a QSM after the large sawtooth collapse at 12 sec and the second without a QSM. The fourth trace from the top shows that the number of particles drops by more than 12% in the first shot, but by less than 5% in the second shot, indicating that about 7% of the total number of particles were lost due to enhanced particle flux by the QSM. In the first shot, \(B_r/B_0\) (wall) \(\approx\) 0.23%. Similar enhanced pumpout rates are observed with QSM's after pellet injection, but the lack of similar shots with and without a QSM makes a quantitative comparison difficult.
Figure 1. Comparison of two similar shots with and without a QSM. a) \( T_e(0) \), b) QSM amplitude, c) diamagnetic stored energy, d) total number of particles, e) ICRH power. The degradation of particle and energy content is evident during the QSM.

**MHD Effects on Momentum Confinement** Neutral beam injection in the direction of the plasma current generally enhances the toroidal momentum of the plasma, but this extra input momentum can be reduced by the onset of low \( m, n \) MHD activity. Unlike the effect on particle and energy confinement, the plasma momentum appears to be affected by any observable amplitude of coherent MHD activity, in that there is almost always a good linear correlation between the MHD frequency of oscillation and the central ion toroidal rotation velocity [3]. This is particularly evident when mode locking occurs and the rotation ceases on the MHD signals as well as on the central ion toroidal rotation velocity, within the error of the measurement. Even in cases when the MHD activity is oscillating rapidly at amplitudes of only \( \frac{d_B}{B_0(\text{wall})} \approx 0.01\% \) (with \( m=3, n=2 \)), there can be a loss of central plasma momentum by more than 50%. Such cases clearly demonstrate that strong toroidal mode coupling is responsible for the degradation in the central momentum, perhaps being lost to the vacuum vessel walls through induced eddy currents.

**MHD Effects on Energy Confinement** Large amplitude MHD activity can also significantly degrade the plasma stored energy. Figure 2 shows that the amount of enhanced degradation in the stored energy, \( \Delta W/W \), depends almost linearly on the relative amplitude of the MHD mode, \( \Delta b_T(n=1)/B_0(\text{wall}) \). This data contains combined NBI, ICRH, and ICRH alone, fast oscillating \( m=3, n=2 \) modes, and slowly rotating or locked \( m=2, n=1 \) modes that arise after a large sawtooth collapse. The contribution to the drop in stored
Figure 2. Relative change in plasma stored energy before and during a QSM versus the relative QSM amplitude.

Figure 3. Effects of QSM on $T_e$ at different radii showing strong mode coupling between $m=1$ and $m=2$.

Figure 4. Thomson scattering electron temperature profiles at nearly opposite phases of the $m=2$, $n=1$ QSM oscillations showing a strongly coupled $m=1$ mode in the plasma center. A typical error bar shows the accuracy of the measurements. The phase of the oscillations on the QSM signal at the times of the two temperature profiles are indicated in the lower figure.
energy due to the sawtooth collapse itself is minimized by taking the measurement during the MHD activity only after more than one confinement time following the sawtooth collapse to allow the plasma to reheat.

The time evolution of this enhanced degradation is shown in Figure 1. The central electron temperature and stored energy remain low during the QSM on the first shot, but increase again with continued heating on the second shot without a QSM. The first shot had about a 17% drop in stored energy while the second shot without a QSM had an initial drop of 10% at the sawtooth collapse that subsequently returned to near the peak value of the stored energy before the ICRH was turned off. Since the drop in the electron temperature in both shots was about the same, the difference in the change in stored energy between the two shots must be due to the QSM. Note that locked modes arising before density limit disruptions also show a similar degradation in confinement [4].

Such evident MHD effects on global confinement parameters also manifest themselves directly in changes in the electron temperature from the edge right up to the center of the plasma. Figure 3 shows the strong toroidal mode coupling of the QSM oscillations near the plasma edge, analyzed with magnetic pick-up coils to be \( m=2, n=1 \), with oscillations on the ECE across much of the plasma. The \( m=1 \) character of the ECE oscillations is evident by the phase inversion across the plasma center. Note that the sawteeth persist despite this large \( m=1 \) mode. Further evidence of the effects of the QSM oscillations on the center of the plasma are shown in Figure 4. The LIDAR Thomson scattering system [5] acquired electron temperature profiles at nearly opposite phases of the QSM oscillations. The difference in the profiles indicates an \( m=1 \) structure across the plasma center with an amplitude of \( \pm 600 \text{ eV} \). Such a large change in temperature is well outside the errors in the measurements.

**Conclusion** The strong coupling between MHD modes observed from the limiter right up to the plasma center demonstrates that these modes can affect most of the plasma. This strong mode coupling must then be at least partly responsible for the observed degradation in particle, momentum, and energy confinement during large amplitude MHD activity. The correlation of the X point of the magnetic island with locally enhanced radiation suggests that particle and energy transport may be increased locally around the X point of the mode. The observed degradation in particle and energy confinement depends on the mode amplitude rather than on whether or not the mode is rapidly oscillating or locked, while the loss of momentum depends on the frequency of oscillation of the mode more than on the mode amplitude.

**References**


THE SAWTOOTH IN JET


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INTRODUCTION Detailed experimental studies [1] of the sawtooth collapse in JET have shown that it occurs very rapidly (~100µs) and X-ray tomography [2] has shown a strong m=1 component, with a spatial dependence in agreement with the predictions of an ideal mhd instability model [3]. However, there are still many remaining problems in our understanding of the sawtooth oscillation and some of these will be discussed below.

INSTABILITY GROWTH RATE The growth rate of the sawtooth instability has been determined by finding the radial centroid, $r$, of the X-ray emission function calculated by tomographic inversion of the data collected with the two orthogonally mounted JET X-ray cameras. This quantity is a sensitive measure of the displacement of the plasma central region and is independent of the relative phase of the m=1 instability present during the sawtooth crash. Fig.1 shows a rather slow movement of $r$ before the sawtooth collapse and the onset of a sudden growth in $r$ during the collapse. The time behaviour of $r$ is in good agreement with the perturbed magnetic field measurements made at the plasma edge [4]. Although the growth rate corresponds to the growth rate of an ideal mhd mode, there is no satisfactory explanation of the sudden increase in the growth rate of the instability or of its trigger mechanism.

The centroid is also a very sensitive monitor of plasma movements between sawtooth collapses and mhd activity and other effects are observed which are too small to be seen by individual X-ray detectors.

MONSTER SAWTOOTH During additional heating experiments it has been observed that there are long sawtooth free periods (~3s), terminated by a large sawtooth called the "monster". Tomographic analysis of these (Fig.2) shows great similarities with normal sawteeth. The core of the initial symmetric state (1) is rapidly (in ≤ 100 µs) displaced radially (2) until it reaches the inversion radius (shown as a dotted line in the Fig.2) where it spreads poloidally (3). The displaced core then collapses with a rapid loss of energy (4) followed by a gradual return to a symmetric but much flattened state (5). The topology of the monster collapse is identical with the normal sawtooth but the inversion radius is considerably larger (by up to 50%). The n=1 nature of the collapse is clearly shown by the X-ray measurements taken with a set of toroidally spaced detectors (Fig.3) and from ECE measurements. This is also confirmed by the magnetic measurements which show a dominant n=1 component with a small n=2 component, similar to normal sawteeth. Although the early stages of the
collapse can probably be explained in terms of the growth of an mhd instability the subsequent rapid loss of energy from the displaced core requires alternative explanations such as enhanced transport for both monster and normal sawteeth.

PARTIAL SAWTEETH Partial sawteeth with a very small collapse amplitude are frequently observed, nearly always accompanied by m=1 mhd activity. The fast part of the collapse has not yet been observed in detail but it seems to involve processes leading to a flattening of the profiles around the q=1 surface with the central parameters remaining relatively unaltered. A slower process has also been observed which leads to an even smaller redistribution of energy at the q=1 surface. The reasons for the different behaviour of the partial sawteeth are unclear.

q-PROFILE It is well known that the q-profile plays a strong role in the sawtooth instability and the sawtooth inversion radius is generally taken to correspond to q=1. This conclusion has been further confirmed on JET by the existence of the m=n=1 "snake" oscillation [5] which also exists at this radius. The identification of normal sawteeth with an interchange instability led to the conclusion that q(0)=1 and that the change in q(0) in the sawtooth crash was ~ 1-2%. This conclusion was supported by analysis of the movement of the snake during the crash.

However, the monster sawteeth show quite a different picture and raise new problems in understanding the sawtooth. Faraday rotation measurements [6], resistive diffusion calculation and magnetic equilibrium calculations all show that q(0) = 0.8 before both normal and monster sawtooth crashes and 6q in the monster crash is ~ 20%. These observations are not in agreement with the earlier understanding of the sawtooth crash and an alternative theoretical approach has therefore been investigated.

THEORY We have taken the point of view that, for normal sawteeth as well as for monsters, q(0) is significantly below one, and that the onset of the collapse is determined by the stability threshold of resistive m=1 modes. The stability plane (Fig.4) is determined by the parameters \[ K = - \left( \frac{r_0}{R_0} \right)^{3/2} \omega / \epsilon - \frac{1}{\omega} \] and \[ \lambda = \omega q / \left( \omega A^2 \right)^{1/2} \], evaluated at the flux coordinate \( r = r_0 \) where \( q(r_0) = 1 \). Here, \( \omega \) is the normalised MHD energy functional, which includes the effects of toroidicity [10] and plasma shaping [11]. The numerical evaluation of \( \omega \) shows that triangular shaping in JET gives a negligible contribution; the elliptical contribution is destabilising, although typically 10% of \( \omega \). The other parameters are as follows:

- \( \sigma = \eta \omega_0^2 \left( 4 \pi r_0^2 \right) \), \( \eta \) is the resistivity, \( \omega_A = s_0 v_A / \sqrt{3} R_0 \) is the Alfvén frequency, \( s_0 \approx 0.7 \), \( q(r_0) \) is the magnetic shear, \( \omega_0 \) and \( \omega_1 \) are the electron and ion diamagnetic frequencies respectively (\( \omega_0 \approx Z \omega_1 \)).

Although uncertainties remain in the exact location of the marginal stability curve, this is expected [7] to lie close to the curve shown in Fig.4. Layer effects near the q=1 surface, including electron thermal conductivity, ion viscosity, and tight aspect ratio, lead to stability on the left of this curve. Ideal MHD internal kink modes correspond to the regime where \( \lambda > 1 \).

In our present approach, the nonlinear evolution of the m=1 mode must leave q(0) below one, in contrast with Kadomtsev's relaxation model, as no significant drop of q(0) can be expected during the ramp of normal sawteeth. The important point here is that stable q profiles with q(0)<1 are found if \( \lambda > 1 \) is negative and sufficiently large. This requirement can be satisfied, for instance, if the local magnetic shear parameter \( s_0 \) is
small. In fact $\Lambda_\text{H}$ is negative and scales mainly with $s_0^{-5/3}$ during a sawtooth ramp in the ohmic phase, since the poloidal beta is small and the global equilibrium parameters cannot evolve significantly. Thus we propose an alternative prescription for the q profile in the relaxed state. We assume that the most significant changes in the q profile after the collapse are localised near the q=1 surface, where the magnetic shear is reduced over a radial interval of width $\Delta r_0$. Outside this interval the q profile is only slightly modified, with q(0) being reset to a prescribed value below unity, while $r_0$ is chosen to be near the observed inversion radius. This prescription is suggested by measurements of the q profile in TEXTOR [12] and is compatible with initial measurements of q(0) in JET [6]. It can be justified in terms of the growth of a magnetic island which triggers the faster late stage of the collapse as a critical width of order $\Delta$ is attained. The pressure profile after the sawtooth crash is taken to be flat in the region where q$>1$.

Using this prescription, the time evolution of $\Lambda_\text{H}$ and $\Lambda_\text{*}$ has been monitored with an equilibrium transport code which assumes neoclassical resistivity. JET plasmas with $q_\text{m}=4.5$ have been considered so far. The trend in the ohmic phase is shown by the solid-line trajectories in the stability plane of Fig. 4. The trajectories start well in the stable region at the beginning of the sawtooth ramp and approach the marginal stability curve after about 90 ms, at the end of the ramp. The first part of these trajectories is affected by the prescribed q profile after the collapse and should be considered with caution. However, for the same ohmic sawtooth simulated with two different initial values of $s_0$, the trajectories become close to each other on a time scale of the order of tens of milliseconds (Fig. 4). A similar sensitivity has been observed with different choices of the initial q(0) ranging from 0.7 to 0.9, and even less to $\Delta$ ranging from $r_0/4$ to $r_0/2$. This is consistent with the requirement that the values of the stability parameters at the onset of the collapse should not depend on fine variations of the prescribed q profile.

The dotted lines in Fig. 4 correspond to the evolution of the stability parameters during the ramp of a typical monster sawtooth for two initial values of $s_0$. For the case shown, the parameter $\Lambda_\text{H}$ becomes positive for $\beta_\phi \approx 0.15$ ($\beta_\phi$ defined in Ref. [4]) after about 300 ms. To account for the observed stability, it has been proposed that the resistive instability threshold is modified by the presence of non-thermal energetic ions [13].

REFERENCES
Fig. 1 (right) Plasma centroid movement during sawtooth crash.

Fig. 2 (below) Tomographic analysis of monster sawtooth crash.

Fig. 3 (right) Contour plot of toroidal X-ray emission during crash. (Chord radius = 38 cm). n=1 is determined as there is only one point of maximum emission, at 110°.

Fig. 4 Trajectories in the stability plane. Solid curves: ohmic sawtooth (period ~ 90 ms); dotted curves: monster sawtooth (duration ~ 2 s). The starting point in the stable region corresponds to the beginning of the sawtooth ramp. Initial values of the local shear: s₀=0.1 (upper solid and upper dotted curves); s₀=0.2 (remaining curves). Other initial values for all curves: q₀=0.8; rₑ=40 cm; Δ=rₑ/2. The local shear increases along the trajectories: s₀=0.1(∗); 0.2(x); 0.25(○); 0.3(Δ); 0.5(□); 0.7(○).
HIGH ELECTRON AND ION TEMPERATURES PRODUCED IN JET
BY ICRH AND NEUTRAL BEAM HEATING

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Introduction

The installation of 8 ICRF antennae on JET has enabled 16 MW of power to be coupled in both monopole and toroidal dipole configurations. With on-axis hydrogen minority heating such power levels have produced electron temperatures in the vicinity of 10 keV in the centre of the discharge. The highest temperatures often occur during sawtooth-free periods, or 'Monster' sawteeth which can persist for over 3 sec and during which the global energy confinement time is about 20% greater than that in plasmas disrupted by sawteeth. Recent experiments have been aimed at achieving both high electron and ion temperatures simultaneously. In this paper we describe the results of two such series of experiments. The first study used ICRF with ³He minority ions in He and deuterium plasmas and investigated the effect of minority concentration on bulk ion heating. The second series combined hydrogen minority ICRH in deuterium plasmas with up to 6 MW D⁰ neutral beam power injected at an energy of 40 keV/nucleon to provide additional ion heating. We also report measurements of the electron heating rates for the RF-only experiments obtained from the initial rate of increase of the electron temperature during sawteeth. These results indicate an increase in the proportion of direct heating as the minority density is increased.

ICRF Heating with He³ Minority Ions

The He³ minority ICRF experiments were carried out using a plasma current $I_p$ of 3.5 MA and an elongation ratio of 1.5. The toroidal field, $B_T$, was 3.4 T in order to make the ratio $I_p/B_T$ (MA/T) close to unity to provide optimum confinement. The RF frequency was 34 MHz to give on-axis heating. Deuterium and He³ plasmas were used, (but with no discernible difference in results) with central electron densities, $n_e(o)$, in the range $3.8 \times 10^{19} < n_e(o) < 4.8 \times 10^{19}$ m⁻³. ICRF power levels of up to 12 MW were coupled using both monopole and toroidal dipole antennae configurations. The plasma $Z_{eff}$ was typically 4.5 (visible bremsstrahlung) and the total radiated power was about 50% of the input power. An example of the time evolution of the central electron and ion temperatures $T_e(o)$ and $T_i(o)$, during heating is shown in Fig 1 for a ³He concentration $n_{He³}/n_e(o) = 0.11$. The values of $T_e(o)$
and $T_i(o)$ were obtained, respectively, from ECE data and Doppler broadening measurements of X-rays from He-like nickel ions. Note that at 3.4 T the ECE diagnostic calibration gives a value of $T_e(o)$ which is about 20% higher than Thomson scattering results. Neutral particle energy spectra confirmed the above $T_i(o)$ values. As the RF power is ramped up the sawtooth amplitude and period increase until at 47.8 sec, a 'Monster' sawtooth is formed. During the Monster sawtooth, $T_e(o)$ rises to 9.1 keV (ECE) and $T_i(o) = 7.8$ keV. At $t = 49.4$ sec one of the RF generators tripped, causing a reduction in the power which, in turn, precipitated the sawtooth crash 0.3 sec later. This delay is routinely observed and has given rise to speculation that the fast minority ions stabilise the sawteeth. No further sawtooth stabilisation occurs beyond $t = 49.4$ sec showing the existence of an RF power threshold between 7 MW and 9 MW for this operating scenario.

The values of $T_i(o)$ obtained with dipole phasing are shown in Fig 2 as a function of the RF power, $P_{RF}$, divided by $n_e(o)$ for a range of $^3$He concentrations. There is clear tendency for the highest values of $T_i(o)$ to be achieved with the highest minority concentrations as expected since the minority ion tail energy will be least under these conditions. Conversely, stronger minority heating of the electrons is expected at low $^3$He density and there is a tendency in this direction in the values of $T_e(o)$ shown in Fig 3. The effect is less pronounced than in the case of $T_i(o)$, possibly because of increased direct electron heating through mode conversion or transit-time magnetic pumping (TTMP) at the higher $^3$He densities. The monopole antennae phasing gave similar results.

Electron Heating Rates

An analysis of the rate of rise of $T_e$ following each sawtooth crash both during the RF ramp phase and following the Monster sawteeth, (see Fig 1) has enabled electron heating rates to be obtained. Results are shown in Fig 4 where the central power density, $P_e$, flowing to the electrons is plotted against $P_{RF}$ for low and 'high' $^3$He densities. In the low density case, $P_e$ shows the non-linear behaviour expected of minority species heating. Note the hysteresis effect, which could be due to the finite minority ion slowing down time, as $P_{RF}$ is reduced from 11 MW to 8 MW. At the higher $^3$He density the linear increase of $P_e$ with $P_{RF}$ implies an increasing proportion of direct electron heating through TTMP, Landau damping, or mode conversion of the fast magnetosonic wave. It should be stressed however that this analysis does not give unambiguously the RF heating of the electron since it does not account for expulsion of the minority ions following the sawtooth crash or changes in the radiated and ohmic power.

Combined (H)DICRF and $D^6$ Neutral Beam Injection

These experiments were carried out in deuterium limiter discharges with $I_p = 3$ MA, and $B_T = 2.8$ T. The RF frequency was 43 MHz so that the $\omega = \omega_{CH}$ resonance intersected the magnetic axis. The central electron density was typically $\sim 6 \times 10^{19}$ m$^{-3}$. Power levels up to 15 MW of ICRF (dipole phasing) and 6 MW of NBI ($E_{beam} = 80$ keV) were used to give a maximum energy content of 6 MJ and a D-D reaction rate of $3.5 \times 10^{15}$ s$^{-1}$. 
The results are shown in Fig 5 where the values $T_i(o) = 7.7$ keV and $T_e(o) = 10$ keV were achieved at the highest power levels. These temperatures are similar to those achieved with the $^3$He minority ICRF experiments but required approximately twice the value of $P_A/n_e(o)$. The precise explanation for this awaits further analysis but is probably due to a combination of factors such as the poorer confinement at the lower plasma current, the broader power deposition of the beam heating and increased radiated power.

Summary

Central electron and ion temperatures in the vicinity of 8 keV have been achieved simultaneously at moderate densities in JET both by He$^3$ minority ICRH and combined ICRH and NBI with total powers up to 22 MW. The ICRF ion heating appears to be optimum for $5\% < n_{He^3}/n_e < 10\%$. Detailed analyses of sawteeth indicate an increase in the ratio of direct to minority ion heating of the electrons as the He$^3$ density increases.

Acknowledgement

We wish to acknowledge with pleasure the assistance of all our colleagues in the JET team. Particular thanks go to the tokamak operating team and to the members of all the diagnostic groups involved in the measurements reported here.

References


Fig 1: $T_e(o)$ and $T_i(o)$ during RF Heating. $I_p = 3.5$ MA, $n_e(o) = 4.7 \times 10^{19}$ m$^{-3}$, $n_{He^3}/n_e(o) = 11\%$. 
Fig 2: $T_i(\omega)$ vs $P_{RF}/n_e(\omega)$. The lines are least squares fits for He$^3$ densities < 2% and 2% - 4%.

Fig 3: $T_e(\omega)$ vs $P_{TOT}/n_e$. The solid line is the best linear fit which passes through the origin.

Fig 4: $P_e$ vs $P_{RF}$ showing minority heating (circles) and direct heating (crosses).

Fig 5: $T_e(\omega)$ and $T_i(\omega)$ vs total power and additional power $P_A$, respectively, divided by central density.
EFFECT OF SAWTEETH AND SAFETY FACTOR q ON CONFINEMENT DURING ICRF HEATING OF JET

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1. Introduction: The ion-cyclotron resonance heating (ICRH) experiments on JET have been carried out in a variety of conditions: Limiter and X-point configuration with H or He^3-minority heating schemes in D, He^3 or He^6-plasmas. Moreover, second-harmonic heating experiments in H-plasmas and combined ICRH and neutral-beam injection heating experiments have also been carried out. RF power exceeding 16 MW for several seconds have been coupled to a 3.3 MA JET plasma (f = 42 MHz) with 8 antennas operating in the dipole configuration [1]. The stored energy reached 6 MJ with T_{eo} = 8 keV, T_{eo} = 6 keV and n_{eo} = 6 x 10^{19} m^{-3}. The radiated power was less than 42 % of the total input power. A significant fraction of the energy is carried by the heated fast ions which are well confined in JET.

In this paper, our efforts are directed to confinement scaling studies with plasma current I_p and safety factor q for the above mentioned heating scenarios and discharges. With the production of "Monster" or ultralong sawteeth (period greater than the energy confinement time) in JET, we are also able to assess the effect of sawteeth on confinement. The confinement data in JET fitted well with an off-set linear law. Therefore, in the following, we present our results in terms of incremental confinement time (\tau_{inc}) and demonstrate that \tau_{inc} improves with I_p at lower currents (1 - 2 MA), saturates and subsequently degrades at higher currents (4 - 5 MA). These results are compared with a model which assumes that \chi (where \chi \propto 1/\tau_{inc}; see below) depends on two competing mechanisms, that is, it decreases with I_p but increases as the size of the q = 1 surface is increased at higher I_p/B_p or low q_a.

2. Results: For the results presented below, we include the JET ICRH data obtained in 1987 both for limiter and X-point discharges. The plasma parameters under ICRF heating are in the following range:

0 \leq P_{RF} (MW) \leq 16 ; 32 < f (MHz) \leq 43 ; 2 \leq B_p (T) \leq 3.5 ; 
1 \leq I_p (MA) \leq 5 ; 0.7 \leq P_{OH} (MW) \leq 4.5 ; 2 \leq n_e \times 10^{19} (m^{-3}) \leq 5 
3 \leq T_e (keV) \leq 10 ; 2 \leq T_{io} (keV) \leq 7.5 ; 0.2 \leq \frac{P_{rad}}{P_{TOT}} \leq 0.8 
2 \leq Z_{eff} \leq 5

A plot of stored energy (from diamagnetic loop measurements) W_d vs (P_T - W_d) is illustrated in Fig.1 for different I_p as indicated where P_T is the total input power and the dot refers to the time derivative.
Similar data that exists for \( I_p = 1.5, 3, 3.5 \) and 4 MA is not shown on the figure for clarity. \( W_0 \) also includes the significant perpendicular energy that is carried by the fast ions created during the minority heating and is estimated to be less than 1 MJ. It is clearly seen that the \( W_0 \) follows an off-set linear law [2,3]. The \( \tau_{inc} \) is determined by the slope of a line which is least square fitted to the data points for the power scan at a given \( I_p \). Since there is little difference in the energy confinement for the various minority species heating scenarios [4] we do not identify them specifically in the data presented below. Except in some specific cases, generally monopole and dipole antennas give similar confinement results, we therefore do not separate the monopole and dipole data.

In order to reduce the scatter of the data, we have subdivided large data sets at a given \( I_p \) into smaller subsets generally regrouping them by a day or consecutive days of operation. We also exclude the data with \( P_{rad}/P_{TOT} > 0.6 \). The radiation can be systematically taken into account but has not yet been done for the data presented here. The off-axis heating data is also excluded and the minority ion-cyclotron layer lies within 20 cm of the magnetic axis. The data at a given \( I_p \) is further identified with respect to \( q \), i.e., operation at 2 different \( B_\phi \) with the higher \( B_\phi \) pertaining to He-like minority operation.

From Fig. 1, we note that \( \tau_{inc}(D) \) increases with \( I_p \) up to 3 MA but it is considerably lower at 5 MA. The \( \tau_{inc} \) values both from \( W_0 \) and \( W_e \) (electron kinetic energy) obtained from such scans at these and other currents are plotted in Fig. 2 as a function of \( I_p \). The \( W_e \) is calculated from \( n_{eo}, T_{eo} \) and integration over the profiles. It is seen that \( \tau_{inc}(e) \) increases with \( I_p \), saturates then decreases. The two low points at 3 and 4 MA correspond to large \( I_p/B_\phi \) or low \( q_{cyl} \)-values. These points merge well into the crowd when the data is plotted as a function of \( I_p/B_\phi \) or \( q_{cyl} \) (see below). Similar remarks apply for \( \tau_{inc}(D) \) data. For the \( \tau_{inc}(D) \) data, we have further identified exclusively non-monster scans which roughly follow the \( \tau_{inc}(e) \) behaviour except that of a point at 3.3 MA. This corresponds to an exceptionally good series in which dipole phasing was used with pellet injection in some shots. The other \( \tau_{inc}(D) \) data which feature monster sawteeth show a marked improvement over the non-monster data and is related to the added contribution of perpendicularenergy of the fast ions. Although occasional monster sawteeth have been obtained at 4 and 5 MA, there are no scans under such conditions to deduce \( \tau_{inc} \) at these currents. Note that the present X-point and limiter ICRH discharges give about the same \( \tau_{inc} \) values. Also, the 2 \( \omega_{CH} \) data point compares well with the other data.

The \( \tau_{inc} \) data shown in Fig 2 is now plotted in Fig. 3 as a function of \( I_p/B_\phi \), where this quantity is inversely proportional to \( q_{cyl} \), if the plasma cross sectional area is held constant (the experimental data fits a relation \( q_{cyl} \sim 3.3 \times (I_p/B_\phi)^{0.8} \)). It is seen that maximum \( \tau_{inc} \) is found when \( I_p/B_\phi = 1 \) or \( q_{cyl} = 3.3 \).

We now compare the \( \tau_{inc} \) deduced from certain scans where we sort the data with and without monster sawteeth. The \( \tau_{inc} \) values thus obtained
are plotted in Fig. 4 as a function of $I_p$ and the data points are identified as shown. It is found that on average, there is about 15% improvement in $\tau_{inc}(D)$ with monster over the non-monster case whereas there is no improvement in $\tau_{inc}(e)$ on a similar comparison. This relates to the fact that fast ions in JET are well confined during a monster sawtooth and result in improvement of stored energy whereas frequent sawteeth may eject them periodically losing their contribution to stored energy. The monster sawtooth thus show their importance in reactor environment in preventing the ejection of $\alpha$-particles due to frequent sawtooth crashes.

3. Modelling: We follow the local transport model such as given in Ref. 2 to predict the variation of $\tau_{inc}$ with $I_p$ and $I_p/B_p$. In Ref. 2, $\tau_{inc}$ is written as $\tau_{inc} = \tau_x \eta$ where $\eta$ is the heating effectiveness and it is equal to unity for central highly peaked power deposition profiles. $\tau_x$ is defined as $\tau_x = 3 \alpha/4 \chi$ where $\chi$ is a spatially averaged heat diffusivity [2]. If $\chi = \chi_0$ inside the mixing radius $r_m$ and $\chi(r) = \chi_0/(1 - (r/a)^2)$ outside $r_m$, we can write as follows:

$$\tau_x = (3 \alpha/8 \chi_0)(1 - r_m/a)^2(1 - (a/2)(1 + r_m/a^2))$$

For simplicity, we take $\alpha = 1$ and $\chi_0 = I_p$ to get:

$$\tau_x(s) = \tau_{inc}(s) = \tau_x(0) s I_p^{0.5} [MA] \cdot (1 - (r_m/a)^2)^2$$

where $\tau_x$ and $\beta$ can be adjusted to fit the data. The experimental data fits a relation $r_m = 1.35/q_{cyl}$. Since $\tau_x$ is a function of $I_p$ and $q_{cyl}$, a single curve cannot be fitted to the data. Instead, a point by point comparison has to be made. For clarity, we show fitted points (diamonds) for $\tau_{inc}(e)$ only in Figs. 2 and 3 which should be compared with triangles and are found to be in fairly good agreement assuming $\beta = 0.5$ and $\tau_x = 0.14$ s. This fit can be further improved by adjusting the parameter $\alpha$.

4. Discussion and Conclusions: The behaviour of experimentally obtained $\tau_{inc}$ of centrally heated ICRF discharges as shown in Figs 2 and 3 is understood in terms of a local heat transport model where $\chi_0 \propto I_p^{-\beta}$ and that $\chi$ increases as the size of the $q = 1$ surface is increased at higher $I_p/B_p$ or low $q_a$. The actual value of $\beta$ could be further improved if power scans are made at constant $n_e$. The effect of off-axis heating profiles on $\tau_{inc}$ has previously been discussed in [2,4]. The loss of fast particles by sawtooth is a significant energy loss which is avoided in a monster sawtooth discharge.

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References

Fig. 1: Plot of $\tau_{inc}(D)$ and $\tau_{inc}(e)$ vs $I_p$ for monster and non-monster sawteeth data.

Fig. 2: Plot of $W_D$ vs $P_T - W_D$ at different $I_p$. The scan at 1.1 MA refers to X-point.

Fig. 3: Plot of $\tau_{inc}(D)$ and $\tau_{inc}(e)$ vs $I_p$. A fit of $\tau_{inc}(e)$ only is shown with $\beta = 0.5$ and $\tau_{\chi\phi} = 0.14$ s.

Fig. 4: Plot of $\tau_{inc}(D)$ and $\tau_{inc}(e)$ vs $I_p$ for monster and non-monster sawteeth data.

For fit, see caption of Fig. 2.
MAGNETIC MEASUREMENTS OF THE SAWTOOTH INSTABILITY IN JET


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Abstract: Magnetic coil signals are used to obtain the toroidal mode number spectrum of the sawtooth instability in a small-aspect-ratio tokamak. The excellent correlation between the edge (magnetic) displacement and the core displacement (soft X-ray tomography) demonstrates the validity of this approach and makes magnetic coil measurements an appropriate tool for studying the sawtooth instability global wavefunction.

Introduction: We have shown earlier that the sawtooth crash is accompanied by a simultaneous magnetic spike, which has been called the "gong" mode [1, 2]. Although the physical nature of this magnetic perturbation was not fully established, it was used to characterize the sawtooth instability. Toroidal mode number analysis, performed then with 4 equally spaced coils yielded a dominance of the n = 1 mode, with some 20–30% of n = 2. The gong mode was shown to be particularly strong at the low field side midplane. Typical growth rates in the range of 10^3 s^-1 were found to compare well with the typical sawtooth collapse times [3]. The sign of the mode's helicity was shown to be identical to the field lines' helicity. These observations showed that the magnetic perturbation was intimately related to the sawtooth instability in the core. In this paper we show a detailed comparison of the gong with soft X-ray tomography. The toroidal mode number spectrum evolution is obtained from 8 equally spaced coils, which allows us to follow the n-spectrum from n=0 to 4.

Plasma displacement: To show the evidence that the gong is in fact a direct measurement of the sawtooth instability, we had to find a parameter which could be measured independently by the soft X-ray diagnostics and the magnetic pick-up coils. The appropriate parameter for comparison was found to be the radial plasma displacement associated with the sawtooth instability. The m = 1 core displacement is accurately represented by the first moment of the soft X-ray emission profile F, which components \( X_i = \int x_i F(x_1, x_2) dx_1 dx_2 / \int F(x_1, x_2) dx_1 dx_2 \) define the centroid [4]. The centroid shift is then \( \xi_{sx} = \left( (x_1-x_1')^2 + (x_2-x_2')^2 \right)^{1/2} \), where the primed symbols refer to the position before the growth of the instability.
An estimate of the edge displacement can be made from the value of the poloidal field perturbation, as \( b = \nabla \times (\xi \times B_0) \) in ideal MHD. In the cylindrical approximation the edge radial displacement can be expressed by \( \xi(a) = (b_0(w)/B_0(a))(w/2m)(1 + nq/m)^{-1}[(w/a)^m - (a/w)^m] \) where \( a \) and \( w \) are the plasma and wall radius, \( m \) and \( n \) the poloidal and toroidal mode numbers of the gong. The edge displacement, which is directly proportional to the gong perturbation, can be analyzed for each \( n \)-component. As the main component of the gong was shown to be \( n=1 \), we will compare \( \xi_{n=1}(a) \) with the centroid shift. This yields edge displacements \( \xi_{n=1}(a) \) in the \( 10^{-3} \)m range. We note that gong-like perturbations in the ion saturation current are equally measured on a limiter langmuir probe. The ratio of edge \( n = 1 \) displacement to core displacement yields in the case of Fig. 1: \( \xi_{n=1}(a)/\xi_{sx} \approx 3.10^{-2} \) with \( q_\psi = 4.7 \) and \( b/a = 1.54 \).

The fast rise of the instability is observed to start some 300 \( \mu \)s before the central soft X-ray emission drop (Fig. 1a). Both displacements exhibit the same temporal behaviour and the two traces overlap almost exactly, as shown in Fig. 1b. The agreement is maintained over the whole growth period and while the value of \( b_{0n=1} \) is large. This demonstrates that the gong is directly linked with the sawtooth instability. The discrepancy between the two displacements noticed in the later phase following the temperature collapse, will be explained in the next section. Apart from this restriction which only applies to the second phase, the good agreement between edge and core information shows that the gong and the sawtooth instability form one single global \( n=1 \) motion. The absence of any phase lag between the edge and the core is typical of a single element oscillating system, which is related to the nature of toroidal coupling. An interesting corollary of this good agreement is the demonstration that the soft X-rays represent a good monitor of the magnetic surfaces, at least during the growth of the instability. Most importantly, the global properties of the sawtooth instability allow the use of the gong as an additional tool for studying the internal disruption instability.

**Toroidal spectrum evolution:** The toroidal spectrum contains important information related to the nature of the sawtooth instability. We have therefore extended the measurements to a toroidal array of eight equally spaced coils. Fig. 2a shows the eight integrated \( B_0 \) traces, from which the \( n = 0, 1, 2, 3 \) and cosine \( n = 4 \) components are calculated. We first note that the \( n = 1, 2 \) and \( 3 \) modes all grow proportional to each other and seem therefore to belong to the same general motion. There is a nearly imperceptible \( n = 4 \) component, which suggests that the eight coils used are sufficient to qualify the \( n \)-spectrum of a sawtooth instability during the growth phase and in the case of JET. Some higher \( n \) perturbations may be present during the later profile reorganisation phase, after the saturation of the \( n=1 \) mode, as is suggested by Fig. 2b.

The \( n = 0 \) mode is related to the inward Shafranov-shift (reduction of poloidal field on all low field side coils, Fig. 2a) following a reduction in \( \lambda = \beta_p + 1/2 \). The \( n=0 \) trace evolves like the integrated \( n=1 \) traces. Thus, the profile modifications seem to be a direct consequence of the \( n=1 \) instability. The \( n = 0 \) growth rate, which is only slightly
smaller than the $n=1$ instability growth rate, indicates that the \( \lambda \) changes occur on the same characteristic time scale.

The presence of this $n=0$ component is sufficient to explain the discrepancy mentioned in section 2 between $b_{0n=1}$ and $\xi_{Sx}$ following the sawtooth collapse. While $b_{0n=1}$ monitors a non-local toroidal component, $\xi_{Sx}$ on the contrary only measures the m=1 displacement in one specific poloidal plane. As a consequence, all the n-components contribute to $\xi_{Sx}^*$, including the $n=0$ Shafranov shift in the late phase.

The decay index $\alpha$ of the n-spectrum of the instability, defined as $S(n) \approx A_n n^{-\alpha}$ for $1 \leq n \leq 3$, yields a value of $\alpha$ of $2.5 \pm 0.5$. The most obvious way to interpret the n-spectrum is to assume that its measurement at the edge represents the n-spectrum on the $q=1$ surface, therefore equivalently the m-spectrum on the $q=1$ surface, as $m = nq$ on a rational surface. Although the contribution from other surfaces cannot be excluded a priori, the fact that the fast gong-like changes observed on the soft X-ray chords are dominant within or close to the sawtooth inversion radius further support this core origin of the n-spectrum measured. However, the question must be posed whether departures from axisymmetry or whether non-linear mechanisms play a role by spreading the n-spectrum. Firstly, only gongs with vanishingly small superimposed helical Mirnov perturbations have been retained for analysis, as Mirnov oscillations not only destroy axisymmetry, but also directly pollute the measured n-spectrum. Secondly, very large amplitude gongs ($2.10^{-3} < b_{0n=1} < 7.10^{-3}$) may even show a reduced $n=2$ component during the growth phase, whereas non-linear effects would show the opposite trend. Both these statements therefore support the thesis of the core origin of the n-spectrum for the presented cases. As a consequence, the n-spectrum may directly represent the m-spectrum on the $q=1$ surface. Finally, and for normal sawteeth, the measurements indicate in addition to the $n=1$ mode non negligible $n=m>1$ mode contributions, growing all in parallel, which is not inconsistent with an interchange mode.

In an attempt to find scaling laws for the growth rate (measured at half maximum amplitude), we note that no trend is obtained for the growth rate as a function of external control parameters. The reason appears evident when considering that two consecutive gongs may show growth rate scattering as much as half an order of magnitude. Let us note that the use of Fourier components enables the measurement to be independent of plasma rotation. This shows therefore that the growth rate is controlled by details of the profiles. The growth rate is, however, found to be grossly proportional to the gong amplitude $b_{0n=1}$.

**Conclusion:** The gong mode has allowed us to measure characteristics of the sawtooth instability. Indirect experimental evidence indicates that the measured magnetic n-spectrum may be interpreted as an image of the m-spectrum on the $q=1$ surface. The n-spectrum, measured up to n=4, therefore indicates that the sawtooth instability has the characteristic behaviour of an interchange mode.

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References


Gong and Centroid Shift

![Gong and Centroid Shift](image)

**Fig. 1b**

**Fig. 1a**

**Fig. 1**: The n=1 gong and m=1 centroid shift overlap

**Fig. 2**: θ-coil toroidal spectrum analysis a) the integrated coil signals b) n-spectrum evolution

![Fig. 2a](image)

![Fig. 2b](image)
RESISTIVE BALLOONING MODES UNDER PLASMA EDGE CONDITIONS

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Introduction: Recently resistive ballooning modes have become the object of intensified theoretical investigation. The reason for this increased interest is that they may be linked to confinement degradation observed in tokamak experiments at high plasma pressure. In previous publications [1,2] we investigated the ballooning stability of ASDEX high-β_p discharges and demonstrated that the hard β_p-saturation observed in ASDEX can be naturally explained within the framework of resistive ballooning modes. Here we concentrate the investigation to the region near the separatrix. Furthermore we include a detailed discussion of resistivity effects on the second stable regime. This analysis is based on an analytic large aspect-ratio model for plasma equilibrium in the neighbourhood of a given flux surface (local equilibrium) [3]. Finally, we comment on the resistive stability of toroidal 2-D ASDEX equilibria with experimentally determined current and pressure profiles.

Localequilibria and second region of resistive ballooning stability: To elucidate the characteristic features of the problem in terms of a limited number of parameters we refer to the theory of a local equilibrium presented in [3]. The shape of the flux surfaces is controlled by assigning particular values to the expression

\[ k^2 = r^2(\sqrt{1+k} - 1)^2[4 + r^2(\sqrt{1+k} - 1)^2 - 4\cos(\theta - \theta_x)(\sqrt{1+k} - 1)]. \]

Here \((r, \theta)\) are polar coordinates with the radius \(r\) normalized by a geometrical parameter \(r_0\). \(\theta_x\) is the angular position of the X-point and \(k\) the distortion of the flux surface; \(k = 1\) is the separatrix case and \(k = 0\) is a circle. The resulting resistive ballooning equations are

\[ \frac{b}{h} \left( \frac{b}{h} \right) \left( \frac{1}{b^2 + P^2} \right) \frac{du}{d\theta} - \alpha K(u + v) - \gamma^2 \left( \frac{1}{b^2 + P^2} \right) u = 0 \]

and

\[ \frac{b}{q^2 h} \left( \frac{b}{h} \right) \left( \frac{dv}{d\theta} \right) + \alpha K \left( \frac{\gamma^2}{\alpha^2} + \frac{\eta n^2}{S_A \gamma} \right) u - \left( \frac{\eta n^2}{S_A \gamma} \right) \left( -\alpha K + \gamma^2 \left( \frac{1}{b^2 + P^2} \right) \right) \gamma^2 \frac{1}{b^2 + P^2} \right) v = 0 \]

where \(W, D, b, h\) and \(f\) are defined in [5] and
\[ K = -\frac{1}{h}[P\left(\frac{dr}{d\theta}\cos\theta - r\sin\theta\right) + \frac{1}{b}\left(\frac{dr}{d\theta}\sin\theta - r\cos\theta\right)] \]

\[ P = -b\int \left[ \alpha\frac{W}{D} - \Lambda + 2\frac{f}{j} - \alpha r\cos\theta \right] \frac{h}{b^3} d\theta \]

\[ L = 1 + \frac{\eta n^2}{S_A\gamma} \left[ \frac{1}{b^2} + P^2 \right], \]

with \( \alpha = -2\mu_0 r_0^2 p'/B_p^2 \), \( \beta = \gamma_T \mu_0 p/B^2 \), the magnetic Reynolds number \( S_A = \tau_R/\tau_A \) and the ratio of the specific heats \( \gamma_T = c_p/c_v \) and the growth rate \( \gamma \) is normalized to the Alfvén frequency \( \gamma_A = \sqrt{\tau_0^2 B_p^2/(r^2 \mu_0 p)} \).

In the limit of circular flux surfaces \((k = 0)\) these equations reduce to the \( s - \alpha \) model, with the shear \( s \) related to the current density parameter \( \Lambda \) by \( s = 2 - \Lambda \). We start the investigation of the second regime in this limit. Fig.1 shows the real part of the growth rate versus \( \alpha \) for ideal ballooning modes (dot-dashed line) and resistive modes (solid lines). The hat-shaped curve shows second stability behaviour. The slight bulge in this curve at \( \alpha \sim 0.6 \) indicates the transition to overstability [6]. This resistive curve closely parallels the ideal stability curve and is reproduced by the \( \Delta' \)-criterion for resistive ballooning modes [6,7] to reasonable accuracy:

\[ \Delta' = \frac{2\gamma^{5/4}(1 + 2q^2)^{1/4}}{S_A(\eta^2 s^2 q^2/S_A)^{3/4}} \left\{ \frac{\Gamma(\frac{1}{4})}{\Gamma(\frac{3}{4})} - \frac{\Gamma(\frac{1}{4} + \frac{4}{3})}{\Gamma(\frac{3}{4} + \frac{4}{3})} \right\}^{-1} \]

with \( \hat{\gamma} = \gamma \gamma_A/\gamma_R \), \((\gamma_R/\gamma_A)^3 = n^2 \eta s^2 q^2 \beta^2 (1 + 2q^2)/S_A \) and where \( \Delta' \) is in general to be evaluated numerically [8]. This dispersion relation (4) is derived in the limit \( \gamma \ll \gamma_S = \sqrt{\beta} \), where \( \gamma_S \) denotes the sound frequency. In the other limit \((\gamma \gg \gamma_S)\) a compressibility mode, not driven by \( \Delta' \), can occur. Its growth rate is governed by the dispersion relation [9]

\[ \gamma^3 + \gamma S^2 (1 + 2q^2) \gamma = \alpha^2 \frac{n^2 q^2}{2 S_A} \]
This mode can also be seen to occur in our numerical results in Fig. 1, where the branch occurring for \( \alpha > 2.6 \) \((\beta = 3 \times 10^{-3})\) is compared with the solution of (5) (long-dashed line). It should be noted that in general \( \gamma << \gamma_S \) for typical tokamak parameters, but in the second stable regime the high \( \alpha \)-values make the region where \( \gamma \sim \gamma_S \) more accessible. From Fig. 1 it can be seen that the dispersion relation (5) continues to apply reasonably well to the regime \( \gamma > \gamma_S \) where \( \gamma \propto \eta/(S_A \gamma_S^2) \). Thus we conclude that the existence of a second stable regime depends on \( \gamma_S \) and \( S_A \) - for large enough values of either parameter a second stable region exists below the critical value of \( \alpha \) at which compressibility modes (5) are destabilized. This is separately demonstrated by the short dashed curve in Fig. 1, representing the compressibility mode for \( \beta = 5 \times 10^{-3} \). The corresponding hat-shaped solutions are essentially unchanged with respect to this increase in \( \beta \), so that effectively a second stable window exists for this case. It should be noted that the picture is in general more complicated, since the region of instability has to be maximized by optimizing the free parameter \( \theta_0 \) which appears in eqs. 3.4. Our conclusions presented above, however, are not affected by the variation of \( \theta_0 \). Furthermore we find similar results on the basis of a model where the Shafranov shift of circular flux surfaces is taken into account.

We now move on to study the non-circular case, i.e. \( k > 0 \). Fig. 2 shows for \( k = 0.885, \Lambda = 0.8 \) the growth rate in the complex \( \gamma \)-plane for a branch entering a second stable resistive ballooning regime. There are, however, branches with \( \gamma > \gamma_S \) which do not show second regime stability as \( \alpha \) is increased; these branches can be suppressed by increasing the sound frequency.

For the ideal ballooning mode coalescence of the first and second stable regime occurs as the separatrix is approached \((k \to 1)\) for a sufficiently high current density parameter \([5]\). It has only proved possible, for the cases examined, to obtain this coalescence for the resistive ballooning mode by increasing the local \( \beta_0 \) to 5%, which is rather larger than the expected values near the separatrix. For \( \beta_0 < 5% \) and large \( k > 0.9 \) an unstable mode, which may be related to the \( \gamma \propto S_A^{-1/2} \) (as \( S_A \to \infty \)) branches of Ref. 10, occurs. The quantitative dependences of these modes on the parameter \( k \) and the question whether and how their behaviour is related to the appearance of compressibility modes is still under investigation.

**Ballooning stability in the ASDEX separatrix region:** We now briefly examine the ballooning stability of a typical ASDEX high-\( \beta_p \) discharge by solving the full resistive ballooning equations on the basis of a separatrix-bounded MHD equilibrium [1.2]. To make parametric studies we alter the value of the pressure gradient from the experimentally observed quantity using a multiplication factor \( C_p \). Results for the real part of the growth rate versus this factor \( C_p \) are shown in Fig. 3 for various flux surfaces from \( r/a \approx 0.87 \) to \( r/a \approx 0.994 \).
In line with our previous results [1,2] we find a transition from purely growing to overstable modes, when compressibility effects become important.

As soon as the bifurcation into the complex $\gamma$-plane occurs, the real solutions form separate branches which move to larger multiplication factors with decreasing growth rates (dashed lines), whereas the overstable branches (solid lines) show the opposite behaviour and are therefore the more unstable solutions. From the observation that the bifurcation to overstability moves to smaller values of $Re\{\gamma\}$ when $r/a$ increases ($\rightarrow \eta$ decreases and the strong compressibility dependence of the bifurcation point we conclude that the overstable branches are associated with the propagation of ion sound waves [6].

Interpreting these results in terms of absolute pressure gradients we find that the highshear surfaces close to the separatrix can support larger pressure gradients.

**Conclusions:** To summarize we have found that in the separatrix region resistive ballooning modes are generally more stable than in the plasma confinement region. However, we have demonstrated that for a second resistive ballooning stable regime to exist sufficiently large values for $\gamma_S$ or $S_A$ are required. Ideal calculations for the local separatrix equilibrium model show Bishop’s results [4,5] for coalescence of the first and second stable regimes as the separatrix is approached ($k \rightarrow 1$). In the resistive case we find that separatrix effects are only strong enough to completely stabilize a ballooning mode at rather large values of $\beta_0(>5\%)$. Consequently the coalescence effect is restricted to this region of parameter space. Whether the small growth rates, which in general accompany resistive ballooning modes near $k = 1$, inhibit the L-H transition predicted in Refs. 4,5, requires additional analysis. Finally we found ASDEX high-$\beta_p$ equilibria to be resistively unstable with small growth rates below the maximum value of $\beta_p$. In the separatrix region larger absolute pressure gradients are supported so that growth rates decrease and corresponding effects on the transport do not appear to be significant here.

3. C.M. Bishop, Culham Report CLM-R249.
Ballooning Instabilities in Tokamaks with Sheared Toroidal Flows

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The use of unbalanced neutral beam injection has induced toroidal plasma rotation that approaches the sound speed in TFTR [1] and in JET [2]. The toroidal rotation profiles measured in the TFTR device also display a significant amount of velocity shear [1]. Because ballooning instabilities can impose a limit on the beta ($\beta$) value that can be achieved in Tokamak devices, it is important to determine what impact plasma rotation and velocity shear can have on the stability properties of the plasma. Previous theoretical studies of ballooning modes in toroidal devices have suggested that the ballooning mode formalism described either through Weyl sequences [3] or with the covering space concept [4] breaks down when the velocity shear in the toroidal flow is finite [5]. In this work, the covering space concept is employed to represent ballooning mode structures that vary slowly along the magnetic field lines but rapidly across them. This eikonal ballooning representation is applied to the linearised magnetohydrodynamic (MHD) equations in axisymmetric systems with toroidal mass flows. In order to resolve the $\nabla \cdot \nabla$ operator acting on perturbed quantities in a toroidal device, the eikonal function $S$ must satisfy not only the usual ballooning condition $B^2 \nabla S = 0$ but also the condition $dS/dt = 0$. This entails a Doppler in $S$, which in the PEST-1 magnetic flux coordinate system $(s, \theta, \phi)$ with straight field lines becomes $S = \psi - Q(s)\Omega(s) + k(s)$. The surface functions $q$, $\Omega$, and $k$ correspond to the safety factor, the plasma rotation frequency and the radial wave number, respectively. The resulting set of initial value ballooning mode equations are:

$$\frac{\partial b}{\partial t} = -\left[ n^2 \eta |\nabla S|^2 + \frac{2 (\nabla S \cdot V) (V \cdot \nabla S)}{|\nabla S|^2} \right] b + \frac{|\nabla S|^2}{B^2} (B \cdot V) V^\perp$$

$$\rho_n |\nabla S|^2 \frac{\partial v}{\partial t} = \rho_n \left[ 2 (\nabla S \cdot V) (V \cdot \nabla S) \right] v + B^2 (B \cdot V) b - 2 (B \times V S \cdot k) \rho$$

$$- [B \times V S \cdot (V \cdot V) V] \rho$$
\[ \frac{\partial p}{\partial t} = - \frac{B \times \nabla_s \cdot \nabla p}{B^2} v_\perp \]

\[ \frac{\partial \rho}{\partial t} = - \frac{B \times \nabla_s \cdot \nabla \rho_M}{B^2} v_\perp \]

where \( b_\perp, v_\perp, p \) and \( \rho \) are the perpendicular magnetic field, the perpendicular velocity, the pressure and the mass density perturbations, respectively. The equilibrium magnetic and velocity fields are \( H \) and \( V \), respectively, the magnetic field line curvature is \( \kappa \), the equilibrium pressure, mass density and resistivity are \( P, \rho_M \) and \( \eta \), respectively, and \( n \) is the toroidal mode number. It is evident from these equations that normal mode solutions can only be constructed when the toroidal rotation is rigid. The previous work on this subject [3,5] was based on the application of the ballooning representation to the Frieman-Rotenberg energy principle [6]. This energy principle, however, imposes apriori that the solutions obtained be normal modes of the system. Consequently the theory that was developed is limited only to the rigid rotor model. The initial value equations we have derived do not constrain the solutions to evolve as \( \exp(i\omega t) \), therefore we are able to investigate the effects of finite velocity shear. The ballooning formalism is still valid when there is shear in the toroidal flow, but the solutions obtained are not normal modes of the system.

Fixed boundary equilibria that model the JET device are obtained with the ATRIME inverse moments code [7]. This code generates axisymmetric MHD equilibria with isothermal toroidal mass flows. The profiles that must be specified are the plasma mass function \( M(s) \), the inverse of the safety factor and the plasma flow function \( U(s) = 0.25M_i \Omega^2(s)/T(s) \), where \( T(s) \) is the plasma temperature and \( M_i \) is the ionic mass. The plasma pressure can then be constructed from \( M(s) \) and \( U(s) \) [7]. The equilibria obtained are mapped from the \((s,\chi,\phi)\) magnetic flux coordinates employed in the ATRIME code to the straight field line coordinates \((s,\theta,\phi)\) using the relation \( \theta=\chi+\lambda(s,\chi) \), where \( \lambda \) is a periodic renormalisation parameter calculated internally in the code. A numerical JET equilibrium with a thermal beta \( \beta_T=4.9\% \), a rotational component of beta \( \beta_R=1.1\% \) and Mach number 0.925 at the magnetic axis is generated by specifying \( M(s)=0.04(1-4s^3+3s^4) \), \( 1/q(s)=1-2s^3/3 \) and \( U(s)=0.023 \). As the initial value equations are evolved in time, periodic bursts of ballooning are observed, as is shown in Fig. 1. The mode structures at the peaks of the bursts labelled as a), b) and c) in Fig. 1 are displayed in Fig. 2. The instability structure that develops has strongly ballooning characteristics, but is displaced by \( 2\pi \) in the extended poloidal angle domain from one burst to the next. An instability growth rate that is remarkably linear can be extracted from the peak values of the ballooning bursts. In addition to the profiles required for the equilibrium calculation, the local stability analysis requires the
specification of either $T(s)$ or $\Omega(s)$. Thus it is possible to investigate the effects of plasma rotation and velocity shear with a single equilibrium which we do in this study by varying the temperature profile. A sequence with $\Omega(s=0.85)$ fixed and variable $\Omega'(s)$, where prime indicates a derivative with respect to $s$, is obtained by choosing $T(s)=1-0.7225s$, $T(s)=1-0.85s^2$, $T(s)=1-s^3$ and $T(s)=1-s^4/0.85$. A sequence with $\Omega'(s=0.85)$ fixed and variable $\Omega(s)$ is obtained with $T(s)=(1-1.026239s^2)^2$, $T(s)=(1-0.804909s^3)^2$ and $T(s)=1-s^3$. At fixed $\Omega'$, we find that the growth rate and frequency of the ballooning bursts are insensitive to variations in $\Omega$. On the other hand, for fixed $\Omega$, the frequency varies linearly with $\Omega'$ and goes to 0 when $\Omega'$ vanishes. The growth rate decreases with increasing $\Omega'$ which indicates that the velocity shear has a stabilising effect. This can be understood by noting that two fluid elements on adjacent flux surfaces become physically separated in space as time evolves which inhibits the formation of an instability structure.

![Graph](image)

**Fig. 1.**
The value of $v_\perp^2$ integrated over the extended poloidal angle $\theta$ domain as a function of time. The time is normalised to the poloidal Alfvén time.
The instability mode structures at three different times that correspond to the peaks of ballooning burst labelled with a), b) and c) in Fig. 1, respectively.
RESISTIVE BALLOONING MODES IN DIFFERENT COLLISIONALITY REGIMES

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The stability of resistive ballooning modes has been the subject of many theoretical investigations [1,4]. The aim of the present paper is to discuss the various stabilizing and destabilizing mechanisms in different collisionality regimes.

In the low-\(\beta\) limit the averaged Braginskij two fluid equations in the cold ion limit reduce to the following system [3]

\[
\frac{d}{dX} \left( \frac{X^2}{\Gamma_1} \frac{df}{dX} \right) + p \frac{d}{dX} - X^2 Q Q_1 \left| \phi - \frac{i \omega_e}{Q_1} (1 + \eta_e) (p + a_1 t) \right| = 0
\]

\[
\frac{d}{dX} \left( \frac{d\rho}{dX} - \frac{1}{\Gamma_1} \frac{d\phi}{dX} \right) - Q^2 G \left( \frac{\alpha X^2}{\nu G} + \frac{\omega_e}{Q} \right) \left| p + a_1 t + \frac{i Q}{\omega_e (1 + \eta_e)} \phi \right| + Q^2 G (t - p) = 0
\]

\[
\alpha_2 \frac{d}{dX} \left( \frac{d\tau}{dX} - \frac{\eta_e}{(1 + \eta_e)} \frac{\Gamma_2}{\Gamma_1} \frac{d\phi}{dX} \right) + \nu G Q \left( p - \frac{5}{2} \right) +
\]

\[
+ i \omega_e (1 - \frac{3}{2} \eta_e) G \left| p + a_1 t + \frac{i Q \phi}{\omega_e (1 + \eta_e)} \right| = 0
\]

where \( Q = -i \omega_e / \omega_\tau \), \( X = r \xi, Q_1 = Q + i \omega_e (1 + (1 + \alpha_1) \eta_e) \), \( \omega_\tau = \omega_\tau \omega_e, \eta_e = d \ln T_e / d \ln n_e, G = \left[ (1 + 2 \alpha_1 \beta_2) \right]^{-\alpha_1}, \Gamma_1 = 1 + X^2 / Q_1, \Gamma_2 = 1 + (1 + a_2 / a_2) \right) X^2 / (\omega_e \eta_e), \alpha_1 = 0.71, a_2 = 1.95, a_3 = 3.2, \nu = \left( \nu_e / \omega_\tau \right) (m_e / m_i), \xi \) is the extended poloidal variable and \( \omega_\tau = \omega_\chi \left( \frac{12}{9} \right)^{1/3}, \chi = \left( \frac{12}{9} \right)^{-1/3} \) with \( I \) the toroidal number and \( S \) the magnetic Reynolds number. Moreover

\[
\phi = \frac{\varepsilon \phi}{T_e} \left| \frac{-i \omega_e (1 + \eta_e)}{Q_1} \right| = \left| \phi \right| \frac{-i \omega_e (1 + \eta_e)}{Q_1}
\]

An optimal ordering can be defined which makes comparable the effect of inertia, line bending, parallel compressibility, parallel thermal conductivity and
perpendicular transport

\[ Q_1 = Q - \omega_* - v - x^{2/3} \quad X = x^{-1/3} \]

Upon using the latter ordering, Eqs (1) to (3) reduce to

\[ \phi'' + p \frac{D}{Q_1} - X^2 Q \left[ \phi - \frac{i\omega_*}{Q_1} (1 + \eta_e) (p + a_1 t) \right] = 0 \]

\[ p'' + Q_2 G_1 t - Q_2 p \left[ G_1 + a_2 \frac{X^2}{Q} \right] - \left[ \frac{a_2 Q^2 X^2}{v} + i\omega_* G_1 \right] \left[ p + a_1 t + \frac{iQ_1}{\omega_* (1 + \eta_e)} \phi \right] = 0 \]

\[ a_3 \frac{d^2}{dx^2} \left[ t - \frac{a_1 a_2}{a_3} \frac{iQ_1}{\omega_* (1 + \eta_e)} \phi \right] - \frac{5}{2} v G_1 Q_1 \left( 1 + \frac{2}{5} \frac{a_2 a_4}{Q} \frac{X^2}{Q} \right) + \]

\[ + p v Q G_1 + i\omega_* \left( 1 - \frac{3}{2} \eta_e \right) G_1 \left[ p + a_1 t + \frac{iQ_1}{\omega_* (1 + \eta_e)} \phi \right] = 0 \]

with the following boundary conditions for \( X \to 0 \)

\[ \phi \sim a - \frac{\Delta'}{X} \frac{X}{Q_1} \quad (9) \]

\[ p \sim 1 \quad (10) \]

\[ t \sim \frac{\eta_e}{1 + \eta_e} a + b - \frac{\eta_e}{1 + \eta_e} \frac{\Delta'}{X} \frac{X}{Q_1} \quad (11) \]

with \( \Delta' \) given in Ref. [2] and \( a \) and \( b \) being arbitrary constants. Equations (9)-(11) are very similar but not identical to the boundary conditions used in Ref. [3].

In the resistive MHD limit \( \nu \to \infty \) Eqs (6)-(8) can be analytically solved for \( a_2 = 1 \) [2]. If both the effect of perpendicular transport and parallel thermal conductivity are neglected in the layer the mode is purely growing with \( \gamma \sim (\Delta')^{4/3} \) and no threshold. If the effect of parallel thermal conductivity is retained the growth rate is unchanged for \( \Delta' \gg \Delta'_c \approx \nu^{5/3} (n^{2/3})^{-1/3} \) while for \( \Delta' \ll \Delta'_c \) the mode is weaker with the growth rate given by

\[ \gamma \sim \gamma_0 \left( \Delta'/\Delta'_c \right)^{2} \]
with $\gamma_0 = \omega_{ce} v^{4/3}$. Again there is no threshold. A threshold $\Delta'$ is obtained by retaining the effect of perpendicular transport. In this case it is convenient to formally consider the limit $\alpha_2 \to 0$ in which case $\gamma = 0$ for $\Delta' \Delta'_c = \alpha_2^{1/2}$. Moreover a new unstable branch exists with $\gamma = \gamma_0 \alpha_2 (\Delta')^{-4}$.

Upon decreasing collisionality the growth rate becomes comparable with $\omega_s$ and for $\Delta' \Delta'_c \leq \omega_s^{5/3} v^{5/3}$ the growth rate of the mode is

$$\gamma \sim \gamma_0 \omega_s^{-2/3} v^{-2/3} (\Delta'/\Delta'_c)^{2} |1 + (1 + \alpha_i) n_c|^{-2/3} v^{1/3}$$

The latter expression is consistent with the smallness of parallel compressibility and parallel thermal conductivity for $\omega_s v^{4/3} \gg 1$ and $\gamma \gg \omega_s^{2} (\Delta'/\Delta'_c)^{2} v^{4/3}$, respectively. This regime can exist only for $\Delta' \Delta'_c \gg 1$. If $(\Delta'/\Delta'_c) < 1$ the mode is stabilized by the effect of parallel compressibility with a damping rate given by $\gamma = -\gamma_0$.

In this regime the destabilizing effect of $\Delta'$ is negligible and equations (6)-(8) reduce to a single second order differential equation in real space for the electrostatic potential, with no driving terms.

A similar regime is obtained also when the effect of parallel thermal conductivity is retained, yielding

$$\gamma \sim -\gamma_0 \omega_s^{4/3} v^{-2/3} e^{i \theta}$$

This result is essentially the same result as that of Ref. [4].

Finally in the semicollisional limit (defined by $(v_{s}/\omega_{s}) (L_s^{2}/L_{ce}) (m_{e}/m_{i}) < 1$) the mode is stable with $\gamma \propto 1/v$.

The conclusion of the present analysis is that the resistive ballooning mode is stable for the collisionality values typical of the experiments. To destabilize the mode requires the inclusion of other effects as e.g. trapped electrons [5].

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ANALYSIS OF SAWTOOTH STABILIZATION IN JET


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Introduction — During experiments with high power additional heating in JET, it is found that the sawteeth may be spontaneously stabilized for periods of up to 3 s. This offers an opportunity to study the behaviour of tokamak plasmas in the absence of mixing due to sawtooth collapses, which normally confuses transport and confinement analysis. In addition, recent measurements of the current density profile in JET have shown that during these periods the central safety factor q0 reaches values of ~ 0.8, a result which raises fundamental questions about the nature and cause of the sawtooth instability in tokamaks. We propose that the stabilization may be due to the presence of a non-thermal ion energy distribution in the plasma.

Observations — Spontaneous stabilization of sawteeth has been observed over a wide range of conditions in JET (Table 1), the principal constraint being the requirement for a minimum level of auxiliary heating power (Paux ≥ 3 MW). This power threshold depends, however, on many parameters, among them: majority and minority gases, RF antenna configuration, density, heating profile and radiated power fraction. Nevertheless, certain heating schemes produce stabilization with high reliability for input powers above the threshold (e.g. ICRH monopole heating on-axis in D(He4) with Ip ≤ 3.5 MA).

Evolution of the principal discharge parameters for one of the longest sawtooth-free periods obtained is shown in figure 1. As indicated in the figure, the plasma current was ramped

<table>
<thead>
<tr>
<th>Table 1: Conditions for sawtooth stabilization</th>
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<tbody>
<tr>
<td><strong>Combined heating (ICRH + NBI):</strong></td>
</tr>
<tr>
<td>- ICRH on-axis, NBI co-injection.</td>
</tr>
<tr>
<td>- Ip = 2 – 5 MA ( (q^\psi = 6 – 3.4) ).</td>
</tr>
<tr>
<td>- Paux &gt; 7.5 MW ( (&lt;n_e&gt; &lt; 4 \times 10^{19} \text{ m}^{-3}) ).</td>
</tr>
<tr>
<td><strong>ICRF heating alone:</strong></td>
</tr>
<tr>
<td>- ICRH on-axis.</td>
</tr>
<tr>
<td>- Ip = 1.5 – 4 MA ( (q^\psi = 4.6 – 9) ).</td>
</tr>
<tr>
<td>- Paux &gt; 3 MW ( (&lt;n_e&gt; &lt; 3 \times 10^{19} \text{ m}^{-3}) ).</td>
</tr>
<tr>
<td>- H, He^3 minorities, mono/dipole antenna.</td>
</tr>
<tr>
<td><strong>NBI heating alone (2 cases in SNX-point):</strong></td>
</tr>
<tr>
<td>- Ip = 2 MA, Paux = 7 MW, ( &lt;n_e&gt; = 2 \times 10^{19} \text{ m}^{-3} ).</td>
</tr>
<tr>
<td>- Ip = 3 MA, Paux = 7 MW, ( &lt;n_e&gt; = 3.5 \times 10^{19} \text{ m}^{-3} ).</td>
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during this period, since there is evidence that this may prolong the stable period. The central electron temperature is usually found to saturate on a timescale of the order of the energy confinement time. However, the density, ion temperature and stored plasma energy can continue to rise throughout the stable period, depending on the heating scheme employed.

Transport and Confinement — It has previously been reported that electron thermal transport in this regime is well represented by a model in which profile or temperature gradient constraints are invoked. Recent analysis of energy confinement during stable periods has confirmed that electron energy confinement is L-mode like (with the exception of stable periods achieved in H-modes), but that total energy confinement is 15–20% higher than in equivalent sawtooth discharges, in part as a result of the central accumulation of fast ions which are normally depleted at \( I_p = 3 \text{ MA}, q_0 = 4.5, \langle n_e \rangle = 2.5 \times 10^{19} \text{ m}^{-3} \) sawtooth collapses. This suggests that sawtooth stabilization may be important in the confinement of \( \alpha \)-particles in the near-ignition regime.

The absence of sawteeth also permits the determination of particle transport coefficients within the \( q=1 \) surface from the analysis of density profile evolution. It is found that the diffusion coefficient at \( r/a = 0.35 \) is \( 0.75 \text{ m}^2\text{s}^{-1} \), which is of the same order as that outside the \( q=1 \) surface in equivalent sawtoothing discharges. These results indicate that the principal advantage deriving from the stabilization of sawteeth will not be due to improved confinement inside the \( q=1 \) surface, but rather to increased fusion reactivity, due to the peaked profiles obtained, which may be further enhanced if non-thermal fusion scenarios are exploited.

Analysis — As reported previously, the sawtooth collapse terminating these periods follows the usual JET behaviour, which led originally to the conjecture that the stabilization occurred due to non-inductively driven or pressure driven current. However, recent analysis of the current profile evolution in this regime has shown that the central safety factor \( q_0 \) attains a value well below unity before the final sawtooth collapse. Figure 2 shows the evolution of \( q_0 \) during a sawtooth-free period as determined by FIR Faraday rotation measurements (polarimetry), calculation of resistive diffusion (TRANSPP?) — which at present excludes bootstrap current — and magnetic equilibrium analysis (IDENTC). The three measurements agree within systematic errors (\( \pm 20\% \)), and show a monotonic decrease of \( q(0) \) during the stable period to values - 0.7—0.9. Since the sawtooth instability in JET strongly resembles an ideal mhd mode, this represents a considerable theoretical problem.
Discharges in which sawtooth-free periods occur are characterized by a population of energetic ions, principally RF-accelerated minority ions, with energies in the range of several hundred keV. The non-thermal $\beta_p$ may be significant therefore, - 0.1, and it is thought that these ions may play a significant role in the stabilization mechanism. Figure 3 illustrates an experiment designed to investigate the role of additional heating. By lengthening the RF pulse it was possible to extend the stable period, up to a maximum of 3 s. When the ICRH was switched off during the stable period, the delay until a sawtooth collapse was a small fraction (~2%) of the central resistive skin-time, but was of the order of the central fast ion slowing-down time $\tau_{FS}$, where,

$$\tau_{FS} (s) = \frac{n_e}{0.06 T_e} \left(10^{19} \text{ m}^{-3}\right),$$  

which was ~0.5 s for these conditions. Although the energy confinement time was also of this order, the decrease in $\beta_p$ accompanying the switch-off was predicted to be stabilizing. This suggests, therefore, that the loss of fast ions was the dominant destabilizing effect.

**Theory** — The influence of an energetic ion distribution on the stability of the resistive and ideal $m=1$ modes has been investigated using the experimental observations that $q_o < 1$ and $\beta_p \approx 0.3$ during sawtooth-free periods. It is found that trapped hot ions introduce a new threshold for the $m=1$ instability. Stabilization arises from the fact that the magnetic drift motion of the hot ions exceeds the phase velocity of the $m=1$ mode, with the result that additional energy is required to perturb this population.

$$\left[-\omega (\omega - \bar{\omega})\right]^{-\frac{1}{2}} = \omega_A [\lambda_H + \lambda_R (\omega)]^{-\frac{1}{2}} \frac{Q^3/2}{8} \frac{\Gamma((Q - 1)/4)}{\Gamma((Q + 5)/4)},$$  

Under these conditions the dispersion relation for the mode is
where \( Q^2 = i\omega (\omega - \omega_i^*) (\omega - \omega_e^*) / \epsilon_\eta \omega_A^3 \), with \( \omega_i^* \) and \( \omega_e^* \) the electron and ion diamagnetic frequencies, \( \epsilon_\eta = \eta c^2 / (4\pi \sigma_0 \omega_A) \) the inverse magnetic Reynolds's number, \( \eta \) the resistivity, \( \omega_A = \delta v_A / \sqrt{3} R_o \) with \( v_A \) the Alfvén velocity, \( s = r_0 q' (r_0) \) and \( q(r_0) = 1 \). The parameter \( \lambda_H \) (see ref 11) is proportional to the mhd energy functional \( \delta W \) \( \lambda_H > 0 \) for ideal mhd instability.

Finally, the parameter \( \lambda_K(\omega) \), which represents the hot ion effects, is defined by
\[
\lambda_K(\omega) = \frac{4\pi^2 i}{(B_0^2) \xi_0^2} \int_{r_0}^{r_0} dr \int_0^{2\pi} \frac{d\theta}{2\pi} \xi \hat{e}_x \cdot \hat{n} \hat{p}_\perp h(\omega)
\]
(4)
where \( \xi_0 \) is the radial displacement, \( \xi \) is the curvature vector and \( \hat{p}_\perp h(\omega) \) is the perturbed transverse hot ion pressure. When the bounce averaged magnetic drift frequency \( \omega_{p,h} \) of the hot ions is larger than the mode frequency \( \omega \), but \( (1 - q_0)^2 < \omega / \omega_{p,h} \), \( \lambda_K \) can be approximated as
\[
\lambda_K(\omega) \approx - (\pi / 3 \xi^2) (r_0 / R_o) \beta_{p,h} \omega / \omega_{p,h}
\]
(5)
where,
\[
\beta_{p,h} = - [8\pi / B^2 (r_0)] \int_{r_0}^{r_0} dr \left( r / r_0 \right)^{3/2} \frac{d}{dr} \left( r / r_0 \right)^{1/2} p_{\perp,h}
\]
(6)
Resistive internal modes are unstable with a growth rate \( \gamma / \omega_A - \epsilon_\eta \xi^{1/3} \) in the limit where \( |\lambda_H + \lambda_K| < \epsilon_\eta \xi^{1/3} \) (and \( \omega_i^* - \omega_e^* \), \( \omega_e^* / \omega_A < \epsilon_\eta \xi^{1/3} \)). However, if \( \beta_{p,h} \) is sufficiently large, the relevant regime satisfies \( |\lambda_H + \lambda_K| > \epsilon_\eta \xi^{1/3} \). If \( (\lambda_H + \Re \lambda_K)^2 < 0 \), the m=1 tearing mode should be stable because of a combination of toroidal and high temperature effects. In the remaining case, \( (\lambda_H + \Re \lambda_K)^2 > 0 \), the dispersion relation reduces to
\[
\omega(\omega - \omega_i^*) = - \omega_A^2 \left| \lambda_H + \lambda_K(\omega) \right|^2
\]
(7)
If also \( \omega_i^* \) and \( \lambda_K \to 0 \), the ideal mhd dispersion relation is recovered. Using the approximate expression (5), it is found that the ideal internal kink mode is stabilized by the hot ions when
\[
\beta_{p,h} > \alpha (r_0 / R_o) \beta_{p} \omega_p / \omega_i^*,
\]
(8)
where \( \alpha \) is a numerical factor, of order unity, primarily determined by the \( q \)-profile. For reasonable estimates of the hot ion distribution, this condition is consistent with experimental values of the relevant parameters. This result also reflects the fact that it is difficult to produce sawtooth stabilization when the RF deposition profile is centered significantly off–axis, since this would produce insufficient peaking in the profile of \( p_{\perp,h} \) within the region where \( q < 1 \) to yield a positive \( \beta_{p,h} \). More detailed information on the energy and phase space distribution of the hot ions is required, however, to enable a more detailed comparison with the experimental results to be undertaken.

We wish to acknowledge R Goldston and D McCune of PPPL for the use of TRANSP, and C Best for supervising its installation at JET.

References

THE DYNAMIC BEHAVIOUR OF THE ELECTRON TEMPERATURE PROFILE IN THE TEXTOR TOKAMAK PLASMA

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1. Introduction

The time evolution of the electron temperature was studied in TEXTOR discharges by means of a 8-channel heterodyne ECE-system of very low noise figure. Typical profile developments during current ramping, gas injection and soft landing phases were detected by continuous recording of ECE-data and off-line computer processing. The present time resolution of 100psec allowed the observation of rapid profile changes during disruptive events of the discharge. The spatial resolution of the $T_e$-measurement in radial direction was $\Delta R = 2$ cm, in vertical direction it was $\Delta Z \leq 5$ cm. The ECE-channels were relatively calibrated by means of an oven with a surface temperature of 600°C and absolutely by a series of soft x-ray temperature determinations.

2. $T_e$-profiles during OH-discharges

Strong dynamics of the $T_e$-profile are most clearly observed during the early phase of the discharge before reaching the current plateau (see fig. 1). After start-up a skin effect is seen in a broad ECE-radiation profile. The onset of the positional feedback adjusts the plasma to a prefixed position and the skin effect disappears soon. A peaked profile then begins to develop. With the beginning of the sawtooth (st) activity the $T_e$-profile broadens and the peak temperature falls. In fig. 2 a time window of 100 msec duration during the plateau phase is shown. It reveals a periodic undulation of the $T_e$-profile. This oscillation is caused by st-activity. In the down slope of the $T_e$-profile at a fixed radial position a cut is made. It visualizes the propagation of heat pulses which are launched from the q=1-region /1/. Details of the st-crash phase are shown in fig. 3. At the end of the st-period the central temperature peaks. The rounded profile gets first eroded by the action of a precursor oscillation. At the end of the crash a flattened profile with shoulders outside of the q=1 surface was observed. The central $T_e$-crash occurred in about 300psec. Then the central $T_e$ slowly begins to rise again. The shoulders move radially outwards and disappear and a new st-cycle has begun.

The st-activity with its profitable influence on plasma performance can be lost by an experimental accident e.g. by an accumulation of impurities. The sawteeth then do not start or even are stopped. The electron temperature on axis can fall far below the normal value usually observed in a st-discharge. A hollow $T_e$-profile was found under these conditions.
and in some cases the discharge ended with a hard disruption. Fig. 4 shows such a profile. The bolometric radiation diagnostic showed a radially strongly peaked profile. It emits the local Ohmic heating power input from the center. Discharges were even observed which began their st-activity only after a small disruption had occurred and probably some impurities had been swept out.

The end of the discharge is generally initiated by a programmed soft landing. The $T_e$-profile shrinks and the central temperature falls rapidly at the end of the current pulse.

Fig. 1 $T_e$-profile evolution during the start-up phase of the TEXTOR discharge.

Fig. 2 Sawtooth activity during the plateau phase of an OH-discharge. Heat pulse propagation is seen at a radial position $R=24$ cm off axis.
3. The density limit disruption

The electron temperature profile development plays a dramatic role in the prehistory of the density limit disruption. With an increasing average density and a slowly peaking density profile the electron temperature profile was observed to contract and the plasma detached from the limiters. The current profile contracted and the current density in the center increased noticeably as detected by polarimetric measurement technique /3/. This process is accompanied by a strong increase in the total radiation from the outer plasma zone. Oxygen and carbon were the most prominent impurities identified in this phase. At this critical stage of the contraction the m=2-mode is triggered. It is observed magnetically and appears as a precursor oscillation in the ECE-channels positioned close to the q=2-surface. Fig. 5 shows a sequence of T_e-profiles at the density limit disruption. During the last st-cycle before disruption precursors in ECE are detected. For a q(a)=5.4 plasma the sawteeth remained to the very end of the discharge. Fig. 5 demonstrates details of the electron energy quench phase. Under the influence of the growing m=2-mode the T_e-profile shows at first in the equatorial plane phases of rapid contraction and relaxation. Then the central temperature begins to drop and a visible radial structure is preserved during the energy decay. Possibly a magnetic flux tube is formed which discharges a major part of the electronic energy from the center into the plasma boundary within less than 1 msec. During this energy quench the plasma current begins to increase to a little tip only. The current quench itself begins later and lasts still for about 40 msec.

![Te-profile graph](image)

Fig. 3 T_e-profile before, during and at the end of the sawtooth crash in an OH-discharge in TEXTOR.
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Fig. 4
Hollow $T_e$-profile of a non-saw-toothing discharge.

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Fig. 5
$T_e$-profile development during a density limit disruption. The time intervals between different profiles are 200 µs. The energy quench lasts less than 1 msec.
ROLE OF SAWTOOTH CRASHES IN \( \beta \) SATURATION AND COLLAPSE IN THE PBX TOKAMAK


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1. Introduction

Because the high \( \beta \) operation of a commercial fusion reactor is essential for its economy, it is important to experimentally study the stability of high \( \beta \) discharge in today's tokamaks. Several tokamaks have in recent years achieved \( \beta \)'s close to the stability limit predicted for the kink and ballooning modes, and they have investigated the MHD activity, energy confinement properties, and energy losses at this limit. Of particular interest is the correlation between the observed MHD activity, deterioration in confinement, and the theoretical predictions for the \( \beta \) limit. In earlier PBX papers, we discussed this correlation for several regions of the Troyon-Gruber parameter \( \beta_c = \frac{\mu_0 I_p}{a_{mid} B_{t}} \). The highest \( \beta \)'s were achieved at highest \( \beta_c \) (>2), lowest \( q \)-edge, and with a very high current ramp (1.5 MA/s). The investigation of \( \beta \) saturation and collapse in PBX at low \( \beta_c \) = 0.5, high \( \beta_p \) and high \( q \)-edge has shown a good correlation between the energy losses causing the saturation and collapse and the MHD activities in the plasma. A similarly good correlation was also found at medium \( \beta_c \) (=1) and \( q \)-edge. This paper will focus on slightly higher \( \beta_c \) (=1.5 to 1.7), medium \( q \)-edge and lower \( \beta_p \).

2. Experiment and Results

The sawtooth characteristics as well as the sawtooth influence on energy confinement can depend very much on some of the discharge parameters. One of the discharge parameters, that seems to considerably change the MHD and sawtooth characteristics of the discharge and the energy loss mechanism, is the current ramp. This is especially true for the discharges near the \( \beta \) limit. In the following we describe two types of discharges whose main difference is the value of the current ramp. Although these discharges have similar \( \beta \) limits, they experience very different energy loss activity during the \( \beta \) saturation and collapse. In our analysis we interpret the change in the fast diamagnetic loop signal as a change in energy.

a) No current ramp case (\( \frac{d I_p}{dt} = 0 \)). The plasma parameters in these discharges were as follows: \( I_p = 320 \text{ to } 350 \text{ kA}, B_t = 2.3 \text{ T}, q \text{-edge} = 4.8 \text{ to } 5, a_{mid} = 30 \text{ cm}, R = 144 \text{ cm}, P_{neut} = 2.8 \text{ MW}, \text{ plasma elongation } \kappa = 1.8, \text{ indentation} = 18\%, \langle \beta_t \rangle = 2.2\%, \text{ and } \beta_c = \frac{\mu_0 I_p}{(a_{mid} B_{t})} = 1.5. \)

The \( \beta \) saturation phase of these discharges is characterized by several MHD loss mechanisms responsible for the \( \beta \) saturation. The continuous, discontinuous or bursting energy losses, which cause the \( \beta \) saturation and collapse, we will call "additional losses" in order to separate them from the baseline energy losses, that are associated with the nonperturbed confinement. The energy losses associated with the sawteeth and measured by a fast diamagnetic loop were normally low, allowing a claim, that in these discharges the sawteeth represented a minor energy loss. These losses for the four sawteeth in Fig. 1. were 1.2% of the total plasma energy (at 517 ms), 1.9% (at 549 and 587 ms), and a large sawtooth induced loss of 22% (at 625 ms) just before the current disruption. The last large energy loss sawtooth is an exception and it happens sometimes just before the major disruption. In Fig. 2, the relative sawtooth modulation amplitudes for the four sawteeth are shown, the highest modulation being associated with the highest energy loss. These sawteeth differ strongly from the ohmic sawteeth in at least one aspect: the position of the sawtooth inversion radius is much larger (almost double: \( z = \pm 28 \text{ cm} \) than the position of the maximum precursor modulation (\( z = \pm 18 \text{ cm} \) for the large sawtooth at 625 ms and \( \pm 15 \text{ cm} \) for the other three smaller sawteeth) both of which are usually associated with the \( q = 1 \) surface. Only in the case of...
first, smaller energy loss sawtooth these two positions almost coincide (z = ±15 cm). The heat pulse propagation time constant is much shorter for the giant sawtooth (80 μs) than for the three smaller sawteeth (~250 μs). At the sawtooth crash time one observes also a 20% drop in neutron flux indicating that the fast beam particles are being ejected from the center of the discharge. Although the sawtooth energy loss per event represents the largest instantaneous loss, the average loss is a smaller part of the total additional energy loss, that causes β saturation and collapse, because of the rather low sawtooth repetition rate.

The energy losses between the sawteeth and not during the sawtooth crash itself are more important in causing the β saturation and collapse. One can identify several MHD phenomena contributing to this: ERP's (or ELM's), a global continuous mode with a strong m = 1, n = 1 component and "minicrashes" - small sawtooth-like crashes on surfaces beyond the q = 1 surface. The inversion radius of these minicrashes is between z = ±30 and ±40 cm, placing these minicrashes at the magnetic surfaces with q larger than 1 but less than 2. A series of many repetitive minicrashes is seen in Fig. 1. between 558 and 577 ms. The energy loss per minicrash is typically several tenths of a percent, but because of the high repetition rate, they can contribute considerably to the total additional losses. A considerable fraction of the additional losses cannot be ascribed to any MHD activity. This can be noted in Fig. 1. between each pair of sawteeth, where, although no MHD activity is seen, one observes a net plasma energy loss in the diamagnetic loop (550 to 553 ms, 549 to 558 ms, etc.) or a stagnation or smaller increase than expected in the plasma energy. The mechanism causing these losses could not be identified.

b) Current ramp case (dI/dt = 0.5 MA/s). Most of the plasma discharge parameters were similar to the "no current ramp" case. The parameters with different values were: I_n = 400 kA, q-edge = 4.5, P_inj = 5 MW, plasma elongation κ = 1.6, <β> = 3%, and β_e = 1.5 to 1.7.

The most important difference to the "no current ramp" case is that the β saturation and collapse in these discharges are caused almost solely by the extreme energy losses during the course of the sawtooth crashes. Only a small part of the additional losses were caused by the ERP's and other MHD activity. In Fig. 3, a typical behaviour of various plasma signals is shown during the β saturation part of the discharge. The dominating events in energy loss are the sawtooth crashes during which one observes huge losses of plasma energy (11% at 610 ms, 2.4% at 633 ms, and again 11% at 653 ms). Although the sawtooth repetition rate is not very high, the large energy loss per event alone forces the β saturation. Although not shown in this figure, the minicrashes can also appear between the sawteeth, but their contribution to the additional losses is small. The relative sawtooth crash amplitudes for these three sawteeth are shown in Fig. 4. These differ considerably from those in Fig. 2. The inversion radius is again very large (z = ±22 cm or even ±27 cm in the case of the first sawtooth, case a). However, the position of the corresponding relative precursor modulation peak is close to z = ±10 cm. This large difference between the sawtooth inversion radius and the position of the of the maximum precursor modulation must be significant and should tell us something about the character of these large energy loss sawteeth. The precursor is again m = 1 / n = 1 mode. The sawtooth with the smallest energy loss of 2.4% has the shallowest crash amplitude (case b). The sawteeth with the higher energy loss of 11% have a deeper crash modulation and the heat pulse amplitude beyond the inversion radius is small. In case a this heat pulse amplitude is extremely small, which could indicate that the plasma confinement even in this region (q > 1) becomes very poor during the crash. There is also an indication of several maxima and minima in the crash amplitude modulation, which could mean, that there are several alternating regions with a better or poorer confinement. In-out crash amplitudes in the major radius direction in Fig. 4, also show, that the heat pulse propagates inward. This could indicate that the outer regions (regions of larger R) suffer more from the loss of confinement during the crash.

3. Discussion

The differences in the MHD and energy loss behavior between two types of discharges allow us to conclude that a strong current ramp stabilizes certain modes. It is not clear which agent stabilizes these modes: is the stabilization caused by the current profile broadening due to the skin effect or is it a consequence of the pressure profile change. What is obviously stabilized by the strong current ramp is the activity causing the "unknown loss"
This decrease in continuous mode and "unknown activity" losses in the strong current ramp case is fully compensated by very large sawtooth losses. These large sawtooth energy losses are observed only at high β's. The nature of these strong sawtooth losses is not understood. One can propose several scenarios, which could explain this sawtooth behavior. Since the sawtooth inversion radius is much farther out than the position of the peak precursor modulation, one can imagine that this large loss of energy is caused by the loss of confinement also in the region beyond q = 1. This loss of confinement in the region for q > 1 can be caused by either simultaneous (with the sawtooth crash) or very fast successive ergodization of this region. The proof that this is possible are the minicrashes, that occur independently of sawtooth on magnetic surfaces with q > 1. The simultaneous or the fast successive ergodizations on various (rational) surfaces between q = 1 and the edge of the plasma would then explain the sudden onset of the large energy loss during the sawtooth crash. A criticism against this thesis is that no mode other than m = 1, n = 1 is seen before the crash. However, one might reply, that no minicrash precursor has been observed even in the case of larger minicrashes. Another mechanism causing a sudden loss of confinement can be invoked from the Fig. 4., where one sees in the bottom part of the figure a large difference in "inboard" and "outboard" confinement. Since the confinement for the large outboard major radius is much poorer than for the inside part of the discharge, one could invoke instabilities, that are less stable on the outboard side. Again, possible candidates are the ballooning modes and drift waves. These modes could be presumably excited by a large heat pulse, resulting from the original sawtooth on the q = 1 surface, accompanied by an increase in the pressure gradient beyond q = 1. However, part of this in-out asymmetry is caused by the inward movement of the plasma column into a new equilibrium position, which occurs because of the large loss of plasma energy. No matter what is the cause of the large energy loss during these sawtooth crashes one can be sure, that it is not a simple case of ergodization or reconnection only around the q = 1 surface.

Some of the statistics for these sawteeth is shown in Figs. 5 and 6 as a function of the heat pulse propagation constant τ_{hp}, as observed in soft x-rays, and energy decay rate τ_{δΑδ}, as observed in diamagnetic loop signal. The larger the energy loss ΔW/W the shorter the heat pulse propagation constant, another indication of decreased confinement between q = 1 and the edge. On the other hand the larger losses cause longer leaking (τ_{δΑδ}) of the energy at the edge of the plasma. The triangles represent the high current ramp case and the full circles the no ramp case. From the statistics is again clear, that the energy losses in the no ramp case are much less severe.

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Fig. 1. Diamagnetic loop and other signals during $\beta$ saturation in no current ramp case.

Fig. 2. Vertical and horizontal relative sawtooth amplitude during the four sawteeth of Fig. 1.

Fig. 3. Various plasma signals during $\beta$ saturation in the strong current ramp case.

Fig. 4. Vertical and horizontal relative sawtooth amplitude for three sawteeth of Fig. 3.

Fig. 5. Relative sawtooth energy loss as a function of the heat pulse propagation constant. $\Delta - \frac{dl_p}{dt} = 0.5 MA/s; \alpha - \frac{dl_p}{dt} = 0$.

Fig. 6. Relative sawtooth energy loss as a function of decay time constant of the diamagnetic loop signal.
CENTRAL ELECTRON POWER DEPOSITION FROM dTeo/dt MEASUREMENTS ON TFTR


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Introduction. The sawtooth, or internal disruption, which is the slow rise and abrupt drop in central electron temperature and density, is a feature of many tokamak discharges. Details of its behavior, especially those which immediately precede the collapse, are quite complicated and still not completely understood.\(^1\)\(^2\)\(^3\) Yet immediately after an internal disruption the central region of the plasma is quite simple. The electron density and temperature profiles are approximately flattened over the central portion of the discharge. The extent of this region depends on safety factor at the edge of the plasma; for example, for \(q(a)\approx 4\), this region, in which \(q=1\) according to the Kadomtsev model, extends to a normalized radius of about 0.25.

Immediately after a sawtooth collapse, the power balance for the electrons at the center of the discharge can be written as:

\[
P_{\text{net}} = P_{\text{oh}} + P_{\text{bf}} + P_{\text{be}} + P_{\text{ie}} - P_{\text{rad}}
\]

(1)

Here \(P_{\text{oh}}\) is the ohmic heating power, \(P_{\text{bf}}\) is the power loss due to beam fueling (-1.5Teo-dNeo/dt). \(P_{\text{be}}\) is beam ion to electron power. \(P_{\text{ie}}\) is the thermal ion to electron power and \(P_{\text{rad}}\) is the radiated power. The power lost by conduction from the plasma center will be ignored since electron temperature profile is approximately flattened by the collapse; similarly the power convected to or from the central region will be ignored since the density profile is approximately flat and particles convected to the center from adjacent regions are at the same temperature as those at the center.

Thus, a measure of the rate of rise of central electron stored energy is a direct measure of the central power deposition:

\[
\frac{3}{2} N_e \Delta T / \Delta t = P_{\text{net}}
\]

(2)

This measurement can be checked against the central power deposition as computed by SNAP, a steady state, one dimensional tokamak simulation code.

The rate of rise of Teo was obtained from a twenty channel grating polychromator\(^4\) which measured second harmonic electron cyclotron
emission from the plasma. The excellent spatial (3 cm radial resolution and 5 cm antenna spot size) and temporal (5 kHz digitizing rate for 2 seconds) resolution make it ideally suited for this experiment. Channel spacing of the instrument is such that there are usually eight channels at or inside the inversion radius (for $q(a)=4$), making it relatively easy to measure the rate of rise at the plasma center.

It is necessary to measure several other plasma parameters in order to calculate the central power deposition in SNAP. Central ion temperature and central plasma toroidal rotation velocity are obtained from x-ray crystal spectroscopy measurements of doppler broadened and doppler shifted Kα resonance lines of He-like ions such as NiXXVII or FeXXV. Ion temperature profiles and plasma toroidal rotation velocity ($V\phi$) profiles are available from CHERS which measures doppler broadened and doppler shifted lines from charge exchange recombination of CVI in the visible region. Radiated power is measured with a bolometer array, $Z_{eff}$ is obtained from visible bremsstrahlung or x-ray PHA measurements and electron temperature and density profiles from Thomson scattering, with the density profile normalized to the line average density from a far infrared interferometer. The grating polychromator was calibrated to Thomson scattering $T_e$ data for these measurements.

Results from Beam Heated Discharges. Beam heated discharges with the following parameters have been studied: $R=2.5m$, $a=0.8m$, $B=4.7T$, $1.2MA<I_p<2MA$, $1.5MW<P_{HEAT}<19MW$, $2.2<Z_{eff}<5.5$. Almost all discharges in this data set were coinjection only.

A comparison of the central electron power deposition measured from the rate of rise of sawtooth with that calculated using SNAP is shown in Fig.1, where the ratio of the measured to calculated central power deposition is plotted versus total heating power. SNAP calculations included rotation and beam driven current effects and used measured $T_i$ values and $T_e$ and Ne profiles from Thomson scattering, which were not the profiles at the bottom of the sawtooth. Thus the $P_{ei}$ calculated by SNAP had to be adjusted using $T_e$ measured at the bottom of the sawtooth from the grating polychromator and $T_i$ values calculated by SNAP.

Agreement between the measured and calculated values of the central power deposition (Fig.1) is rather poor for reasons which are not entirely clear at present. One possible explanation is that there is also a large sawtooth feature on the central ion temperature and central density. (The $T_e$ sawtooth amplitude is exceptionally large - $\Delta T_e/T_e=0.35$ - for these discharges). In high power beam heated discharges, $P_{ei}$ is the dominant term in the power balance at the bottom of the sawtooth, and decreases by a factor of two (assuming nothing else changes) between the bottom and the top of the sawtooth due mostly to the $T_e^{-1.5}$ dependence of
Pei. This is evidently very difficult to take into account given the 100 ms time resolution on Ti measurements and the problem in determining Neo with high precision.

In the single case in which the Thomson scattering density profile is available at the bottom of the sawtooth (P=11.4MW), the measured central power deposition was 262 mW (determined from dTeo/dt measurements and Neo from Thomson scattering). Calculated central power deposition from SNAP was 352 mW for a measured average central ion temperature of 13.5 keV. An ion temperature sawtooth of ΔTio/Tio=0.2 [or for Tio=13.5keV, Tio,min=10.5keV] is required to account for this discrepancy between the measured and calculated central power deposition. This is not unreasonable: in this discharge ΔTio/Tio=0.3. It must, however, be verified with detailed Tio measurements.

Shown in Fig. 2 is a plot of measured central power deposition versus heating power normalized to central density. There is a distinct saturation in central power deposition, as might be expected, due to rising plasma density and a corresponding broader beam power deposition.

Results from Ohmic Discharges. Central power deposition for ohmic discharges has also been investigated using SNAP. The ohmic power deposition is computed from Te and Zeff measurements using the formula for the classical Spitzer conductivity as given in Ref. 5, and assuming that q=1 inside the inversion radius. Using Pei and Prad from SNAP, the net power input (Eqn. 1) was compared to that determined from the rate of rise of Teo. The results are shown in Fig.3, where the ratio of the measured central power deposition to that calculated assuming q(0)=1 is plotted for a number of TFTR discharges. In all cases the measured central power deposition is larger than the calculated value. The most interesting and unambiguous cases are for those discharges labeled with crosses (+) in Fig.3. These are low density (Neo<1.8E19m⁻³), high Teo (Teo=5keV), high Zeff discharges. They are nearly pure carbon plasmas in which Prad and Pei are small (<15% of the measured input power). The measured input power is clearly much larger than what is calculated assuming q(0)=1.

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Fig. 1 Ratio of measured to calculated central power deposition vs total heating power.

Fig. 2 Measured central power deposition vs heating power normalized to central density.

Fig. 3 Ratio of measured to calculated $P(0)$, assuming $q(0)=1$, for ohmic TFTR discharges. The discharges indicated by + are high Zeff, low density.
MHD Stability in High $\beta_T$ DIII-D Divertor Discharges*


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DIII-D has achieved a volume averaged toroidal beta $\beta_T$ of over 6% for values of $I_P/aB_T$ varying from 1.8 to 2.5 in H-mode divertor plasmas [1], as summarized in Fig. 1. The safety factor at 95% of the enclosed poloidal flux $q(0.95)$ ranged from 2.2 to 2.8. Stable, high $\beta_T \geq 5\%$ discharges, with $\beta_N = aB_T\beta_T/I_P \leq 2.3$, have also been sustained for many energy confinement times $\tau_E$. As $\beta_N$ is increased toward the Troyon limit ($\beta_N \simeq 3.5$), a rich variety of MHD activity appears [2]. High toroidal mode number $n=3$, 4 magnetic fluctuations first occur, preceding the $n=1$, 2 modes which may grow, stop rotating, and cause beta collapse and lead to plasma disruptions. This is illustrated in Fig. 2, where the power spectrum obtained from the measured Mirnov signals for a high $\beta_T$ DIII-D divertor discharge is shown. This particular discharge has a maximum $\beta_T$ value of 5.2% and $\beta_N$ of 2.9. Secondly, when $\beta_N$ is raised, the relative fluctuation amplitudes tend to increase as the unstable plasma region propagates radially outward. The poloidal mode number $m$ generally varies as $n+1$, which is consistent with the results from the ideal MHD stability calculations.

The high $n=3$, 4 modes do not seem to have any apparent effects on $\tau_E$, whereas the $n=1$, 2 modes can degrade $\tau_E$ significantly. In the absence of the $n=1$, 2 modes, high $\beta_T$ itself does not seem to have any apparent effect on $\tau_E$. However, $\tau_E$ in these low $q$ ($q(0.95) \leq 3$) high $\beta_T$ H-mode discharges are lower than those found in high $q$ ($q(0.95) \geq 3$) low $\beta_T$ high quality H-mode discharges. These low $q$ high $\beta_T$ plasmas usually have frequent giant edge localized modes (ELMs) with periods as short as a few milliseconds and a large sawtooth inversion radius ($0.3 - 0.5a$). The interaction of these two effects may play a role in the degradation of $\tau_E$ in these plasmas.

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Detailed ideal kink and ballooning stability analysis has been carried out for a $\beta_T = 6.2\%$ and $\beta_N = 2.5$ discharge. The plasma MHD equilibrium is reconstructed self-consistently using measured kinetic profile data as well as magnetic data. The results show that the plasma is stable to the $n = 1, 2, 3$ ideal modes without a conducting wall and is stable to the infinite $n$ ballooning mode over most of the plasma region. This is shown in Fig. 3, where the pressure gradient determined from the experimentally measured data is compared against the theoretically computed first ideal ballooning marginal pressure limit for various plasma flux surfaces labeled by the enclosed normalized volume $V_N$. Near the plasma center, the result depends on essentially unmeasurable details in the central shear. Near the edge, the plasma is approaching the first marginal ballooning mode pressure limit, particularly just before a giant ELM. This is further illustrated in Fig. 4, where the experimentally measured normalized pressure gradients $2\mu_0 R_0 \partial P/\partial z/\partial B_n^2$ near the plasma edge along the line of sight of the Thomson measurements are compared against the theoretically computed first ideal marginal ballooning pressure limits for a number of high $\beta_T = 4-6\%$ DIII-D divertor plasmas with various values of $S(0.95)/q(0.95)^2$. Here, $S(0.95)$ is the plasma global shear at 95% of the enclosed poloidal flux. Before the occurrence of a giant ELM, the edge pressure gradient is near or exceeding the ballooning limit, whereas after an ELM, the edge pressure gradient drops below the ballooning limit. The details of the analysis are given in Ref. 3. These results are similar to those observed in low $\beta_T$ H-mode DIII-D divertor plasmas [4]. The correlations between the observed edge pressure gradients and the theoretically computed ballooning marginal pressure limits suggest that the ideal ballooning mode limits the achievable pressure gradient built up in the edge region, although not the total volume averaged beta $\beta_T$ in these discharges, and may be responsible for the appearance of the giant ELMs.

In summary, DIII-D has achieved a $\beta_T$ of over 6% in H-mode divertor plasmas and has not exceeded the beta limit as allowed by ideal MHD instabilities.

References


Fig. 1. Toroidal beta $\beta_T$ achieved in DIII-D divertor plasmas as a function of $I_p/aB_T$(MA/mT).

Fig. 2. Power spectrum from Mirnov measurements for DIII-D divertor discharge 55771, which has a maximum $\beta_T$ of 5.2% and $\beta_N = aB_T\beta_T/I_p$ of 2.9.
Fig. 3. Comparison of measured pressure gradient (solid line) against theoretically computed marginally stable first regime ideal ballooning stability pressure gradient limit (dashed line) for various plasma flux surfaces labeled by the enclosed normalized volume $V_N = V(\rho)/V(\alpha)$.

Fig. 4. Comparison of the measured normalized pressure gradient ($T_i = T_e$) along the vertical Thomson line of sight against the theoretically computed marginally stable first regime ideal ballooning stability pressure gradient limit for a number of high $\beta_T$ DIII-D plasmas before and after the occurrence of giant ELMs.
CHARACTERISTICS OF LOW-q DISRUPTIONS IN PBX


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The objective of the Princeton Beta Experiment (PBX) was to push into the second regime of stability to high-n ideal ballooning modes/1/. To accomplish this goal, the PDX vessel was modified in order to produce plasmas indented on the inboard side; the indentation increased the plasma stability to the high-n ballooning/2/ and low-n internal kink modes/3/. However, to insure against extreme fast ion loss associated with the internal kinks, two of the four PDX neutral beamlines were re-oriented from perpendicular to a tangential injection direction. The results of approximately two years of operation indicated that, although PBX achieved then world record values of volume averaged toroidal beta, \( \langle \beta_t \rangle = 5.5\% \), the plasmas still resided in the first stability regime. Various mechanisms were responsible for this limitation; at high-q, the plasma stored energy was seen to saturate in time and then decrease over the course of tens of milliseconds at constant plasma current and beam power/4/. This saturation and collapse was associated with several types of MHD activity: Edge Relaxation Phenomena (ERPs), a saturated amplitude continuous mode that locked as the plasma rotation velocity went to zero, large sawteeth, and both parallel and perpendicular fast ion loss associated with internal kink oscillations. At low-q, the plasma \( \langle \beta_t \rangle \), and the plasma itself, were limited by a hard disruption during the neutral beam injection and strong current ramp phase of the discharge. The strong current ramp (in excess of 1.5 MA/sec) was employed to aid in broadening the plasma current profile. The subject of this paper is to examine the disruptions observed in discharges with \( q_{95} \leq 4.5 \), where \( q_{95} \) is the q-value at the 95% flux surface as determined from a flux-based, between shots equilibrium calculation. Typical discharge parameters for these types of shots were \( I_p = 450 \) to 600 kA, \( B_t = 0.8 \) to 1.2 T, \( P_{inj} = 3.6 \) to 6.0 MW, and densities up to 6 x \( 10^{19} \) cm\(^{-3}\). While the mode responsible for the disruption may have been, to zero order, an \( n=1 \) ideal external kink, as concluded in

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previous studies/5,6/, we will report considerable evidence that suggests that non-ideal effects were important.

Figure 1 is a plot of the measured \( \langle \phi \rangle \) vs. the Troyon parameter \( \beta_c = n_0 L / a B_i \) for all plasmas with indentations \( \geq 5\% \). The closed circles indicate plasmas that were disruption-free through the neutral beam injection phase, while the open circles indicate discharges that disrupted during this time. The first feature to note is that the PBX plasmas achieved \( \beta_c \) values up to 2.7, corresponding to \( q_a \) values down to 2.2, and equivalent circular \( q \) values of down to 1.0. This low value of equivalent circular \( q \) indicates the high current density that the indented plasmas were able to sustain. Also seen in the figure is that the plasmas were limited to \( <\beta_p> \) values \( \leq 2.5-3.0 \beta_c \), the first regime boundary determined empirically by Troyon. The third important feature to note is that all discharges with \( \beta_c \geq 1.9 \) disrupted during the neutral beam injection phase, regardless of the value of \( <\beta_p> \) or how close to the first regime stability boundary the plasma was. The one exception is a pellet fuelled discharge which disrupted a few msec after the beams were turned off. Although the relation between \( \beta_c \) and \( q_a \) is not one-to-one, the value of \( \beta_c \geq 1.9 \) approximately corresponds to \( q_a \leq 3.3 \). The probability that a plasma would disrupt increased dramatically as \( q_a \) decreased. At \( q_a = 5 \), the probability of disruption was 25%, increasing to 80% and 100% at \( q_a = 4 \) and 3 respectively.

Figure 2 shows the sequence of events leading up to a typical disruption. Shown are the displaced toroidal flux as measured by a diamagnetic loop, the neutron flux, the time integrated signal from an outer midplane Mirnov coil, the raw signal from this Mirnov coil, and a near central chord from the vertically viewing soft X-ray (SRX) array. Note the existence of a 14 kHz signal which grew on a time scale of 200 \( \mu \)sec just prior to a sawtooth-like crash. At the crash, both the neutron flux and displaced flux dropped; as much as 30% of the plasma energy, could be lost at the time of the crash. The crash was followed by a delay period of two msec before the plasma underwent a discharge ending disruption. The plasma current remained constant through both the crash and the delay period. As will be seen, the growth time of the crash precursor ranged from tens of \( \mu \)sec to a few msec, and the duration of the delay period was seen to range from zero to ten msec. The structure of the precursor (whose frequencies ranged from a few to 25 kHz) in all the discharges studied was predominantly \( m=3/n=1 \), although some even-\( m \) poloidal component was observed in conjunction with the \( m=3 \) component. At most, the amplitude of this even-\( m \) component was 40% of that of the \( m=3 \) component.

The inversion radius of the SXR emissivity profile across the crash varied dramatically from shot to shot. Some shots exhibited a structure change that indicated an inversion radius close to the center of the discharge (\( \pm 10 \) cm off the mid-plane in a plasma that had approximately a 40 cm vertical extent off the midplane). These discharges were associated with
a relatively small energy loss (<10%) at the time of the crash. As the positions of the inversion radii increased, more energy was lost at the crash. Some plasmas exhibited SXR structures that had no inversion radius at the crash; a drop in emissivity was seen in all channels.

The delay period is interpreted to be a period during which a locked mode grew. The locked mode typically had a strong n=1 component, but also exhibited non-insignificant n=0 and n=2 components. It is believed that field errors were responsible for the mode lock, as the mode locked at the same toroidal location in all discharges studied, both at high-q and low-q. It was determined that a sufficient amount of energy was lost at the crash as to cause inward motion of the plasma center on time scales faster than the position control system could correct, and the plasma ultimately went vertically unstable. However, this sequence of events was initiated by the mode growth and crash, believed to be the manifestation of the n=1 ideal internal kink/5/.

Figure 3 shows the growth time, \(\tau_G\), plotted as a function of a quantity related to the ratio of the edge to central fluctuation amplitude. This quantity, \(A_{pe}\), is in arbitrary units, but is constructed from the normalized fluctuation amplitudes measured near the edge (outer midplane Mirnov coils) and near the center (SRX array). As \(A_{pe}\) increases, the amplitude of the edge fluctuation increases relative to that in the center. Several features are of note; the first is that the growth times ranged from 20 \(\mu\)sec to several msec, spanning the range expected for ideal and resistive time scales. As was discussed in an earlier paper, growth times on the order of 100 \(\mu\)sec can be accounted for within the framework of ideal theory/6/; however, the majority of the discharges shown in the figure exhibited growth times in excess of this value and therefore cannot be readily explained on the basis of ideal theory alone. Figure 3 also shows that the faster growing modes were associated with higher values of edge to central fluctuation amplitudes. A dependence not shown in this graph is that the discharges exhibiting the higher values of \(A_{pe}\) were those at the highest \(<\beta_p>\) and \(<\beta_c>/\beta_c\). This was due primarily to the increase in the relative edge fluctuation amplitude with \(<\beta_p>\) and \(<\beta_c>/\beta_c\). Little change in the relative central fluctuation amplitude was seen with increasing \(<\beta_p>\) and \(<\beta_c>/\beta_c\); however, the central fluctuation amplitude was seen to decrease strongly as \(q_0\) decreased. Nevertheless, it appears that it was the discharges with the highest relative level of edge fluctuations that behaved most ideally in terms of the time scales for mode growth, and these discharges occurred at the highest \(<\beta_p>\).

Figure 4 shows the delay time, \(\tau_D\), plotted as a function of the central or near central rotation velocity, as measured from the Doppler shift of the O\(^{8+}\) line. A complimentary plot (not shown) is a plot of the growth time as a function of rotation velocity. The relation inferred from this latter plot is that the fastest growth rates were associated with the fastest rotating plasmas. The relationship to be inferred from Figure 4 is that the highest rotation velocity plasmas were clearly associated with discharges that had zero delay. The range of rotation velocities seen in the figure was not due to differences in the amount of perpendicular to tangential power in the various cases; all discharges had essentially the same value of this ratio. Therefore, if we take the pre-disruption crash for non-zero delay shots to be evidence for some resistive process, then it was the slowly rotating plasmas that seem to be the most resistive (e.g., slowest growth rates and non-zero delays). This result may indicate some relationship between plasma resistivity and plasma viscosity.
In conclusion, while ideal MHD may be an adequate zero order description of the growing mode that led to disruptions in PBX, the details of the mode indicate that non-ideal effects must play an important role. This was shown first in a previous paper /5/ where, in order to match the phase of the measured signal with that of a signal produced from ideal theory, some approximation to the plasma resistivity was needed. Here, we present additional evidence for these non-ideal effects. Among them are the existence of a sawtooth-like crash leading up to the disruption, the existence of growing locked modes after this crash, and the relatively long growth times of the disruption precursor mode for most of the discharges. In addition, the relation between the growth and/or delay times and the measured plasma rotation speed may indicate a relationship between plasma viscosity (momentum diffusion) and plasma resistivity.

Acknowledgements

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TRANSPORT CODE STUDIES OF M=2 MODE CONTROL BY LOCAL ELECTRON CYCLOTRON HEATING IN TFR

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INTRODUCTION
It has been realized for a long time that low m tearing modes can be stabilized by a suitable tailoring of the current density profile [1, 2]. One possibility that has been suggested to achieve such a profile tailoring is local Electron Cyclotron Heating (ECH) [3, 4]. Experimentally this method was successfully applied in T-10 [5] and, more recently, in TFR [6, 7]. For future applications of ECH to mode control it is important to analyse whether these positive results were indeed a consequence of profile tailoring. The aim of the transport code simulations of the experiments in TFR presented here is to answer this question. As the m=2, n=1 tearing mode plays a dominant role in the occurrence of major disruptions, it is particularly important to understand the experimental results on the stabilization of this mode properly. Therefore, we have concentrated on this mode.

The mode control experiments on TFR were performed mostly in a plasma regime that exhibited strong MHD activity [6, 7]. Discharges in this regime were characterized by a safety factor at the edge q_a ≤ 3.2 [6, 7]. We simulated discharges in this regime with the following basic plasma parameters: major radius R = 0.96 m; minor radius a = 0.18 m; toroidal field B_t = 2.45 T; plasma current I = 130 kA and volume-averaged density \( \langle n_e \rangle = 1.3 \times 10^{19} \text{ m}^{-3} \). For these parameters q_a = 3.0, and the electron cyclotron resonance for 60 GHz (\( \nu_e = 0.14 \text{ m} \)) is slightly outside the q=2 surface (\( r = 2 \approx 0.13 \text{ m} \)). The anomalous electron heat conductivity, \( \chi_e \), was chosen in such a way that an Ohmic target plasma, unstable to the m=2, n=1 tearing mode, was obtained with an energy confinement time in reasonable agreement with the experiment. The anomalous particle diffusivity was assumed to be given by \( D = 1/4 \chi_e \). Further, the neo-classical ion heat-conductivity and the classical 'Spitzer' resistivity were used.

The evolution of the m=2, n=1 tearing mode was evaluated for cylindrical geometry using quasi-linear theory [8]. Disregarding the effects of island heating and of the increased transport across the island, the evolution of the tearing mode with poloidal and toroidal mode numbers m and n, resp., located at the radius \( r_0 \) with \( q(r_0) = m/n \), is well approximated by [8]

\[
\frac{d}{dt} w_{m,n} = 1.66 \frac{\eta(r_0)}{\mu_0} \Delta'_{m,n}(w_{m,n}),
\]

where \( \eta(r_0) \) is the resistivity at \( r_0 \) and \( \Delta'_{m,n}(w_{m,n}) \) is the jump in the logarithmic derivative of the disturbed helical flux function \( \psi_{m,n} \) [9] over the entire width \( w_{m,n} \) of the magnetic island.
SUMMARY OF SIMULATIONS

A moderately unstable target plasma (Δ'(0)=4.7 m⁻¹, w_sat = 2.3 cm) was obtained with

\[ \chi_e = 0.63 \times 10^{19} \frac{T_e^{1/2}}{n_e} e^{1.6 ((r/a)^2 - 1)} \text{ m}^2/\text{s}. \]  

The TABLE summarizes the results of the simulation without ECH.

For this discharge the effect of a 100 ms ECH pulse on the island size was evaluated. To study the dependence of this effect on the location of the heating, the wave frequency was varied between 61.0 GHz (x_r = 12 cm) and 59.4 GHz (x_r = 15 cm). For these resonance positions single-pass absorption is very low, typically 15 to 20%. During the experiments, a mirror which reflected the transmitted waves obliquely and changed their polarization to the X-mode, allowing for enhanced absorption, was mounted opposite to the wave launchers [7]. As the effect on the mode depends critically on the exact power deposition profile, full ray-tracing calculations including absorption along the reflected rays were performed.

Figure 1 shows the evolution of the width of the m=2, n=1 island during the injection of 300 kW of ECH (~180 kW absorbed) for five different positions of the electron cyclotron resonance x_r. The effect on the evolution of the magnetic island is seen to be very sensitive to the position of the heating. The optimum effect, a complete suppression of the magnetic island, is obtained only for x_r = 13.5 cm, i.e. for heating almost exactly on the q=2 surface. As little as 0.5 cm difference in x_r already leads to a significantly decreased effect. The stabilizing effect of ECH, however, is only temporary: 20 to 40 ms after the initial decrease of the island width, the island starts to grow again. Heating on the inside of the q=2 surface (x_r = 12 cm) even leads to an increase in island width well above its initial size. A strong dependence on the injected (or, equivalently, absorbed) power is also observed in the simulations. For the injection of 100 and 200 kW (60 and 120 kW absorbed, resp.) at x_r = 13.5 cm only a small effect on the island is found and complete suppression of the island is never obtained.

DISCUSSION AND CONCLUSION

Comparing the results of the simulations with the experimental data the following differences are found. First, the optimum results in the experiment were obtained by heating approximately 2 cm outside the q=2 surface [7], while the best results in the simulations were obtained for heating almost exactly on the q=2 surface. Secondly, in the simulations the injection of 300 kW is required for complete stabilization while in the experiments suppression of the MHD activity was easily obtained with one gyrotron, i.e. ~170 kW. This discrepancy could be explained by the fact that only around 60% of the power is accounted for in the simulations while it is likely that, in the experiment, after multiple reflections or conversion to Bernstein waves at the upper-hybrid resonance, the remaining 40% was also absorbed close to the cyclotron resonance, not changing significantly the deposition profile, but leading to a higher total absorption. Finally, in the simulations the island suppression is only temporary while in the experiment the MHD activity could be suppressed for a full 300 ms pulse of ECH [7]. This discrepancy is quite serious since there is a simple reason for it, which is directly related to the nature of the profile changes that can be obtained by local heating. This is illustrated by Fig. 2 which gives the temperature and the current density profile before ECH,
10 ms after the start of ECH and at the end of the 100 ms ECH pulse for the injection of 300 kW at \( x_r = 13.5 \) cm. After 10 ms of ECH a clear reduction of the current density gradient around \( r_{q=2} = x_r \) is found such that \( \Delta'(0) \) is well below zero (stable case), while further away from \( x_r \) the current density profile is hardly changed. On a longer timescale the changes of the current density profile are more global and a strong broadening of \( j(r) \) takes place, induced by the broadening of the temperature profile. This broadening of the current density profile is destabilizing to the \( m=2, n=1 \) tearing mode. These results agree with model calculations of the effect of profile tailoring by local heating on the \( m=2, n=1 \) tearing mode [10] which have shown that stabilization, in steady state, requires a narrow and accurately localized power deposition profile, typically \( \Delta r/a = 1\% \). In TFR, however, the width of the deposition profile was around 10\% and the requirements for steady state stabilization, hence, were not met.

We therefore are led to the conclusion that the suppression of the MHD activity with ECH obtained in the experiments on TFR cannot be due to current profile tailoring alone. The most probable cause of the stabilization as obtained in the experiments, in fact, is the effect of ECH on the plasma position: the application of ECH almost always led to a horizontal outward shift of the plasma and a consequent increase in \( q_a \) and, thereby, to stabilization of the mode (see e.g. Ref. [7]).

ACKNOWLEDGEMENTS.
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REFERENCES.

**TABLE**

<table>
<thead>
<tr>
<th>Simulation</th>
<th>Experiment</th>
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<td>( T_e(0) ) (eV)</td>
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<tr>
<td>( T_i(0) ) (eV)</td>
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<tr>
<td>( n_e(0) ) ( (10^{19} \text{ m}^{-3}) )</td>
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<td>( \tau_e ) (ms)</td>
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<td>( V_{loop} ) (V)</td>
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Typical experimental data are given for comparison.
FIG. 1  The effect of ECH on the size of the \( m=2, n=1 \) magnetic island. The evolution of the island during a 100 ms, 300 kW ECH pulse is shown for five different positions of the EC resonance.

FIG. 2  The effect of ECH on the temperature (a) and the current density (b) profiles. The profiles before ECH, after 10 ms ECH, and at the end of the ECH phase are given for the case with 300 kW of ECH at \( x_r = 13.5 \) cm.
SIMULATION OF PLASMA CONTROL IN THE TCV TOKAMAK WITH HIGH FREQUENCY STABILIZATION


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Abstract:

The Tokamak Simulation Code (TSC) is used to study evolution, control, and growth rates of resistive, axisymmetric instabilities in the TCV (Tokamak à Configuration Variable). TCV is currently under construction. Its parameters are: $B=1.43$ T, $R=0.87$ m, $a=0.24$ m, $b/a$ up to 3/1, $I_p$ up to 1200 kAmp. The axisymmetric resistive growth rate is calculated for various equilibria passively stabilized by the TCV vacuum vessel. There exists a range of plasmas which are ideally stable, but which have resistive growth rates greater than those which can be controlled with 100 Hz, 12 pulse thyristor power supplies. We are able to control these high growth rate plasmas by producing a fast, relatively low power radial flux using coils inside the vacuum vessel. Detailed TSC evolutions showing the response of the plasma to these combined control and shaping coils are presented.

Introduction:

It has been demonstrated (1,2) with the Tokamak Simulation Code TSC (3) that a plasma in the TCV Tokamak can be evolved from a near-circular shape to a highly elongated plasma with a relatively broad current profile. Because of the limited bandwidth and voltage of the thyristor shaping power supplies and vessel attenuation, plasmas with peaked current profiles and high elongation having growth rates above 1000 / sec for the axisymmetric resistive vertical instability cannot be controlled. A wide range of highly elongated plasmas with peaked profiles are ideally stable(2), but they have resistive growth rates in the range 1000 - 5000 / sec, which is much larger than the inverse of the vessel time constant of 6.7 msec. We investigate methods of extending the operation range to include plasmas with peaked current profiles.
Fast Stabilization Coils:

A plasma is stabilized against vertical motion by image currents induced in the vacuum vessel. The resistive decay of these currents, and hence the loss of stabilizing flux, causes vertical resistive instability. The role of the active coil system is to provide supplementary flux so as to maintain the plasma at its desired position. In order to achieve this, we consider using additional coils inside the vessel in the region near the outer vessel corners and compare the results with using coils outside the vessel. The coils are placed in the corners (Fig. 4) to reduce interference with the plasma inside the vessel.

Results:

-- The vacuum radial flux produced by a square voltage pulse of a given amplitude and 0.2 msec duration is an order of magnitude larger, for coils inside the vessel compared with coils outside. The flux produced by the inside coils is proportional to the linearly increasing coil current.

-- With coils inside the vessel, it has been demonstrated that it is possible to stabilize both a 2/1 Dee plasma in the upper half and a symmetric 3/1 Dee plasma, with growth rates in the range of 1000 to 5000/sec. The plasmas are stable with no visible distortion. The interesting result is that even though the radial flux pattern of the corner coils is highly distorted by vessel image currents, and symmetric about the midplane, on the time scale of less than 1-2 msec, plasmas both symmetric and nonsymmetric about the midplane can be stabilized. Also, higher order axisymmetric modes are not observed to be dangerous on this time scale.

-- We conclude that it is sufficient to compensate on average the lost radial flux due to wall resistivity with radial flux from local sources in the corners on this time scale. On longer time scales than 1 msec, the outer 16 shaping coils precisely maintain the plasma shape, and limit the power requirements on the inner coils, as shown below.

Examples:

As an example, we choose a 3/1 elongated Dee shaped plasma with $B=1.43T$, $I_p=1200\, \text{kAmp}$, $q_{\text{lim}}/q_{\text{axis}}=3$, plasma-wall gap=0.04 m, growth rate of 3600/sec, initially at rest. The applied coil voltage is proportional to a linear combination of the flux difference between two flux measurement points at the top and the bottom of the plasma and its time derivative, with $P=400000\, \text{volts/weber-radian}$, $D=80\, \text{volts-sec/weber-radian}$. The inside coils have a resistance of 0.005 ohms, and the power supplies for these coils are ideal amplifiers with a nonlinear limit of +50 volts. This system limits the inside coil currents to $+10\, \text{kAmps}$ at low frequency. The lower plasma boundary is commanded to move upwards at constant speed (0.02 m in 2 msec). The command velocity of 10 m/sec represents a very approximate upper limit for three types of occurrences which can induce plasma motion: 1. normal programmed motion during evolution to high elongation, 2. vertical motion due to a sawtooth or minor internal disruption in a non up/down symmetric plasma, 3. fluctuations at high beta.

A simulation is made with only the inside coils active, and the outside shaping coils short-circuited (with shaping currents present), for an interval of 0.002 seconds. The voltage initially jumps to -50 volts due to the $D$ term in the feedback, in order to start the plasma moving vertically. The coil current also goes negative. Once the plasma starts moving upwards at 10 m/sec, the voltage changes sign to restrain the plasma, which is trying to go vertically unstable. The current also changes sign. During the first msec, the plasma moves smoothly upwards, following the command. However, during the period 1.0 to 1.5 msec, the coil voltage reaches its maximum, +50 volts, but the coil current increase is limited by an $L/R$ approach towards a maximum of 10 kAmp. After 1.5 msec, because of this saturation, the control of the plasma is lost, and the plasma moves upwards with a growth rate of 3600/sec.
In order to avoid this saturation, it is necessary for the fast inside coils and the outer shaping coils to work together, with the outer coils providing control and limiting the inside coil currents on a time scale of greater than 1 msec. The method adopted gives the same voltage per turn for the fast upper coil (no. 17, Fig. 4) and the four upper, outer shaping coils (nos. 9-12). The negative of this voltage is applied to the fast lower coil (no. 17) and the four lower, outer shaping coils (nos. 13-16). By this method, since there are four times as many coils with a much lower resistance, the outer coils reduce the current required in the inner coils on the long time scale.

This result is shown in Figs 1-4. In Figs 1 and 3, the currents in coils 18 (inside) and 16 (outside) are shown. The applied voltage, which is the same for both coils, is given in Fig. 2. The evolution of the limiting flux surface of the plasma (expanded) and the coil and vessel geometry are shown in Fig. 4. The outside coils are activated at 0.2 msec, after the plasma has started moving. Apart from the turn-on transient (the outside coils have no voltage limit in these simulations), the voltage and current of the coils inside the vessel and the plasma position evolve for the first msec similarly to the previous case (inside coils only). However, after the first msec, the inside coil current and voltage reach a maximum and then decrease, while the current in the outer coils continues to increase, which is the desired effect. The plasma lower boundary evolves linearly as programmed, demonstrating that control of the plasma can be continuously maintained. The actual thyristor shaping supplies on the outside coils should be able to act similarly to the linear model used here on a 1-2 msec timescale. The behavior of the magnetic axis reflects the non-rigid nature of the plasma, since it initially moves upwards at the same speed as the plasma lower boundary, then moves more slowly, since the plasma is scraping off on the upper limiter.

Based on calculations of radial flux due to coils inside or outside the vessel, control by outside coils of high growth rate plasmas would require high bandwidth and unrealistically high voltages. As an example, repeating the above evolution simulation with only coils outside the vessel, the same PD coefficients and no voltage limit, large plasma oscillations and control voltages an order of magnitude larger than in Fig. 2 are produced. With limiting conditions of 50 volts per turn on coils just outside the vessel at the corners, the voltage saturates for long periods, and control is not achieved.

Conclusions:

It has been demonstrated that high growth rate plasmas can be controlled with a simple additional coil system with high frequency response and reduced power. Inside and outside coils can operate in different bandwidth regions, and therefore reduce the total power requirements.

References:


Fig. 1-4. Controlled evolution of 3/1 elongated Dee shaped plasma, showing lower inside coil current and voltage, lower outside coil current, and evolution of limiting flux surface in TCV geometry.
INTRODUCTION

Sawtooth oscillations of central temperature [1] occur in almost all tokamaks. Experimentally, the sawteeth are observed as a sudden crash in the central temperature, or in the central X-ray emission, followed by a slow rise until the next crash. Kadomtsev [2] first proposed a theoretical model of the sawtooth cycle. Numerical simulations have been carried out by several groups. The first non-linear simulations were performed by Waddell et al. [3] using the reduced MHD equations. Sykes and Wesson [4] followed repeated oscillations by assuming Spitzer resistivity $\eta \propto T^{-3/2}$ and introducing an equation for the temperature evolution including ohmic heating and perpendicular thermal diffusion. However, within this model, they found that the oscillations were decaying in time. Denton et al. [5] found periodic oscillations by introducing a strong thermal conductivity along the field lines. All computations carried out so far have been made with plasma parameters far from those characteristic of current experiments. In this paper we shall present simulations performed with plasma parameters closer to those of the experiments, in particular regarding the value of the Lundquist number $S = l/\eta = \tau_{\text{R}}/\tau_{\text{A}}$ and the ratio of the energy confinement time $\tau_{\text{E}}$ and the resistive time $\tau_{\text{R}}$ (here $\tau_{\text{A}}$ is the Alfvén transit time).

THE MODEL

The simulations presented in this paper are based on the standard, straight cylinder, low-$\beta$, reduced MHD equations [6]. The electron temperature is evolved self-consistently with highly anisotropic thermal diffusivity and ohmic heating. The code is similar to that used by Bondeson [7], but the present simulations incorporate neoclassical effects (neoclassical resistivity and bootstrap current) and the perpendicular thermal conductivity $\kappa_p$ is one inferred from the power balance of the Frascati Tokamak (FT). In the normalized units used in the code, the temperature corresponds to the poloidal beta, which is related to the ratio of the energy confinement time $\tau_{\text{E}}$ and the resistive diffusion time $\tau_{\text{R}}$ by the relation $T = 4/3 \tau_{\text{E}}/\tau_{\text{R}}$ (with $\tau_{\text{R}} = 3/2 nT/\eta^{3/2}$). The code uses finite differences in the radial direction (with a radial grid approximately uniform with 200 points), and Fourier expansion (with both sine and cosine components) in the azimuthal and toroidal directions. We restrict ourselves to single helicity perturbations with $m/n = 1$, retaining mode numbers up to $m/n = 4/4$. The $q$-value at the edge has been kept fixed in time.

SIMULATION RESULTS AND CONCLUSIONS

In the following we shall refer to a typical scenario of the ohmically heated FT machine with $B_T = 6$ T, $a = 0.2$ m, $R = 0.83$ m, $q_a = 2.6$, central density $n(0) \sim 1.8 \times 10^{20}$

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m⁻³ and central temperature T(0) ~ 500 ± 1000 eV. With these parameters one has $t_A \approx 0.085$ μs, $S \approx 10^7$ and normalized $T \approx 0.02$. We have assumed $\kappa_l = \kappa_l(0) \left[1-(r/r_0)^2\right]^{-2}$, which well represents the results of the power balance on FT. Moreover the parallel thermal conductivity $\kappa_l$ and the viscosity $\nu$ have been taken as constant over the whole cross-section. We present two simulations: one with rather small $S$ number similar to the

![Figure 1](image1)

**Fig. 1:** a) Normalized central temperature $T(0)$, b) central safety factor $q(0)$ and c) energies of the various Fourier components (the $+$ is the $m/n = 1/1$, the $\Delta$ is the $m/n = 2/2$, the $+$ is the $m/n = 3/3$ and the $\Delta$ is the $m/n = 4/4$) vs time (in Alfvén time units) for the first simulation ($S = 10^4$, $\kappa_l(0) = 7 \times 10^{-6}$, $\kappa_l = 40$, $\nu = 10^{-4}$, $T \sim 1.4$)

![Figure 2](image2)

**Fig. 2:** a) Two-dimensional contour plots of the helical flux function $\psi$, (only in the central region of the plasma), b) electron temperature and c) stream function $\phi$ at different times for the first simulation ($S = 10^4$, $\kappa_l(0) = 7 \times 10^{-6}$, $\kappa_l = 40$, $\nu = 10^{-4}$, $T \sim 1.4$)
simulation of Denton et al. [5], and a second one with higher $S$, to emphasize the role of the plasma parameters. For the first simulation we choose $S = 10^4$, $\kappa_1(0) = 7 \times 10^{-6}$, $\kappa_4 = 40$, $v = 10^{-4}$ (and hence $T \sim 1.4$) and neoclassical effects neglected. The time dependence of $T(0)$ and $q(0)$ is presented in Fig. 1a,b. Regular sawteeth are obtained with a large variation of $q$ in the centre $\Delta q \sim 0.25$. The energies of the various Fourier components are presented in fig. 1c, and show a regular repetitive sequence of growth followed by

![Fig. 3: Same as Fig. 1 for the second simulation ($S = 5 \times 10^5$, $\kappa_1(0) = 2 \times 10^{-5}$, $\kappa_4 = 26.67$, $v = 0.5 \times 10^{-6}$, $T \sim 0.02$)](image)

![Fig. 4: Same as Fig. 2 for the second simulation ($S = 5 \times 10^5$, $\kappa_1(0) = 2 \times 10^{-5}$, $\kappa_4 = 26.67$, $v = 0.5 \times 10^{-6}$, $T \sim 0.02$)](image)
a period of decaying MHD activity before the next crash. Fig. 2a shows two-dimensional contour plots of the helical flux function $\psi = \psi - (a^2 - r^2) B_r / 2 q R$ with $q = 1$ ($\psi$ is the flux function and the magnetic field $B$ is given by $B = V\psi \times \hat{z} + B_T \hat{z}$) during the time of one crash, and Fig 2b,c, show the electron temperature and the stream function $\phi$ respectively (the velocity $v$ of the fluid is given by $v = V\phi \times \hat{z}$). We emphasize that in this simulation $\kappa_0 / \kappa \sim 1$, almost two order of magnitude smaller than the value obtained in FT, allowing large changes in the current profile during one sawtooth cycle. For the second simulation we choose $S = 5 \times 10^5$, $k_1(0) = 2 \times 10^{-5}$, $k_2 = 26.67$, $v = 0.5 \times 10^{-6}$ (and hence $T \sim 0.02$). In Fig. 3a,b,c we show the evolution of $T(0)$, $q(0)$ and of the energies of the Fourier components. Very little variation of the central $q$ is observed, corresponding to small-scale, local variations in the current density, and MHD activity persists during the whole cycle. The small-scale variations can be seen more clearly in the two-dimensional plots of $\psi$, $T$ and $\phi$ (Fig. 4a,b,c). Sawteeth still occur, but show very little resemblance with the experimental behaviour.

In conclusions we have shown that rather distinct sawteeth are produced for $S < 10^5$, but simulations with $S > 10^5$ tends to produce profiles that are all the time very close to marginal stability, $q = 1$ over the entire central region. Thus this study suggests that toroidal and finite pressure effects must be accounted for to properly describe the sawteeth.

FOOTNOTE AND REFERENCES

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SAWTOOTH STABILIZATION BY ENERGETIC TRAPPED PARTICLES

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Experiments on JET using high power ICRF on-axis heating of a minority ion species\(^1\) show efficient electron heating with the occurrence of long sawtooth-free periods of up to 1.6 sec. In this work we show that a high energy trapped ion population such as that produced by the heating in JET can significantly decrease the growth rate of the resistive internal kink mode, leading to an increase of the sawtooth period.

Recently, the effect of an energetic trapped particle population on magnetohydrodynamic modes in a tokamak has been explored with the use of a variational formalism.\(^2-6\) The usual branch of the ideal internal kink, unstable for plasma \(\beta\) (the ratio of plasma pressure to magnetic pressure) greater than a threshold value, is stabilized by a trapped particle population as long as the average toroidal precession rate of the particles is greater than the mode growth rate.

On the other hand, it was found that for \(\beta\) near the internal kink threshold value the trapped particles resonantly destabilize a second branch of the internal kink mode, with a real frequency given by the average precession frequency of the particle distribution, provided that the trapped particle beta, \(\beta_t\), exceeded a threshold value. This branch is responsible for the fishbone oscillation. The dispersion relation describing both branches of this mode was generalized to include resistive effects in Ref. 6. The dispersion relation takes the form

\[
\delta W_c + \delta W_k + \frac{8S^{-1/3} \Lambda^{-9/4} \Gamma(\Lambda^{3/2}+5/4)}{\Gamma((\Lambda^{3/2}-1)/4)} \left[ \frac{\omega_{\Lambda}}{\omega_R} \right]^{1/2} \Omega(\Omega+i\omega_{\Lambda}) \Gamma(\Omega+i\omega_{\Lambda}) = 0, \tag{1}
\]

where \(\Omega = -i\omega/\omega_R\), \(\Lambda = (\Omega+i\omega_{e}/\omega_R)(\Omega+i\omega_{\Lambda}/\omega_R)^{1/3}\), \(\omega_R = S^{-1/3} \omega_A\) is the resistive frequency, \(S\) is the magnetic Reynolds number, \(\omega_A\) is the shear Alfvén frequency

\[
\omega_A = \frac{V_A}{\sqrt{3} Re q'}
\tag{2}
\]

with \(V_A\) the Alfvén velocity, \(R\) and \(r\) the major and minor radii, respectively, and \(q' = dq/dr\) with \(q\) the safety factor. The \(\omega_{\Lambda}\) terms are diamagnetic frequencies with \(\omega_{\Lambda} = -(c/neBr)(dp/\sqrt{dr})\), \(\omega_{e} = (c/neBr)(dp_{e}/dr)\), and \(\omega_{\Lambda} = \omega_{\Lambda} + 0.71 (c/eBr)(dT_e/dr)\). The term in Eq. (1) involving the \(\Gamma\) functions arises from the inertial layer, so all
expressions are evaluated at the \( q = 1 \) surface. The inclusion of the

diamagnetic terms was carried out by Bussac et al.\(^8\) and Ara et al.\(^9\) generalizing the work of Coppi et al.\(^{10}\) The expression \( \delta W_k \) is the

minimized ideal variational energy for the internal kink, first calculated
by Bussac et al.\(^7\) and \( \delta W_k \) is the kinetic contribution coming from the

trapped particle distribution \( F \),

\[
\delta W_k = \frac{2^{3/2}}{B^2} \m_\perp m^2 \int d(\alpha B) \int \frac{dEE}{k_b^2} \frac{\omega(\alpha/E + \omega_d^* / \omega_d) F}{\omega_d^*} 
\]

with \( \gamma = (2 \int y rdr / r_S^2) \), \( r_S \) the \( q = 1 \) radius, \( \alpha = v_T^2 / v^2 \), \( \omega_d^* \) a
differential operator associated with the diamagnetic drift frequency, and

\( K_2 \) and \( K_b \) are elliptic functions arising from bounce averaging. Details
of the derivation of this expression are found in Refs. 3 and 5.

We have examined the solutions to Eq. (1) using a numerical code
developed for the investigation of the fishbone mode.\(^7\) Numerical values
of temperatures, densities, diamagnetic frequencies, etc., were chosen to
approximate the experiments done on JET. The kinetic contribution \( \delta W_k \) is
generated by a Monte Carlo procedure. For minority species ion cyclotron
heating experiments the hot trapped particle distribution is well
approximated by

\[
F(E, \mu, r) = n(r) e^{-E/T} \delta(u/E-\alpha) 
\]

The value of \( n \) is typically of the order of a few percent of the average
plasma density, and the temperature \( T \) ranges from 70 to 150 keV depending
on the ICRF power.

Results of a Monte Carlo simulation are shown in Fig. 1 for hydrogen
minority species in JET. Shown is the growth rate as a function of

trapped particle density for an approximate JET equilibrium with \( R = 296 \)
cm, for two different values of the magnetic Reynolds number, \( S = 3 \times 10^6 \)
and \( 10^7 \). The trapped particle density ranges from zero to \( 10^{12} \) cm\(^{-3}\). The particle distributions were of the form given by Eq. (13), and had


temperatures of 50 to 150 keV. The toroidal field was \( B = 24 \) kG and the

average trapped particle precession rates were \( \langle \omega_d \rangle = 2 \times 10^6 \) sec and \( 6 \times 10^6 \) sec, respectively. The shear Alfvén frequency was \( \omega_A = 2 \times 10^6 \) sec, and the
diamagnetic frequencies were \( \omega_d^* = 3 \times 10^4 \) sec and \( \omega_d^* = 2 \times 10^4 \) sec. From these curves it is clear that a significant lengthening of
the sawtooth period is possible, assuming that the sawtooth period scales
as the inverse of the linear growth rate. We note that the trapped
particles also stabilize the ideal internal kink\(^5\) and thus the

qualitative features of this result are unchanged if a short time scale
ideal model of the sawtooth is used.
Fig. 1: Growth rate as a function of trapped proton density, for two values of the magnetic Reynolds number $S$, and different energies. Equilibrium parameters and Alfvén and diamagnetic frequencies were chosen to approximate JET.

We find that to stabilize the tearing mode it is necessary to produce a trapped particle density which is large enough to destabilize the fishbone. The occurrence of the fishbone is however also contingent upon proximity to the internal kink threshold. Depending on plasma $\beta$ and other equilibrium parameters we must in general expect either sawtooth stabilization or fishbone oscillations, with fishbone oscillations dominating for high $\beta$ operation. The two could exist together only if the fishbone failed to eject the trapped particles.

During neutral beam heating of JET at lower toroidal field values (21 kG) and higher beta, magnetic signals associated with the fishbone mode were observed. The threshold value of $\beta_p$ for the fishbone was calculated to be $2 \times 10^{-3}$. Because of the efficient ejection of trapped ions by the fishbone the trapped particle population would be effectively limited to approximately this value, which corresponds, for the neutral beam energy of 70 keV, to a trapped particle density of $3 \times 10^{11}$/cm$^3$, too low a value to produce stabilization of the sawtooth. In general since the stabilization condition implies that the fishbone threshold be exceeded, it appears, perhaps unfortunately, that at high beta operation the occurrence of the fishbone will not allow sawtooth stabilization with trapped particles. To our knowledge, fishbone oscillations have not yet been produced during ICRF heating, so the nature of the transition from sawtooth stabilization to fishbone is still unexplored experimentally.
The signature for this transition would be the occurrence of ion bursts with an associated drop in neutron production during the sawtooth period.

In conclusion, we find that the presence of a high energy trapped ion population introduces a stabilization of the sawtooth in a tokamak. Numerical calculations of the effect are in reasonable agreement with experiments on JET, giving almost an order of magnitude increase in the period. Although the stabilization of the sawtooth mode in low beta discharges is encouraging, at higher $\beta$ the same trapped particle population should destabilize the fishbone branch, as has been observed to happen with neutral beam injection. The fishbone should be then expected to limit the trapped particle population to a value too low to provide sawtooth stabilization.

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References

INTERNAL KINK MODES IN THE ION-KINETIC REGIME

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Abstract. The m=1 kink mode is investigated in the high temperature regime where the width of the singular layer is determined by the mean ion gyroradius. A dispersion relation is presented that contains the full ion dynamics. The growth rates are larger than the corresponding ones obtained from fluid theory. Diamagnetic stabilisation is weaker than in the fluid case. If \( n_i \) is sufficiently large an unstable mode is found for all values of \( \omega \) even at zero resistivity.

Introduction. We refer to a low-\( \beta \) axisymmetric toroidal plasma configuration. The radial profile of \( q = rB_\phi / RB_\theta \) is assumed to be monotonic with \( q \) on axis below one and with finite shear at the \( q=1 \) surface at \( r = r_0 \). In this geometry, the structure of modes with toroidal mode number \( n=1 \) and dominant \( m=1 \) component is characterised by a narrow boundary layer near the \( q=1 \) surface where non-ideal effects become important \([1,2]\). Outside this layer the mode amplitude matches onto the relevant ideal-MHD solution. In this paper we consider plasmas with temperatures such that the mean ion gyroradius exceeds the width of the singular layer as obtained from fluid theory. We adopt a Vlasov description for kinetic ions which includes ion gyroradius effects to all orders. The electron response, which is isothermal for the modes under consideration, is obtained from resistive fluid equations. We focus our attention on kink regimes where the ideal-MHD driving force \( \lambda_H (-\delta W, \text{ the ideal-MHD energy functional}) \) is positive or small which corresponds to ideally unstable or marginal stable modes. These regimes are relevant to the explanation of the disruption phase of the sawtooth oscillation in tokamaks. The main toroidal effects are included in \( \lambda_H [3] \).

Basic Equations. We adopt the formalism of Ref.[4] where the singular layer equations for collisional electrons and kinetic ions are derived in Fourier space, with \( z \) the variable conjugate to \( z = (r-r_0)/r_0 \). This formalism allows us to obtain the kinetic ion response for arbitrary \( \rho_i \) in algebraic form, as opposed to the integral form in coordinate space. The normal mode equation for the perturbed parallel current density \( J_\parallel \) is

\[
\frac{d}{dz} \left[ \frac{1}{\lambda D} \right] \frac{d}{dz} \left[ \frac{1}{\tau \rho_i^*} \right] J_\parallel - \left( \frac{\lambda}{\tau \rho_i^*} \right) \left( \frac{\lambda - i \rho_i^*}{z^2} \right) J_\parallel = 0
\]  

(1)
where \( \lambda_* = -\omega \tau_A \), \( \lambda = -i \omega \tau_A \), \( \omega_* \) is the electron diamagnetic frequency, \( \tau_A \) is the Alfvén time with respect to the shear field, \( \rho_i^2 = T_i / (m_i \Omega_i \tau_A^2) \), \( \tau = T_e / T_i \), \( \lambda_i = -\omega \tau_A (1 + \eta_i) / \tau \), \( \eta_i = \text{d} \ln T_i / \text{d} \ln n \), and \( \varepsilon_1 \) is the inverse magnetic Reynolds number. The function

\[
D = \tau (1 - \Gamma_0) (1 + \frac{i \lambda_2}{\lambda}) + i \eta_1 \frac{\lambda}{\lambda} \rho_i^2 z^2 (\Gamma_0 - \Gamma_1) \tag{2}
\]

where \( \Gamma_0, \Gamma_1 = \Gamma_0, 1 \) (\( \rho_i^2 z^2 \) \( \exp (\rho_i^2 z^2) \)), \( \Gamma_0, 1 \) being modified Bessel function, represents the finite gyroradius terms in the perturbed ion density \( n_i / n = (i \lambda_2 / \lambda - D) \varepsilon/\tau \).

Kink modes in the large ion gyroradius regime are found for

\[
\frac{\lambda}{\tau \rho_i^2}, \quad \frac{\lambda}{\tau \rho_i^2} (\lambda - i \lambda_2) \ll \frac{1 + \tau}{\tau + i \lambda_2 / \lambda}, \tag{3}
\]

i.e. when the gyroradius is larger than the resistive and inertial scale lengths.

In Refs.4 and 5 we have used the model \( D^{-1} = (\tau + i \lambda_2 / \lambda)^{-1} + (1 + i \lambda_2 / \lambda)^{-1} (\tau \rho_i^2 z^2)^{-1} \), which is an interpolation formula between the fluid \( (\rho_i^2 z^2)^{-1} \) and the large gyroradius \( (\rho_i^2 z^2) \ll 1 \) response. When \( \rho_i^2 z^2 \ll 1 \) the ion density perturbation is governed by the E×B and polarization drifts, and by the density and temperature gradients. In the opposite limit, the ions can move freely across the magnetic field on the scale of the wavelength, and the density perturbation is a Boltzmann response. The model is consistent with these two limiting cases. It can be shown that except for some spurious roots, the model gives not only qualitatively but also quantitatively correct results for small values of \( \eta_i \). For larger values of \( \eta_i \) the equilibrium gradients enter differently in the model and in the full response for \( \rho_i z = 0(1) \) and the model breaks down.

Inequalities (3) allow us to solve (2) with the full ion dynamics [6]. The resulting dispersion relation is

\[
\lambda (\lambda - i \lambda_2) (1 + \tau)^{1/2} \int_0^\infty \frac{dt}{t^2} \frac{1}{1 + g/D} = (1 + \tau)^{1/2} \lambda_H \phi \frac{1}{\lambda - i \lambda_2} + \frac{(1 + \tau) \varepsilon_1}{\lambda - i \lambda_2} \rho_i^2 z^2 \left( \frac{\lambda}{\lambda - i \lambda_2} \right)^{1/2} \tag{4}
\]

The corresponding solutions have regular amplitudes if \( \text{Re}(\lambda + i \lambda_2 / \tau)^{1/2} > 0 \).

For positive values of \( \lambda_H \) in the limit where resistivity is negligible \( (\varepsilon_1 = 0) \) we obtain two roots for small values of \( \eta_i \). The real part of \( \lambda \) is plotted in Fig.1 for the unstable root as a function of \( \lambda_2 \) for \( \eta_i = 1 \). The variables on both axis are normalized on the growth rate \( \lambda_0 = [(2^{1/2} / \pi) \lambda_H \rho_i^2]^{1/2} \) at \( \lambda_2 = 0 \). We will take \( \tau = 1 \). The unstable root acquires an oscillatory
character for increasing values of $\lambda_X$. Complete stabilization occurs for $\lambda_X^2$, i.e. for $\lambda_X^2(\lambda_H/\rho_i)^2$, which can also be written as $82\lambda_H/\rho_i$ where $\lambda = (L_s/L_\eta)/\eta$ and $L_s$ and $L_\eta$ the shear and density scale lengths respectively. This stabilization occurs for larger values of $\lambda_X$ than in fluid theory.

The same plot is given for $\eta_i=2$ in Fig.2. For this value of $\eta_i$ three roots are found [6]. The mode that corresponds to the one of Fig.1 becomes stabilized at a somewhat lower value of $\lambda_X$. The third root emerges at the origin with an extremely low value of $\lambda/\lambda_X$ and exists when $\eta_i$ is larger than $\eta_{ic}=1.67$. A window in $\lambda_X$ exists where all three modes are purely oscillatory. For increasing values of $\lambda_X$, an almost purely growing mode emerges. These features cannot be obtained from the model for D. A small resistivity removes the stable window and one of the roots become unstable for all values of $\lambda_X$ (thin line).

![Fig. 1](image1.png) ![Fig. 2](image2.png)

In Fig.3 the unstable root is plotted for $\eta_i=5$. Although the growth rate decreases at first faster than $\lambda_X$ than at lower values of $\eta_i$, one of the modes remains unstable for all values of $\lambda_X$ even without resistivity. In Fig.4 the real part of the frequency for the unstable root is plotted at marginal ideal stability ($\lambda_H=0$) for finite resistivity and for different values of $\eta_i$. The frequencies are normalized at the growth rate

$$\lambda_0 = \rho_i [2\epsilon_i/\pi^2 \rho_i]^1/7$$

at $\lambda_X = 0$. For increasing values of $\lambda_X$ the growth rate decreases but the mode remains unstable. Increasing values of $\eta_i$ tend to reduce the growth rate for small values of $\lambda_X$ and to increase it for large values of $\lambda_X$ where, however, the growth rate remains small.
REFERENCES

THE $m=1$ INTERNAL KINK MODE IN A TOROIDAL PLASMA WITH A FLAT $q$-PROFILE NEAR $q=1$ \\

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ABSTRACT. The energy principle is used to study the ideal MHD stability of the $m=1$ internal kink mode in a toroidal plasma. The equilibrium configuration is allowed to have a broad region with $q$ close to unity. The minimization of the energy functional yields an implicit equation for the growth rate that can be solved by simple numerical means. The examples that we present retain the essential features of experimentally expected $q$-profiles. We find that the growth rate depends very sensitively on $q$ when $q$ is very close to unity. The highest growth rates are of the order $\varepsilon/\tau_A$, where $\tau_A$ is the poloidal Alfvén time.

INTRODUCTION. One of the instabilities that is crucial to the confinement of tokamak plasmas is the internal disruption or sawtooth instability. In general the sawtooth collapse is very fast, a typical timescale being $10^{-3}$s. The disruption is characterized by an $m=n=1$ displacement of the central plasma column. On the basis of soft X-ray tomography on JET it is suggested that a cold bubble enters the central plasma and that the hotter plasma spreads around it [1]. Such a flow would be consistent with a flat $q$-profile [2,3]. Experimental results [4,5], however, also show that in a number of circumstances $q$ is well below unity in the centre. In our treatment $q$ is allowed to be flat and close to unity in a layer I of arbitrary width, where the constant $\xi$ approximation is not valid (see Fig. 1). The case that this layer reduces to a single $q=1$ surface, and also the case that $q=1$ in the entire central plasma, are included.

ENERGY PRINCIPLE. We consider the stability of a low beta MHD equilibrium with circular flux surfaces. The minimization of $E(\xi)$, which is the sum of the kinetic and the potential energy associated with a perturbation $\xi$, yields the growth rate and the spatial structure of the mode. In the minimization we use an expansion to powers of the inverse aspect ratio. Minimization with respect to the component of $\xi$ parallel to B shows that plasma compression is small. The remaining part of $\xi$ consists of a poloidal component $\chi$ and a radial component $\xi$. A first minimization yields $(\partial^2/\partial \zeta^2)(\xi r)+\partial \chi/\partial \xi)$. \\

Fig. 1. The considered $q$-profiles and corresponding eigenfunction.
=O(\varepsilon^2), \text{ i.e., the compression of field lines is small. We restrict ourselves to the } m=1 \text{ mode. Since the toroidal coupling is } O(\varepsilon), \text{ the } m=0, 2 \text{ harmonics are } O(\varepsilon). \text{ Other poloidal harmonics are of higher order and thus negligible.}

The leading order Euler-Lagrange equation for the } m=1 \text{ harmonic } \xi_1(r) \text{ is }

\begin{align}
\left[\frac{1}{\varrho} - \frac{1}{4}\right]^2 + M r^2 \right] \xi_1(r) &= O(\varepsilon^2),
\end{align}

where } \Gamma \text{ is the growth rate normalized to the poloidal Alfvén time } \tau_A = \frac{B}{R} \rho^{-\frac{1}{2}}, \text{ and the inertial factor is }

M(r) = 1 + \left(1 + \frac{1}{2} \Gamma^2/\beta Y\right)^{-1} + \left((\frac{2}{q} - 1)^2 + \frac{1}{2} r^2/\gamma Y\right)^{-1}.

The first term in Eq. (1) arises from the bending of magnetic field lines, the second term is due to inertia and plasma compression. The above steps minimize the energy to } O(\varepsilon^2). \text{ After minimization with respect to } \chi_m=0, E \text{ is a function of } \xi_m (m=1, 2) \text{ only.}

**METHOD AND MAIN RESULTS.** We consider a q-profile that is close to unity in region I in the plasma, while in parts II [q-1] is not small (see Fig. 1). The minimization of } E \text{ with respect to } \xi_1 \text{ and } \xi_2 \text{ is performed in regions I and II separately. In regions II, Eq. (1) implies that } \xi_1(r) \sim O(\varepsilon^2) \text{ so that } \xi_1 \text{ is constant to leading order. In region I the constant-} \xi_1 \text{ approximation is not valid since the coefficient in Eq. (1) is } O(\varepsilon^2). \text{ In this region the energy functional is minimized by means of an additional expansion in the small parameter } [q-1]. \text{ We match the solutions in regions I and II at their boundaries } r=r_1. \text{ Thus we obtain an implicit equation for the growth rate, }

\begin{align}
\left[&S_2 - S_1^2 / S_o - A(r_+) + A(r_-) \right] \\
&\left[ -1/S_o + \frac{9}{16} A(r_-) - \frac{1}{4} r^2 \xi(r_-) \right] \\
&= \left[ -S_1/S_o + \frac{3}{4} A(r_-) \right]^2.
\end{align}

Here, the quantities } S_i (i=0, 1, 2) \text{ are integrals over the layer I, }

\begin{align}
S_i(r) &= \frac{1}{R^2} \left[ \frac{1}{q-1} r^4 \right]^{1/2} + M r^2 \right] \beta_p(r) \left[ \frac{1}{q-1} r^4 \right]^{1/2} + M r^2 \right], \quad S_1 \equiv S_1(r_+),
\end{align}

where the poloidal beta and the internal inductance are given by

\begin{align}
\beta_p(r) &= -2 \frac{R^2}{B_o} r^{-n} \left[ \frac{1}{q-1} r^4 \right]^{1/2} + M r^2 \right] \beta_p(r) \left[ \frac{1}{q-1} r^4 \right]^{1/2} + M r^2 \right], \quad \lambda(r) \equiv 2 r^{-n} \left[ \frac{1}{q-1} r^4 \right]^{1/2} + M r^2 \right] \lambda(r) \left[ \frac{1}{q-1} r^4 \right]^{1/2} + M r^2 \right] \lambda(r) \left[ \frac{1}{q-1} r^4 \right]^{1/2} + M r^2 \right].
\end{align}

The quantities } A(r_+) \text{ are the boundary values of the function } A(r) \text{ which is defined as }

\begin{align}
A(r) &\equiv r^2 \frac{1-a}{3+a} \text{; } a(r) \equiv \frac{2-q}{q-1} r \xi^2 / \xi,
\end{align}

where } \xi(r) \text{ is the solution of the leading order } m=2 \text{ Euler equation }

\begin{align}
\frac{d}{dr} \left[ r^3 \left( \frac{1}{q-1} \right)^2 \xi \right] &= 3 \left( \frac{1}{q-1} \right)^2 r \xi,
\end{align}

in the respective regions II. Equation (3) can be written as an equation for } A(r) \text{ with the property } (dA/dr) \sim O(q-1), \text{ so that } A(r_+) \text{ are invariant to small shifts of } r_- \text{ and } r_+. \text{ The approximations applied in regions I and II are not justified in thin layers between I and II. When these layers are taken into}
account, we obtain O(ε) corrections to the integrals $S_i$ that make Eq. (2) insensitive to small redefinitions of the boundary positions $r = r_\pm$. In addition, we have calculated the energy in the interval $[r_-, r_+]$ to one higher order in $(q-1)$. This $O(ε)$ contribution becomes significant if $q=1$ in the entire central plasma and $β_p$ is small, and should then be included in Eq. (2). The radial eigenfunction $ξ_1(r)$ decreases continuously to zero in the interval $[r_-, r_+]$. It is obtained to leading order and is given by

$$ξ_1(r) = \begin{cases} ξ_1(0), & r < r_-; \\ 0, & r > r_+; \\ ξ_1(0) \left[1 - \frac{S_0(r)}{S_0} \right] - (S_1(r) - \frac{1}{S_0} \frac{1}{S_0} A(r_-) - \frac{1}{4} r_-^2 A(r_-)) \left[\frac{1}{S_0} + \frac{3}{4} A(r_-) \right]^{-1}, & r_- < r < r_+. \end{cases}$$

**EXAMPLES.** We will show how the growth rate depends on the pressure gradient and on basic characteristics of the $q$-profile by solving Eq. (2) for simple model profiles. We consider cases where $q$ is approximately constant in the central plasma, with $q(o)$ either above or below unity. This region is surrounded by a layer $[r_-, r_+]$ where $δq ≡ q-1$ is small and taken to be constant, file for Fig. 3. see Fig. 2.

We adopt a parabolic pressure profile in this layer so that $β_p(r)$ is constant. For these profiles, the integrals $S_i$ are easily evaluated, and from the $q$-profile in the centre and in the outer region we deduce that $A(r_-) = -r_-^2 q_0 = 1/2 r_-^2 (1-q(o)^2)$ and $A(r_+) = 1/4 r_+^4$ are good approximations. Equation (2) expresses the quantity $f^2 \equiv δq^2 + Mt^2$ in terms of the parameters $q(o)$, $r_+$, and $β_p$. If $f$ is large, $O(ε/τ_A)$, the growth rate strongly depends on $δq$ near $δq = f$.

In the first example we take $r_- = 1/2 r_+$. Figure 3 shows lines of constant $f(q(o), β_p)$. The growth rate increases with $β_p$ and $q(o)$. $f^2$ is always positive and vanishes with $β_p$ when $q(o)<1$. In the case of vanishing $δq$, we have $f^2 = M^2 t^2$ so that the mode is always unstable. In our second example we take $r_- = 0$, and $q-1=δq$ is a small constant in the entire central plasma. Then Eq. (2) reduces to $S_2 = A(r_+)$ [6], or approximately,

$$f^2 = \frac{1}{4} \frac{r_+^2}{R^2} β_p^2 .$$

First we note that Eq. (4) is accurate only for sufficiently high values of
$\beta_p$, for $\beta_p > 0.05$ say, since otherwise $O(\delta q)$-terms become significant. Large growth rates $-\varepsilon/\tau_A$ occur for $|\delta q| < 1/2 \beta_p r_+/R$. The mode is stable for $\delta q > 1/2 \beta_p r_+/R$. For negative $\delta q < 1/2 \beta_p r_+/R$, on the other hand, a $q$-1 surface is present at $r = r_+$. In that case the limit $\lim r_+ - r_+$ should be considered. Moderate growth rates $-\varepsilon^2/\tau_A$ are found since Eq. (2) becomes

$$\frac{1}{S_o} = \frac{9}{16} A(r_-) - \frac{1}{4} r_+^2 \delta (-) + \left[ \frac{3}{4} A(r_-) - r_+^2 (\beta_p + \frac{1}{2} r_+) \right] [A(r_+) - A(r_-)]$$

which is equivalent to the expression for the growth rate that has been obtained in the approximation of a singular layer [7,8,9]. In Fig. 4, lines of constant growth rate are plotted in the $\beta_p$-$q(0)$ plane for the indicated $s$-profiles. For $q(0) = 1$, Eq. (4) has been used while for $q(0)$ below unity Eq. (5) has been applied. In the figure these limit cases are connected by sketched dashed lines. In conclusion, the ideal $m=1$ mode is very sensitive to the part of the $q$-profile that is close to unity. Non-monotonic $q$-profiles with a minimum close to $q=1$ are unstable. On the other hand, if $q(0)$ is below unity, the mode is stable for a sufficiently low pressure gradient. However, an evolving $q$-profile has to pass an unstable regime to get below unity.

**ACKNOWLEDGEMENT.** This work was performed under the Euratom-FOM association agreement with financial support from NWO and Euratom. Fig. 4. Curves of constant growth rate $\Gamma$ for the displayed $q$-profiles. The full lines are given by Eq. (4) for $q(0)=1$, and by Eq. (5) for $q(0)$ below 1.

**REFERENCES**

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EIGENVALUES OF RELAXED TOROIDAL PLASMAS

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According to the theory of relaxed states /1,2/, a turbulent plasma relaxes to a free-force configuration of minimum energy subject to the constraint of constant magnetic helicity \( K = \int A \cdot B \, d\tau \). These configurations are described by:

\[
\text{curl } \vec{B} = \gamma \vec{B} \quad \gamma = \text{constant.} \tag{1}
\]

On a conductive boundary the containment magnetic field has to satisfy

\[
\vec{B} \cdot \hat{n} = 0 \tag{2}
\]

In order to select the correct minimum energy solution of (1),(2) one has to consider the eigenvalues of this system /3/. There exist eigenvalues at which the toroidal magnetic flux vanishes ("zero-flux eigenvalues"). It has been shown /2,3/ that in these relaxed states \( \gamma \) cannot exceed the smallest eigenvalue \( \gamma_{\text{min}} \) and that for toroidal discharges there is a maximum toroidal current which is fixed. For these reasons the knowledge of \( \gamma \) in (1),(2) is of importance.

Since the analytical solution of (1),(2) in toroidal coordinates seems to be impossible, the system (1),(2) has been solved /3/ in the infinite aspect-ratio limit \( R/a \to \infty \) (cylindrical approximation) and also for a torus with rectangular cross section. In the latter case (1) is separable in cartesian coordinates. Toroidal containers with other cross sections than a rectangle have, however, not been calculated.

In this paper we apply on (1),(2) a method to solve equations of this type for a torus of arbitrary aspect-ratio and arbitrary cross section /4/. We use cylindrical coordinates \( r, \phi, z \) and assume axisymmetry, i.e., \( \partial / \partial \phi = 0 \). The meridional cross section in the \( r,z \) plane of the toroidal metallic vessel wall shall be described by an arbitrary curve \( z = z^*(r) \) along which (2) has to be satisfied. Along such an arbitrary curve (2) takes the form:

\[
\frac{dr}{B_r(r,z)} = \frac{dz}{B_z(r,z)} \quad , \quad z = z^*(r) \tag{3}
\]

since the field lines are tangential to this curve \( z = z^*(r) \). This may be written in the form

\[
- B_r(r,z^*) \frac{dz^*}{dr} + B_z(r,z^*) = 0 \tag{4}
\]

(+) On leave from Atomic Energy Authority, Plasma Phys. Unit, Cairo, Egypt.
Now for axisymmetry \( \partial \phi = 0 \) equation (1) takes the form

\[
D_r = -\frac{1}{y} \frac{\partial B_\phi}{\partial z}, \quad B_z = \frac{1}{y} \frac{1}{r} \frac{\partial}{\partial r} (r B_\phi)
\]  

(5)

and

\[
\frac{\partial B_r}{\partial z} - \frac{\partial B_z}{\partial r} = \gamma B_\phi
\]  

(6)

Inserting (5) into (4) we obtain

\[
\frac{\partial (r B_\phi)}{\partial z} + \frac{\partial (r B_\phi)}{\partial r} dr = d(r B_\phi) = 0 \quad \text{or} \quad r B_\phi = \text{const.}
\]  

(7)

Here and in the following we omit the asterisk.

The solution of (5), (6) is given by

\[
B_r = \sum_{\ell=1}^{N} k_\ell \sin(k_\ell z) \left\{ a_\ell J_1(\sqrt{y^2-k_\ell^2} r) + b_\ell Y_1(\sqrt{y^2-k_\ell^2} r) \right\}
\]  

(8)

\[
B_z = \sum_{\ell=1}^{N} \sqrt{y^2-k_\ell^2} \cos(k_\ell z) \left\{ a_\ell J_0(\sqrt{y^2-k_\ell^2} r) + b_\ell Y_0(\sqrt{y^2-k_\ell^2} r) \right\}
\]  

(9)

\[
B_\phi = \sum_{\ell=1}^{N} \gamma \cos(k_\ell z) \left\{ a_\ell J_1(\sqrt{y^2-k_\ell^2} r) + b_\ell Y_1(\sqrt{y^2-k_\ell^2} r) \right\}
\]  

(10)

Here \( J_0, J_1 \) are Bessel and \( Y_0, Y_1 \) are Neumann functions (Bessel function of the second kind).

From (7) the boundary condition (2) reads now

\[
\sum_{\ell=1}^{N} r_i \gamma \cos(k_\ell z_i) \left\{ a_\ell J_1(\sqrt{y^2-k_\ell^2} r_i) + b_\ell Y_1(\sqrt{y^2-k_\ell^2} r_i) \right\} = C_i
\]  

\( i = 1, \ldots, p \)  

(11)

When we insert (8), (9) into (4) we obtain the boundary condition in the form

\[
-dz \sum_{\ell=1}^{N} k_\ell \sin(k_\ell z_i) \left\{ a_\ell J_1(\sqrt{y^2-k_\ell^2} r_i) + b_\ell Y_1(\sqrt{y^2-k_\ell^2} r_i) \right\} +
+ dr \sum_{\ell=1}^{N} \sqrt{y^2-k_\ell^2} \cos(k_\ell z_i) \left\{ a_\ell J_0(\sqrt{y^2-k_\ell^2} r_i) + b_\ell Y_0(\sqrt{y^2-k_\ell^2} r_i) \right\} = 0
\]  

(12)

We see that for each partial solution \( \ell = 1, 2, \ldots, N \) one has 2 partial amplitudes \( a_\ell, b_\ell \). Thus we have altogether 2N partial amplitudes. The P boundary conditions (11) or (12) have to be satisfied for all continuously varying \( r, z \) along \( z = z (r) \). This can be done exactly for \( N \to \infty \). For practical purposes \( N \) must be finite. We therefore assume that (11), (12) be satisfied in P collocation points...
1: At first sight, (11) or (12) contains 3N+P+1 unknowns namely 2N amplitudes, N separation constants \( k_\theta \), P constants \( C_i \) and \( \gamma \). We thus need to solve 3N+P+1 transcendental equations. The accuracy of the result depends apparently on P. For fixed P we must assume that the accuracy is also constant. Instead of solving the transcendental system (11) or (12), according to experience it is easier to choose arbitrary but evenly distributed values of the \( k_\theta \) between 0 and \( \gamma \) and to use (12) instead of (11). Then one has 2N+1 unknowns and equations (12) become a system of linear equations for \( a_\phi \), \( b_\phi \) and for \( \gamma \). In order that this system possess as a nontrivial solution, its determinant of order 2N must vanish. The transcendental equation thus obtained yields the eigenvalue \( \gamma \). Since the system (12) is linear and homogeneous with a vanishing determinant, we may choose one variable, e.g., \( a_1 = 1 \) and one equation may be omitted. Then the remaining 2N-1 inhomogeneous equations may be solved for \( a_2, \ldots, a_N, b_1, \ldots, b_N \). The fixed number \( P \) of collocation points determines the number \( N \) of partial solutions by \( N = P/2 \).

The method may be applied to tori of arbitrary meridional cross section and arbitrary aspect ratio \( \alpha = R/a \), where \( R \) is the major and a some sort of a minor radius. For a circular cross section we have

\[
z_i = z_i(r) = \sqrt{a^2 - (r_i - R)^2}, \quad i = 1, \ldots, P
\]

for the meridional curve. Thus \( z_i \) in (12) may be expressed by (13). In this way we obtained table 1.

We see that the (lowest) eigenvalue \( \gamma \) depends on the aspect ratio \( \alpha \).

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Table 1. Dependence of \( \gamma \) from \( \alpha \) for \( N = 6 \)

One might argue, that the choice of the \( k_\theta \) is highly arbitrary. Table 2. demonstrates that \( \gamma \) does not depend strongly on the choice of the \( k_\theta \).

In order to check the accuracy, we plotted the circle defined by the \( P \) collocation points (13), but using (12). As fig. 1. (for \( \gamma = 2.570507 \) ) shows, the boundary condition (12) describes realy a continous circle. The collocation points \( r_i, z_i \) are marked by small squares in Fig. 1.
<table>
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Table 2. Dependence of $\gamma$ from the $k_p$ for N=6
( $a=0.35$, $R=2$. and $a=.35$, $R=.95$ (Φ) )

References.

EFFECT OF THE ELECTRON ENERGY TRANSPORT COEFFICIENT ON THE STABILITY OF THE TEARING MODES.

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Tokamak electron temperature profiles since 1980 /1/ are observed to follow some form of profile consistency which, at present, is widely accepted to be represented by a gaussian shape /2,3/.

Such a property of the temperature profile produces the paradoxical effect on the models describing plasma in terms of energy transport coefficient that the electron temperature profiles appear to be determined by non-local processes restoring a unique temperature profile independently of the position and of the power of an additional heating /4/.

One seemingly natural solution of the problem relays on the obvious observation that current density profiles in a tokamak discharge must be stable against the principal MHD instabilities thus being limited to a narrow class of allowed magnetic topologies.

These ideas could suggest that the current density rather than the temperature profile is the key of the problem. A Kadomtsev-like model, usually linked to transport codes to limit the central safety factor q in the order of unity, is a clear example of how transport and MHD constraints have to be simultaneously fulfilled to fit experimental profiles /5/.

Tearing modes have been addressed as candidate mechanism for the profile consistency /6/ by observing that computed and measured /7/ current density profiles are stable against such instabilities and that in the ECR-Heated Stellarator WVII-A a profile consistency is not observed /8/. We remind that the driving mechanism of the tearing modes are the current density gradients, absent in stellarators.

Moreover it can be shown that the gaussian temperature profiles used to fit experimental data together with reasonable assumptions on the resistivity produce current density profiles stable against tearing modes /9/. A possible conclusion is that the "anomalous transport coefficients" describing experimental data (i.e. gaussian temperature profiles) include the effects of the tearing modes in a hidden way. If this hypothesis is accepted it becomes relevant to investigate on the physical ground of local transport coefficients that produce temperature and current density profiles unstable against, and affected by, tearing modes.
A steady state transport model, embodying a principle of $\Delta'$ minimization together with an enhanced transport in the vicinity of the rational surfaces $q=m/n$, has been implemented to study the effect of the radial dependence of the electron energy transport on the stability of the tearing modes.

The stability parameter $\Delta'$ is defined as $10$

$$\Delta' = \frac{1}{\psi} \frac{d\psi}{dr} \bigg|_{r=s^\pm} \quad \epsilon \to 0$$

where $\psi$ is the perturbed helical flux function (with indexes $m,n$) obtained solving the equation

$$\frac{1}{r} \frac{d}{dr} \left( \frac{r \frac{d\psi}{dr}}{} \right) - \frac{m^2}{r^2} \frac{\psi}{\psi} - \frac{B_\theta}{B_\bot(1-qn/m)} = 0$$

A simplified transport equation is assumed

$$\frac{d}{dr} (r n_x \frac{dT}{dr}) = -r P$$

where $n$ is the electron density 
$x$ the heat transport coefficient 
P the radially uniform power source
and $T$ the electron temperature.

The radial behaviour of the diffusion coefficient has been chosen of the form 

$$n_x = n_0 x_0 (1 - \lambda r/a).$$

In this way a high value of $\lambda$ corresponds to steep (unstable) temperature-current density profiles at the plasma boundary.

Solving eq (1) we obtain

$$T(r) = \left( \frac{r}{r/a - 1} \right) + \frac{1}{\lambda} \log \left( \frac{1 - \lambda r/a}{1 - \lambda} \right) \frac{n}{2\lambda}$$

where $\Pi = Pa^2/n_0 x_0$.

Eq. (2) describes the steadystate electron temperature for different $\lambda$ i.e. for a different slope of the diffusion coefficient. Fixing a toroidal field and imposing a safety factor $q(a)$ the current density and $q$ profiles are easily obtained assuming a Spitzer resistivity. The electron temperature in the central region is limited in order to prevent current density profiles with $q<1$.

The marginally stable $T$- $J$ profiles are obtained flattening $T$ on both sides.
of the rational surfaces until the instability parameter $aA'$ is lower than a prefixed value ($\sim 10^{-3}$).

We consider 8 major modes ($m/n=2/1, 3/1, 3/2, 5/2, 4/3, 6/3, 7/3, 8/3$). Starting from the $m/n=2/1$ mode the temperature is flattened on both sides of the rational surface until the corresponding $J$ profile becomes stable. The process is then repeated for higher order tearing modes and for the $q=1$ region and iterated to reach stable and uniquely determined $T(r) - J(r)$ profiles.

**RESULTS**

The efficiency of the described mechanism for shaping temperature profile appears clearly by comparing Fig.1 and Fig.2. Fig.1 represents temperature profiles parametrized in $\lambda$ ($0.1<\lambda<0.9$) with $q(a)=3.1$ and Fig.2 shows temperature profiles parametrized in the same way when the effect of the tearing modes is included. The profile consistency-like behaviour shown in Fig.2 is summarized in Fig.3 where the ratio between the mean temperature and the peak temperature, $<T>/T(0)$, is plotted versus $\lambda$ for $q(a)=3.1$. In the same figure the ratio $<T>/T(0)$ of temperature profile affected by tearing modes; the shape of the temperature profile is kept constant by tearing modes and, for a wide choice of $\lambda$ (i.e. $n_x$ slope), the final shape resembles a $n_x$-constant temperature profile. The increased transport induced by tearing modes masks the real local transport leading to the observed transport profiles.

In Fig.4 $<T>/T(0)$ is plotted for different $\lambda$ in function of $q(a)$. The simplified model that we adopted and the restricted number of rational surfaces taken into account limit the validity of the results above described for $q(a)$ ranging from 3 to 3.5. Nevertheless the physical model described in this work suggests a possible interaction between MHD instabilities and transport, and it could be used in conjunction with those transport coefficients that vanish for $r \to a$. /11/

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Fig. 1 Normalized Temperature Profiles without Tearing Modes.

Fig. 2 Normalized Temperature Profiles with Tearing Modes.

Fig. 3 Ratio between mean temperature and peak temperature vs. lambda

Fig. 4 Ratio between mean temperature and peak temperature vs. q(a)
FINITE–BETA MINIMUM ENERGY STATES ARISING FROM A MULTIPLE SCALE APPROACH TO TAYLOR'S MINIMUM ENERGY PRINCIPLE

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INTRODUCTION

In Tokamak's and RFP's a plasma relaxes on a timescale, short compared with the resistive diffusion time, to characteristic equilibrium states largely independent of initial and external conditions (e.g. selfreversal of RFP's, profile consistency of Tokamaks).

In formulating his famous variational principle, J.B. Taylor has assumed a 'slightly imperfect' plasma and minimized the magnetic energy, subject to the global constraints of helicity and flux conservation. This principle actually corresponds to the relaxation of a pressureless plasma to the minimum energy state on an intermediate timescale.

The relaxation is usually considered to be mainly due to tearing modes, each resonant on some rational magnetic surface. The timescale of plasma relaxation is then the one of the resistive tearing modes and the corresponding length scale is that of the thickness of a resistive layer, \( \Delta \). The large difference between the different scales offers an ideal opportunity for the application of a multiple scale method.

In our multiple scale approach Taylor's heuristic constraints of flux and helicity conservation emerge for the relaxation time scale as true invariants for the zero order quantities. The prediction which can be drawn from the resulting finite–\( \beta \) minimum energy states allow for a considerable improvement of Taylor's theory.

THE MULTIPLE SCALE APPROACH

We first consider the three characteristic timescales \( \tau_1 = \tau \) (ideal MHD) \( \tau_2 = \tau \) (relaxation) and \( \tau_3 = \tau \) (resistive diffusion) with the ordering \( \tau_1 \ll \tau_2 \ll \tau_3 \), and introduce the three dimensionless times \( \tau_1 = t/\tau_1 \), where

\[
\tau_1 = a \sqrt{\frac{\rho t_0}{B_0}}, \quad \tau_2 = \frac{\mu \Delta^2}{\eta} = \tau_3 \left( \frac{\Delta}{a} \right)^2, \quad \tau_3 = \frac{\mu_0 a^2}{\eta}. \tag{1}
\]

As the dominant length scale we consider the thickness of a resistive layer \( \Delta \), and the plasma radius \( a \), and introduce the dimensionless space variables

\[
x^i_1 = \frac{x^i}{\Delta}, \quad x^i_2 = \frac{x^i}{a}. \tag{2}
\]

In our multiple scale approach we then employ the expansion parameters \( \epsilon = \tau_1/\tau_2 \) and \( \delta = \Delta/a \) (both of the order \( 10^{-2} \) to \( 10^{-3} \)) and obtain

\[
\frac{\partial}{\partial x_i} = \frac{1}{\tau_1} \left( \frac{\partial}{\partial x_1} + \epsilon \frac{\partial}{\partial x_2} + \epsilon \delta Q_3 \frac{\partial}{\partial x_3} \right), \quad \nabla = \frac{1}{\Delta} \left( \nabla_1 + \delta \nabla_2 \right). \tag{3}
\]

For all physical quantities \( Q \) we now choose a multiple scale representation of the form

\[
Q = \overline{Q} \left( Q_0 + \epsilon Q_1 + \delta Q_2 + \epsilon^2 Q_3 + \epsilon \delta Q_4 + \epsilon^2 Q_5 + \epsilon^3 Q_6 + \ldots \right)
\]

where \( \overline{Q} \) is a characteristic dimensional value of \( Q \) and \( Q_i = Q_i \left( t_1, t_2, t_3, x^i_1, x^i_2 \right) \). This multiple scale approach then is applied to the full set of resistive MHD equations. To solve the resulting (dimensionless) equations on all five scales and for all quantities \( Q_1 \) is...
practically unfeasible, even in the first orders. However, since we are primarily interested in the static (or stationary) equilibrium states of the zero order quantities $Q_0$ on the two first scales, the problem simplifies substantially if one considers the asymptotic time limits $t_{1,2} \to \infty$.

**Time limit** $t_1 \to \infty$ and ideal MHD equilibrium states

From our multiple scale method we obtain in the leading order 1 just the ideal MHD equations containing only zero order quantities $Q_0$, $\partial / \partial t_1$ and $\nabla_1$.

The stable static equilibrium states an ideal MHD plasma relaxes towards in the limit $t_1 \to \infty$ are obtained by minimization of the total potential energy, subject to the constraints of the ideal MHD equations, as shown by Kruskal and Kulsrud /2/. Thus in the leading order 1, the result of the limit $t_1 \to \infty$ is

$$b \nabla_1 p_0 = (\nabla_1 \times \vec{B}_0) \times \vec{B}_0, \quad b = 2\beta (axis)$$

(4)

along with the requirement of ideal MHD stability. The asymptotic limit $t_1 \to \infty$ in the higher orders of $\epsilon$ and $\delta$ involves the higher order quantities $Q_1$ and is not considered.

**Time limit** $t_2 \to \infty$ and turbulent relaxation

In analogy with the limit $t_1 \to \infty$, the stable static (or stationary) equilibrium state a turbulent plasma relaxes to in the limit $t_2 \to \infty$, is again obtained by minimization of the total potential energy, but now subject to the constraints which determine the dynamics of the relaxation processes on the second timescale. From our multiple scale ordering it follows that

$$\frac{\partial}{\partial t_2} \vec{A}_0 = \nabla_1 \times (\nabla_1 \times \vec{A}_0)$$

(5)

so that the zero order helicity $K_0$ is an invariant of motion. The changes in $K$ are hence of the order of $\epsilon^2$. It further follows that also the zero order toroidal flux is conserved.

The existence of other invariants relating temperature, density and pressure has not been investigated so far. Therefore we consider the specific heat ratio $\gamma$ of the relaxation processes as a free parameter, which is then in the resulting equilibrium solution determined by other quantities.

Due to the calculus of the multiple timescale ordering the result of the asymptotic limit $t_1 \to \infty$, i.e., the equilibrium equation (4) has to be included as a necessary constraint in the calculation of lines $t_2 \to \infty$, i.e., in the minimization of the potential energy.

The assumption that the relaxation is mainly due to tearing mode activity further implies an invariant of the form $\int \vec{c} \cdot \vec{A}_0 \, d\tau = \text{const.}$, $\vec{c} = \text{const}$.

**THE VARIATIONAL PRINCIPLE**

Based on a multiple scale approach we now arrive for the second timescale at the following minimum energy principle:

$$W_0 = \int \left\{ \frac{1}{2} (\nabla_1 \times \vec{A}_0)^2 + \frac{b p_0}{\gamma - 1} \right\} d\tau \to \text{min.} \quad K_0 = \int \vec{A}_0 (\nabla_1 \times \vec{A}_0) d\tau = \text{const.}$$

(5)

$$P = b \nabla_1 p_0 - (\nabla_1 \times \vec{A}_0) \times (\nabla_1 \times \vec{A}_0) = 0, \quad L_0 = \int \vec{c} \cdot \vec{A}_0 d\tau = \text{const.}$$

with boundary conditions: $\delta \vec{A}_0 |_{\partial \Omega} = 0$, $\delta \left( b p_0 + \frac{1}{2} \vec{B}_0^2 \right) |_{\partial \Omega} = 0$.

(6)

With the help of a Lagrangian multiplier technique we obtain from the first variation
\[
\delta W_0 = \delta W_0 - \frac{1}{2} \lambda \delta K_0 + \delta \int \dot{\omega} \left( \chi_1 \right) \, d\tau + \delta L_0 = 0
\]

the Euler equations: 
\[E_1 = V_1 \cdot \ddot{\omega} - \frac{1}{\gamma - 1} = 0, \quad E_2 = \dot{P}^0, \quad E_3 = V_1 \times B_0 - \lambda B_0 - V_1 \times \left( (\ddot{\omega} \times (V_1 \times B_0) - V_1 \times (\ddot{\omega} \times B_0) \right) + \ddot{\omega} = 0 \quad (7)
\]

The stability of the resulting equilibrium solutions with respect to the relaxation modes is ensured by the minimum of \( W_0 \), i.e. by the positive-definiteness of the second variation
\[
\delta^2 W_0 = \int \left\{ \frac{-1}{\kappa} (\nabla \times \ddot{y})^2 - \lambda \ddot{y} (\nabla \times \dot{y}) - (\nabla \times \ddot{\omega}) (\nabla \times \dot{y}) - 2 (\nabla \times \ddot{\omega}) (\nabla \times \dot{y}) \right\} (\nabla \times \dot{y}) \cdot \ddot{\omega} \right\} d\tau \geq 0 \quad (8)
\]

for all permissible variations \( \delta A_0 = \dot{y} \)

**THE STABLE ONE-DIMENSIONAL CYLINDRICAL SOLUTION**

From the Euler equation (7) we obtain the one-dimensional cylindrically symmetric solution as
\[
B_r = 0, \quad B_\theta = \frac{\lambda B}{\mu} \mu_1 (\mu x), \quad B_z = B \left[ J_0 (\mu x) - C \right],
\]
\[
\mu_0 p = \frac{B^2}{2 |\kappa|} \left\{ \frac{1}{(\kappa + 1)} \left[ J_0^2 (\mu x) - J_0^2 (\mu) \right] - 2 \kappa C \left[ J_0 (\mu x) - J_0 (\mu) \right] \right\}
\]
\[
\mu = |\lambda| \sqrt{|\kappa|}, \quad x = \frac{r}{a}, \quad \kappa = \frac{\gamma - 1}{2 - \gamma} < 0. \quad (9)
\]

The constants contained in this solution now depend on the diffusion time \( t_3 \) and at the second space scale \( r_2 \), a dependence which is not discussed here. For convenience we have therefore the original argument of the Bessel functions \( \nu t_1 = \nu \Delta \) replaced by \( \frac{\nu t_1}{\nu \Delta} = \nu x \).

The investigation of the stability requirement \( \delta^2 W_0 \geq 0 \) of eq.(8), which is not quite trivial and discussed elsewhere, yields the stability criterion
\[
\kappa < 0 \quad \text{and} \quad \mu < 3.11 \quad (10)
\]

It further turns out that the dominating mode is an \( m = 1 \) mode with poloidal mode number \( n = (A.1.23/\sqrt{|\kappa|}) \), where \( A \) is the aspect ratio and \( (\cdot) \) means the nearest integer.

**COMPARISON WITH EXPERIMENT**

First we have applied our cylindrical solution (9) to a comparison with experimental magnetic field and pressure profiles of RFP's and Tokamaks. The free parameters in our solution we have determined by prescribing the toroidal values of current and magnetic flux, the pressure on axis and a vanishing pressure and toroidal current density at the wall.

In Fig.1 we have compared the resulting profiles with measurements of a typical \( \eta-\beta-II \) RFP discharge, where the found agreement of the magnetic field profiles is for most measurements within the experimental error. In Fig.2 the comparison is done for a typical TORTUR tokamak discharge. The toroidal magnetic field turns out to be almost constant and the pressure and toroidal current density profiles are bell shaped. The only internal measurement corresponds to the pressure.

In Fig.3 the F-O diagram is shown for typical tokamak values of \( Q_0 = q_{ax} A = 4 \)
and 6 and \( \beta \)-values of 2.5 and 10%. Over the total F–Θ diagram only this very narrow window of stable solutions (solid line) exists containing, however, the experimental parameter range of current Tokamaks (note, that \( \Theta^{-1} = \lambda_{q \text{wall}} \)). The bold-faced part of the solid line corresponds to equilibria for which the pressure is everywhere positive and vanishing at the wall.

In Fig.4 we have plotted the F–Θ diagram for typical RFP values with \( Q_0 = 0.8 \) and values of \( \beta = 2; 10 \) and \( 20\% \). The solid line again represents the stable solutions where for larger negative values of F the equilibrium again becomes unstable. The bold faced part represents the solutions with positive definite pressure. The predicted operational regime for RFP's in our F–Θ diagram well agrees with the experimental data.

The poloidal mode number \( n \) for the dominating \( m = 1 \) relaxation mode strongly depends on \( \beta \) and lies for RFP's in the range \( n = (3;20) \) and for Tokamaks at \( n = (20;500) \). Note, that in the F–Θ diagrams we have only taken into account the stability criterion (16) for relaxation modes. The stability requirement with respect to ideal MHD modes will essentially lead to a limitation of the high-\( \beta \) equilibria.

![Fig.1](image1.png) ![Fig.2](image2.png) ![Fig.3](image3.png) ![Fig.4](image4.png)

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MODE LOCKING IN THE TOKAMAK AND RFP

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Introduction

In large tokamaks (e.g. JET) coherent MHD activity is often observed to cease oscillating, but to continue growing in amplitude [1]. Such locked mode behaviour generally occurs for $m = 2$ (or 3) pre-disruption MHD activity. Locked modes have also been seen in tokamaks when resonant helical fields are applied, using external conductors (e.g. TOSCA [2]). In the HBTX1-B RFP locked mode activity has been observed for modes which are approximately resonant at the magnetic axis [3]. It appears likely that this mode locking is caused by the interaction of the MHD activity with either, the liner/wall, or with resonant helical fields. In this paper we examine these possible causes for mode locking.

Linear resistive wall effects

In both the tokamak and RFP allowing for resistive wall effects can destabilise ideal non-resonant modes with growth rates $\gamma / \tau_w$ ($\tau_w$ = liner time constant). For sub-Alfvenic rotation velocities, linear theory shows these ideal modes remain locked to the wall [4]. In the tokamak such locked modes can occur for $q_a < 2$ and are $m = 2$, $n = 1$ modes. In considering modes which are resonant near the resistive wall careful treatment of the boundary conditions is necessary. Considering the situation of a plasma surrounded by a thin vacuum annulus, which is in turn surrounded by a resistive wall, the constraints applied at the plasma-vacuum interface are pressure balance ($P + B^2/2$ continuous) and that the interface remains a flux surface. At the resistive wall a thin shell approximation is made [4] yielding the boundary conditions

$$\left[ \frac{\partial b_r}{\partial r} \right] = \tau \frac{\partial b_r}{\partial t}$$

where $\left[ \cdot \right]$ indicates the jump across the wall. To study resistive stability with these boundary conditions a set of reduced MHD equations [5] are solved using an implicit initial value approach.

Figure 1 shows how the linear $m = 2$, $n = 1$ growth rate (normalised to the Alfven time $\tau_A$) varies with the edge $q_a$ for a tokamak profile with an axial $q_o = 1.1$; the wall time $\tau_w = 10^3 \tau_A$ and the magnetic Reynolds number, $S = 10^8$. The rapidly
growing mode for \( q_a < 2 \) is an ideal external kink with a growth rate, \( \omega_1/\tau_w \). Calculations confirm the analytic result that plasma rotation has little effect on the growth rate for the ideal \( (q_a < 2) \) mode and that the mode remains locked to the wall. The properties of these modes correlate well with the current limit \( (q_\phi < 2) \) pre-disruption MHD activity which occurs in many tokamaks (e.g. JET), where locked \( m=2 \) modes (often with a very short rotating phase) occur with growth rates \( \sim 0.5/\tau_w \).

In the case of the RFP both internal and external non-resonant ideal modes can occur, when the effects of a resistive liner are considered. However, when a conducting wall is placed just outside the liner (as in HBTX-1B) the external modes are stabilised but the ideal on-axis modes \([4]\) can remain unstable for RFP configurations which are marginally stable with an ideal wall. These modes appear likely candidates for explaining the on-axis locked modes which have been observed in HBTX-1B \([3]\).

Non-Linear Resistive Wall Effects

For resonant MHD activity (tearing modes) linear theory shows resistive wall effects do not lead to mode locking \([4]\). There is however a non-linear effect which causes mode locking \([6,7]\). This arises because the eddy currents driven in the liner by the rotating mode produce a net torque which opposes the plasma rotation. By approximating the region between the resonant surface \( (r_s) \) and the edge as a vacuum, we may calculate the reduction in mode frequency \( (\Omega) \) due to this torque as

\[
\frac{d\Omega}{dt} = -\frac{2r_s(\Delta'_V-\Delta'_W)\Omega Q_\tau}{4(m^2+n^2\varepsilon^2)/A + m^2\Omega^2\tau_w^2/(m^2+n^2\varepsilon^2)}
\]

where \( W \) is the island width, \( 'A' \) is an equilibrium dependent constant, \( \varepsilon \) is the inverse aspect ratio, and \( \Delta'_V, \Delta'_W \) are the jumps in \( d \log b_r/dr \) at \( r_s \) with \( \tau_w=0 \) and \( \infty \), respectively. In the tokamak limit \( (\varepsilon<<1) A^1 \) \([6]\), while for an internally resonant RFP mode, \( A^2 \). The main difference in the mode lock rates between the tokamak and RFP is caused by the term \( (\Delta'_V-\Delta'_W) \) in Eq (2). Apart from near \( q_a \sim 2 \), a conducting wall has little effect on the \( m=2 \) mode stability in the tokamak, and typically \( (\Delta'_V-\Delta'_W) \sim 1 \), while in the RFP the conducting wall has a strong stabilising influence and for resonant modes, typically \( (\Delta'_V-\Delta'_W) \sim 5 \) to 10. This difference in \( (\Delta'_V-\Delta'_W) \), and the differences in profiles means, that (for the same \( \Omega, W \) and \( \tau_w \)) the mode locking rates in the RFP are an order of magnitude faster than the tokamak.

**Fig 2.** Frequency \( (\Omega) \) and \( d\Omega^2/dt \) v time showing mode lock in RFP.
An exception is near \( q_a \) \( \approx \sqrt{2} \) where \((\Delta V - \Delta W)\) can become large in the tokamak and cause rapid mode locking \([8]\). These RFP results are confirmed by non linear calculations. Figure 2 shows how \( \Omega \) and \( d\Omega^2/dt \) vary in time for the RFP profile \( J_B/ZB^2 = 3.6(1 - r^2 \cdot \rho^2) \) with \( m=1, n=11 \) (internal resonant) and \( \tau_w = 10^3 \tau_A \); the rapid mode locking (\( \Omega + 0 \)) is evident. The solid curve for \( d\Omega^2/dt \) is computed from Eq(2) and is in reasonable agreement with the numerical results (broken line). Using the measured values of \( \dot{E} \), at the liner in HBTX and assuming that the broad spectrum of modes will stop the plasma rotation across the entire minor radius we find a mode locking time of order 10\( \mu \)s. This is consistent with the observed decorrelation time in HBTX \([9]\).

For tokamak profiles numerical simulations of mode locking are also in reasonable agreement with Eq(2). The characteristic mode locking time calculated from Eq(2) for JET like parameters is \( \approx 40\)ms when \( f=1kHz \) and \( W=5\% \); which is much faster than the experimentally observed time of \( \approx 40\)ms. However, in comparing these numerical and analytic models with the experiment it must be remembered that the simple MHD model used only causes the plasma rotation in the vicinity of the island to stop. In practice, viscous and toroidal coupling may require that a larger fraction of the plasma stops rotating, and also neo-classical effects may inhibit poloidal rotation. We may nevertheless use Eq(2) to scale the expected island widths at mode lock between various tokamaks. Solving Eq(2) we find the island width \( (W_c) \) to reduce an initial frequency \( (\Omega_o) \) by a factor of 10 is, \( W_c^2 = C_1 \varepsilon^{-2} \gamma \tau_A^2 (0.99 \Omega_o^2 \tau_w + 9.2/\tau_w) \), where \( \gamma \) is the growth rate and \( C_1 \) is an equilibrium dependent constant. Taking typical values we find \( W_c \approx 10\% \) for JET, \( \approx 25\% \) for DITE, and \( \approx 70\% \) for TOSCA. Thus mode lock is likely to occur in JET but unlikely to occur in TOSCA, where \( W_c \) exceeds the disruption threshold. As noted above however care must be exercised in applying Eq(2) directly to the experiment and these values for \( W_c \) are most appropriately interpreted as a relative scaling between the various experiments.

Resonant Helical Field Effects

Resonant helical fields (whether applied intentionally or generated from coil errors) will interact with a rotating MHD instability to affect its frequency. Again assuming a vacuum to exist between \( r_S \) and the plasma edge \( (r=1) \), where the helical coil boundary condition is applied, we find in the tokamak limit that

\[
\tau^2 \frac{d\Omega}{A dt} = \frac{m^2}{32} \frac{I_C}{aB_z} \frac{r_S^{m-2}}{a} \frac{W}{r_S q^2} \sin m\Theta
\]

where \( I_C \) is the applied helical coil current, \( 'a' \) is the minor radius, and \( \Theta_o \) is the poloidal angle between the intrinsic magnetic island the helical coil island. We may solve Eq(3) to find

\[
\Omega^2 = \Omega_o^2 - \frac{m}{16} \frac{I_C}{a} \frac{r_S^{m-2}}{a} \frac{W}{r_S q^2} \frac{\tau_A^2}{\varepsilon} \cos m\Theta
\]
where $\Omega_0$ is the frequency before the helical field is applied. This shows that during one oscillation period an increase and corresponding decrease in frequency will occur as the torque between the rotating island and helical coil acts to enhance and then reduce the rotation velocity. These results are again reproduced by numerical solutions of the reduced MHD equations. Figure 3 shows for a tokamak q-profile ($q_a=3$) the effect of applying an $m=2$, $n=1$ helical field (at $t=0$) to an $m=2$, $n=1$, 2% island rotating at an initial frequency, $\Omega_0=10^{-2}T_A^{-1}$.

As the $m=2$ island grows towards its final saturation width ($\sim 30\%$), the force on the rotating island increases, and the strong deformations to the $B_\theta$ waveform, as the frequency increases and decreases are evident. There are two other noticeable effects in Fig 3. Firstly an increase in the period with time. This seems to be caused by the back EMF of the penetrating helical field. Secondly, for $t<9\times 10^3T_A$ the oscillations change character. This occurs because the peak torque becomes so large that the plasma island is no longer able to make a complete revolution and goes into an oscillation about the static helical coil (i.e. $\Omega^2<0$ in Eq(4) for $\theta_0=0$). These strong deformations to the $B_\theta$ wave forms have been observed in experiments where resonant helical fields are applied (e.g. DITE, CLEO) but there is no experimental evidence of any reverse direction oscillations near the mode lock. These reverse oscillations would however be strongly inhibited by toroidal coupling to $m=3$, and damped if resistive wall effects are considered.

Fig 3. $B_\theta$ at two poloidal locations for rotating island interacting with coil

OBSERVATION OF NONLINEAR RESISTIVE MODE STRUCTURE ON JET TEMPERATURE PROFILES

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INTRODUCTION
The nonlinear evolution of resistive modes in a tokamak plasma is expected to develop finite size magnetic islands around the rational values of the safety factor q/1-2/. The resulting perturbed equilibrium profiles of temperature and pressure could become flatter across these islands. Direct experimental observation of the internal structure of MHD perturbations has been so far provided essentially by X-ray tomography /3/. In one case /4/ visible continuum profile measurements have been successfully used and in another electron cyclotron emission measurements /5/. Recent developments of the LIDAR electron temperature diagnostic /6/ in JET and of the techniques of solution of the equilibrium identification problem /7/ has provided good evidence of strict correlation between the location of low order rational q values computed from the magnetic data, and the localised regions of reduced slope in the temperature (Te).

The correlation is also clearly found with the results of magnetic mode analysis and soft X ray tomography.

DATA ANALYSIS
In Fig. 1 a typical LIDAR Te profile in the equatorial plane of JET is presented, for shot 14010 at t=11 s, and in Fig. 2 the Te profile is associated to that of q obtained from the magnetic equilibrium reconstruction code/7/.

It is possible to identify flatter regions of Te close to the surfaces with q=1,2 and other rational values.

The results of soft X-ray analysis are shown in Figs. 3 to 6. Fig. 3 shows the time evolution of a perturbation rotating at 400 Hz, at t=11 s, and Fig. 4 shows the radial structure of the relative amplitude of the modes present. The peaks corresponding to m=1,2,3 correlate well with Figs. 1 and 2. Fig. 5 shows the reconstruction of the plasma cross section with the m=1 flattening coinciding with that of Fig. 1. The X ray tomography of the central plasma zone (Fig. 6) shows an m=1 island-like structure related to the q=1 surface. Consistently with the fact that in a noncircular toroidal case the X and O points of an island are not equidistant poloidally sometimes the flattening of modes with even parity is seen only on one side of the profile in the equatorial plane. The same circumstance occurs naturally for odd parity modes which do not appear on both sides when even parity modes do.

Figs. 7a and 7b show the position of the “plateaus” of the Te
profile correlated with the q profiles at two different times t=10 s and t= 12 s for shot 12579. This shot exhibits magnetic activity identified in Fig. 8 as a slowly rotating mode. The top trace shows the number of the octant where the '0' points of the n=1 rotating islands are in the midplane, on the low field side. The bottom traces are the sine and cosine components of the radial magnetic field perturbation.

The soft X-ray tomography of Fig.9 confirms the existence of flattening perturbations or island like structures of order m=1,2, located at the identified q rational surfaces. The LIDAR measurements error bars do not smear the systematic geometrical correspondence with rational q locations observed on a random sample of JET discharges. However the radial resolution of the apparatus (~12cm) may affect the width and slope of the flatter regions as well as their location relative to the position of the rational q surfaces deduced magnetically/6/. It is generally observed that the LIDAR profiles are wider on the high field side(Fig.1 and 7). This is consistent with what is observed in the soft X-rays profile reconstruction. Fig 10a and 10b show the q=2 region for shot 12579 on the inner and outer side. On the inner side the perturbation is clearly wider. As the modes rotate the flattening may be seen or not. The structures seen on these profiles appear to evolve slowly, as expected of nearly saturated resistive modes, in non disruptive shots. Although no conclusions are to be drawn at present on the exact nature of all the perturbations observed the experimental facts shown support the interpretation that an identification of saturated resistive modes has been made.

CONCLUSIONS
We have presented evidence of a definite correlation between independent experimental measurements, which support the picture of tokamak equilibrium profiles which reach a state of macroscopic (marginal) stability in presence of saturated (or slowly growing) helical MHD resistive modes.

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Fig. 1 - LIDAR Te profile (shot 14010 t=11 s) vs major radius

Fig. 2 - Te and magnetically identified q profile

Fig. 3 - Central soft X-ray flux amplitude vs time

Fig. 4 - Relative mode amplitude vs R

Fig. 5 - X-ray midplane emission profile vs R

Fig. 6 - X-ray tomography
Fig. 7a - LIDAR Te profile (shot 12579 t=10) and magnetically identified q profile

Fig. 7b - LIDAR Te profile at t = 12 and magnetically identified q profile

Fig. 8 - a) Octant number of n=1 island 'O' point
b) and c) sine and cosine component of radial magnetic field perturbation

Fig. 9 - X-ray tomographic cross section

Fig. 10a - High field side blow up of q=2 region
Fig. 10b - Low field side blow up of q=2 region
XV
FUSION
DUBROVNIK
1988

First Author
Index
<table>
<thead>
<tr>
<th>Author</th>
<th>Page</th>
<th>Author</th>
<th>Page</th>
</tr>
</thead>
<tbody>
<tr>
<td>Murphy T. J.</td>
<td>I–91</td>
<td>Škorić M. M.</td>
<td>II–1283</td>
</tr>
<tr>
<td>Nagatsu M.</td>
<td>III–1187</td>
<td>Snipes J. A.</td>
<td>I–346</td>
</tr>
<tr>
<td>Nagornyj V. P.</td>
<td>II–522</td>
<td>Soldner F. X.</td>
<td>III–874</td>
</tr>
<tr>
<td>Naito O.</td>
<td>I–159</td>
<td>Start D. F. H.</td>
<td>I–354</td>
</tr>
<tr>
<td>Navarro A. P.</td>
<td>III–1103</td>
<td>Sugisaki K.</td>
<td>II–625</td>
</tr>
<tr>
<td>Nave F. M.</td>
<td>I–441</td>
<td>Tabares F. L.</td>
<td>III–1163</td>
</tr>
<tr>
<td>Neilson G. H.</td>
<td>I–486</td>
<td>Tanaka H.</td>
<td>III–1031</td>
</tr>
<tr>
<td>Neudatchin S. V.</td>
<td>III–1147</td>
<td>Tanga A.</td>
<td>I–235</td>
</tr>
<tr>
<td>Neudatchin S. V.</td>
<td>III–1003</td>
<td>Tartari U.</td>
<td>III–1107</td>
</tr>
<tr>
<td>Noterdaeme J. — M.</td>
<td>II–762</td>
<td>Taylor P.</td>
<td>II–573</td>
</tr>
<tr>
<td>O’Rourke J.</td>
<td>I–155</td>
<td>Tibone F.</td>
<td>II–709</td>
</tr>
<tr>
<td>Ohybu N.</td>
<td>I–227</td>
<td>Tokar’ M. Z.</td>
<td>II–675</td>
</tr>
<tr>
<td>Okabayashi M.</td>
<td>I–171</td>
<td>Tsuboi F.</td>
<td>III–1195</td>
</tr>
<tr>
<td>Orsitto F. P.</td>
<td>III–1183</td>
<td>Tsui H. Y. W.</td>
<td>II–585</td>
</tr>
<tr>
<td>Ortolani S.</td>
<td>II–569</td>
<td>Van Milligen B. Ph.</td>
<td>I–318</td>
</tr>
<tr>
<td>Pan C. H.</td>
<td>III–920</td>
<td>Van Niekerk</td>
<td>III–1221</td>
</tr>
<tr>
<td>Pasini D.</td>
<td>I–251</td>
<td>Van Nieuwenhove R.</td>
<td>II–778</td>
</tr>
<tr>
<td>Pavlenko V. P.</td>
<td>III–1209</td>
<td>Vannucci A.</td>
<td>I–203</td>
</tr>
<tr>
<td>Pešić S.</td>
<td>II–858</td>
<td>Vasin N. L.</td>
<td>I–59</td>
</tr>
<tr>
<td>Pfirsch D.</td>
<td>III–1229</td>
<td>Vdovin V. L.</td>
<td>III–1027</td>
</tr>
<tr>
<td>Potapenko I. F.</td>
<td>III–1019</td>
<td>Vinogradov N. I.</td>
<td>I–71</td>
</tr>
<tr>
<td>Pozzoli R.</td>
<td>II–866</td>
<td>Vlad G.</td>
<td>I–409</td>
</tr>
<tr>
<td>Puri S.</td>
<td>II–754</td>
<td>Voronov G. S.</td>
<td>I–463</td>
</tr>
<tr>
<td>Puri S.</td>
<td>III–952</td>
<td>Wagner F.</td>
<td>I–207</td>
</tr>
<tr>
<td>Pustovitov V. D.</td>
<td>II–490</td>
<td>Waidman G.</td>
<td>I–381</td>
</tr>
<tr>
<td>Pustovitov V. D.</td>
<td>II–502</td>
<td>Wang Long</td>
<td>II–819</td>
</tr>
<tr>
<td>Radeztsky R. H.</td>
<td>I–79</td>
<td>Wang Z.</td>
<td>II–835</td>
</tr>
<tr>
<td>Rebut P. H.</td>
<td>I–259</td>
<td>Ward D. J.</td>
<td>I–330</td>
</tr>
<tr>
<td>Rebut P. H.</td>
<td>I–247</td>
<td>Westerhof E.</td>
<td>I–401</td>
</tr>
<tr>
<td>Riviere A. C.</td>
<td>II–811</td>
<td>White R. B.</td>
<td>I–413</td>
</tr>
<tr>
<td>Romanelli F.</td>
<td>I–374</td>
<td>Wurden G. A.</td>
<td>II–533</td>
</tr>
<tr>
<td>Rubel M.</td>
<td>II–683</td>
<td>Yamaguchi N.</td>
<td>II–593</td>
</tr>
<tr>
<td>Ryan P. M.</td>
<td>II–795</td>
<td>Yoshida H.</td>
<td>I–163</td>
</tr>
<tr>
<td>Sadler G.</td>
<td>I–131</td>
<td>Zaki N. G.</td>
<td>III–1260</td>
</tr>
<tr>
<td>Sato M.</td>
<td>II–470</td>
<td>Zarnstorff M. C.</td>
<td>I–95</td>
</tr>
<tr>
<td>Sauter O.</td>
<td>II–758</td>
<td>Zhao Hua</td>
<td>II–601</td>
</tr>
<tr>
<td>Scharer J.</td>
<td>II–787</td>
<td>Zurro B. G.</td>
<td>II–699</td>
</tr>
<tr>
<td>Schild P.</td>
<td>III–1139</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Schoch P. M.</td>
<td>I–191</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Scott S. D.</td>
<td>I–103</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Sesnic S.</td>
<td>I–385</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Shi X. H.</td>
<td>II–526</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Shinohara S.</td>
<td>II–541</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Shishkin A. A.</td>
<td>II–478</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Shoji T.</td>
<td>I–219</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Shukla P. K.</td>
<td>III–916</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Shukla P. K.</td>
<td>III–1205</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Simonen T. C.</td>
<td>I–223</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Sinman A.</td>
<td>II–621</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Sinman S.</td>
<td>II–617</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Sitenko A. G.</td>
<td>III–1233</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Skladnik — Sadowska E.</td>
<td>II–633</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>