20th EPS Conference on Controlled Fusion and Plasma Physics

Lisboa, 26–30 July 1993

Editors: J. A. Costa Cabral, M. E. Manso, F. M. Serra, F. C. Schüller

Contributed Papers, Part II

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CONTRIBUTED PAPERS, PART II

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PREFACE

The 20th EPS Conference on Controlled Fusion and Plasma Physics is organized, on behalf of the European Physical Society (EPS), by "Sociedade Portuguesa de Fisica" (SPF) and "Centro de Fusão Nuclear" (CFN) of "Instituto Superior Técnico" (IST) of the Lisbon Technical University.

The Programme, Format and Schedule of the Conference were determined by its International Programme Committee, which also selected the Plenary and Topical Invited Lectures.

The International Programme Committee has also made the selection of the submitted one-page abstracts. Some of these abstracts, of outstanding quality, have been selected, for both poster and oral presentation of the corresponding four-page papers.

In the odd years the Conference is essentially related with Controlled Fusion Research and it has a reduced format. Therefore, the IPC has only been able to accept for conference presentation about 435 abstracts from the almost 620 received.

The Conference Format is: 9 Review Lectures of 45 minutes, 18 Topical Lectures of 30 minutes, 24 Oral Presentations of Contributed Papers of 20 minutes, 4 poster sessions with about 110 posters each and a Special Evening Public Lecture.

Lisboa, June 1993

The Editors

Acknowledgements:

The Conference Organizers acknowledge the financial support of the following Agencies and Institutions:

- Junta Nacional de Investigação Científica e Tecnológica
- Commission of the European Communities
- Fundação Calouste Gulbenkian
- Instituto Superior Técnico
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Topic 3

Alternative Confinement Schemes
Analysis of Reversed Field Pinch plasmas in RFX

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Introduction

An analysis of plasmas produced in the RFX Reversed Field Pinch (RFP) experiment [1] (minor radius a=0.457 m and major radius R=2 m) is presented. Global energy confinement time, poloidal beta and radiation losses are reported in different I/N regimes, where I is the plasma toroidal current and N is the average line density. The study refers to discharges with plasma current between 0.5 and 0.65 MA, while the I/N parameter varies in the range 2 ≤ I/N ≤ 6 (10^{-14} \text{ A m}) (assuming a flat density profile). In order to compare the measured plasma resistivity with the classical one the helicity balance equation has been used, the plasma internal magnetic field profiles being reconstructed from externally measured global magnetic quantities. Finally the transport properties are briefly discussed in terms of stochastic thermal conductivity.

Results

In RFX the plasma current is sustained in the flat-top phase by external power supplies for about 100 msec [2]. Typical electron density (at 0.5 MA) is 3-4 \times 10^{19} \text{ m}^{-3}, while electron temperature is in the range 300-450 eV with ion temperature being comparable. Density profiles [3] are rather flat, since the first wall of RFX is completely covered by graphite tiles and density sustainment is practically achieved by wall loading; while temperature profiles [3] are relatively more peaked. Using these measurements, the global energy confinement time defined from a stationary power balance as:

\[ \tau_e = \frac{3 \langle n k (T_e + T_i) \rangle}{2 \; V \; I} \]  (1)

where \langle > \) represents the volume integral, V is the resistive loop voltage and I the plasma current, can be easily calculated. If radiation losses are negligible with respect to the input power, eq. (1) estimates the energy confinement time due to transport processes.
In many RFP devices [4,5,6,7] it has been observed that the confinement is strongly dependent on the parameter I/N with the best performances obtained at low I/N, although the electron temperature shows an opposite behaviour with I/N. These features are also confirmed in RFX as shown in Fig.1. \( \tau_e \) scales like \((I/N)^{-0.5}\). Note that the points in Fig.1 (and in all the figures reported in this paper) are obtained considering time averages over an interval of about 20 msec during current flat-top. The experimental points refer to more than 200 discharges and the scatter is due to the fact that a selection of homogeneous conditions has not been done. As shown in ref. [2] for example the matching of the field errors at the gaps gives a clear reduction in the applied loop voltage.

In Fig.2 the radiated power, \( P_{rad} \), normalized to the input power is plotted vs. I/N. It is quite low if I/N \( \geq 2 \times 10^{-14} \) (A m) and is around 10% of the input power. Some uncertainty is present in the measurements of \( P_{rad} \), since they are presently obtained with a single chord bolometer and the absolute value of \( P_{rad} \) could be affected by systematic errors due to profile effects. The points plotted in Fig.2 refer to an assumed annular distribution of the plasma radiation profile. The uncertainty due to the precise form of the profiles should not affect the relative scaling of the experimental data with I/N, if the profiles do not vary with I/N.

The decrease of \( \tau_e \) with I/N can be discussed writing eq. (1) as:

\[
\tau_e \propto \beta_0 \; \tau_{e0}^{3/2} / Z_e
\]
where \( Z^*_k = \eta^k / \eta^{cl} (Z=1) \), is the so-called helicity balance resistivity anomaly factor [8]. As shown in Figs.3,4 both the poloidal beta (\( \beta_\theta \)) and \( Z_k \) scale unfavourably with I/N, a feature previously observed in other RFPs [4,5,6,7]. Note that \( \beta_\theta \) has been calculated using a quite flat density profile, \( n \sim (1-r^6) \), and a temperature profile, \( T_e \sim (1-r^3) \), in agreement with experimental measurements [3]. The trend of \( \beta_\theta \) with I/N is due essentially to the fact that the decrease of density is not compensated by a corresponding increase in electron temperature. The scaling of \( Z^*_k \) with I/N however is not so easily interpreted.

A possible explanation may be found, as suggested in ref. [9], in the dependence of the anomaly factor on the quantity \( E_\phi / E_D \), \( E_\phi \) is the applied electric field and \( E_D \) the Dreicer field for runaway of electrons. The dependence of \( Z^*_k \) on this ratio is predicted by the Kinetic Dynamo Theory (KDT) [9]: the increase of \( Z^*_k \) with decreasing \( E_D \) (or decreasing density) can be interpreted as due to the diffusion of electrons in the stochastic magnetic field and to the loss of momentum to the wall. This interpretation is also consistent with the measurements of an asymmetry between the heat fluxes on the electron and ions drift sides at the plasma edge [10]. In RFX, however there are presently other indications that \( Z^*_k \) is just in the same range (see Fig. 4) of the spectroscopic \( Z_{eff} \) [11] in the explored interval of I/N, so that the helicity balance resistivity agrees fairly well with the Spitzer's one. If this is the case the increase of \( Z^*_k \) with I/N can be more easily interpreted as related to the increase in the impurity content.

From the measured confinement time it is possible to calculate the stochastic magnetic diffusion coefficient [12]:

\[
D_{st} = \frac{a^2}{\tau_e v_{th}}
\]
Using the measured values of electron temperature we find $D_{st} = 2 - 3 \times 10^{-5}$ m. Considering the fluctuation level at the wall (of the order of 1% of the total field) this does not fit a quasi-linear estimate [12], $D_{st} \geq 6 \times 10^{-4}$ m. This may be due to the fact that the stochasticity parameter $s$ [12] in the RFP is larger than unity so that the quasi-linear estimate is not suitable. Another possible explanation could be that the quantity $R = (b \cdot L/B \delta)$ [13], where $\delta$ and $L$ are respectively the transverse and the parallel correlation lengths of the magnetic field, is larger than unity in this way a turbulent regime results and the magnetic diffusivity should be calculated as $D_{st} = (b/B) \delta$. However in that case a very short transverse correlation length $\delta$ of the order of few millimetres can be estimated. Despite these difficulties, the stochastic transport model is believed to be also for RFX a good candidate to explain the observed energy confinement.

Conclusions

Global energy confinement, poloidal beta and resistivity anomaly factor have been measured in the RFX device. The best confinement has been achieved at low values of $I/N$, as previously observed in other RFPs, due to unfavourable scaling of poloidal beta and resistivity anomaly with $I/N$. The calculated values of global energy confinement (1 msec) and poloidal beta ($\leq 10\%$) are not completely satisfactory for the achieved current level (0.5 to 0.65 MA). On the other hand, the present results suggest that a significant improvement should be obtained just by reducing the impurity level, which is presently quite high [11], e.g. by conditioning the wall or by using a pellet injector to peak the density profiles.

References
MEASUREMENTS OF ELECTRON DENSITY FLUCTUATIONS ON RFX

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1. INTRODUCTION

RFP discharges are known to support a wide variety of fluctuations which may be related to the turbulent processes taking place in the plasma during relaxation and which play an important role in anomalous transport [1]. In ZT-40, density fluctuations have been measured by interferometric methods [2] and significant fluctuations are found in the range 10-150 kHz. On the RFX reversed field pinch experiment [3] the electron density is measured by a multichord two-color infrared interferometer [4] with a resolution of 1.2 x 10^{17} m^{-2} for the line density and a maximum bandwidth of 1 MHz which is thus adequate for the fluctuations measurement. In this paper RFX density fluctuations below 20 kHz are analysed in terms of spectrum, RMS and correlation between the chords during the RFP phase and for different I and I/N regimes.

2. FLUCTUATIONS CHARACTERISTICS AND SIGNAL TREATMENT

The RFX interferometer [4] is two-color, with a CO₂ laser beam for the plasma density measurement and a He-Ne laser for vibration compensation. It uses heterodyne detection at 40 MHz. Presently, eight interferometric chords supply the density measurement and provide a reasonable picture of the spatial profiles. Figure 1 shows the interferometer chordal positions and angles. The resolution is 1.2 x 10^{17} m^{-2} and the maximum bandwidth is of 1 Mhz.

The density fluctuation spectra (more details are found in figure 4) has the highest amplitude in the 1-6 kHz interval, and is nearly constant and higher than the background noise towards the 20 kHz Nyquist frequency (fig 3). The spectrum of fluctuations sampled at 100 kHz is similar to that at 40 kHz and the fluctuation amplitude at the Nyquist frequency 3 times higher than noise. Thus, the frequency interval between 0.5 and 20 kHz seems to contain sufficient information on the density fluctuations at low frequency.

Some density fluctuations general features can be inferred just from the density measurements:
- The fluctuation level is larger for the outermost chords (1 and 8 in figure 1);
- The fluctuation level of chord 1 is greater than that of chord 8 in most cases;
- Zero time lag correlation is found at all frequencies in the range of study.

Figure 2 illustrates the former characteristics. In our analysis we have window filtered the density signal with a 6th order Butterworth filter having a cutoff frequency at the 95% of the amplitude.

3. ANALYSIS TECHNIQUES: POWER SPECTRUM, RMS AND CORRELATION

In order to compute the fluctuations power spectrum of the signal \( S_x(f) = \langle X^*(f)X(f) \rangle \), (where \( X(f) \) is the FFT of \( x(t) \) and the brackets represent average over data subsets), the density signal is previously filtered for the chosen frequency window and the time interval
of interest is selected. That time interval is divided into a number of subsets which are chosen with a 50% overlapping. A Hanning window is applied to each subset in order to reduce the effects of leakage before applying the FFT algorithm. Averaging the power spectrum over the mentioned subsets and over different shots reduces the statistical fluctuations and highlights the underlying mean behaviour of the fluctuations [5].

The RMS has been calculated for a given time interval after filtering the density signal for the chosen frequency window.

The correlation between two interferometric signals has been calculated in the frequency space from the cross power spectrum $S_{xy}(f) = \langle X(f)Y(f) \rangle$ with the above windowing and overlapping techniques in order to obtain its phase, which gives information on the time shift between $x$ and $y$ as a function of frequency. We have calculated the correlation coefficient in the time space, defined as $C_{xy} = \langle x(t)y(t) \rangle / \sqrt{\langle x^2 \rangle \langle y^2 \rangle}$, for different frequency windows of interest and for all pairs of interferometric signals.

4. RESULTS

POWER SPECTRUM: Power spectra decrease exponentially, with some peaks around 1-2 kHz, being from 6-8 kHz nearly constant up to 20 kHz as can be seen from the continuous line in figure 3. Its amplitude is much larger than background noise showed in fig. 3 as a dashed line. The background amplitude is constant from 2 kHz onwards and at the Nyquist frequency is of an order of magnitude below the density fluctuation amplitude. The shots chosen for this example belong to a low density shot group and hence give a pessimistic estimate of the fluctuation/background noise RMS rate because, as we will see later, the RMS increases with density. The power spectrum amplitude of the outermost chords is always higher than that of the inner chords but the spectrum has always the same general features.

RMS: The RMS of chords 1 and 8 is higher than that of the inner chords which present very similar values among them. The normalized RMS/<n> is about 2% for the inner chords and 5-10% for the outer chords. In addition, significant differences between the RMS value of chord 1 and 8 are found from shot to shot, chord 1 having a higher RMS in most cases while chord 8 may have RMS values similar to those of the inner chords. Figure 4a shows two representative shots in which these differences are observed. It is seen that shot 2006 with a more symmetric density profile (fig. 4b) has also a more symmetric RMS profile. The opposite happens for shot 2003 with a density profile which is shifted towards the external part of the torus.

The absolute RMS is found to increase with the density. Averages of the normalized RMS (RMS/<n>) have been calculated for the same I/N intervals described before. The normalized RMS shows a weak increase with I/N. Averages taken for two data sets with different peak current (550 and 650 kA) have been also computed. RMS/<n> does not scale with current.

CORRELATION: The computed phase difference between the fluctuation signals is zero and the equal time correlation computed for all pairs of interferometric chords is very high, ranging from 0.4 to 0.9. It generally decreases with frequency and with chord separation as seen in figure 5. Correlation between chords 1 and 8 is seen to increase for shots in which the density profile is highly symmetric.

The correlation coefficients have been calculated for two data sets at different plasma current (550 and 650 kA). No substantial differences are found between the two cases. Just a higher correlation for the [2,6]kHz interval in the 650 kA set is encountered.
Correlation averages for different I/N sets at [0.5,20]kHz give no significant differences.

5. SUMMARY AND DISCUSSION

We have analyzed the RFX density fluctuations with the interferometric technique for the [0.5,20]kHz interval. The outer chords RMS is substantially higher than that of the inner chords which have a similar fluctuating level. The correlations between the chords have zero time lag and the correlation coefficient is high, ranging from 0.4 to 0.9. This RMS and correlation behaviour gives us information on the spatial distribution. A high correlation at zero time lag suggests the presence of a structure with \( m=0 \) symmetry. On the other hand, the RMS is much higher in the outer chords. The picture that comes out is that of a fluctuating annulus localized in the outer plasma \( (r/a>0.62) \) which is 'seen' mainly by the outer chords (1 and 8) and often only by chord 1. The differences which are often encountered between the RMS of chord 1 and 8 may be due to differences in the toroidal shift of the density profile as fig. 4 suggests. Chord 1 and 8 are nearly symmetric as seen in fig. 1. The fact that a small toroidal shift which is \( \pm 2 \) cm may change the fluctuating zone 'seen' by chord 8 supports the idea of a very localized and external annulus. We can hypothesize a small width for this plasma annulus supported by the fact that the differences of RMS among the inner chords are not significant. In an RFP the external zone corresponds to the reversal region, in which the toroidal field changes direction and where many resonant modes are present. These could be the source of the observed large fluctuations.

REFERENCES


![Figure 1. The eight RFX interferometer chords. Positions starting from chord 1 are: \( r=0.338,0.233,0.14,0.053,0.053,0.141,0.226,0.403 \) m. The angles respect the horizontal are 90,90,80,16,90,90,79,89,90,90 degrees.](image)
a) density fluct. background

$\rho \sim 10^{-10} \text{m}^{-1}$

Continuous line: typical power spectrum obtained by averaging 16 shots at high $I/N$ ($7 \times 10^{-14}$) for [20,40]ms and a low cutoff frequency at 500 Hz. Dashed line: background spectrum for the same data set, taken for [100,120]ms.

b) dataset, taken for [100,120]ms.

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**Figure 2:**

- **a)** 8-chord density measurements for shot n. 2038.
- **b)** 10 ms interval from the same shot filtered with a 0.5-6 KHz frequency window.

**Figure 3:**

Continuous line: typical power spectrum obtained by averaging 16 shots at high $I/N$ ($7 \times 10^{-14}$) for [20,40]ms and a low cutoff frequency at 500 Hz. Dashed line: background spectrum for the same data set, taken for [100,120]ms.

**Figure 4:**

- **a)** Density profiles as obtained from chordal density measurement for two different shots (#2006 and 2003).
- **b)** The corresponding RMS calculated for a frequency window of 0.5-2 KHz and a time interval of (30-60 ms).

**Figure 5:**

Correlation coefficient matrix $C_{jk}$, plotted for each $j$ and $k\neq j$, calculated for [30,60]ms during the RFP phase, taking averages over seven shots from a set at 650 KA for 3 frequency intervals:

- [0.5,2] KHz
- [2,6] KHz
- [6,10] KHz.
Magnetic field errors and non-axisymmetric behaviour of the plasma in RFX

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Introduction
RFX is a RFP experiment operating in Padua with R=2 m and a=0.457 m /1/. The device is equipped with a thick shell to stabilize MHD modes in the plasma. On the other hand, the shell equilibrium currents are forced to deviate at the two poloidal insulating gaps causing a local distortion of the magnetic field configuration. The experience of other RFP devices has clearly shown the importance of minimizing magnetic field errors at the gaps to improve discharge parameters and to increase its duration /2,3/. In RFX, sixteen coils, making up the so called Field Shaping Winding, have been located around the torus in order to generate part of the equilibrium field and produce a m.m.f equal and opposite to the plasma m.m.f. An accurate programming of the currents in this winding has been sought to get as close as possible to an axisymmetric magnetic configuration. A detailed analysis of the magnetic configuration is allowed by the RFX magnetic measurement system /4/. In particular, for the analysis presented in this paper, the measurements of the B13 component of the magnetic field have been provided by 8 poloidal arrays of pickup coils and by 2 toroidal arrays of 72 pickup coils, lying at 22.5° from the equatorial plane of the machine, on its inner and outer side; moreover, poloidal flux loops, Rogowski coils, saddle coils mounted at the gaps have been used.

Field error effect on plasma discharge performance
Several pulses (from 1770 to 1880) have been dedicated to study the horizontal equilibrium and to minimize the field error at the poloidal gaps throughout the discharge by properly programming the 8 Field Shaping Winding (FSW) thyristor amplifiers. All the pulses were executed with the same initial applied loop voltage (360 V) and the same magnetic energy initially stored in the magnetizing winding (29 MJ). Moreover, the flat-top poloidal amplifiers were always active in the examined cases and a steady-state vertical field was applied to reduce the plasma outward shift /5/. To characterize the quality of the magnetic configuration, for each pulse we evaluated the poloidal field error by taking the average of the absolute values of the measurements provided by the couples of saddle coils spanning the two poloidal gaps of the machine. This result has then been normalized with respect to the zero-order poloidal field at the plasma edge and time averaged during the flat-top in order to allow the comparison among different shots.

Typical results in terms of plasma current are shown in fig 1: pulse 1774 has been obtained controlling only the current of the outermost coil, while pulse 1860 is a significant example of an effective open loop programming of all the eight FSW amplifiers. The reduction of the time averaged relative field error from 15% to 6% resulted in an increase of duration of about 30 ms.
As an example of incorrect programming, pulse 1853 is presented, where the relatively short duration is due to an excess of external vertical field with respect to the ideal configuration. Fig. 2 shows the general trend of the pulse duration vs. average relative field error, for shots with 2.5-10^{-14}<I/N<5-10^{-14} \text{ Am}, 1.5<\Theta<1.6 and steady state vertical field Bv=6 to 8 mT. The pulse duration is defined by the loss of field reversal.

As a matter of fact, duration is also affected by the switching off instant of the flat-top amplifier; this is the case for pulses longer than 85 ms.

**Toroidal distortions of the poloidal field configuration**

The toroidal arrays of pickup coils have been used to investigate the toroidal distortion of the configuration due to the penetration of the magnetic field through the gaps. In order to get rid of systematic uncertainties, the calibration constants have been corrected by imposing a condition of toroidal axisymmetry when the field error at the gaps, measured by the saddle coils, vanishes (usually at about 15 ms). In fig.3 the poloidal field measured by the outer array of probes is shown as a function of the toroidal angle at different times (shot 1860). The position of the gaps is characterized by two peaks of growing amplitude and constant width. The toroidal extension of the gap effect, defined as the region where the poloidal field distortion is not lower than 10% of its maximum value, has been estimated to be 40° (20° on each side) and does not vary significantly throughout the discharge. A periodicity of order 24 is also present, possibly due to the effect on the plasma of the toroidal field ripple, which is produced in RFX by 48 coils.

In order to compare the field penetrated through the gaps with the vertical field diffused across the shell, the non-axisymmetric component of the vertical field has been evaluated in the gap region on the basis of the measurements of the outer and inner poloidal pickup coils. A measure of the average vertical field across the whole shell is then obtained by taking the difference between the signals of the outermost and the innermost poloidal flux loop. The results for the two fields, starting from the axisymmetry condition instant, are shown in fig.4 (shot 1860). The field through the gaps accounts for a substantial part of the total field; the difference, which is growing in time, can be associated with the decay of the shell equilibrium currents.

**Evaluation of the local plasma shift near the gaps**

The local horizontal plasma shift at the gap has been estimated on the basis of a first order approximation of the equilibrium magnetic configuration /6/. The m=1 B_\parallel cosine component has been calculated as the difference between the measurements of the inner and outer pickup coils closest to the gap, while the m=1 B_\parallel sine component has been derived from the measurements of the saddle coils across the gaps. These signals are measured on the shell inner surface, where the fields present maximum amplitude, therefore this is the highest estimate of the horizontal shift.

In fig.5 the shifts for the same pulses shown in fig.1 are presented. As expected, in shot 1774 the plasma column is significantly shifted outwards due to the insufficient equilibrium field, while the inward shift of shot 1853 is consistent with an excess of external equilibrium field. In
shot 1860, characterized by a better programming of the FSW amplifiers, the local shift turned out to be much smaller, even if larger than the average plasma column shift, about 1.2 cm for this discharge. The slight equilibrium field defect still occurring at the gaps accounts for this discrepancy.

The average radial field, measured by the saddle coils at the gaps, is shown for a typical shot in fig.6 (shot 1823). A remarkable feature of this graph is the different behaviour observed at the two gaps. This is confirmed by fig.7, where time averages of the absolute values of the relative radial fields at the two gaps are correlated. The data show that the radial field can be larger than 1% either at one gap or at the other. The radial field can be related to a local vertical shift of the plasma column, but the relationship depends on the plasma horizontal shift.

In the shots where either gap exhibits a high radial field, evidence of enhanced local interaction with the first wall at the same toroidal position is given by CCD camera recordings.

Conclusions
Analysis of RFX shots confirmed the importance of an accurate control of the field errors at the gaps to improve discharge performance. The presence of the poloidal gaps has proven to be a major cause of non axisymmetric behaviour of the plasma column. A reduction of the local horizontal shift at the gaps can be obtained with an effective programming of the FSW amplifiers.

References
/5/S.Martini et al., this conference, 3/25.
IMPURITY BEHAVIOUR IN RFX

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Introduction
Impurity line emission has been routinely measured in RFX in the spectral range 20-7000Å by means of three vacuum spectrometers and a multichord visible one, in order to monitor the impurity content in the discharge under different operational conditions. The simulation of the experimental data by a 1-dimensional model has allowed the evaluation of the plasma effective charge and of the impurity diffusion.

Experimental results.
Carbon lines dominate the RFX spectrum: in fact the RFX inconel vacuum vessel is entirely covered by graphite tiles, in order to minimise the influx of metallic high-Z, highly radiative impurities. Indeed no metal lines have been detected; high values of the effective charge may anyway be reached, due to the Carbon high content in the discharge (typically about 5-7% of the electron density). Lines from Oxygen (about 1-2% of electron density), Nitrogen and Helium are also observed, the latter being present during some pulses after the He glow discharges which generally follow the cleaning procedures in Hydrogen. The emission spectrum measured in a 650 kA pulse in the range 100-1000Å by a flat-field survey spectrometer is shown in fig.1. The resonant lines from C V (40.26Å) and C VI (33.7Å) and the O V (151.5Å) and O VI (150.1Å) are currently monitored by a high resolution grazing incidence spectrometer [1]. For C V and C VI the absolute content has also been evaluated, the grazing incidence spectrometer being absolutely calibrated at 40Å via the branching ratio technique. To infer the ion content from the absolute line intensity it is necessary to know the radial position of the ion and the corresponding electron density and temperature. Space-resolved measurements of C V (2271Å) and C VI (5291Å) have been performed by a 9-chord spectrometer, showing that both these ions emit essentially from a rather external plasma shell. The electron temperature of C V has been measured by the line intensity ratio of the resonant 40.26Å and intercombination 40.73Å lines [2], while the interferometric measurements show a very flat density profile.

The behaviour of O V+O VI (arbitrary units) and C V+C VI (absolute values) as a function of the electron density during the plasma current flat-top phase is shown in fig. 2. The Oxygen ion concentration
increases with the electron density, indicating that the Oxygen production mechanism might be associated to chemical sputtering [3]. The correlation of Carbon with the electron density is much less clear and this behaviour suggests that for Carbon production also other mechanisms must be considered, i.e. self sputtering, physical sputtering or local phenomena. Strong local plasma-wall interaction phenomena have been observed by two CCD cameras equipped with a C I interference filter. The strongest interaction is generally toroidally localised at the region where the largest magnetic field perturbations exist due to the presence of one of the two poloidal gaps in the aluminium shell surrounding the vacuum vessel[4], and a toroidal asymmetry is related to the toroidal field ripple. Also poloidal asymmetries are observed also by the CCD cameras and by the multichord spectrometer, revealing a higher emission level from the chords viewing the external region of the poloidal section of the plasma. A stronger interaction with the external part of the vacuum vessel is due to the external shift of the plasma column.

Carbon densities measured at two different plasma current levels are shown in fig.3; at 550kA and 680kA the Carbon density is almost the same, but the contribution of the stripped C VII is not considered. On the other hand the electron temperatures in these discharges are very similar and therefore the contribution of C VII may be considered to be the same. In terms of effective charge, higher current discharges have higher density, and therefore the Carbon Z_{eff} is lower. However Carbon density shows a slightly increasing behaviour with the fraction of ohmic power dissipated during the start-up phase of the discharge. This dependence will be better investigated in the near future by performing pulses with a slow current ramping.

Impurity transport simulation.

To evaluate the effective charge it is necessary to know the contribution of the completely stripped ions (C VII and O IX) that are not directly measured. They have been estimated by a 1-dimensional impurity collisional-radiative model [5], where the diffusion fluxes are varied in order to simulate the experimental time behaviour of Carbon and Oxygen ions and their measured radial profiles. Input to the model are the electron temperature and density profiles and the impurity influxes. Experimental estimates of the temperature and density profiles are available [2], but a reliable measurement of the Carbon influx is very difficult: in fact, due to the relevant plasma-wall interaction asymmetries already discussed, a local measurement is not representative of the average influx. Therefore the impurity influx value has been varied up to the best simulation of the measured
absolute line intensities; in particular C V 2271Å and O VII 1623Å - measured by a vacuum Czerny-Turner spectrometer absolutely calibrated by means of a D2 lamp - have been used as reference lines. Carbon and Oxygen ion distributions obtained from the simulation of the 550kA discharges are shown in fig.4 and 5, respectively. The figures refer to a time during the plasma current flat-top. The corresponding radial profiles of C V 2271Å and C VI 5291Å lines are in good agreement with the experimental inverted profiles; as an example, in fig.6 the calculated and inverted C V profiles are shown. The impurity diffusion fluxes used for this simulation are the classical fluxes enhanced by a factor which increases with the radius from 5 on the axis to 100 at the wall, corresponding to a diffusion coefficient in the external region of about 10-30 m²/sec. The corresponding effective charge and plasma dilution profiles are drawn in fig.7. On axis Z_{eff} values of about 5 are obtained, corresponding to a plasma dilution of about 0.25. Slightly lower Z_{eff} values (~4 on the axis) with a similar profile are obtained for higher current discharges: in fact the electron density is higher, while the impurity content and the ionization degree are very similar, with only small differences in the ion distribution.

References
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PLASMA HEATING STUDIES IN RFX

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Electron and ion temperature measurements performed during the initial operational phase of RFX\(^1\) are presented. RFX is a large (a=.457m, R=2m) Reversed Field Pinch machine (RFP) whose aim is to verify at high plasma current - up to 2 MA - the validity of the scaling laws found so far on smaller experiments. The first wall is made up with more than 2000 graphite tiles which protect almost entirely the inconel vacuum vessel. A 6 mm thick aluminium shell, presenting two poloidal and two toroidal insulated gaps, surrounds the vacuum vessel for stability purposes. Since RFP's could reach ignition with ohmic heating only, no additional heating systems are installed on RFX.

The first period of operation was spent experimenting a variety of operational modes and discharge parameter regimes. This entailed a large variation of the plasma characteristics corresponding to different levels of pulse optimization.

The electron temperature measurements relied on the x-ray pulse height analysis carried out by means of four Si(Li) detectors\(^2\) organized in two systems, one single and one triple, placed 15° apart on the equatorial plane. The triple system was also able to be tilted of +/- 18° in order to look at different plasma positions up to 20 cm from the center. Particular care was taken in examining the x-ray spectra to avoid pile-up distortions and to check the absence of impurity lines. Generally a 125 µm thick Be filter was used and the temperature was evaluated from the slope of the spectrum in the 1.5 - 2.6 keV range. When the photon flux per diode was adequate (20-30 kHz) a temperature value was provided every 10-15 ms. Fig.1 reports the central electron temperatures as a function of I/N, N being the linear density. \(Te\) increases with I/N but the variety of discharge conditions reflects in a data spread. The temperature values range between 200 and 450 eV and refer to the time comprised between 30 and 50 ms during the shot. However in many discharges the photon flux was such that the average had to be taken over the whole pulse. It should also be
mentioned that the Si(Li) detectors were placed 15° and 30° respectively from one of the two poloidal insulating gaps of the shell, where large field errors may be present during the shot and where indeed strong plasma wall interactions together with higher magnetic fluctuation amplitudes are observed. In fact the x-ray spectra from the single system, which is lacking of a beam dump on the opposite side of the vessel, were sometimes affected by distortions which were interpreted as due to processes related to plasma wall interactions and as such rejected. Subgrouping the data according to plasma current and experimental conditions, the trend of $T_e$ increasing with $I/N$ is sometime more clear, but at the expense of a reduced statistics.

The technique based on the impurity line intensity ratio was also used to provide further $T_e$ points along the minor radius. The intercombination to resonant line ratio$^3$ of C V was routinely monitored, while the same ratio for O VII and the I(215)/(I)(193) ratio$^4$ of O V was sometime measured. The ion temperature was measured by two neutral particle analyzers (NPA) - a time-of-flight and a charge exchange system - and by the Doppler broadened impurity lines of O VII (1623 Å -II order), CVI (5291 Å), C V (2271 Å), C IV (1548 Å - II order), O V (1371Å - II order) and C III (2296 Å) detected by a vacuum high resolution Czerny-Turner spectrometer$^5$. Equipartition times between impurities and Hydrogen in the temperature range characteristic of RFX are lower than the energy confinement time and therefore impurity and Hydrogen temperatures are comparable. The pertaining positions along the radius of the spectroscopically determined ion and electron temperatures were deduced by inverting multichord observations and making use of a 1-D impurity transport code$^6$. Also the NPA data required a high degree of modelling and a Montecarlo code was used to deconvolve the opacity effects on the escaping neutrals; the result being the determination of the plasma radius to be associated to the ion temperature value directly computed from the slope of the energy spectrum and an extrapolation of the on axis one.

In Figs 2 and 3 the ion and electron temperature profiles are shown for a 550 kA and a 680 kA discharge respectively. These profiles are the result of an averaging process, being the spectroscopic data taken over different pulses. Electron and ion temperatures show to be comparable, thus suggesting the presence of an anomalous ion heating process; in fact the energy equipartition times with electrons (5 ms for H and about 4 ms for C in steady state) exceed several times the global
energy confinement time which is approximately 1 ms. It is worth mentioning that in particular circumstances, characterized by high a Θ value during the reversal phase - Θ = B0(a)/<Bz> - ion temperatures up to 800 eV have been observed. The 550 kA temperature profiles are fittable with a function of the type 1-(r/a)α with α between 2 and 3, while α = 4 appears quite suitable for the 680 kA ones. It appears that for the two reported plasma conditions the effect of increasing the current was mainly that of flattening the temperature profile rather than increasing the central value.

Reporting nₑTₑ(0) versus I, as in Fig.4, shows how the kinetic energy content increases with current. Within the experimental uncertainties the rate of the plasma energy enhancement with current is enough to keep β approximately constant as shown in Fig.5. In the poloidal β computation - β₀ is the ratio between kinetic and magnetic pressure in the poloidal section - the experimental profiles were considered. The electron density profile measured by the eight chord interferometer was included, with a radial dependence of the type 1-(r/a)α with α = 6 for both the cases. The constancy of beta with the current in these type of discharges is however maintained at the expense of the central energy confinement: the flattening of the temperature profile is in fact an evidence of a central thermal conductivity enhancement. Qualitatively this is in agreement with the observation at the edge of an enhanced energy flux asymmetry between electron and ion drift sides, which is currently interpreted as an increase in the energy flux of fast electrons escaping the central and hotter plasma core due to their high mobility and to the stochasticity of the magnetic configuration.

References

1 L. Fellin et al. Proc. 17th SOFT Rome (1992), paper V15
4 T. Kato, J. Lang, K.E. Berrington, Atomic and Nucl. Data Tab., 44(1990) 133
6 L. Carraro et al., “Impurity behaviour in RFX”, this conference
7 V. Antoni et al. “Characterization of the RFX edge plasma”, this conference
Optimization of RFP formation and sustainment in RFX


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Introduction

In the first months of 1993, the RFX experiment (R=2m, a=0.46 m) [1] has operated at reduced volt-second (6 V·s out of 15) to study the formation and sustainment of the RFP in a relatively safer power input regime, before increasing the parameters to reach the design value of 2MA plasma current.

At present the RFP configuration is obtained, similarly to the ETA-BETA II experiment [2], in the aided mode (see fig.1): a capacitor bank is discharged into the toroidal winding to produce an initial toroidal flux, \( \Phi_T \), then the plasma current, \( I_T \), is induced by varying the poloidal flux stored in the magnetising winding [1]; the free oscillation of the toroidal circuit continues and \( \Phi_T \) decays during the initial plasma current rise, until the toroidal field at the wall, \( B_T(a) \), reverses and the toroidal circuit is crow-barred. The overall performance of the plasma during the RFP sustainment phase is strongly influenced by the control performed on density, toroidal field and plasma position during the formation phase.

As soon as the RFP is obtained, a clear improvement of confinement is seen and the plasma current increases again until the applied toroidal voltage \( V_T \), which decreases exponentially, no longer matches the resistive drop. In RFX it is also possible to insert a pre-programmed flat-top power amplifier by which \( V_T \) can be sustained and controlled in the range (0-60 V). In this way quasi-steady RFP current flat-top phases lasting \( 90 \) ms can be obtained which terminate only when the amplifiers are switched off and \( V_T \) is no longer sustained (see fig.1 pulse 1860).

On the other hand, the RFX vacuum vessel is presently surrounded by a stabilizing aluminium shell with a time constant for diffusion of horizontal magnetic field of 0.4 s; hence, during the \( 60-90 \) ms of the RFP phase, toroidal equilibrium is determined by the shell and by a steady state vertical field (bias \( B_v \)) applied previously to the start-up by a set of toroidal coils (Field Shaping Winding, FSW). The FSW is also used to compensate the total ampere-turn produced by the plasma, to match the vertical field at the shell gaps and to control the equilibrium at late times of the pulse. Therefore the optimum performance during the RFP phase is obtained by a proper combination of bias \( B_v \) and dynamic FSW programming.

Optimization of density behaviour

Active density control is performed by a pre-programmed set of piezoelectric and electromagnetic valves with a maximum capability of 900 and 7000 mbar-l/s respectively. Since operation at high wall temperature (\( T_{wall}<350 \) C) has not yet been possible, the full graphite tiles armour behaves a large H reservoir which determines the density behaviour through the discharge in most cases. In fact a single pulse sustained by gas puffing at a density of \( 3\times10^{19} \) m\(^{-3} \), i.e. in the optimum I/N range of a few \( 10^{14} \) A·m\(^{-3} \) (I/N = plasma current / line density), is sufficient to load the wall with an amount of H which will keep the density in the same range for the next pulse with no additional feed except the initial filling gas. Operation at low density can be obtained only after the wall has been 'emptied' by He glow discharge cleaning or by a series of pulses performed with little filling gas only, whereas a moderate gas puff during the initial phase allows to operate at medium or high density with I/N approximately constant through the pulse.
A strong gas feed or a substantial influx from the wall (recycling >1) has been found to be very important in the RFP formation phase. This is shown in fig. 2 where 3 discharges obtained with the same poloidal and toroidal field programming and different density control are compared. In pulse 1429 the recycling was low and the RFP could be achieved only with a high filling pressure, in pulse 1554 a low filling pressure and a higher recycling allowed a better performance and, finally, in pulse 1600, a moderate gas puff in the setting-up phase and high recycling allowed to obtain the highest current.

The magnetic configurations of the above cases are very similar if compared in terms of $F - \Theta$ diagram ($F \equiv B_T(a)/\langle B_T \rangle$, $\Theta \equiv B_\Theta(a)/\langle B_T \rangle$), although the time of $B_T(a)$ reversal is slightly different. Therefore the observed behaviours cannot be explained by differences in the current density profile, but rather indicate a lower effective plasma resistivity for the higher density cases, as shown in fig.2 by the resistivity on axis computed from a global helicity balance [4], $\eta_K(0)$. Only at later times, when $V_T$ decreases in the 40 V range, the pulses at lower density show lower resistivity. Conversely the electron temperature, $T_e$, is higher at lower density [5] (370 for pulse 1554 and 280 eV for 1600) even soon after the $B_T$ reversal time. This suggests that both in the RFP formation and in the early current ramping phase, when the input power is high, the impurity influx from the wall at low density causes an increase in $Z_{eff}$ which more than compensates for the higher $T_e$ in terms of electrical resistivity. In fact the resistivity anomaly factor defined as $Z^* \kappa = \eta_K(0)/\eta_{cl}$ (where $\eta_{cl}$ is the classical Spitzer resistivity computed with $Z=1$ from the $T_e$ values) is ≈5 for pulse 1554 and ≈3 for 1600.

Fig. 1: Comparison between two pulses with (1860) and without (1733) sustainment of the $V_T$ by the flat-top amplifiers.

Fig. 2: Comparison between three pulses with different density control.
Study on toroidal field circuit control during setting-up

The effect of the control of the toroidal circuit on the RFP formation phase has been studied by varying the period and the amplitude of the $\Phi_T$ oscillation. Comparing cases with different $B_T(a)$ amplitudes and otherwise similar parameters (see fig.3) it is seen that the initial slope of the plasma current, $dI/dt$, is independent from the value of $B_T$. Indeed $dI/dt$, except for the breakdown delay which is influenced by the initial toroidal field and filling gas pressure, depends in general only on the applied $V_T$. On the other hand, at higher $B_T$ the critical values of the safety factor $q(a)$ (at rational values in the range 0.1-0.4, see also [1]) are reached later, allowing a higher current to be achieved before the plasma resistance increases (see fig.3 pulse 1866). As a consequence, since the period of the toroidal circuit is the same, $B_T(a)$ then goes to zero in a shorter time, thus reducing the losses connected to the highly turbulent reversal phase. Conversely, if the period of the toroidal circuit is increased, although higher currents are initially achieved, then the longer field reversal phase causes larger losses and the RFP may not even be reached. The difference with the MST experiment [6], where the RFP is obtained in the self-reversal mode, could be related to the relatively high distance between the toroidal windings ($r=0.62$ m) and the plasma, similarly to the ETA-BETA II experiment [2].

The influence of $B_V$ in the setting-up phase is shown in fig. 4, where two cases with zero (pulse 1870) and with 6 mT (pulse 1865) vertical field are compared. Again no influence on plasma current slope after breakdown was found. The higher current reached in pulse 1865 is due to the better plasma centering relative to the vessel which entails a larger plasma cross-section and lower plasma-wall interaction.

Control of equilibrium and field errors in flat-topped discharges

The bias vertical field improves the performance of the experiment also during the current flat-top phase as already observed in RFX [1] and in other RFPs [7,8]. For the present current level of 0.5-0.7 MA, a broad optimum range of $B_V$ is found between 5 and 8 mT. As longer pulses are obtained using the flat-top poloidal amplifier, it becomes more and more important to control the currents in the FSW to limit the magnetic energy stored outside the plasma and to control the field errors at the gaps [9]. By properly preprogramming the FSW amplifiers a compensation of the plasma ampere-turn within 3% and a reduction of the vertical field errors at the shell gaps within 10% during the flat-top have been achieved as shown in fig. 5 (pulse 2074). When the field errors are well controlled the mean horizontal shift of the plasma $\Delta_H$ remains also close to the low value due to the bias $B_V$. For comparison an outward drift of the $\Delta_H$ and oscillations of the mean vertical shift $\Delta_V$ observed in the presence of high field errors is shown in fig. 5 (pulse 2077). The control of the field errors and of the associated enhanced plasma-wall interaction in the gap regions causes also a reduction of the resistive loop voltage as shown in fig.6.
Conclusions

The study of the control on density, toroidal field programming, toroidal equilibrium and field errors at the gap led to improve the overall plasma performance in RFX. Higher currents and lower plasma resistance are obtained at the end of the RFP formation with the same applied $V_T$ by optimizing the control of density, toroidal field and steady state vertical field. Quasi-steady state pulses in terms of plasma current, $F$ and $\Theta$ parameters, electron density, temperature, $B_\Theta$ and $T_E$ can be maintained for times of the order of 0.1 s as long as enough $V$'s are available when the proper control procedures are applied also in the RFP phase.

REFERENCES


Fig. 4 Plasma current and safety factor in the setting-up phase with $B_\psi = 6$ mT (1865) and without $B_\psi$ (1870)

Fig. 5 Comparison between two pulses with low (2074) and high (2077) field errors at the shell gaps

Fig. 6 Toroidal voltage at maximum $I_T$ versus normalized average field error at the gaps for a set of pulses with $I/N = 2.5 \times 10^{-14}$ A·m
HYDROGEN RECYCLING AND IMPURITY PRODUCTION IN RFX

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Introduction The performance of the RFX experiment [1], whose first wall is almost entirely covered by graphite tiles, is found to be affected by strong hydrogen recycling and light impurities (mainly carbon and oxygen) content. In this paper preliminary analysis of the impurity production mechanisms is presented.

Experimental set up and methods H\textsubscript{\alpha} emission has been currently monitored over 9 vertical chords on the same poloidal section. Uncalibrated CCD cameras with CI (9077Å) interference filters have been used to look at the first wall with the aim of monitoring also heating effects on the graphite tiles. Two of them have been placed on the same equatorial port looking along the two opposite toroidal directions. A third one has monitored the graphite tiles at the bottom of the vessel together with a Ge pyrometer for surface temperature measurements. Single Langmuir probes have been negatively polarized up to 150V to detect the ion saturation current, whereas heat sensors mounted on an horizontally insertable limiter have measured the surface temperature. Silicon samples have been also exposed to the plasma on the shadow of the graphite tiles and then analyzed by nuclear techniques.

Results CCD cameras have shown strong toroidal asymmetries which have been found to occur preferentially in a region close to one of the insulated gaps of the aluminium shell. Power loading effects have been found on some tiles of this region but more often on the tiles close to the bottom or top diagnostic ports. These effects appear as strong emission on the electron drift side of the border of the tiles lasting long after the pulse. The corresponding surface temperature cannot presently be specified, but certainly exceeds 600°C, which is the lower typical threshold for front illuminated Silicon CCD's. The pyrometer (which views vertically) and the energy sensors (see fig.1) show a wall temperature increase 10 times larger than what would be expected (ΔT=20°C) for a uniformly distributed input power. Toroidal periodicity associated to the toroidal field ripple has been observed: it is more evident during the RFP formation phase, i.e. when the toroidal field reverses at the edge, but may be observed also during the current flat top. Poloidal asymmetries are also observed, mainly associated to an external shift of the plasma column and in many situations a vertical asymmetry reveal that the plasma is drifting downwards. In particular the H\textsubscript{\alpha} data reveal that differences of an order of magnitude in the influx may be
present between internal and external chords (see fig.2). Spectroscopic measurements show that the content of oxygen and carbon increase with density, though C behaviour is less evident /2/. Electron temperature at the edge is found to be constant with line average density whereas the edge density tends to increase more than linearly /3/.

Discussion The recycling factor $R$ can be conveniently expressed by the simple relationship:

$$ R = \frac{\Gamma_{\text{Hin}}}{\Gamma_{\text{Hout}}} = 1 - \frac{\tau_p}{\tau^*} $$

where $\tau_p$ is the hydrogen confinement time and $\tau^*$ is the characteristic time of density variation, $\Gamma_{\text{Hin}}$ and $\Gamma_{\text{Hout}}$ are the H influx and outflux respectively. In RFX $\tau_p = 1 \text{ ms}$ /1/ and, when density is not sustained by puffing during the discharge, $\tau^*$ is of the order of several 10 ms. In this case $\Gamma_{\text{Hout}}$ can be reasonably identified with $\Gamma_{\text{Hin}}$ as given by $H_\beta$ intensity, so that $\Gamma_{\text{Hin}} = \Gamma_{\text{Hout}} = \Gamma_H$. This conclusion agrees with the consideration that the measured hydrogen flux to the RFX wall, $\Gamma_H = 10^{22} \text{ m}^{-2}\text{s}^{-1}$, is large enough to give saturation of the carbon surface even in a few discharges with pulse length of $\approx 0.1 \text{ s}$. In fact RFX operates mainly with an H loaded wall, making a satisfactorily density control rather difficult to achieve. Electron density is controlled to some extent by using gas puffing with fast piezoelectric valves. The particle outflux $\Gamma_{\text{out}}$ is deduced by the ion saturation current measured by Langmuir probes. In fig.3 the time evolution of $\Gamma_H$ and $\Gamma_{\text{out}}$ are reported for a shot where a clear correlation is observed. It should be mentioned that the two measurements refer to different toroidal positions and therefore the ratio $\Gamma_{\text{out}}/\Gamma_H$ is affected by toroidal asymmetries. However generally $\Gamma_{\text{out}}/\Gamma_H$ is found to increase with $n_e$ as it can be deduced in fig.4, in agreement with what is found for impurities /2/.

Carbon density in the bulk plasma, $n_C$, in steady state can be related to the hydrogen density $n_H$ by the simple equation:

$$ \frac{n_C}{n_H} = \frac{\tau_c}{\tau_p} \eta \left( Y_H + Y_C + \frac{\Gamma_C}{\Gamma_H} Y_{CC} + \frac{\Gamma_{ox}}{\Gamma_H} Y_{ox} \right) $$

where $\tau_p$ and $\tau_c$ are the confinement times of hydrogen and carbon ions respectively. The factor $\eta$ describes the ability of the carbon neutrals released from the wall to penetrate the plasma. $\Gamma_{ox}$ and $\Gamma_C$ are the flux densities of oxygen and carbon ions, $Y_H$ and $Y_{CC}$ are the physical sputtering yields for hydrogen on carbon and for carbon self-sputtering, $Y_C$ and $Y_{ox}$ are the yields for chemical reaction involving production of hydrocarbons or CO respectively. A significant contribution of radiation enhanced sublimation and thermal sublimation, occurring at temperatures greater than 1200 °C, can be excluded, since only seldom wall temperature in correspondent of hot spots exceeds 600 °C. The ratios $\Gamma_{ox}/\Gamma_H$ and $\Gamma_C/\Gamma_H$ have been therefore evaluated by the relative concentrations measured on silicon samples
exposed to the plasma. The values found, as averages in 10, 30 and 50 discharges, are respectively $\Gamma_\text{ox}/\Gamma_\text{H} = 3\%$ and $\Gamma_\text{C}/\Gamma_\text{H} = 30\%$.

Physical sputtering yields are estimated with the formulas given in ref. /4/ and with ion energies corresponding to $T_\text{e} = 10$ eV /3/, assuming that carbon ions return to the wall with charge state 3. By using enhancement factors to account for energy distribution, assumed maxwellian, and isotropic angular distribution, they can be estimated as $Y_\text{H} = 1.5\%$ and $Y_\text{CC} = 25\%$. The chemical sputtering of low energy (=50 eV) hydrogen ions is reported to achieve a maximum of $Y_\text{C} = 4\%$ at wall temperatures in the range 500 to 800 °C for high fluxes /5/, thus allowing only some local contribution to carbon erosion. An additional contribution of oxygen to chemical sputtering is also expected, since values of $Y_\text{OX}$ up to 1 have been quoted /6/.

The global carbon yield, corresponding to the term within brackets in eq.(2) can be estimated to be in the range 10-16%. This figure is consistent with the relative carbon concentration $n_\text{C}/n_\text{e} <10\%$ experimentally found by spectroscopy /2/, thus a value of $\eta \tau_\text{P}/\tau_\text{e}=1$ is expected, i.e. no significant shielding of carbon neutrals released from the wall should occur if $\tau_\text{P}/\tau_\text{e}=1$, consistently with numerical simulations /7/.

Among the impurity production mechanisms the dominant one may be then identified in the C self—sputtering, which accounts for more than 50% of the global yield. This is consistent with little dependence on the electron density found for the edge electron temperature and for the wall temperature. The weak dependence of C content on $n_\text{e}$ may be however associated to the O chemical sputtering since O is found to increase with density /2/.

References
/2/ L. Carraro et al., Impurity Behaviour in RFX, This Conference Proceedings.
/3/ V. Antoni et al., Characterization of the RFX Edge Plasma, This Conference Proceedings.
/6/ J. Roth, in Data Compendium for Plasma-Surface Interaction, Nuclear Fusion special issue, 1984, p. 72.
Confinement studies of high current density reversed-field pinch plasmas in Extrap-T1

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The experiments described in this paper were carried out on the Extrap-T1 device [1] after modifications for RFP operation. The device has a high aspect ratio $R_0/a = 0.5$ m/0.057 m = 8.8. In spite of the small cross-sectional area, discharge currents of up to 100 kA are achieved, with current densities of about 40 MA/m$^2$ on axis. For these discharges, the line average electron density is above $0.5 \cdot 10^{20}$ m$^{-3}$, and $I/N$ is maintained less than $2 \cdot 10^{-18}$ Am. A typical discharge is shown in Fig. 1.

We have performed two series of discharges, one series with $\Theta$ fixed around 1.7 and $I_\phi$ varied in the range 40 kA to 100 kA, and another series with $I_\phi$ fixed around 55 kA with $\Theta$ varied from 1.6 to 2.6. The variation of $\Theta$ was achieved by adjusting the level of aided reversal. The measurements were performed during flat-top $I_\phi$ and $\Theta$ conditions. For each working point in the $I_\phi - \Theta$ plane, 20 consecutive discharges were performed while maintaining nearly identical conditions within each series, such as wall conditioning, base vacuum and filling pressure. In the following, we study the ensemble average of the parameters at each working point. All calculated quantities were evaluated for each individual discharge, and the final results were then averaged for each ensemble.

In addition to the controlled variation of the discharge parameters $I_\phi$ and $\Theta$, there was a variation in the parameter $T_e/n_e$ within the current scan. The data falls into two groups, which are referred to as "high $T_e/n_e$" series and "low $T_e/n_e$" series in the following. The cause of the variation is related to wall conditioning history, which included glow discharge cleaning, experimental run history and base vacuum conditions.

The study wherein discharge current was varied with constant $\Theta$ is presented first. In Fig. 2 a summary of both observed and calculated parameters is shown. We can distinguish clear trends within each of the two series. For both high $T_e/n_e$ and low $T_e/n_e$ series, the electron temperature and electron density increase with current. The scaling within the high $T_e/n_e$ series is $2.4$ eV/kA. $Z_{eff}$ is roughly constant for all groups of discharges, yet somewhat lower for the low $T_e/n_e$ discharges. The radiated power amounts to 15–30% of the Spitzer input power, and this ratio is in general larger for the
Figure 1: Example of pulse data for discharge #25954 ($I_\phi \approx 90$ kA, $\Theta \approx 1.7$) From top to bottom: Plasma current $I_\phi$. Total loop voltage $U_{tot}$. Average toroidal field $\langle B_\phi \rangle$ and edge toroidal field $B_\phi(a)$. Pinch parameter $\Theta = B_\phi(a)/\langle B_\phi \rangle$ and field reversal ratio $F = B_\phi(a)/\langle B_\phi \rangle$. Line average electron density $n_e$ measured by interferometry. Hydrogen influx $\Gamma$ derived from Ha radiance.

low $T_e/n_e$ series. The fraction of Spitzer input power $P_S/P$ decreases with current within the high $T_e/n_e$ series, and is significantly higher in the low $T_e/n_e$ series. At constant $\Theta \approx 1.7$ the fluctuation level $\delta B/B$ increases with current. An important observation is that $\delta B/B$ is lower for all currents within the low $T_e/n_e$ series. In addition, the non-Spitzer loop voltage is correlated with the fluctuation level. Again there is a clear distinction between the two series. The high $T_e/n_e$ series exhibits a significantly larger non-Spitzer component of the loop voltage for all currents. $\beta_\phi$ decreases with current from 13 % at 40 kA to about 8 % at 95 kA. Here both series have almost the same scaling with $I_\phi$. We note that the particle confinement time is significantly lower for the low $T_e/n_e$ series, while the energy confinement time is slightly higher. It should be noted, that the discharges with low $T_e/n_e$ at 90 kA have a $+20$ % inductive component of the total loop voltage, compared to the other data where the inductive component was in the range $\pm10$ %.

In Fig. 3 the results for the discharges with varied pinch parameter at constant plasma current and low $T_e/n_e$ are presented. The temperature shows no clear trend with $\Theta$. The density however increases with $\Theta$. $Z_{eff}$ remains at 2 over the range of
Figure 2: Overview of results obtained at fixed pinch parameter, $\Theta \approx 1.7$, as a function of plasma current. The different symbols indicate low $T_e/n_e$ ($\Box$) and high $T_e/n_e$ discharges ($\bullet$). Left from top to bottom: Line average electron temperature $T_e$ from VUV spectroscopy. Line average electron density $n_e$ from interferometry. $Z_{\text{eff}}$ from VUV spectroscopy. Ratio of radiated power to Spitzer input power $P_{\text{rad}}/P_S$. Center from top to bottom: Fraction of Spitzer input power $P_S/P$. Non-Spitzer loop voltage $I_\phi(R-R_S)$. Edge toroidal field fluctuation level. Right from top to bottom: Poloidal beta. Particle confinement time $\tau_p$ (above the horizontal line) and energy confinement time $\tau_E$ (below the horizontal line).

The ratio of radiated power to Spitzer input power remains in the range 15–30%, and the fraction of Spitzer input power $P_S/P$ decreases with $\Theta$. The most striking trends within the $\Theta$-scan are the increase of the fluctuation level $\delta B/B$, which reflects enhanced tearing mode activity, and the equally clear increase of the non-Spitzer component of the loop voltage. At $I_\phi \approx 55$ kA the fluctuation level increases from $\leq 1\%$ at low $\Theta$ to $\geq 2\%$ at high $\Theta$, whereas $I_\phi(R-R_S)$ increases from 25 V to 100 V. In spite of the favourable scaling of $\beta_\theta$, which increases from $11\%$ at $\Theta \approx 1.7$ to $\geq 20\%$ at $\Theta \approx 2.5$, $\tau_E$ is degraded due to the significant increase of the non-Spitzer component of the loop voltage. The particle confinement time, on the other hand, shows an improvement with $\Theta$.

Three observations are common to both the current scan and the $\Theta$-scan. First, the plasma pressure appears to be mainly a function of $I_\phi$ and $\Theta$, since the result for $\beta_\theta$ is essentially unaffected by the change of $T_e/n_e$. Second, the series with high $T_e/n_e$ ratio has a higher level of non-Spitzer input power. As a result, the total input power and
Figure 3: Overview of results obtained at fixed plasma current, $I_0 \approx 55$ kA, as a function of pinch parameter.

The energy confinement time are comparable for both high $T_e/n_e$ and low $T_e/n_e$ series for a given plasma current. The magnetic fluctuation level correlates well with this non-Spitzer loop voltage within current scan and $\Theta$-scan and for different $T_e/n_e$ ratios. Third, the particle confinement time is anticorrelated with the energy confinement time, which is observed in the $\Theta$-scan and also by a comparison of the two current-scan series with different $T_e/n_e$. This is not an artifact of the Hα analysis procedure, since the dependence on electron density of the measure of ionisations per photon is considered in the analysis. The simultaneous increase in particle confinement time and magnetic fluctuation level is a notable result.

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References

PROBE MEASUREMENTS ON THE EDGE PLASMA OF RFP DISCHARGES IN EXTRAP T1.

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Introduction. The T1 device is used for experiments on reversed field pinch and EXTRAP discharges. We have reported previously on measurements with passive probes in RFP discharges with 40 kA plasma current[1,2]. A major objective has been to provide data on plasma edge conditions in order to improve the understanding of plasma-surface interactions. In this contribution we present data which are extended up to 90 kA plasma current, together with additional measurements with Langmuir probes and heat flux probes.

Experimental. The dimensions of T1 are R/a = 0.5 m/0.057 m. The wall consists of 316L stainless steel bellows and is operated at ambient temperature. The inner radius of the bellows is 57 mm, there are six fully poloidal constrictions at 56 mm and the 18 mm wide portholes are welded at r = 59 mm. Edge probe measurements have been performed in RFP discharges in hydrogen and deuterium with plasma current from 40 to 90 kA and a pinch parameter θ ≈ 1.7. The 40 kA discharges in deuterium have already been discussed in detail[1]. Typical time evolutions of plasma current and density in 90 kA discharges in hydrogen are shown in figure 1. Reversal is maintained roughly from 0.1 to 0.45 ms. The passive probes and the geometry of the wall structure have been described before[1,2]. In a new series of 90 kA discharges in deuterium, passive graphite probes have been exposed in similar manner. For the current and heat flux measurements we have used a probe head which has previously been used on ETA-BETA II [3]. It consists of three tungsten pins which are mounted in graphite protection in such a way that they collect particles arriving mainly in directions tangential to the wall, a pair of pins collecting in the same direction and a single pin in the opposite direction, all three at the same minor radius. The collection area of each pin is 1.2 mm². In the experiments described here, all probes were movable along a minor radius at 20° below the outer midplane. To sample current and heat flux in different directions, the active probe head was turned around its axis between discharges. A number of different operation modes have been employed. The two probe elements facing the same direction have been used as a Langmuir double probe, or one of them has been used as a single probe, applying bias with respect to the vacuum vessel. Either the current has been measured throughout the discharge at constant applied voltage, or the voltage has been swept at 10 kHz using a Kepco Bp-1000 bipolar current amplifier. The average heat flux density throughout a discharge was calculated from the temperature excursion of the probe elements, measured with thermocouples. The Langmuir and heat flux measurements have been made in hydrogen plasmas. The central electron temperature from soft X-ray measurements and the line averaged density do not appear to have been significantly influenced when the probes were inserted, but there was a correlation of $H\alpha$ intensity at the same toroidal position with probe insertion depth.

Results. As in [1,2] the trapping of deuterium and stainless steel components was measured on graphite probes exposed at radii from r = 50 mm to r = 80 mm, both on the flat surface which had been exposed parallel to the wall and on the radially extended cylindrical sides. In the 40 kA discharges, the trapping rate of deuterium at
the wall position in short exposures, time averaged over complete discharges, was $4 \times 10^{22}$ m$^{-2}$ s$^{-1}$, and deuterium trapping saturated at $7 \times 10^{20}$ m$^{-2}$ on samples exposed to 200 discharges [1,2]. In the 90 kA discharges the corresponding numbers were found to be $6 \times 10^{22}$ m$^{-2}$ s$^{-1}$ and $1.7 \times 10^{21}$ m$^{-2}$, respectively. The collection rate of stainless steel components was 1-2% of the D trapping rate in the 40 kA discharges and roughly 3% in the 90 kA discharges. The trapping rates on the side of the probes which faces the ion drift direction are shown in figures 3 and 4.

A time trace of Langmuir double probe current at 60 V applied bias is shown at the bottom of figure 1. By averaging over $0.3 < t < 0.4$ ms and changing the applied bias between discharges, Langmuir characteristics can be constructed as shown at the top of figure 2. At the bottom of figure 2 a characteristic is shown which has been obtained by ramping the voltage throughout the same time interval in a different discharge. Both characteristics indicate an electron temperature $T_e \approx 12$ eV. Due to limitations of the bias supply, full ramped characteristics could only be obtained at large radii, where the plasma density was low. However, the slope of the characteristic around zero applied voltage could often be used together with ion saturation current data from operation in the stationary voltage mode, in order to estimate the electron temperature. This procedure results in $T_e = 10 - 15$ eV, fairly independently of plasma current and probe position in the range $50 < r < 60$ mm. Figures 3 and 4 show radial distributions in the edge region of poloidal heat flux $q_\theta$ and ion flux density $\Gamma_\theta = j_{sat}/e$ (thus presented under the assumption of negligible impurity current) in the ion- and electron drift directions for 40 kA and 90 kA discharges.

Discussion. We may assume that there is no immediate isotope dependence, so that the passive probe measurements in deuterium plasmas and the active probe measurements in hydrogen plasmas are directly comparable. In a preliminary analysis we may also assume that the perturbation of the edge plasma due to the presence of probes can be neglected.

In the 40 kA deuterium discharges an averaged radial flux $\Gamma_r \approx 7 \times 10^{22}$ deuterium m$^{-2}$ s$^{-1}$ with energies similar to a $kT_i \approx 100$ eV Maxwellian distribution were inferred from the low fluence trapping rate and the saturation level [1,2]. By the same arguments,
\[ \Gamma_r \approx 9 \cdot 10^{22} \text{ m}^{-2} \text{ s}^{-1} \] and energies corresponding to \( kT_i \approx 200 \text{ eV} \) can be derived for the 90 kA discharges. These radial fluxes can be compared to the value \( \Gamma_r = 1 - 2 \cdot 10^{23} \text{ m}^{-2} \text{ s}^{-1} \) which is obtained from \( H_\alpha \) in similar conditions [1,4], and the complications involved in such a comparison have been discussed elsewhere [1].

In agreement with observations on other RFP experiments [3,9] a high and asymmetric parallel heat flux is detected in the plasma edge, being strongly peaked in the electron drift direction, where values as high as 1 GW/m² occur in the 90 kA discharges. This phenomenon has been attributed to hot electrons which escape from the core plasma along stochastic field lines and may be indicative of the relaxation mechanism which dominates in RFP plasmas.

If the ion temperature is similar to the electron temperature, it is common practice [7] to calculate the local electron density from the ion saturation current density \( j_{sat} \) on a Langmuir probe as

\[ n_e = 2 \cdot j_{sat} / e \sqrt{2kT_e/m} \]  

If we apply this formula we obtain e.g. for the 90 kA discharges \( n_e \approx 8 \cdot 10^{18} \text{ m}^{-3} \) at \( r = 57 \text{ mm} \), about 10% of the line averaged density. Obviously if \( T_i \gg T_e \) the density can be lower.

Given \( n_e, T_e \) and typical ion energies at the edge it is possible to calculate the mean free paths for charge exchange and electron impact ionisation of neutrals leaving the wall, using rate coefficients from the literature [8]. Using the numbers stated for 90 kA discharges we get a total mean free path \( \lambda_n \approx 5 \text{ mm} \) for thermal hydrogen atoms, roughly 75% suffering charge exchange before being ionised by electron impact. Energetic neutrals due to reflection or ion impact desorption are likely to penetrate many cm through the plasma.

From the neutral path length and the arguments below about molecular reemission we see that the edge region is not a strong source of ions, hence if the radial particle loss from the plasma is taken as diffusive, we have \( \Gamma_r \approx -D_r \cdot dn_e/dr \). From figure 4 we can estimate \( |dn_e/dr| \approx 5 - 10 \cdot 10^{20} \text{ m}^{-4} \) at \( r = 51 \text{ mm} \), so that \( D_r \approx 100 - 200 \text{ m}^2/\text{s} \).
This can be compared to the coefficient of Bohm diffusion \( D_B = \frac{kT_e}{16eB} \approx 4 \text{ m}^2/\text{s} \), if evaluated at the edge.

The edge plasma conditions are closely linked to plasma surface interactions, e.g. impurity production and recycling. As discussed previously, the measured collection rate of metals at wall radius in the 40 kA shots is consistent with uniform sputtering at the wall by \( T_e \approx 100 \text{ eV} \) deuterium. The higher metal collection rate in the 90 kA discharges is still consistent with sputtering by more energetic deuterons. In a simplified analytical model of diffusion and surface recombination, Doyle [5] treated the molecular reemission of hydrogen from a surface as diffusion limited if \( W > 1 \) and recombination limited if \( W \ll 1 \), where the dimensionless parameter \( W = R \sqrt{\Phi K_r} / D \), \( R \) is the implantation range, \( \Phi \) is the implantation flux density, \( K_r \) is the surface recombination rate coefficient and \( D \) the diffusion coefficient of hydrogen in the material. He also gave the expression \( \tau = (R^2/D)(1+1/W^2) \) for the response time for molecular reemission. In our case of 90 kA discharges, we may take \( \Phi = 10^{23} \text{ m}^{-2} \text{s}^{-1} \). From the passive probe data the typical implantation energy can be estimated to \( E \approx 300 \text{ eV} \), corresponding to \( R \approx 40 \text{ Å} \). For the stainless steel wall we may take \( K_r \approx 10^{-35} \text{ m}^4/\text{s} \) and \( D \approx 10^{-14} \text{ m}^2/\text{s} \) [6]. This gives \( W = 0.4 \) and \( \tau = 12 \text{ ms} \). Thus molecular reemission is probably negligible in the 0.5 ms discharges discussed here.

Conclusions. Measurements in RFP plasmas in the T1 device with passive probes have been extended to high current operation. The edge plasma has also been investigated with Langmuir probes and heat flux probes. As on other RFP experiments, a strongly asymmetric heat flux is observed in the edge region. Examples are given of edge electron temperature, edge density and the flux and energy of particles hitting the wall. By characterising the edge plasma conditions this work provides a basis for modelling plasma surface interactions in T1. These aspects will be developed further in another publication.

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References.
Global Confinement and Edge Transport Measurements in the MST Reversed-Field Pinch


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Several approaches have been used to improve global confinement in the MST reversed-field pinch ($R = 1.5$ m, $a = 0.5$ m, $I < 750$ kA) including solid-target boronization, circuit modifications and hydrogen pellet injection. Of these, boronization has been the most dramatic, resulting in long periods free of major sawteeth in many discharges during which global energy confinement improves by over a factor of two. The sawtooth-free periods usually terminate in strong relaxation events. A large confinement database confirms that $\beta$ increases with density at fixed current up to a maximum of 10% near the radiation limit.

An extensive series of edge fluctuation-induced transport measurements have been made in low-current discharges. Electrostatic fluctuations can account for much of the edge particle flux with $D_{eff} \sim 50$ m$^2$/s, but account for less than 10% of the energy flux ($x_{eff} \sim 60$ m$^2$/s); however, the latter measurement is insensitive to ions and fast electrons. Magnetic fluctuations are found to result in small, ambipolar particle flux and corresponding small energy flux carried by fast electrons.

Introduction. Over the past year, attempts have been made to improve confinement in the MST reversed-field pinch by several means, including boronization, circuit improvements, and hydrogen pellet injection. During the same period, detailed studies of edge fluctuation-induced transport have been carried out in low current (200 kA) discharges with insertable probes in order to determine the underlying mechanisms of RFP transport near the plasma edge. We summarize the results of these studies here.

A large confinement database has been accumulated for a wide range of plasma current and density; one of the main results is reaffirmation of the increase of $\beta$ with density at constant plasma current, as indicated by Fig. 1 in which central electron temperature is shown along with lines of constant $\beta_e \equiv nT_e/\beta_0^2(a) = 5\%$. $\beta_e$ values of 5% (corresponding to $\beta = 10\%$ for $T_i = T_e$) are achieved only near the high-density limit given by $I/N \sim 2 \times 10^{-14}$ A-m. Ohmic power input and radiated power also increase near the high-density limit; this should improve in higher current-density devices such as RFX.

Boronization and Sawtooth-Free Periods. The inner aluminum wall of the MST vacuum vessel is boronized by ablating a solid boron carbide target in either RFP or pulsed discharge cleaning (PDC) discharges [1]. The boron layer acts as a low reflux coating and getter of light impurities, reducing, for example, edge CIII and OIII radiation to 40% of pre-boronization levels. After boronization was begun, we observed long sawtooth-free periods with reduced plasma resistance and global energy confinement improved over comparable discharges in which sawteeth or discrete dynamo events occur every 5 ms or so as a result of sudden nonlinear coupling of $m = 1$ internal tearing modes [2].

Fig. 2 displays a high-current discharge with a long sawtooth-free period ($t = 11-31$ ms). These periods are usually terminated by a violent dynamo event accompanied by a large sudden density rise which returns the plasma to typical sawtoothing discharge parameters. Table 1 shows a comparison of parameters before and after the dynamo event for the discharge of Fig. 2 — in all, the energy confinement time is almost a factor of three better during the sawtooth-free phase, although it does not reach the 2 ms confinement times of high-density discharges. Impurity radiation and edge $H_\alpha$ emission are observed to drop to record low levels during this time, implying an increase of $\tau_{e\alpha}$ and reduction of $Z_{eff}$ which partly accounts for the fourfold reduction of plasma resistance. The current profile is highly peaked during the sawtooth-free
<table>
<thead>
<tr>
<th>Parameter</th>
<th>Before Sawtooth</th>
<th>After Sawtooth</th>
</tr>
</thead>
<tbody>
<tr>
<td>$I_p$ (kA)</td>
<td>520</td>
<td>410</td>
</tr>
<tr>
<td>$\Theta$</td>
<td>1.90</td>
<td>1.75</td>
</tr>
<tr>
<td>$F$</td>
<td>-0.25</td>
<td>-0.25</td>
</tr>
<tr>
<td>$\bar{n}$(cm$^{-3}$)</td>
<td>0.75 x 10$^{13}$</td>
<td>0.85 x 10$^{13}$</td>
</tr>
<tr>
<td>$T_{e0}$ (eV)</td>
<td>350 ± 40</td>
<td>300 ± 50</td>
</tr>
<tr>
<td>$T_i$ (eV)</td>
<td>300 ± 20</td>
<td>200 ± 20</td>
</tr>
<tr>
<td>$P_{OH}$ (MW)</td>
<td>5</td>
<td>12</td>
</tr>
<tr>
<td>$P_{rad}$ (MW)</td>
<td>1</td>
<td>3</td>
</tr>
<tr>
<td>$R_p$ ((\mu\Omega))</td>
<td>19</td>
<td>73</td>
</tr>
<tr>
<td>$\tau_E$ (ms)</td>
<td>1.4</td>
<td>0.5</td>
</tr>
</tbody>
</table>

Table 1: Comparison of global plasma parameters before and after the major discrete dynamo event which ends the sawtooth-free period for the discharge shown in Fig. 1. $\tau_E$ is calculated by setting \(\langle n_e T_e + n_i T_i \rangle = \bar{n}(T_{e0} + T_i)\).

phase, as $\Theta = 1.90$ is quite high for $F = -0.25$; the current profile broadens suddenly with the sawtooth, with $\Theta$ dropping to 1.75. The current broadening and decrease of central electron temperature are additional factors which increase plasma resistance following the sawtooth; anomalous resistivity may also contribute. Unlike typical sawteeth, which often produce sudden increases of $T_i$ [3], the ‘monster’ sawtooth leads to a radical drop of $T_i$ as measured by a charge-exchange analyzer; this may be due in part to the increased charge-exchange flux from outer regions of the plasma.

MHD activity is still present in small, rapid bursts during the sawtooth-free period, producing small, regular ‘mini-sawteeth’ in many signals. The $m = 1$ toroidal mode spectrum is similar to the case of normal sawteeth, except the amplitudes are reduced. This activity is a new manifestation of the discrete dynamo which evidently produces little current broadening and plasma-wall interaction.

**Circuit Improvements.** Several improvements to the poloidal and toroidal field circuits have been implemented, allowing low-current flat top durations of 40 ms, radial field error below 10% at all current levels, and plasma current as high as 700 kA. These changes have also greatly reduced the occurrence of mode locking [4] observed to degrade confinement in the past.

An external inductor has been added to the toroidal field crowbar circuit which allows additional control of reversal. Lower external inductance results in improved volt-second consumption but delayed reversal due to the higher initial toroidal field required for self-reversal to the same steady-state $F$ value.

**Hydrogen Pellet Injection.** Hydrogen pellets have been injected into MST discharges with the hope of peaking the density profile and improving confinement as in tokamaks. However, the pellet flight time (0.5–1.0 ms) is comparable to the particle confinement time, and density profile relaxation takes place during pellet ablation. As a result, high-density pellet-fueled discharges are observed to have flattened density profiles similar to gas-fueled discharges of the same density. Pellet injection does allow the transient achievement of high density discharges ($I/N < 2 \times 10^{-14}$ Am) without wall loading and radiation collapse.

**Edge Fluctuation-Induced Transport.** An extensive series of probe measurements have been made in 200 kA discharges in order to measure the contributions to fluctuation-induced transport. A summary of these measurements is shown in Table 2.

Particle and energy transport produced by electrostatic fluctuations is measured with triple Langmuir probes [5], which are insensitive to the fast electron component present in the edge [6]. Electrostatic fluctuations can account for much of the edge particle flux, with an effective diffusion coefficient $D_{eff} \equiv \Gamma/|\nabla n| \sim 50$ m$^2$/s; they account for a small part of energy flux,
consistent with the 20 eV bulk electron temperature measured by the probes. The effective thermal diffusivity for the bulk electrons is \( \chi_{\text{eff}} \approx \frac{q_r}{n|\nabla T_e|} \sim 60 \text{ m}^2/\text{s} \).

Fast electron transport produced by magnetic fluctuations is measured by an electron energy analyzer (EEA) and magnetic probes while the energy flux produced by magnetic fluctuations, dominated by fast electrons, is measured with a pyrobolometer probe [7] and magnetic probes. Magnetic fluctuations account for a small fraction of total edge particle and energy transport.

The undetected radial energy flux may be in the form of electrostatic fluctuation-induced transport of fast electrons and ion energy flux, neither of which has been measured. Efforts are currently underway to estimate the relevant quantity for fast electron energy transport, \( \langle n_{\text{fast}} B_\theta \rangle \) as well as the radial ion flux using new probe techniques.

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Table 2: Radial power and energy flow at the edge of MST for \( I_p = 200 \text{ kA} \), \( \bar{n} = 1 \times 10^{13} \text{ cm}^{-3} \).

<table>
<thead>
<tr>
<th>Mechanism</th>
<th>Particle</th>
<th>( q_r/\Gamma )</th>
<th>Energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>All Particles</td>
<td>(~ 2 \times 10^{21} \text{ m}^{-2} \text{s}^{-1} = \Gamma)</td>
<td>( 180 \text{ eV} \sim T_{\text{fast}} )</td>
<td>( 60 \text{ kW/m}^2 = q_r )</td>
</tr>
<tr>
<td>Global</td>
<td>( H_\alpha, \bar{n} )</td>
<td>( 20 \text{ eV} \sim T_e )</td>
<td>( 6 \text{ kW/m}^2 \sim 0.1 q_r )</td>
</tr>
<tr>
<td>Bulk Electrons</td>
<td>(~ 2 \times 10^{21} \text{ m}^{-2} \text{s}^{-1} = \Gamma)</td>
<td>( D_{\text{eff}} \sim 50 \text{ m}^2/\text{s} )</td>
<td>( \chi_{\text{eff}} \sim 60 \text{ m}^2/\text{s} )</td>
</tr>
<tr>
<td>ES Fluctuations</td>
<td>( \langle \bar{n} \bar{E}_\phi \rangle )</td>
<td>( 150 \text{ eV} \sim T_{\text{fast}} )</td>
<td>( 3 \text{ kW/m}^2 \sim 0.05 q_r )</td>
</tr>
<tr>
<td>Fast Electrons</td>
<td>(~ 1 \times 10^{20} \text{ m}^{-2} \text{s}^{-1} \sim 0.05 \Gamma)</td>
<td>( \langle \bar{n} \bar{E}_\phi \rangle )</td>
<td>( \langle \bar{n} \bar{E}_\phi \rangle )</td>
</tr>
<tr>
<td>Mag Fluctuations</td>
<td>( \langle \bar{n} \bar{E}_\phi \rangle )</td>
<td>( \langle \bar{n} \bar{E}_\phi \rangle )</td>
<td>( \langle \bar{n} \bar{E}_\phi \rangle )</td>
</tr>
</tbody>
</table>

Figure 1: Central electron temperature, measured by Thomson scattering, decreases weakly with increasing density for a given plasma current, resulting in increase of $\beta_e$ with density.

Figure 2: A high current discharge with a sawtooth-free period ($t = 11-31$ ms): (a) applied poloidal (upper trace) and toroidal voltage, (b) plasma current, (c) central line-averaged density, (d) reversal and pinch parameters, (e) radiated (lower trace) and Ohmic power, (f) edge $H_\alpha$ emission, (g) central soft x-ray emission (dominated by OVIII and AlXII line radiation) (h) C$X\alpha$ ion temperature with central electron temperature indicated at two times.
Controlled Wave Form of the Plasma Current on ATRAS-RFP Experiment

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Y. Osanai; Department of Electronics Engineering Tokyo Metropolitan College of Aeronautical

Abstract: The current wave form of a reversed field pinch (RFP) plasma is controlled by superposing an additional loop voltage. Many types of the wave form are obtained and the current durations are also elongated in some case. Then F and Θ values are not affected by the variation of the wave form. This means that we can control the wave form but not control the spatial distribution of the magnetic field except its magnitude. A necessary condition to maintain constantly the F, Θ values are discussed and need a radial diffusion velocity.

Experimental device: ATRAS (Approach to Toroidal Reactor by Alternative System) is a RFP machine which has the major radius of 50 cm and the minor radius of 9 cm limited by thirty two stainless-steel limiters. The vacuum vessel which is made of stainless steel with the thickness of 0.4 mm is covered by Aluminum double shell of 10 mm thickness. It has not the vertical field coils to make the equilibrium. The characteristic of ATRAS device is that it has two independent Ohmic heating (OH) circuits. One of which called OH-1 is used to produce a basic RFP plasma. It has eighteen coils arranged as shown in fig. 1 and connected in series. This arrangement does not make out stray field around the plasma region except small vertical field. The quarter period of this circuit is 1.5 ms. The maximum volt-second is 0.21 Vs. Another one called OH-2 has six coils arranged as shown in fig. 1. The outside shells divided toroidally to two sections are used as the basic coils. We can change the period of this circuit by variable connection methods of these coils. Here, we connected these coils to series and obtained the longest quarter period of 550 μs. The wave form of the plasma current is controlled by the superposition of the toroidal loop voltage induced by the discharge of OH-2 bank. The toroidal field coil is constructed by thirty-six elements shown in fig. 1. These elements are arranged toroidally at regular intervals and connected in series. To make small the toroidal stray field, the axis of the vacuum vessel is set 2 cm inner from the axis of the toroidal field coil. This coil is

Fig. 1. Poloidal cross-section of the ATRAS-RFP machine which has two independent Ohmic heating circuits.
connected to three banks making the bias, reversal and sustainment field. By these discharge, we can obtain the various driving modes for toroidal magnetic field.

**Basic plasma:** Basic plasma is made by using only OH-1 circuit. Plasma current rises linearly and reaches the maximum value of 60 kA after 960 μs of discharge. Then it decreases convexly and abruptly disrupt with exponential curve. The current duration time is about 1.4 ms. The toroidal magnetic field on the plasma surface reverse the direction at 100 μs after the beginning of discharge. F and Θ values become almost constant after 200 μs from the start of discharge. The central position of the plasma current moves violently in this formation phase of the RFP configuration, then it is almost static with toroidal shift of 2cm. These mean that the RFP configuration is constructed in early stage and maintained until the plasma current disrupt in spite of the increasing and decreasing of the plasma current. These situations are confirmed by using of the magnetic probe inserted in the plasma column.

**Wave form control:** When the current of the basic plasma begin to decrease, the loop voltage produced by the discharge of the OH-2 circuit is superposed on the plasma surface to control the wave form of the plasma current. Many kinds of the wave form are obtained according to the loop voltage induced by the discharge of the OH-2 circuit. These wave forms are shown in Fig. 2(a) with a toroidal loop voltage on the plasma surface, where labels attached to each curve show the charging voltages of the OH-2 circuit. The plasma current increases with the charging voltage and flat top of 600 μs is obtained when the charging voltage is 2 kV. The plasma current increase than that of the start point of the OH-2 discharge when the charging voltage is higher than 2 kV, as seen in fig. 2(a). In this case, the onset time of the current disruption could be retarded at most about 600 μs than that of the basic operation by controlling the induced voltage. The F and Θ values is almost constant through the discharge, keeping the values before the OH-2 discharge in any operation of OH-2 discharge, and also the central position of the plasma current is not affected by the OH-2 discharge, as shown in fig. 2(b).

![Fig. 2(a). Controlled wave forms of the plasma current for each charging voltage of OH-2 bank and the loop voltage in case of 3 kV charging to OH-2 bank. The dotted line shows the loop voltage in case of vacuum shot with the 3 kV charging to OH-2 bank. Fig. 2 (b). The time variations of the F and Θ values for each shot corresponding to the current wave forms shown in Fig. 2 (a). The dotted line show the toroidal shift of the central position of the plasma current.](image-url)
In any operation of OH-2 discharge, the plasma current begin linearly to decrease at the time when the loop voltage of the corresponding vacuum shot would have gone through the null point, as long as the plasma current lasted till this time. Then, the loop voltage become almost constant from that time till the disruption. The relation between the loop voltage and the plasma current is shown in fig. 3.

Things take place the way similar to the above when OH - 2 discharge is started just before disruption of the plasma current. It seems that the condition occurring the current disruption grows in very short time from this result and is overcome by superposing the additional loop voltage.

From these facts, we can see that the spatial distributions of the plasma current and the magnetic field do not change but only these magnitudes changes on superposition of the OH - 2 discharge. This means that the RFP configuration is very hard and we could use this property to control the output energy of the reactor.

**Discussions:** Necessary condition which maintain the RFP configuration with a constant F - value is considered. For the simplicity, the radial variations of the applied loop voltage in the plasma are neglected, which is valid provided the skin depth of the plasma column is sufficiently long compared with the minor radius of the plasma column. The plasma current induced by the loop voltage and diffusion flows according to the Ohm's law, \( \mathbf{j} = \mathbf{E} + \mathbf{v} \times \mathbf{B} \).

The poloidal and toroidal components of the plasma currents are written as

\[
\mathbf{j}_p(t) = \frac{B_p B_\theta V_{\text{loop}}}{2\pi R \eta_\parallel (B_\phi^2 + B_\theta^2)} \left( 1 - \frac{1}{\eta_\|} \right) \frac{v_p B_\parallel}{C_n \eta_\parallel} B_\phi^2 V_{\text{loop}} \left( \frac{B_\phi^2 + B_\theta^2}{C_n B_\phi^2} \right) + \frac{v_p B_\parallel}{C_n \eta_\|} \frac{v_p B_\parallel}{C_n \eta_\parallel} \]

where \( R \) is major radius of the torus, \( V_{\text{loop}} \) is loop voltage on plasma surface, \( \eta_\parallel \) is resistivity along the magnetic field line, \( C_n \) is defined as \( C_n = \eta_\parallel / \eta_\perp \) where \( \eta_\perp \) is resistivity perpendicular to the line of the magnetic field, \( v_r \) is radial velocity of the plasma particles, and another symbols are used as usual meanings. The first term of the equation for the poloidal current is the current which is induced by the loop voltage and is flowing along the magnetic line of force. This term changes the sign at the reversal surface because the toroidal magnetic field changes the sign. If the poloidal current change the sign in the plasma region, the reversal surface move toward the axis. The second term prevent this because it has the opposite sign.

The condition of which the poloidal currents do not change the sign in the plasma region is written as
This means that some diffusion is needed to sustain the RFP configuration and the position of the reversed surface is dependent on the plasma surface condition of the toroidal field. The value of \( v_r \) must be sufficiently small to maintain sufficiently long the energy confinement time for fusion reactor. That is,

\[
\tau_B \lesssim \frac{a}{v_r}
\]

where \( a \) is the minor radius of the plasma column. This condition is not so hard to satisfy. Because, \( V_{\text{loop}} \) may be several volts in the high temperature plasma and the confinement field is several teslas in the reactor condition. And, if we estimate that \( C_\eta \) is 2 - 3, the diffusion velocity which need to maintain the RFP configuration becomes \( \sim 0.1 \text{ m/sec} \). In our experiment, these critical values of \( v_r \) are calculated on each radius of the plasma column under the assumption that the field configuration is Modified Bessel Function Model \(^{(1)}\) and \( C_\eta \) is 3, and shown in fig. 4. Where, the magnitude of the toroidal field on axis is 0.4 T, \( \Theta = 2.0 \), and \( V_{\text{loop}} = 100 \text{ V} \). The skin depth for the loop voltage produced by the OH-2 circuit is 2.3 cm when we assume the plasma temperature of 100 eV and the density of \( 10^9 / \text{m}^3 \). The skin depth becomes longer in our experiment because the self inductance on each magnetic surface is large in the outer region of the plasma column. From these facts, the diffusion time, \( a/v_r \), becomes same order of the plasma current duration.

Conclusions: The wave form of the plasma current is controlled by the superposition of the additional toroidal loop voltage and the current duration of the discharge is also elongated by addition of the sufficient high loop voltage. But we can not control the \( F \) and \( \Theta \) values. This means that the RFP configuration is very hard and the things controlled are magnitudes of the current and magnetic field. The diffusion velocity, \( v_r \), is necessary to maintain the static RFP configuration. But, the diffusion time decided by \( a/v_r \) is longer than the energy confinement time in the reactor condition.

REFERENCES

ION TEMPERATURE GRADIENT DRIVEN INSTABILITY AND ANOMALOUS ION HEATING IN RFP

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I. INTRODUCTION

The experimental observations in Reversed Field Pinch (RFP) suggest the existence of anomalous ion heating. It has been commonly assumed that the heating results from damping of the velocity fluctuations associated with current driven MHD instabilities through parallel ion viscosity. In the present work, we suggest a possible mechanism of the existence of collisionless perpendicular viscosity due to ion gradient driven instability (the so called $\eta_i$-mode, $\eta_i = \frac{d \ln T_i}{d \ln n_i} = \frac{L_B}{L_T}$). The excess of electromagnetic energy associated with tearing mode fluctuations is dissipated through the anomalous viscosity, resulting in direct ion heating. The measurements in RFX have shown that the ion temperature profile is more peaked than the density profile, while the ion temperature is observed anomalously high. This suggests that the $\eta_i$-instability may exist and play a role in the direct heating of ions. It is found that the growth rate of the $\eta_i$-instability is much larger than that in a Tokamak because both the destabilizing effects, bad magnetic curvature and negative parallel compressibility, are enhanced in RFP configuration. Furthermore, the ion anomalous viscosity is estimated by employing the $\gamma/k_{\perp}^2$ approximation. The corresponding ion heating power density has been calculated from the flux velocity fluctuations associated with tearing modes, and found to be sufficient to balance the ion energy confinement losses.

II. EIGENMODE EQUATION AND SOLUTIONS

The eigenmode equation for instability is derived from gyrokinetic equation in cylindrical coordinates. We assume $\left(\frac{\nabla \cdot \nabla}{\omega}\right)^2 = \frac{\omega_d}{\omega} << 1$ for ions, and the adiabatic responses of electrons. Here $\omega_d$ is the magnetic drift frequency. It has been found that in order to justify the self consistency of above approximations, the condition $\varepsilon_T / \varepsilon_B << 1$ and $\varepsilon_T / q << 1$ must assumed. The quasineutrality equation retaining full finite Larmor radius effects is given by

$$\frac{d^2 \phi}{d \chi^2} + F \frac{d \phi}{d \chi} + G = 0$$

(1)
In order to solve the eigenmode equation analytically, we consider two asymptotic parameter regions: \( b_0 << 1 \) and \( b_0 >> 1 \).

1. In the \( b_0 << 1 \) limit, the mode has a broad structure in \( \chi \), but \( b_0 s^2 \chi^2 << 1 \) remains valid. On expanding function \( \Gamma(b) \) with small argument we obtain a simplified equation having the eigenfunctions of the form \( \phi = H_n \exp(-\sigma^2 \chi^2) \), where \( \sigma^2 = i \varepsilon_n / \Omega q b_0 t s (1 + \varepsilon^2 / q^2)^{1/2} \) and \( H_n \) are Hermite polynomials. The corresponding eigenvalues are

\[
\Omega = \frac{-b_0 \eta_i}{2} + \left[ \frac{2 \varepsilon_n \eta_i + i s \varepsilon_n \eta_i}{\tau E} (1 + \varepsilon^2 / q^2) \right]^{1/2} \quad \text{for} \quad b_0 << (\varepsilon_T / E_B)^{1/2}. \]

And

\[
\Omega = -b_0 \eta_i + \left[ \frac{2 \varepsilon_n \eta_i + i s \varepsilon_n \eta_i}{\tau E b_0} (1 + \varepsilon^2 / q^2) \right]^{-1/2} \quad \text{for} \quad 1 >> b_0 >> (\varepsilon_T / E_B)^{1/2}. \]

2. In the \( b_0 >> 1 \) limit, the mode structure is very localized in \( \chi \). Upon expanding Eq.(1) about \( \chi = 0 \) for \( \chi^2 << 1 \) and further expanding \( \Gamma( b_0 ) \) for large \( b_0 \), we obtain
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\[ \Omega = \frac{\eta_i}{2(1 + \tau)\sqrt{2\pi b_\theta}} + \frac{3\varepsilon_B + \varepsilon_\tau}{2\varepsilon_B\tau} g_2 \left( \frac{\varepsilon^2}{a + \tau} \right)^{\frac{1}{2}}, \text{ and } \sigma^2 = \frac{q^2 s}{q^2} b_\theta \sqrt{1 + \frac{\varepsilon^2}{q^2}} \]

for \( \varepsilon_B^2/2\varepsilon_\tau^2 \gg b_\theta \gg 1 \). When \( b_\theta \gg \varepsilon_B^2/2\varepsilon_\tau^2 \), the kinetic effects, both ion Landau damping and magnetic drift resonance, become important and must be taken into account. This will be studied elsewhere.

From above results, we can see that in the long wavelength region \((b_\theta < (\varepsilon_\tau/\varepsilon_B)^{1/2})\), the magnetic drift \( \omega_{di} \) plays an important role for the instability. Instead, in the short wavelength region \((b_\theta > (\varepsilon_\tau/\varepsilon_B)^{1/2})\), the growth rate is mainly determined by the parallel compressibility. The numerical results of Eq.(1) is shown in Fig.1, where the mode frequency \( \Omega_r \) and growth rate \( \Omega_i \) are plotted versus \( b_\theta \) for \( \varepsilon_\tau = 0.08 \), \( \eta_i = 10 \), \( \varepsilon_B = 0.4 \) and \( q = 0.2 \).

### III. ANOMALOUS ION VISCOSITY AND ION HEATING

The anomalous ion viscosity coefficient \( \nu_{\eta_i} \) induced by \( \eta_i \)-instability can be estimated as \( \nu_{\eta_i} = m_i \rho_{\eta_i} / v_k \), where \( \eta \) is the growth rate, \( k_\perp = \sigma^2 \) is the mode width in radial direction. According to this model, \( \nu_{\eta_i} \) has been calculated in following different parameter regions:

a) \( \nu_{\eta_i} = m_i \rho_{\eta_i} / \sqrt{2\pi} b_{\perp} \) \( \varepsilon_\tau \) \( v_s \) \( \sqrt{1 + \frac{\varepsilon^2}{q^2}} b_{\theta} \) for \( b_\theta << (\varepsilon_\tau/\varepsilon_B)^{1/2} \),

which is a monotonically increasing function of \( \sqrt{b_\theta} \).

b) \( \nu_{\eta_i} = m_i \rho_{\eta_i} / 2\sqrt{2\pi} L_{\tau} \) \( v_{\tau,\perp} \) \( \rho_\perp^2 \) \( \sqrt{b_{\theta}} \) for \( 1 > b_\theta \gg (\varepsilon_\tau/\varepsilon_B)^{1/2} \),

which is also proportional to \( \sqrt{b_\theta} \).

c) \( \nu_{\eta_i} = m_i \rho_{\eta_i} / 2\sqrt{\pi} L_{\tau} \) \( \rho_\perp^2 \) \( \sqrt{b_{\theta}} \) for \( 1 < b_\theta < \varepsilon_B^2/2\varepsilon_\tau^2 \),

which is independent of \( b_\theta \). Here \( \rho_\perp \) is the ion Larmor radius.

For \( b_\theta \gg \varepsilon_B^2/2\varepsilon_\tau^2 \), we expect that the kinetic effect would damp the instability and cause the decreasing of \( \nu_{\eta_i} \). Taking \( \nu_{\eta_i} \) resulting from the Eq.(4) as upper limit of the anomalous viscosity value, in the RFX case, \( T_i = 2500 \text{ eV} \), \( T_e = 500 \text{ eV} \), \( \eta_i = 2 \times 10^{13} / \text{cm} \), \( L_T = 20 \text{ cm} \), \( B = 0.2 \text{ T} \), \( Z_{eff} = 3 \), and \( \tau_E = 1 \text{ ms} \), we obtain \( \nu_{\eta_i} = 4.5 \times 10^{-6} \) (poise). In Fig.2, both the anomalous and classical perpendicular viscosity versus \( T_i \) are plotted for the RFX parameter, showing that the collisionless viscosity is several orders of
magnitude larger than the collisional one. It is obvious that, for higher temperature plasma, the collisionless viscosity becomes more important.

During the relaxation process in a RFP, the excess potential energy is released from the equilibrium magnetic field through instabilities which generate plasma motion. The ion viscosity damps the velocity fluctuation associated with the instabilities, leading to direct ion heating. Several measurements support the interpretation of the low frequency magnetic fluctuations in RFP as tearing modes[6]. Here we employ the tearing mode magnetic fluctuation $\delta B_r$ as an energy source to estimate ion heating power density, which can be approximated as[8]

$$ P_{\text{vis}} = \eta_{ni} \left( \frac{\partial v_{\theta}}{\partial \tau} \right)^2 = \frac{\varepsilon_{\text{vis}}}{w^2} \frac{\varepsilon_{\text{vis}}}{\varepsilon_{\text{vis}}} \frac{S^2}{\eta_{A}^2} \frac{v_A^2}{m^2 R^2 q^2} \left( \frac{\delta B_r}{B_0} \right) $$

(5)

where $\delta v_{\theta}$ is the amplitude of velocity fluctuation of tearing mode; $W$ is the magnetic island width, $S$ is the Lundquist number, $v_A$ is the Alfvén velocity. By Taking above RFX parameters and $m=1$, we obtain: $S=1.5 \cdot 10^6$, $\delta v_{\theta} = 19.3 v_A \delta B_r / B_0$, and $P_{\text{vis}} = 7.89 \cdot 10^9 (\delta B_r / B_0)$ (erg/cm$^3$sec).

Upon balancing $P_{\text{vis}}$ with ion energy confinement loss $n_i T_i / T_E$, it turns out that the amplitude of magnetic fluctuation should be above the level of $\delta B_r / B_0 = 1 \cdot 10^{-3}$, which is reasonable as compared with experimental results[9].

Finally, we should point out that, in this model, $\eta_{ni}$-instability also induces ion anomalous thermal conductivity, which results in energy transport with the power density $P_{\chi i}$. On comparing the $P_{\chi i}$ with $P_{\text{vis}}$ we have:

$$ \frac{P_{\text{vis}}}{P_{\chi i}} = 2 \frac{S^2}{\eta_{A}^2} \frac{v_A^2}{m^2 R^2 q^2} \frac{S}{B_0} \left( \frac{\delta B_r}{B_0} \right) $$

(6)

In RFX case, by taking $\delta B_r / B_0 = 2 \cdot 10^{-3}$, Eq.(6) leads to $P_{\text{vis}} / P_{\chi i} \approx 6$. These two power density can balance with each other only when the ion temperature higher than 1.5 KeV. This indicates that the $\eta_{ni}$-instability induced anomalous thermal diffusion is not the main mechanism of energy confinement losses in present RFP's.

REFERENCES
(1) R.B.Howell et al; P.F. 28 743 (1985)
(2) K.Ogawa et al; N.F. 25 1295 (1985)
(3) P.G.Carolan et al; 14th EPS Vol.II 469 (1987)
(4) C.Gimblett; Europhys. Lett. 11 541 (1990)
(7) L.Carraro et al; to be published in 20th EPS conf. Lisboa,(1993)
(8) P.H.Rutherford; P.F. 16 1903(1973)
(9) The RFX team;14th IAEA Conf. CN-56-H-1-1,Wurzburg,(1992)
MHD studies of stationary turbulent dynamics

in a reversed-field pinch

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We present results obtained from a number of long-time simulations runs of stationary Reversed-Field-Pinch (RFP) dynamics. The dissipative compressible MHD equations are solved in cylindrical geometry using a 3D semi-implicit spectral code with appropriate selection of modes \((m,n)\). The boundary conditions are: rigid wall \(v_r(a)=0\), constant toroidal current \(I\) and constant toroidal magnetic flux \(\psi_t\), i.e. \(\Theta = B_\phi (a)/< B_z > = \text{const.} \) and \(d\bar{B}/dr|_r=a=0\), where the bar indicates the average over the wall surface and \(<\>\) the volume average. In most cases a resistivity profile constant for \(r<0.95a\) and sharply increasing in the edge region is chosen, but some runs were performed with uniform resistivity. The aspect ratio of the system is \(R/a=4\). More details concerning the model are found in Refs. [1],[2]. Runs are followed for several thousands Alfvén times, which allows good statistical averages. Lundquist numbers \(S\) are varied between \(10^3\) and \(10^5\).

Except for low values \(S \sim 10^3\), where a single helicity behaviour tends to dominate, we always find turbulent states. Boundary conditions and resistivity profile seem to determine the (average) behaviour uniquely. (The possibility of high-viscosity quasi-single-helicity states discussed recently [3] is eliminated by choosing magnetic Prandtl number \(\eta/\nu=1\)). There is, however, a rather strong dependence of the dynamical behaviour on the numerical resolution, i.e. the number of modes included, if the latter is too small.

Fluctuation spectra are dominated by \(m=1\) modes, \((m,n)=(1,n)\) with \(-15 \leq n \leq -8\). While the \(n\)-spectrum drops sharply for \(n > -10\) \((n > -10m\) for higher \(m\), corresponding to nonresonant modes \(q_{nm} > q(0)\), the decrease at high \(|n|\), \(n \leq -20\), is much more shallow indicating substantial excitation of modes close
to the field reversal radius. Figure 1 illustrates the magnetic modal energies \( \langle \delta E_{mn}^H \rangle \) for two typical runs with \( S=3 \times 10^3 \) and \( S=10^5 \). In the latter case the total number of modes included (107) is only marginally adequate. On the other hand the long exponential tail in the energy spectrum in

![Graph showing magnetic modal energies for two typical runs](image)

**Fig. 1:** Magnetic modal energies \( \langle \delta E_{mn}^H \rangle \) for two typical runs

Typical pseudo-spectral simulations (e.g. Ref. [4]) clearly indicate that a large fraction of the modes included in such computations is redundant. The scaling of the average magnetic fluctuation amplitude \( \delta B \) in the plasma with Lundquist number is of crucial importance for the energy confinement in the RFP. As is clearly seen in Fig. 1 the fluctuation energy decreases with increasing \( S \), in contrast to the findings in Ref. [4].

Figure 2 gives the volume averaged energy \( \langle \delta E^H \rangle \) for different values of \( S \). The simulation results closely fit the scaling \( \langle \delta E_{mn}^H \rangle = S^{-0.44} \), or \( \delta B = S^{-0.22} \). This result is consistent with the scaling expected for a dynamical state dominated by current sheet reconnection, \( \delta v_\perp / \delta B = S^{-1/2} \), which using Ohm's law at the
reversal point \( \eta J_0 = \langle \delta v \times \delta B \rangle \) yields \( \delta B = S^{-1/4} \). In fact formation of intense current sheets is found to occur intermittently. Such current sheets lasting about \( 10^2 \tau_A \) are rather localized in toroidal azimuth, \( \nabla \phi = \pi/3 \). As seen in Fig. 3 sheet currents are flowing more in poloidal than toroidal direction, which is to be expected since \( \mathbf{j} \) is roughly parallel to \( \mathbf{B} \). The total contribution of such current sheet reconnection to the average dynamo effect is at present difficult to assess.

Measurement of \( \delta B \) within the plasma have been performed only on low current discharges in some experiments. In most experiments \( \delta B (a) \) is
usually measured yielding a stronger $S$-dependence, $\delta B = S^{-1/2}$ OHTE [5] and $\delta B = S^{-2/3}$ TPE IRM15 [6], in an $S$-range comparable to ours. We therefore give the ratio of the volume average fluctuation energy of mode number $m$, $\langle E_m \rangle$ to the surface averaged value $\overline{E_m}(a)$, $R_m = \langle E_m \rangle / \overline{E_m}(a) : R_0 = 1, R_1 = 2, R_2 = 100, R_3 = 500, R_4 = 10^4$. The $m=1$ component $\overline{E_1}(a)$ dominates the $m$-spectrum like the volume average $\langle E_1 \rangle$ and the $S$-scaling is similar.

Considering toroidal and poloidal components we have $\delta B_\varphi(a)/\overline{B_\varphi}(a) = 0.135 (=0.063)$ and $\delta B_\theta(a)/\overline{B_\theta}(a) = 0.05 (=0.022)$ for $S=3\times10^3 (S=10^5)$, indicating a well-defined polarization of the fluctuations, which is also observed experimentally. Extrapolating the numerical values of the $m=0$ and $m=1$ components to $S=10^5$ we find that for the poloidal component $m=1$ remains dominant, while for the toroidal component $m=0$ becomes larger than $m=1$, which is also in agreement with experiments. The absolute values are, however, significantly larger than observed experimentally.

On the other hand the field reversal is generally less pronounced in the simulations, $F = \overline{E_\varphi}(a)/\langle B_\varphi \rangle \approx -0.12$ for $S=3\times10^3$ and $-0.08$ for $S=10^5$ as compared to typically $F=-0.2$ in experiments. Since a possible reason for this discrepancy could lie in the resistivity profile, several runs were performed with uniform resistivity. While the fluctuation amplitudes indeed decrease preserving approximately the $S$-scaling, the field reversal becomes even weaker. The origin of the deeper reversal in the experiments may be caused by non-MHD effects, such as the kinetic dynamo effect [7].

References

TRANSPORT IN A PARAMAGNETIC PINCH WITH NON-CIRCULAR CROSS-SECTION

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1. Introduction. In order to reach fusion-relevant temperatures in pinch experiments, such as ULQ and RFP, it is necessary to strongly reduce the plasma-wall interaction. To do this, poloidal divertor configurations have been introduced, where the plasma cross-section inevitably becomes non-circular. So far, no study of how this non-circularity affects the transport properties has been made, and this is the subject of the present paper. Using a simple Ohm's law and two different forms for the resistivity, we can calculate the changes in particle- and energy confinement times due to the non-circularity.

2. Basic equations. Assuming that Ohm's law takes the simple form $E + \mathbf{v} \times \mathbf{B} = \eta \mathbf{J}$, the classical particle confinement time is given by [1]

$$\tau_p = \frac{N_p(\psi_s)}{\Gamma_p(\psi_s)}, \quad N_p(\psi_s) = \int n dV, \quad \Gamma_p(\psi_s) = \int \frac{n \eta |\mathbf{J}|^2}{|\nabla \psi|} dS,$$

where $N_p(\psi_s)$ is the total number of plasma particles, $\Gamma_p(\psi_s)$ the flux of particles leaving the plasma, $n$ the number density of particles, $\eta$ the total pressure, $\eta$ the plasma resistivity, and the plasma boundary is given by the poloidal flux $\psi = \psi_s$. For given pressure- and current profiles the density profile $n(\mathbf{r})$ is then chosen such that the particle flux $\Gamma_p$ is independent of $\psi$, corresponding to a situation where all particles are born along the magnetic axis. This procedure [2] is useful because it provides a simple and well-defined method for estimating $\tau_p$ as function of various profile- and geometrical parameters. For the classical energy confinement time $\tau_E$ we adopt the definition
\[ \tau_c = \frac{W(\psi_s)}{P_\Omega(\psi_s)}, \quad W(\psi_s) = \int_{\psi < \psi_s} \frac{3}{2} p \, dV, \quad P_\Omega(\psi_s) = \int_{\psi < \psi_s} \eta |J|^2 \, dV. \]

\(W(\psi_s)\) and \(P_\Omega(\psi_s)\) are the total thermal energy and ohmic energy input, respectively. In calculating these integrals we use the same profiles as in calculating \(\tau_D\) above. For an axisymmetric plasma the condition \(dI_D/d\psi = 0\) leads to the density profile

\[ \frac{n(\psi)}{n_0} = \left( \frac{k(\psi)}{k_0} \right)^{\frac{1}{3}} \left( \frac{1 dp/d\psi}{1 dp/d\psi_0} \right), \quad k(\psi) = \int \frac{\eta |J|^2}{B_p} \, dl_p, \]

where the integral is to be taken once around the flux surface in the poloidal direction and index 0 denotes values on the magnetic axis \(\psi = 0\). With given pressure- and current profiles, all what remains is to specify the resistivity profile. Here we use two different profiles; the first is the classical Spitzer scaling

\[ \frac{\eta_{cl}}{\eta_0} = \left( \frac{T}{T_0} \right)^{-3/2} \left( \frac{p}{p_0} \right)^{-3/2} \left( \frac{n}{n_0} \right)^{3/2} \]

and the second, \(\eta_{st}\), is derived from the condition of stationarity, \(V \times E = -\partial B/\partial t = 0\). It should be noted, however, that for general non-circular equilibria the confinement times defined above must be computed via a numerical equilibrium code.

3. Non-circular constant current pinch. For weakly non-circular, straight, constant current pinch equilibria, however, the confinement times can indeed be calculated analytically and expressed in terms of integrals. The corresponding equilibrium theory is developed by Wahlberg et al.[3]. The plasma boundary is assumed to be given by \(r_s(\theta) = a[1 + AN\cos(N\theta) + ...]\), where \(\theta\) is the poloidal angle. With \(B = V \times B + B_z \hat{z}\), \(B_s\) is given by [3]

\[ \frac{\psi}{\psi_s} = \left( \frac{r}{a} \right)^2 - 2A_N \left( \frac{r}{a} \right)^N \cos(N\theta) + \left( N - \frac{1}{2} \right) A_N^2 \left( \frac{r}{a} \right)^{2N} \cos(2N\theta) + ... . \]

The non-circularity coefficient \(A_N\) is formally assumed to be small, \(|A_N| \ll 1\). The linear functions \(p(\psi)\) and \(B_z^2(\psi)\) are characterized by the two parameters

\[ \lambda = \frac{B_{z0}^2 - B_{zN}^2}{B_{z0}^2}, \quad \beta_0 = \frac{p_0}{B_{z0}^2/2\mu_0} \]

measuring the degree of paramagnetism and strength of the axial magnetic field, respectively. \(\eta_{cl}\) becomes a strongly increasing function of \(\psi\) while
\( \eta_{st} \) is more flat (since \( J_z \) is constant). The resulting particle confinement time can be written \( \tau_p = \tau_p^{(0)} + \Delta \tau_p \), where \( \tau_p^{(0)} \) corresponds to a circular pinch. Similarly, the energy confinement time is written \( \tau_E = \tau_E^{(0)} + \Delta \tau_E \). The changes \( \Delta \tau_p \) and \( \Delta \tau_E \) are both of order \( A_N^2 \). If the non-circularity is due to an applied magnetic multipole field \( |A_N| \) is limited due to the fact that the plasma boundary must be located inside the separatrix [4]. For \( N >> 1 \) this limit is given by \( \max(|A_N|) = \Lambda/N \), where \( \Lambda \) is of order unity. Fig. 1 shows the relative changes in confinement times divided by \( A_N^2 \) and \( N^2 \). These values should then correspond to those found in experimentally relevant geometries. They are plotted as functions of \( N \) and for the limiting values of \( \lambda \) and \( \beta_0 \). (For \( \{\lambda = 1, \beta_0 \rightarrow \infty\} \), however, \( \Delta \tau_p \rightarrow -\infty \).

4. Conclusions. Although other shaping methods may lead to larger changes we concentrate upon the case where the non-circularity is due to an applied multipole field. Also we are mainly interested in the experimentally relevant parameter range where \( \lambda = 1 \) and \( \beta_0 = 0 \). We find that for \( \eta = \eta_{el} \) a triangular cross-section \( (N = 3) \) is optimal, resulting in improved particle- and energy confinement. For \( \eta = \eta_{st} \), on the other hand, both particle- and energy confinement are made worse by the non-circularity. Apparently the exact form of the resistivity profile is of great importance. However, being second order effects in the non-circularity, the changes in confinement times are usually rather small (\( \leq 15\% \)).

References
Fig. 1 Relative changes in confinement times (divided by $A_N^2$ and $N^2$) for (a) classical resistivity scaling and (b) resistivity consistent with a stationary state.
MHD Equilibrium Generation by Partial RF Current Drive in Reversed Field Pinch Plasma

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1. Introduction

Fast Magnetosonic (FM) wave \( \omega_{m} > \omega > \omega_{e} \) has been found to be an useful candidate for the current drive in high beta plasma as RFP plasma \( V \). The radial profile of the driven current density was theoretically investigated as a function of initial refractive index parallel to the magnetic line of force \( Z_{n} \). Here, theoretical investigation on MHD equilibrium generation by the partial current drive is reported, where the noninductive driven current is combined with ohmic current.

2. RF Current Drive

As the current driver, fast magnetosonic wave \( \omega_{m} \) is used. According to the solution of the dispersion relation, \( G(k, \omega, r) = 0 \) and the wave equation, this wave has the propagation characteristics accessible to high density region with a wide band of \( N_{n} \) (parallel refraction index) spectrum and the main absorption mechanism due to transit time magnetic pumping (TTMP), which is the strong in the higher beta plasma and has a relatively high current driving efficiency.

We note that at this frequency and ion temperature the power loss to the deuterons is negligible small because of \( \omega/\omega_{e} = 10 \).

3. Calculation Scheme of Wave MHD Equilibrium

3.1 MHD Equilibrium Model

We are particularly in RF current drive for TPE-RX RFP device under design study. Here the plasma is modeled as cylindrical plasma and the equilibrium state is described by a partially relaxed state model (PRSM) /3/ and Suydam criterion; MHD equilibrium equation \( dp/dr = j_{B_{z}} - j_{p} \), PRSM \( j_{p} = \lambda B_{0} / \mu_{0} = B_{0} \lambda_{s} \left[ 1 - (\psi_{s} - \psi_{s} - \psi_{s} + \psi_{s}) \right] / \mu_{0} \), and Suydam criterion \( dp/dr = S_{B} / r B_{0} \left( dq/dr \right) / \left( q B_{0} \right)^{2} \). The equilibrium equation combined with PRSM and Suydam criterion is solved numerically by using the Runge-Kutta-Gill method. The electron and ion density profiles are specified as given function of the poloidal flux \( \psi \) where \( \psi \) is the solution to the MHD equilibrium equation with an initial guess for the diamagnetism \( \lambda(\psi) \). The electron and ion temperature profiles are known from the effective charge number \( Z_{e} \) and the pressure profile: \( n_{e} = n_{e} \left[ 1 - (\psi_{e} - \psi_{e} + \psi_{e} - \psi_{e}) \right] \); \( n_{i} = n_{i} \left[ 1 - (\psi_{i} - \psi_{i} + \psi_{i} - \psi_{i}) \right] \); \( T_{e} = T_{e} \left[ 1 - (\psi_{e} - \psi_{e} + \psi_{e} - \psi_{e}) \right] \); \( T_{i} = T_{i} \left[ 1 - (\psi_{i} - \psi_{i} + \psi_{i} - \psi_{i}) \right] \). The wave is launched near the plasma edge, at a density just sufficient to permit wave propagation.

The wave frequency \( \omega/2\pi \), the initial parallel index of refraction \( N_{n} \), and the initial poloidal index \( N_{p} \) are specified at the start. The ray's progress is calculated thereafter from a WKB analysis of the spatial variation in the dielectric tensor, \( K(k, \omega, r) \).
Along the ray trajectory the imaginary part of the k vector is determined in order to compute wave damping: \[ P = p_0 \exp \{-2\gamma (l_1) d_1 \}, \quad \gamma = \text{Im}(k_1), \] where \( P_0 \) is the injected wave power and \( l_1 \) the perpendicular component of the ray path. At each time step, the wave power decrement \( dP \) and the change in flux \( d\psi \) are computed and, with \[ V = \frac{dV}{d\phi} = \frac{f d\phi}{B_0}, \] the electron absorbing power density is found: \[ Q = \frac{dP}{V} \frac{d\psi}{d\phi}. \]

Likewise, at each time step the value of \( k_1 \) is computed and, from the local value of \( T_e \), the ratio of \( v_i \) to electron thermal speed \( v_1 \) is found: \[ u = \frac{\omega}{k_1 v_1}, \] where \( v_i = (T_e/m_e)^{1/2} \). From this we can evaluate the normalized ratio of current density to power density in a uniform medium: \[ \frac{j_1}{P} = \frac{\Omega}{\omega} \left( \frac{Z}{\omega v_1} \right) \left( \frac{k_1}{n_p} \right)^2 \left( \frac{v_1}{\omega} \right)^2. \]

Contributions to \( j_1 \) accumulate as the ray turns and makes multiple passages through the flux surface. We typically terminate the ray tracing once \( F/P_0 = 1 \times 10^{-4} \). It is desirable that the wave power is almost completely absorbed before the ray returns to the wall, so reflection is not a factor.

### 3.3 Convergence

The function \( \lambda(\psi) \) is checked for convergence: it must be accurately known in order to solve the equilibrium equation. Typically, it converges to relative error \( \Delta \lambda/\lambda \) of \( 10^{-3} \) very fast, needing seven iterations between the MHD computations and the ray tracing computations. By this point the parameters of \( I_0 \) (total toroidal current), \( I_0 \) (total poloidal current), \( \beta_N \) (averaged poloidal beta) and \( \Theta \) (theta value) have also converged to three significant figures.

### 4. Fast-Wave Equilibrium by Partial Current Drive

All the equilibria described here have \( 2xR_0 = 9.42 \text{m}, a = 0.3 \text{m} \). The diamagnetism and the toroidal field at the center, \( \lambda \) and \( B_0 \), are determined from \( F \) value, \( a \), and \( l \).

The input variables include the electron peak density and temperature, \( n_e \) and \( T_e \), the density profile exponents, \( m_e \), the Suydam factor, \( S_e \), and the diamagnetism exponents, \( m_n \) and \( n_n \), as well as \( \omega/2\pi \), \( N_i \), and \( P_i \).

A reference case has \( n_e = 2.3 \times 10^{19} / \text{m}^3, T_e = 4 \text{.Kev} \), \( m_e = 5.0 \), \( n_e = 2.0 \), \( S_e = 0.4 \). If the initial guess for \( \lambda(\psi) \) uses \( m = 5.0, n = 2.0 \), it corresponds to an initial equilibrium with \( I_0 = 2.0 \text{MA} \), \( I_0 = 36.1 \text{MA} \), \( \beta_\theta = 8.0 \text{.4\%} \), \( F/\Theta = 0.20/1.7 \) and a monotonic \( q(\psi) \)-profile with \( q(0) = 0.145 \) and \( q(\psi) = -0.027 \). We select a wave with \( \omega/2\pi = 300 \text{MHz} \), \( N_i = 9.9 \) and \( P_i = 5.0 \text{MW} \). By the seventh iteration, the solution converges to an equilibrium with \( I_0 = 2.2 \text{MA} \), \( I_0 = 28.6 \text{MA} \), \( \beta_\theta = 8.0 \text{.4\%} \), \( F/\Theta = 0.20/1.71 \) and a higher shear \( q(\psi) \)-profile with \( q(0) = 0.16 \) and \( q(\psi) = -0.027 \). The global current driving efficiency is \( I_0 / F_0 = 5.1 \times 10^{-3} \) and \( I_0 / P_0 = 4.0 \times 10^{-3} \), where \( I_0 \) and \( I_0 \) are the total driven poloidal and toroidal current, respectively, and \( I_0 > I_0 \).

In order to achieve strong damping, with single pass absorption, the ratio \( u = \omega/k_1 v_1 \), which determines the strength of electron TTMP damping, must be approximately unity. However, this limits the current driving efficiency. For this reference case, the electron power density, \( Q \), peaks off axis, at \( r = 0.20 \text{m} \), where \( u = 0.7 \). We find \( j/p = 17 \). In comparison, the normalized driving efficiency \( j/p \) is larger near the wall \( (j/p = 13 \text{ with } u = 2.0 \) at \( r = 0.75 \text{a} \), but \( Q \) is much smaller by orders of magnitude. We note that the condition \( u < 1 \) makes the driven current decrease because of trapped particle effect.

Fig. 1 shows the converged diamagnetism profile \( \lambda(\psi) \), which determines the final equilibrium. The profile is flatter in core plasma region than in the initial guess. Fig. 2 shows the \( q(\psi) \) profile in the initial and the final equilibrium: the dotted curve is for the initial one and the solid
curve for the final one. The comparison of q profiles indicates that the magnetic shear is enhanced by the wave injection and suggests that the plasma stability is improved over a wide region, especially at r=0.2a.

In the survey of single Ni, we independently vary the input parameters according to our problem, first the target plasma properties and then the wave characteristics.

We first consider a variation of beta value by changing the Suydam parameter S.. All other parameters are held fixed at the values of our reference case, except S=0.8, m., n., T, and P=10MW. One might expect the beta value to rise as the peak pressure is increased.

It is striking, however, that the ratio of total driven current I to injected wave power P., I/P., decreases as beta value increases. This is mainly due to the inverse relation of j, and density. Another effect of varying the beta value is shown by the j, and q profiles. Most of the power is absorbed while the ray reaches r=0.5a because TTMP is strong at high beta. The j, and j profiles have peaks near the boundary. The enough poloidal driven current in the outer region reverses the toroidal field, which could reduce the need for a dynamo mechanism within the plasma. The global current driving efficiency is I,P.,=2.1X103 and I,P.,/P.,=8.4X10-3, which are lower than those at low beta (S=0.4). An important inference drawn from this result is that increases in equilibrium beta value do not necessary require increases in I/P.. In fact, for a reactor it is most encouraging that I may be minimized while, for given P., beta value might be pushed to high values.

Fig.3 shows the converged diamagnetism profile \( \lambda (r) \), indicating to be very flat except the boundary region and approach to Taylor's minimum energy state. Fig.4 shows the q profiles in the initial and the final equilibrium: the dotted curve is for the initial one and the solid curve for the final one. The comparison of q profiles indicates that the magnetic shear is enhanced by the wave injection and suggests that the plasma stability is improved in the outer region.

The wave damping becomes strong with the increase of frequency, and the strongest when \( N \approx 5.5 \) for f=300MHz. The smaller \( N \approx 5.5 \) wave has both the better accessibility to high density region and higher current driving efficiency, however the weaker damping. As the plasma beta value becomes higher, the damping effect due to higher harmonic ion cyclotron resonances is observed when \( N \approx 5.5 \) is w/\omega_0c=10. Therefore the use of higher wave frequency is desired to avoid these resonances.

5. Conclusion

A more stable RFP equilibrium having a higher shear and a flatter \( \lambda \) profile is found with a single ray fast–wave partial current drive calculation. Furthermore the enough poloidal driven current in the outer region leads to the hope that the need for a dynamo mechanism could be reduced, thus the energy confinement time be improved.

The increase in equilibrium beta value does not necessarily require increase in the global current driving efficiency, and makes the plasma approach to Taylor's minimum energy state.

6. Discussion

The spatial dependence of current generation is a function of the wave power profile, the spatial variation of \( u=\omega/k_1v_\perp \) and the electron temperature and density profiles. With a single \( N \approx \) in the wave spectrum, \( j_\parallel \) \( (\psi) \) peaks fairly strongly around one flux surface, and it is difficult to obtain a very broad \( j_\parallel \)-profile which is large at both the magnetic axis and the boundary. In consequence, \( j_\parallel \) tends to be quite small either at r=0 or at r=a. The MHD equilibrium is difficult to obtain for these extreme cases of nearly zero density. The equilibrium by full current drive which does
not need ohmic current may not be achievable unless a broader, properly selected spectrum can be employed. Calculation with multiple N-rays, with the goal of generating the wave-equilibrium, will be reported at a later time.

Fig. 1 shows the profile of converged diamagnetism \( \lambda (r) \) and driven parallel current density \( j_{||} \) (\( f=300 \text{ MHz}, N_i=9.9, S_s=0.4 \)).

Fig. 2 shows the \( q \) and \( q' \) profile in the initial (dotted curve) and final (solid curve) equilibrium (\( f=300 \text{ MHz}, N_i=9.9, S_s=0.4 \)).

Fig. 3 shows the profile of converged diamagnetism \( \lambda (r) \) and driven parallel current density \( j_{||} \) (\( f=300 \text{ MHz}, N_i=9.9, S_s=0.8 \)).

Fig. 4 shows the \( q \) and \( q' \) profile in the initial (dotted curve) and final (solid curve) equilibrium (\( f=300 \text{ MHz}, N_i=9.9, S_s=0.8 \)).

Decay Rate and Particle Confinement Time in SPHEX

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Introduction
The plasma in the SPHEX spheromak is made up of the central column in which the current is driven directly by the Marshall gun which serves as the plasma source, and the surrounding toroidal annulus in which the current is driven indirectly by the rectification of a large-scale $n=1$ oscillation at about 20 kHz which is a constant feature of the system [1]. This mode is generated in the central column and radiates outward across the interface between column and annulus; it is absorbed in the first few cm of the annulus, and the current drive is transferred throughout the annulus by the "MHD dynamo" effect [2]. There is evidence that the column and annulus are separated by a good flux surface, with transport across it due only to the effects of the mode. In this paper we shall examine the balance of particles and angular momentum and show that estimates of the overall particle confinement time and of the recycling rate can be obtained on this basis.

Transport between column and annulus
The plasma in the annulus is supplied by a flux of particles across the interface from the column, and we shall show that recycling from the wall must also be taken into account. The interface flux is associated with the $n=1$ mode and has been determined using a multiple Langmuir probe to measure $\Gamma = \langle nE_\perp \rangle / B$, giving a value of $6 \times 10^{22}$ ions $m^{-2} s^{-1}$ [3]. We assume that in the steady state the interface flux and the recycling together supply the annulus completely; then if $\tau_p$ is the particle confinement time, the particle balance is expressed by
where $A$ is the interface area, $N = nV$ is the total number of ions in the annulus volume $V$, and $R$ is the relative recycling rate defined as the ratio of the total recycling input rate to the particle content of the annulus.

Angular momentum balance

Since the column is rotating, as is shown by the plasma potential profile [1], the interface flux will also carry angular momentum, and this exerts a torque on the annulus given by $\tau = \Gamma A \frac{r_c^2 \omega_c}{N}$, where $r_c$ and $\omega_c$ are respectively the radius and angular velocity of the column. (This is about an order of magnitude greater than the torque due to the Larmor gyration of the ions carried into the annulus.) The annulus is observed to rotate toroidally [4], but much more slowly than the column, while after the discharge termination the angular velocity decreases over a characteristic "spin-down" time of about 300 $\mu$sec. Angular momentum is lost from the annulus when particles reach the wall; however, particle loss by itself cannot slow down the observed rotation, and we must distinguish between the loss of angular momentum which is due to particle loss, and the reduction of angular velocity which can only be due to an input of non-rotating plasma, i.e. recycling from the wall. Thus we can estimate the recycling rate from the observed spin-down time:

$$R = \frac{1}{\tau_\omega} \equiv -\frac{1}{\omega} \frac{d\omega}{dt}$$

The angular momentum of the annulus may be estimated as $L = N m_i r_m^2 \omega$ where $r_m$ is the radius of the magnetic axis, and the relative input rate from the column is then given by $\tau/L = (\Gamma A/N)K$ where $K = (r_c/r_m)^2 \omega_c/\omega$. This is balanced in the steady state by the transfer of angular momentum to the wall due to particle loss,
giving a relative output rate equal to \((Q/\tau_p)\) where \(Q = \langle r_1^2 \rangle / r_m^2 \); \(r_1\) is the major radius at which a particle reaches the wall, and brackets \(<\rangle\) imply an average over all particles lost to the wall. Thus we have

\[
\Gamma A N = \frac{Q}{\tau_p} \quad (3)
\]

Equations (1 - 3) can be solved to lead to the following results:

\[
\frac{1}{\tau_p} = \frac{K}{K-Q} R \quad (4)
\]

\[
\frac{\Gamma A}{N} = \frac{Q}{K-Q} R \quad (5)
\]

where equation (4) may be written in the equivalent form

\[
\frac{1}{\tau_p} = \frac{r_c^2 \omega C}{r_c^2 \omega C - \langle r_1^2 \rangle \omega} R \quad (6)
\]

On the assumption that particles are lost to the wall uniformly over its surface outside the central column, but not to the central column itself, an approximate calculation leads to

\[\langle r_1^2 \rangle^{1/2} \approx 0.37 \text{ m}, \quad \text{which with } r_m = 0.28 \text{ m gives } Q = 1.75. \]

If \(r_c\) and \(Q\) can be assumed to be constant, and if we assume that the potential at the centre of the column is a constant fraction of the gun voltage, these results will allow the determination of the scaling of \(\tau_p\) with gun current using only non-intrusive diagnostics. We shall now show that they are consistent with the known results in the standard operating conditions with a gun current of 60 kA.

**Experimental results**

With \(r_c = 0.10 \text{ m}, r_m = 0.28 \text{ m}, \omega_c/2\pi \approx 40 \text{ kHz and } \omega/2\pi \approx 1 \text{ kHz}\) we obtain \(K \approx 5\). From the results given in ref [4] we have \(1/R = \tau_\omega \approx 300 \mu\text{sec}, \) and from equation (4) we have \(\tau_p \approx 200 \mu\text{sec}. \) Equation (5) then gives \((\Gamma A/N)^{-1} = 520 \mu\text{sec}, \) which is to be compared with the result obtained directly from the experimental values of \(\Gamma, N\) and \(A\), measured in the standard operating conditions with a gun
current of 60 kA. We have \( N = nV \) with \( n \approx 4 \times 10^{19} \text{ m}^{-3} \) and \( V = 0.18 \text{ m}^3 \), and \( A \approx 0.25 \text{ m}^2 \), giving \( (\Gamma A/N)^{-1} = 480 \mu \text{sec} \). The agreement is sufficient to confirm the internal consistency of the model and gives support to the estimate of \( \tau_p \). We conclude that about 60% of the particle influx is due to recycling and 40% is derived from the central column.

The input from the central column is given by \( \Gamma A = 1.5 \times 10^{22} \) ions/sec; for comparison, if we assume the plasma in the central column to be flowing at the ion sound speed we find the total throughput to be about \( 6 \times 10^{22} \) ions/sec, so about 1/4 of the total throughput is diverted into the annulus. These figures may also be compared with the total gas input, which is derived from a set of 6 gas puff valves; we estimate the input rate to be about \( 1.5 \times 10^{23} \) atoms/sec.

Conclusions

Measurements of the toroidal rotation rate of the plasma and the time-constant for its decay after the gun is shorted have allowed us to estimate the particle confinement time in the SPHEX discharge, and to show that the annulus plasma is sustained by recycling and transport from the central column in the ratio 3:2.

References

MINIMUM ENERGY STATES IN SPHEROMAKS WITH EXTERNAL DRIVING

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Introduction

Relaxation theory [1,2], which postulates that a magnetised plasma minimises its magnetic energy whilst conserving magnetic helicity, has been applied to explain field configurations in Reverse Field Pinches and spheromaks. The relaxed state is described by

\[ \nabla \times \mathbf{B} = \mu \mathbf{B} \quad (1) \]

where

\[ \mu = \frac{\mathbf{j} \cdot \mathbf{B}}{B^2} \quad (2) \]

is spatially constant [1]. For spheromaks with some flux embedded in external conductors, relaxed states are possible for all values of \( \mu \), with resonance as \( \mu \) approaches the first eigenvalue \( \mu_1 \) [2].

In a gun-injected spheromak such as CTX [3] or SPHEX [4], current is electrostatically driven on open flux which connects the electrodes, and this external driving causes departures from simple relaxed states. This process of helicity injection sustains the spheromak against resistive decay. Fields are approximately force-free, described by (1) with \( \mu \) varying across the field. Measurements confirm that \( \mu \) is higher on the driven flux and lower near the magnetic axis during sustainment [5] although the profile can be quite complex [6]. Spheromak equilibria with non-constant-\( \mu \) have been modeled, but assumptions such as linear dependence of \( \mu \) on flux, must be made [5,7]: a model is required to predict the \( \mu \) profile for given driving and dissipation. Diffusive models [8] with helicity transport proportional to \( \nabla \mu \) give monotonic \( \mu \) profiles. A new approach, retaining the fundamental assumptions of relaxation theory [1] and yet predicting non-constant \( \mu \) profiles, is outlined here.

The model

During sustainment of a gun-injected spheromak, the current on the open field is determined by the gun, so that

\[ \mu = \mu_\text{gun} = \mu_0 I_\text{gun}/\psi_\text{sol} \quad \text{on the open flux} \quad (3) \]

where \( I_\text{gun} \) is the gun current and \( \psi_\text{sol} \) the open flux from the solenoid. If dissipation and driving are slow compared with relaxation, a relaxed state should be established; according to Taylor [2], this is a constant-\( \mu \) field with the value of \( \mu \) equal to \( \mu_\text{gun} \) given by (3). However, in practice \( \mu_\text{gun} \) is limited by the plasma ejection criterion and always exceeds \( \mu_1 \) (e.g. for SPHEX, \( \mu_1 = 11.1 \, \text{m}^{-1} \).
and \( \mu_{\text{gun}} \) is typically 24 m\(^{-1}\); such fields with \( \mu > \mu_1 \) both have an unrealistic topology [9] and are not absolute energy minima. A more serious difficulty is that the helicity is then determined only \textit{a posteriori} since \( K = K(\mu) \), and cannot, in general, equal the prescribed value. For example, in a spheromak, \( K \) is given by time-integrating the helicity balance equation [3].

We propose that the field is in a minimum energy state subject both to the constraint that the global helicity has a given value \( K [1] \), and that the value of \( \mu \) on the open flux matches that of the helicity source (3). As argued above, this state cannot, in general, be constant-\( \mu \), and thus the minimum energy field must have spatially varying \( \mu \). Intuitively, one might expect \( \mu \) on the closed (undriven) flux to be constant, in general having a different value from the open flux (a piecewise constant \( \mu \) profile); this would indeed be the minimum energy state if the helicity of the closed flux alone were conserved, but because rather the total helicity is constant, which depends on flux linkages between closed and open field regions, there is no reason to expect a constant-\( \mu \) state.

**Large aspect ratio results**

In order to determine the nature of the minimum energy state in such a driven system, we first investigate an idealised 1D model, a large-aspect ratio approximation to the true geometry. Consider the field within an infinite-length periodic cylinder, with \( r = 0 \) corresponding the magnetic axis and a conducting boundary at \( r = a \), so that total toroidal flux \( \phi \) is conserved. Driven field, representing the spheromak open flux, is in the outer layer \( \phi > \phi_1 \), where the amount of driven flux, \( \phi_1 \), is imposed but the spatial location of this surface \( (r = r_1) \) is to be determined. We impose the total helicity \( K \) (per unit length), where

\[
K = \int_{a}^{\infty} A_{z} 2\pi r \, dr
\]

with \( A_{z}(r = a) = 0 \) for gauge invariance; also the value of \( \mu \) on the driven flux \( (\mu_{\text{gun}}) \) is chosen, so that

\[
\nabla \times B = \mu_{\text{gun}} B \quad \text{on} \quad \phi > \phi_1.
\]

The profile \( \mu(r) \) on the undriven flux \( (\phi < \phi_1) \) is then found by minimising the energy subject to these constraints.

A numerical minimisation approach is used. A polynomial for \( \mu \) is chosen, of degree up to 5; \( \mu = \sum a_i r^i \). The force-free equation is integrated from \( r = 0 \) using a 4th order Runge-Kutta until the toroidal flux equals \( \phi_1 \), and then the fields are matched to the analytical Bessel function solution of (6) describing the field in the outer region \( (r_1 \leq r \leq a) \). The values of \( B_z(r=0) \) and \( a_0 \) are then
iterated until the correct helicity \( K \) and toroidal flux \( \psi \) are attained, and the total energy \( 0 \leq r \leq a \) is calculated. The polynomial coefficients \( a_i \) are varied until an energy minimum is found, using a quasi-Newton algorithm.

A typical \( \mu \) profile is shown in Fig. 1. The helicity \( K \) is chosen to be that of a constant-\( \mu \) field with \( \mu = 2/a \). Note that the energy of this field is 242540 J whereas that of the piecewise-constant profile with the same helicity is 242887 J. Thus the minimum energy state is clearly not a constant \( \mu \) field in the undriven flux, but rather has a \( \mu \) profile in this region. The field with \( \mu = 2 \) everywhere has of course the same helicity and lower energy (242309 J) - but this does not satisfy the current constraint in the driven region. A sharp dip in \( \mu \) compensates for the "too large" \( \mu \) on the driven flux, with \( \mu \) flattening towards the Taylor value near the magnetic axis. Thus, the field \( B \) is close to the Taylor field over the undriven flux. The \( \mu \) profile is not monotonic over the inner region, but exhibits a distinct bump. These results have been verified both by using a sum of Chebyshev polynomials for \( \mu(r) \) and by applying a partial analytical minimisation procedure, leading to a set of ODEs.

Results in SPHEX geometry

The alternative minimum energy principle has also been used to find equilibria in SPHEX geometry. The methodology is similar, except that the equilibrium is found as a numerical solution to the Grad-Shafranov equation with given \( \mu(\psi) \) and gun flux [7]. On the open flux \( (\psi < \psi_{sat}) \mu = \mu_{gun} \) whilst the \( \mu \) profile on the closed flux is found by choosing the coefficients of a polynomial in \( \psi \) to minimise the total energy. The profiles found are qualitatively similar to Fig. 1, in particular having a dip in \( \mu \) near the edge of the closed flux and thus giving good agreement with the experimentally measured \( \mu \) profile [6]. These fields have significantly lower energy than piecewise constant \( \mu \) profiles. For example, with a total helicity of 0.254 mWb², a minimum energy state is found with energy 0.765 kJ whereas if \( \mu \) is constant over the closed flux, the energy for this helicity is 0.784 kJ.

Conclusions

A previous inconsistency in the application of relaxation theory, in that the constraint of fixed helicity is imposed in order to derive the constant-\( \mu \) condition but that the helicity of the calculated field cannot be set equal to the known helicity content of the system, has been resolved. In driven systems such as gun-injected spheromaks, the minimum energy state subject to the two constraints of helicity conservation and applied current on the driven field has a
spatially varying $\mu$ profile. This profile has a rather complex form in spheromak geometry, which corresponds closely to experimental measurements of $\mu$. Whilst more accurate numerical energy minimisation might produce a further fine-tuning of the $\mu$ profile, it has been shown conclusively that the calculated "bumpy" profiles have significantly lower energy (for the given constraints) than fields with constant-$\mu$ on the undriven flux.

The model may be similarly applied to fields with resistive wall layers, where $\mu$ is effectively constrained to be low, such as in the decay phase of spheromaks. In this case, a peak in $\mu$ in the outer layers of the closed flux is expected, giving fields with quite different stability properties from the linear $\mu$ profiles which have been the focus of previous studies. These results also have implications for helicity injection current drive of tokamaks, since we have shown, in contrast to diffusive models [8], that such current drive is not necessarily associated with monotonically decreasing $j/B$ from axis to edge.


**Fig 1** The $\mu$ profile in a cylinder (radius $a$ = 1) for the minimum energy state with helicity equal to that of the constant-$\mu$ field with $\mu = 2$ and driving ($\mu = 4$) in the outer 10% of flux.
The response to a fast changing toroidal magnetic field in ULQ plasma

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1. Introduction
The Ultra-low-q (ULQ) configuration is a toroidal current carrying system with the safety factor $q_a$ of $0 < q_a < 1$. The equilibrium of ULQ configuration is characterized by the profile with a pitch minimum of the magnetic shear in the bulk plasma region. Similar to RFP plasmas, dynamo effect plays an important role in ULQ plasmas, and strong ion heating has been observed, accompanied by the resistance anomaly$^1$.

Ohm's law is given by $\eta \mathbf{v} = \mathbf{E} + \langle \mathbf{v} \times \mathbf{B} \rangle$, where a dynamo electric field is represented by the term of $\langle \mathbf{v} \times \mathbf{B} \rangle$. A poloidal current is, in general, sustained by the poloidal component of the dynamo electric field. To investigate the role of dynamo electric field on the plasma performance, the externally induced poloidal electric field has been applied on ULQ plasmas, by changing the vacuum toroidal magnetic field rapidly.

2. Experimental results of fast changing toroidal magnetic field
1) Experiments with a fast decrease of the toroidal magnetic field
The toroidal magnetic field is rapidly decreased during the ULQ discharge. The time evolutions of plasma parameters are shown in Fig. 1, where the ULQ configuration with $q_a \sim 0.8$ is initially set up, and the toroidal magnetic field $B_t$ is decreased quickly with a time constant of $\sim 500 \mu$sec. Synchronized to the decrease of the toroidal magnetic field, the increase of the toroidal plasma current is observed, accompanied by the decrease of loop voltage. This result is not explained by the empirical scaling that in ULQ plasmas the resistance anomaly becomes larger, as the toroidal magnetic field is reduced$^2$). Figure 2 shows the reduction rate of the resistance as a function of the change rate of the toroidal magnetic field, where the negative value of $dB/dt$ corresponds to the fast decreasing experiments of the toroidal magnetic field. The decrease of the resistance anomaly becomes larger, as the change rate of the toroidal magnetic field
is increased. The improvement of the plasma performance can be accounted for the poloidal electric field induced by the decrease of the toroidal magnetic field. It is inferred that a poloidal component of dynamo electric field, which is strongly coupled with plasma fluctuations, is partly replaced with the induced poloidal electric field, because the direction of this induced field is the same with that of dynamo electric field to sustain the poloidal current.

The time evolution of the safety factor profile is shown in Fig. 3. The configuration with pitch minimum is kept during the decreasing phase of the toroidal magnetic field. The change of the safety factor at the plasma surface is relatively continuous function of time. While around the plasma center, the response of the safety factor seems to be stepwise. The safety factor around the plasma center does not change so much during $t=1.47 \sim 1.92$ ms, and decreases quickly at $t=1.95$ ms. The stepwise decrease of the plasma current has also been observed in constant toroidal field ULQ experiments when the safety factor passes through the rational surface$^3$.

2) Experiments with a fast rise of the toroidal magnetic field.

During ULQ discharge the toroidal magnetic field is rapidly increased, just similar to the adiabatic compression experiments in tokamak plasmas$^4$. The time evolutions of plasma parameters are shown in Fig. 4, where the toroidal magnetic field is increased from $\sim 0.8$ kG to $\sim 2.4$ kG. During the increasing phase of the toroidal magnetic field, the increase of the loop voltage is observed, accompanied by the enhancement of magnetic fluctuations. While, the plasma current is sustained almost constant. The increase of the resistance anomaly due to the increase of the toroidal magnetic field is depicted in Fig. 2, and these results also support that the poloidal electric field induced by the fast change of the toroidal magnetic field plays an essential role for the deterioration of the plasma performance.

Just after the increase of the toroidal magnetic field is terminated, the plasma current begins to increase slowly, accompanied by the decrease of the loop voltage, as shown in Fig. 4, and the maximum plasma current might be limited by large MHD activity. During the increasing phase of the toroidal magnetic field, the safety factor at the wall increases, and has a maximum value $q_{a,max}$ at the end of the increasing phase.

Subsequently, the safety factor begins to decrease, and has a minimum settled value $q_{a, settle}$. Figure 5 shows the correlation between $q_{a,max}$ and
It seems that the values of $q_{a,\text{settle}}$ are spreading around rational numbers of $2/3$, $1/2$ and $2/5$. The elevated toroidal magnetic field may account for the improvement of the plasma performance, as observed in constant toroidal field ULQ experiments.

3. Conclusions

To investigate the role of dynamo electric field in ULQ plasmas, the toroidal magnetic field is rapidly changed during ULQ discharges. When the toroidal magnetic field is decreased/increased), the improvement (deterioration) of the resistance anomaly is observed. This might be account for the poloidal electric field induced by the change of the toroidal magnetic field, because the dynamo electric field is partly replaced/enhanced) with the externally induced poloidal electric field. In the case of the decrease of the toroidal magnetic field, the stepwise change of the safety factor is observed around the plasma center. When the toroidal magnetic field is increased, gradual improvement of the plasma performance is observed just after the increase of the toroidal magnetic field is terminated, and the safety factor seems to settle to rational numbers.

REFERENCES


Fig. 1
ULQ discharge with a fast decrease of the toroidal magnetic field; time behavior of the plasma current ($I_p$), loop voltage ($V_{\text{loop}}$), the toroidal magnetic field ($B_t$), and safety factor at the wall ($q_a$).
Fig. 2
The reduction rate of the resistance as a function of the change rate of the toroidal magnetic field.

Fig. 3
Time evolution of safety factor $q(r)$, as a function of the minor radius $r/a$.

Fig. 4
ULQ discharge with a fast rise of the toroidal magnetic field; time behavior of the plasma current ($I_p$), loop voltage ($V_{loop}$), the toroidal magnetic field ($B_t$), safety factor at the wall ($q_a$) and magnetic fluctuations.

Fig. 5
The safety factor at the wall just after the increase of the toroidal field, $q_{a,max}$, versus the safety factor at the flat top of the plasma current, $q_{a, settle}$. 
Anisotropic Stabilization of Internal Tilting FRC

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Abstract

A preliminary study on internal tilting stability of an anisotropic model for field-reversed configuration (FRC) is presented. Using a trial function approach together with the corresponding anisotropic Energy Principle a stabilizing influence of anisotropy is shown.

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After the prediction of Rosenbluth and Bussac [1], that field-reversed configurations (FRC) should be unstable to internal tilting when the elongation of their separatrix (ratio of half length to maximum radius) is greater than one, the mode was confirmed within ideal MHD theory, first numerically [2], and shortly after that analytically [3]. Predicted e-folding times for typical experimental data (for a general review of FRC experiments see Ref. [4]) are of the order of one to a few microseconds. However, a paradoxical situation, which puzzled researchers over a decade, arose, since practically all FRC experiments showed a robust stability to such mode. As the exception that confirms the rule, Tuszewski et al. reported evidence of the presence of the tilting instability and its destructive consequences on FRX-C/LSM machine, when operating at low filling pressures and with relatively low separatrix elongations ($\approx 3.4$) [5].

Several mechanisms have been proposed in order to explain this discrepancy between theory and experiments: gyroscopic effects due to plasma rotation [6], kinetic effects (via Vlasov-fluid numerical codes) [7], finite Larmor radius effects and Hall terms in Ohm’s law in two-fluid models [8, 10]. However, none of these mechanisms has been completely successful in explaining all the evidence of stability to tilting and the general consensus was that the ideal MHD model is inadequate for this kind of problem.

Here, we want to draw attention to a possible new interpretation of experimental results in terms of an anisotropic plasma pressure within an ideal fluid model. Looking at typical experimental data it can be seen that predicted tilting e-folding times are of the same order of, or smaller than, typical classical FRC ion-ion collision times. In such conditions, it is hard to assume that the plasma obeys a simple adiabatic law and it should be better to assume, at least, double adiabatic equations, allowing for anisotropicity. Starting with an anisotropic equilibrium state, double adiabaticity allows for stabilizing mechanisms in the special case of tilting modes.

For internal modes and anisotropic equilibrium states, the Energy Principle gives the following expression for the potential energy associated to the modes [11, 12]

$$\delta W = \frac{1}{2} \int d^3x \left[ |Q|^2 - \tilde{Q} \cdot \tilde{\xi} + (\nabla \cdot \tilde{\xi}) (\tilde{\xi} \cdot \nabla_{P,1}) + \frac{5}{3} P_{1} (\nabla \cdot \tilde{\xi})^2 \right. $$

$$\left. + \frac{\tilde{B} \cdot (\nabla \tilde{\xi})}{B^2} \nabla \cdot (\sigma_{-} B^2 \tilde{\xi}) - \sigma_{-} \left[ \tilde{Q} + \tilde{B} \cdot \tilde{\xi} + \frac{\tilde{B} \cdot \tilde{\xi}}{B^2} \right) \right] \nabla B^2 \right]$$

(1)
\[
\left\{ [\dot{Q} + \dot{B} \nabla \cdot \dot{\xi} + \nabla (\dot{\xi} \cdot \ddot{B}) - \ddot{\xi} \times \dddot{B}] + 4\sigma_{-} \left[ \frac{\dot{B} \cdot (\ddot{B} \cdot \nabla \dddot{B})}{B^2} \right] + \frac{p_{\perp}}{3} \left[ \nabla \cdot \dddot{\xi} - 3 \frac{\dot{B} \cdot (\ddot{B} \cdot \nabla \dddot{B})}{B^2} \right]^2 \right\},
\]

(2)

where

\[
\begin{align*}
\dot{Q} &= \nabla \times (\dddot{\xi} \times \ddot{B}), \\
\sigma_{-} &= \frac{p_{\parallel} - p_{\perp}}{B^2}, \\
B^2 &= \dddot{B} \cdot \dddot{B}.
\end{align*}
\]

For axisymmetric FRC \( \dddot{B} \) can be written in terms of a stream function as \( \dddot{B} = \nabla \psi \times \dddot{\phi} / r \) and \( \sigma_{-} \) can be assumed to vanish at the separatrix and at the O-point. In such case the fourth and last terms in expression (2) suggest that a background \( p_{\parallel} \) at the separatrix should always allow to stabilize (i.e., \( \delta W > 0 \)). However, when \( \nabla \cdot \dddot{\xi} = 0 \) and \( \dddot{B} \cdot (\dddot{B} \cdot \nabla \dddot{\xi}) = 0 \) the background is ineffective and this suggest that the most dangerous displacement should be divergence free (as is the case in all previous tilting studies) and satisfying also \( \dddot{B} \cdot (\dddot{B} \cdot \nabla \dddot{\xi}) = 0 \). This last condition corresponds to tilting mode like rigid axial shift of the magnetic surfaces plus an azimuthal component that makes \( \nabla \cdot \dddot{\xi} = 0 \). Modes of this kind can be represented by [13]

\[
\dddot{\xi} = \xi_{1}(\psi) \cos \phi \dddot{e}_z - r \frac{\partial \dddot{e}_z}{\partial \phi} \sin \phi \dddot{e}_\phi.
\]

(3)

In order to do a trial function study of the tilting mode in anisotropic FRC, we consider an anisotropic equilibrium which is a natural extension of the Hill's vortex model [14]:

\[
\frac{\psi}{\psi_{0}} = \frac{1}{\sigma_{0}} \left[ 1 - \left( 1 - \frac{3}{2} \sigma_{0} U_{hv} \right)^{2/3} \right],
\]

(4)

where

\[
\begin{align*}
U_{hv} &= U_{0} \left( \frac{r^2}{a^2} \left( 1 - \frac{r^2}{a^2} \right)^{2/3} \right), \\
U_{0} &= \frac{2}{3\sigma_{0}} \left[ 1 - (1 - \sigma_{0})^{3/2} \right].
\end{align*}
\]

being \( \sigma_{-} = \sigma_{0} \frac{\psi}{\psi_{0}} \), with \( |\sigma_{0}| \) constant smaller than unity and \( \psi_{0} \) the maximum value of \( \psi \) inside the FRC, \( a \) and \( b \) are the semiaxes of the elliptical separatrix.

As trial functions for \( \xi_{2} \) we have considered:

\[
\begin{align*}
\xi_{11} &= \left( \frac{U_{hv}}{U_{0}} \right)^2, \\
\xi_{12} &= \left( \frac{U_{hv}}{U_{0}} \right)^2 \left[ 1 - \left( \frac{U_{hv}}{U_{0}} \right)^2 \right].
\end{align*}
\]

(5)

(6)

For \( \sigma_{0} = 0 \), isotropic case \( \xi_{11} \) makes \( \delta W < 0 \) when \( b/a > 2.357 \) and \( \xi_{12} \) when \( b/a > 3.85 \). When \( \sigma_{0} \neq 0 \) the critical elongation for \( \delta W > 0 \) is a function of \( \sigma_{0} \). In fig. 1 we plot the critical \( b/a \) as a function of \( \sigma_{0} \) for the trial function \( \xi_{11} \). For \( \xi_{12} \) a different curve is obtained. As in fig. 2. \( \xi_{12} \) represents a mode concentrated away from the O-point, and shows a stabilizing influence of \( p_{\perp} > p_{\parallel} \) in the FRC. \( \xi_{11} \) represents a mode concentrated near the O-point; in this case both \( p_{\perp} > p_{\parallel} \) and \( p_{\perp} < p_{\parallel} \) have stabilizing
influence. This suggests that $p_1 > p_\parallel$ should have a stabilizing influence far from the O-point while $p_\perp < p_\parallel$ should stabilize near the O-point.

These are preliminary results and a more detailed examination of the function approach is needed. The fact that $\vec{B} \cdot (\vec{B} \cdot \nabla \xi) = 0$ and $\nabla \cdot \xi = 0$ ensure that since at equilibrium $p_\perp \approx p_\parallel$ near the O-point, the $p_\parallel$ and $p_\perp$ evolution equation converge continuously to the isotropic model in that region (as it should be desirable). What kind of anisotropy ($\sigma_0 > 0$ or $\sigma_0 < 0$) could exist in the bulk of an FRC still remain open question but we think that the observed tilting stability in experiments should be an evidence of its existence.

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**References**


Fig. 1 - Critical elongation as function of $\sigma_o$ for $\xi_{z1}$

Fig. 2 - Critical elongation as function of $\sigma_o$ for $\xi_{z2}$
Experiments of Ultra Low q Equilibrium with $q_a < 0.1$

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1. Introduction

We have studied the regime intermediate between the tokamak and the RFP, with $0 < q < 1$ which is called the ultra low q (ULQ) regime. The ULQ experiments with $q_a > 0.1$ have been carried out in REPUTE-1 [1], HBTX-1C [2], and OHTE [3]. The ULQ equilibrium with $q_a < 0.1$ is observed first on SWIP-RFP device [4]. The experimental results of ULQ configuration are given with a passive crowbar, the plasma current duration $\tau_p$, which is over $300\mu s$, is obtained.

2. Experimental results

The ULQ configuration can be established with a passive crowbar and under following conditions: filling pressure $p_0(H_2) = 1$ to $2p_a$, charging voltage of poloidal field bank $V_p = 12.6$ to $20$KV, charging voltage of initial toroidal field bank $V_i = 0.68$ to $2.5$KV and charging voltage of reversed field bank $V_r = 0.56$ to $1.56$KV. The ULQ experiments have been carried out with the surface safety factor $q_s$ less than 0.1.

Fig. 1 shows the toroidal voltage of plasma $V_r$, plasma current $I_p$ and the value of toroidal magnetic field averaged over the minor cross section $\langle B_\phi \rangle$, poloidal signal $\frac{dB_\phi}{dt}$ and $\frac{dB_p}{dt}$. Toroidal voltage $V_r$ is $353.5$ V and the pulse lengths of plasma current is longer than $360\mu s$. The toroidal field of plasma surface $B_{\phi r}$ is $0.019$ T, the poloidal field $B_{pz}$ is $0.055$T. The hydrogen filling pressure is $1pa$.

Fig 2 shows the time evolution of the safety factor $q$ for three different locations of radius $r$ and the time evolution of the
poloidal field near the wall of vacuum chamber \( B_w \) and the reverse parameter \( F = B_w / \langle B_p \rangle \) and the pinch ratio \( \Theta = B_w / \langle B_p \rangle \).

Typical results of ULQ are shown in Table 1

Table 1 The distribution of ULQ in different conditions and other values for \( t = 210\mu s \)

<table>
<thead>
<tr>
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<th>1</th>
<th>2</th>
<th>3</th>
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<tr>
<td>( V_i, KV )</td>
<td>1.70</td>
<td>2.04</td>
<td>2.04</td>
<td>2.55</td>
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<tr>
<td>( V_r, KV )</td>
<td>0.73</td>
<td>0.92</td>
<td>0.92</td>
<td>1.10</td>
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<tr>
<td>( V_p, KV )</td>
<td>19.76</td>
<td>17.96</td>
<td>17.96</td>
<td>20.65</td>
</tr>
<tr>
<td>( P_o, Pa )</td>
<td>1.0</td>
<td>1.5</td>
<td>1.75</td>
<td>2.0</td>
</tr>
<tr>
<td>( \langle B_p \rangle, T )</td>
<td>287.9</td>
<td>305.5</td>
<td>393.5</td>
<td></td>
</tr>
<tr>
<td>( \langle B_w \rangle, T )</td>
<td>0.035</td>
<td>0.046</td>
<td>0.038</td>
<td>0.035</td>
</tr>
<tr>
<td>( q_1 )</td>
<td>0.056</td>
<td>0.022</td>
<td>0.056</td>
<td>0.028</td>
</tr>
<tr>
<td>( q_2 )</td>
<td>0.293</td>
<td>0.130</td>
<td>0.158</td>
<td>0.151</td>
</tr>
<tr>
<td>( q_3 )</td>
<td>0.531</td>
<td>0.376</td>
<td>0.224</td>
<td>0.278</td>
</tr>
<tr>
<td>( \tau_p, \mu s )</td>
<td>360</td>
<td>380</td>
<td>360</td>
<td>380</td>
</tr>
</tbody>
</table>

The profiles of safety factor \( q \) with minor radius at three different times \( t_1 = 80\mu s, t_2 = 170\mu s \) and \( t_3 = 290\mu s \) are shown in Fig. 3. Fig. 2 and Fig. 3 indicate that \( q \) is decreased as the minor radius increases. The measurements show a \( q \)-profile with \( \frac{dq}{dr} < 0 \). Fig. 3 (a) and (c) show \( q \)-profiles, when \( t = 80\mu s \) and \( t = 290\mu s \), it can be seen that there is no a pitch minimum. Fig. 3 (b) shows \( q \)-profiles when \( t = 170\mu s \) with a pitch minimum characterizing Suydam instability near the plasma edge. The distribution of ULQ configuration is shown in Fig. 4. The distributions of \( B_\phi \) and \( B_z \) are given for \( t_1 = 80\mu s, t_2 = 170\mu s \) and \( t_3 = 290\mu s \).

In our experiments the typical plasma parameters are: plasma
current $I_p = 40$ KA, toroidal voltage $V_\varphi = 353.5$ V, the edge safety factor $q_s$ is about 0.056, $P_0(H_2) = 1$ Pa. The electron temperature $T_e \sim 28$ eV and the electron density $n_e \sim 5.7 \times 10^{13}$ cm$^{-3}$ is measured by electrostatic probes. The horizontal displacement of the plasma column is about 8.2 to 9.6 mm.

3. Conclusion

The ULQ discharges with $q_s < 0.1$ are obtained in SWIP-RFP device when the filling gas pressure is sufficiently high [$P_0(H_2) = 1$ to 2 Pa]. Typically, the edge safety factor $q_s$ is $0.05 < q_s < 0.1$. The experimental results show that the distribution of $q$ is decreased with the minor radius increase. The ULQ stable state last for more than $300\mu$s is observed. The profiles of safety factor near the plasma edge with a pitch minimum characteristic is measured at $t = 170\mu$s but when $t = 80\mu$s and $t = 270\mu$s the profiles of $q$ are without a pitch minimum.

Robinson has shown that a $dq/dr < 0$ profile is stable against the global MHD mode. The experiment results are consistent with ideal MHD theory$^{[5]}$ and similar to characteristic of simulation study$^{[6]}$ [The former see Fig. 2, Fig. 3(a) and (c), the later see Fig. 3(b)].

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References

Nonlinear stability of the internal $m = 1$ modes in a pure Z-pinch

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1. Introduction. According to linear, ideal MHD theory the internal region of a pure Z-pinch (i.e. a pinch having no axial magnetic field) is strongly unstable with respect to $m = 1$ (kink) modes [1]. However, various fusion devices based on the Z-pinch often show experimental lifetimes much longer than the characteristic Alfvén time [2]. While numerous attempts have been made to explain this anomalous stability in terms of various non-ideal phenomena (see, e.g., the references listed in Ref. 3), the nonlinear evolution of these instabilities does not seem to have been calculated previously. In the present work the result of such a calculation, based on bifurcation theory, is presented.

2. Linear theory. We consider a cylindrical Z-pinch with uniform current- and mass-density. The details of the analysis summarized below can be found in Ref. 3. For shorter notations we use dimensionless quantities, and express all lengths in units of the plasma radius and the time in units of the Alfvén time $\tau_A = 2\pi r(\mu_o\rho_o)^{1/2}/B_0$ (independent of the radius). Then, using the linearized, incompressible ideal MHD equations, the dispersion relation for the internal (fixed boundary) $m = 1$ modes is obtained in the form [3]

$$\tau k J_1(\tau k) + \frac{2}{1 - \omega^2} J_1(\tau k) = 0$$

where $k$ denotes the axial wave number of the perturbation, $\tau^2 = 4/(1 - \omega^2)^2 - 1$, and $J_1(x)$ is the Bessel function of the first kind, and of order one. Figure 1 shows a few dispersion curves corresponding to the (infinitely many) unstable roots to Eq.(1).

Fig. 1
Dispersion curves for the unstable, internal (fixed boundary) $m = 1$ modes in an incompressible, ideal MHD Z-pinch with uniform current distribution.
It is seen that each mode in the figure has a stable regime, $0 \leq k < k_c$, an unstable regime, $k_c < k \leq \infty$, and one marginal point, $k = k_c$, where $\omega^2 = 0$. Thus, a possible analytical approach to the nonlinear evolution of these modes is to calculate the bifurcated equilibria for the nearly marginal wave numbers $k = k_c$.

3. Nonlinear theory. The equation describing the nonlinear evolution of any nonresonant, ideal MHD mode near marginal stability is given by [4]

$$\eta_{tt} + D_1(k) \eta + D_3(k) \eta^3 = 0$$

(2)

where $\eta(t)$ stands for some suitable quantity representing the amplitude of the mode. Here $\eta$ is chosen as the amplitude of a helical displacement of the magnetic axis. The equation above is obtained by expansion of the ideal MHD equations up to order $\delta^3$, with $\eta$ assumed to be a quantity of order $\delta$. Furthermore, it is assumed that $k - k_c = O(\delta^2)$, where $D_1(k_c) = 0$, and $D_3 = O(\delta)$. Thus, within the range of validity of the equation above, the linear coefficient $D_1(k)$ has the form $(k - k_c)D_1(k_c) = O(\delta^2)$, and the the nonlinear coefficient is given by $D_3(k_c) = O(1)$.

The linear coefficients $D_1(k)$ for the various modes in Fig. 1 are given implicitly through the dispersion relation, Eq. (1) $[\omega^2 = D_1(k)]$. Furthermore, the calculation of the nonlinear coefficients $D_3(k_c)$ can be based on the existence of the equilibrium solutions

$$\eta_{eq} = \pm \sqrt{-\frac{D_1(k)}{D_3(k)}}$$

(3)

to Eq. (2). Since $D_1(k)$ changes sign at $k = k_c$, such solutions always exist for $k = k_c$, irrespective of the sign of $D_3(k_c)$. Moreover, it is easily realized that the nonlinearity is stabilizing if the equilibria (3) exist on the unstable side of the marginal points [where $D_1(k) < 0$, i.e. $k > k_c$ in Fig. 1], and vice versa.

By use of the technique described in Ref. 5 for calculating helical, ideal MHD equilibria consistent with a straight equilibrium with respect to the constraints of i) axial flux conservation, ii) azimuthal flux conservation, and iii) the equation of state (incompressibility in this case), the coefficients $D_1(k)$ and $D_3(k)$ in Eq. (2) are [apart from a common factor that is of no importance as regards the value of the amplitude (3)] obtained in the form [3]

$$D_1(k) = \frac{4}{3\sqrt{3}k} J_1(\sqrt{3}k) + \frac{2}{3} J_1'(\sqrt{3}k)$$

(4)
\[ D_3(k) = \frac{2}{3} P(0) - \frac{1}{3} P_\rho(1) - \frac{4}{9\sqrt{3}k^3} [J_1(\sqrt{3}k)]^3 (k^4 + 6k^2 + 5) \]  

where

\[ P(\rho) = \frac{\pi}{2} \int_0^\rho W(y,y',\rho') \left[ Y_1(\sqrt{3}k\rho)J_1(\sqrt{3}k\rho') - J_1(\sqrt{3}k\rho')Y_1(\sqrt{3}k\rho) \right] \frac{d\rho'}{\rho'}, \]  

\[ W = -\frac{8}{27\sqrt{3}k^2} \left[ \rho^{-1} y_\rho^2 (5k^2\rho^2 + 15) + \rho^{-2} y_\rho y_\rho' (27k^4\rho^4 + 4k^2\rho^2 - 21) \right. \]
\[ + \left. \rho^{-3} y_\rho^2 (33k^4\rho^4 + 4k^2\rho^2 - 15) + \rho^{-4} y_\rho^2 (9k^6\rho^6 + 30k^4\rho^4 - 22k^2\rho^2 + 21) \right] \]  

and \( y(\rho) \equiv J_1(\sqrt{3}k\rho) \). The bifurcated equilibria (3) resulting from the expressions above, with \( D_1 \to (k - k_c)D_1(k_c) \) and \( D_3 \to D_3(k_c) \), are shown by the solid curves in Fig. 2.

\[ \text{Fig. 2} \]

The bifurcated equilibria around the marginal points in Fig. 1, calculated with the exact forms of \( D_1 \) and \( D_3 \) in Eq. (3) (full curves), and by means of the asymptotic expressions (8) and (9) (broken curves), valid for large \( k \).

It is seen that all bifurcated equilibria exist on the unstable side of the marginal points, indicating that all internal \( m = 1 \) modes in the pure Z-pinch are stabilized nonlinearly. In the case of short axial wavelengths, i.e. \( k \gg 1 \), the expressions for \( D_1(k) \) and \( D_3(k) \) above can be simplified by use of the asymptotic expressions for the Bessel functions \( J_1(x) \) and \( Y_1(x) \) with large arguments. In this limit one finds that

\[ k_c = \frac{\sqrt{3}\pi}{4} + n \frac{\pi}{\sqrt{3}} \quad (n \text{ is a large integer}) \]  

\[ \eta_{eq}^2(k, k_c) \approx \frac{\pi}{4} \frac{3\sqrt{3}}{k_c^2} \frac{k - k_c}{\ln[k_c] + \alpha} \quad (k > k_c \gg 1) \]  

where \( \alpha = 1.84 \). The approximation above is shown by the broken curves in Fig. 2.
4. Summary and conclusions. The nonlinear evolution of the internal m = 1 kink instability in a pure Z-pinch (no axial magnetic field) with uniform current distribution has been calculated by means of bifurcation theory, using an incompressible, ideal MHD model of the plasma. It is found that, within the range of validity of this theory, all internal m = 1 modes are stabilized nonlinearly. Although the special case of a uniform current distribution was considered here, it can be argued that [3] in the regime of short axial wave lengths, the result is probably true also for more general profiles, due to the localization of the eigenmodes close to the pinch axis.

In spite of the simple model used, we think that the present work illustrates a phenomenon that could be of importance for the anomalous plasma stability observed in various fusion configurations based on the Z-pinch [2]. In particular, the results obtained indicate that nonlinear MHD effects might be able to explain the absence of destructive internal instabilities, as predicted by linear theory [1], in these experiments. However, we realize that, since the conditions for validity of the ideal MHD model are at best only marginally satisfied in many cases, various non-ideal phenomena are probably also of importance. Furthermore, since the bifurcation theory is limited to the nearly marginal modes, the nonlinear evolution of the strongly unstable modes has to be investigated by other methods, presumably numerical. Thus, a complete picture of the internal Z-pinch dynamics can probably only be arrived at via numerical simulations, involving both non-ideal as well as nonlinear effects.

References


INVESTIGATION OF THE ENHANCED RATE OF MAGNETIC FIELD PROPAGATION ALONG THE ANODE IN Z-PINCH

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The behavior of a plasma near the electrodes in various plasma systems with strong electric current is appreciably influenced by the Hall effect. This phenomenon appears to be of great interest for Z-pincher-like systems.

Consistent description of this phenomenon proves to be possible only within the framework of two-dimensional MHD modeling with overall account of the Hall effect. The results of such a modeling of Z-pincher plasma sheath dynamics for cylindrical geometry are presented here.

We consider dynamics of Z-pincher plasma within the volume bounded by the electrodes. Assuming the equality of ion-electron temperatures, we have the following set of eq [1,2] for plasma density n, hydrodynamic velocity V, magnetic field H and temperature T:

\[ \frac{\partial n}{\partial t} + \nabla \cdot (nV) = 0, \quad (1) \]

\[ \frac{\partial H}{\partial t} = \nabla \times (V \times H) - \nabla \times \left( \frac{c}{4\pi n_m} \nabla \times H \right), \quad (2) \]

\[ \frac{\partial V}{\partial t} + \nabla \cdot (V \otimes V) = - \frac{2}{m_i} \nabla T + \frac{1}{4\pi n_m} [\nabla \times H, H], \quad (3) \]

\[ j = \frac{c}{4\pi} \nabla \times H \quad (4) \]

\[ \frac{3}{2} \left[ \frac{\partial (nT)}{\partial t} + \nabla \cdot (nTV) \right] = - nT \nabla V + j^2/\sigma, \quad (5) \]

where \( \sigma \) is classical conductivity, \( j \) is current density, and \( m_i \) is ion mass. The boundary conditions for Eqs. (1-5) are follows:

\[ V(V_r, V_z) \mid z=0, z=Z = 0; U(U_r, U_z) \mid z=0, z=Z = 0; \quad (1/\sigma \partial H/\partial z) \mid z = 0, \]

where \( U=j/en \) is current velocity. The boundary plane \( z=Z \) correspond to anode and cathode, respectively.

At time \( t=0 \) plasma has the form of the column of radius \( a \) surrounded by magnetic field \( H_\varphi \): \( H = H_0 R_0/r \) for \( r>R_0 \) and \( H=0 \) for \( r< R_0 \), with homogeneous initial plasma temperature \( T_0 \) and density \( n_0 \). The modeling of Eqs. (1-5) is carried out for various values.
of parameter $\Pi_{1}^{-1/2} = \left(\frac{mc^2}{4\pi n_0 e^2 R_0^2}\right)^{1/2}$, which is the ratio of Hall electric field $1/(\text{sec})[j,H]$ to inductive electric field $1/c[V,H]$, namely $\Pi_{1}^{-1/2} = 0.065$, 0.079, and 0.161, and for magnetization parameter $\omega_{eh}\tau_e = 100$.

Figures 1a-1c present the computer-produced plots of the plasma density $n(r,z)$ at the left part and magnetic field $H(r,z)$ at the right part at the moment when magnetic field reaches the axis of the system. The grey-levels of the distributions in Fig.1 differs by a factor 2. The results of computation show that the front of the magnetic field overtakes plasma sheath for $\Pi_{1}^{-1/2} > 0.065$. This condition appears to be a criterion for the penetration of the magnetic field into a nonperturbed plasma near the anode which, in turn, is a reason for the slipping of the current sheath along the anode surface. The radial component of the current gives rise to a "magnetic snowplough" pushing the plasma away from the anode. As the plasma moves away from the anode the sheath mass in the region near the anode decreases and the radial velocity increases here. That behavior of Z-pinch plasma near the anode was experimentally discovered in Numerically, from our computation we obtained the condition such a phenomenon to occur:

1) $(\omega \tau) > 1;$
2) $n_0 R_0^2 < Z_i^2 M/M_H \leq 10^{18}$ cm$^{-1}$,

where $M$ is effective (i.e. averaged over ion species) mass of atoms(ions) in plasma sheath, $M_H$ is hydrogen mass atom and $Z_i$ is effective charge of plasma ions.

Fig.1 show the breaking of the ideal-MHD symmetry of magnetic field dynamics near the electrodes, which is caused by the transfer of magnetic field by electrons with velocity along the current line from cathode to anode the effect was found numerically by A.I.Morozov [1].

Comparison of experimental data and numerical results presented in Fig.2 and Fig.3. In Fig.2 the distribution plasma density at the moment of Z-pinch column formation at the anode is shown (here the anode is on the bottom, cathode absent). This picture was obtained by laser interferometry method in [4]. In Fig.3, plasma Z-pinch column is shown for entire electrode gap (anode is on the bottom, cathode is on the top ). Fig.3 presents the photograph from [5] for Z-pinch scattered visible light. The experimental and numerical are in qualitative agreement.

REFERENCES

2. V.V.Vikhrev, K.G.Gureev // Nuclear Fusion. 1977, v..
Fig. 1a $\Pi_1^{-1/2} = 0.065$

Fig. 1b $\Pi_1^{-1/2} = 0.079$

Fig. 1c $\Pi_1^{-1/2} = 0.161$

Fig. 1a-1c. Distribution of plasma density at left part and magnetic field at right part. The lower boundary of the calculation region corresponds to the anode and upper boundary to the cathode.
Two-dimensional modelling of thermonuclear combustion wave propagation in a z-pinch.

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The development of sausage-type instabilities in initially homogeneous z-pinch plasma column lead to the appearance of dense plasma which temperature substantially higher than the average plasma temperature in the column (~2 keV [1,2], ~10 keV [3]). This fact lead to the idea of using this hightemperature areas for thermonuclear combustion wave initiation along a z-pinch axis [4].

Calculate solution of MHD-equations [5] was made for the case of large radiative energy losses and thermonuclear heat release. The influence of thermonuclear heat emission on the dynamics of sausage instability growth is seen most obviously in the slowing-down of \( \alpha \)-particles in the plasma. In the calculations we assumed local emission of energy by the \( \alpha \)-particles in the plasma.

To fulfill the condition for thermonuclear combustion wave propagation in the axial direction, it is essential to have \( \rho r > A \) in the pinch, where \( \rho \) is the density of the material compressed by the magnetic field and \( r \) is the characteristic transverse dimension of the region occupied by that material (for example the radius), and \( A \) is a constant determined by the type of thermonuclear fuel and the compression conditions.

In solving this problem we considered initially a homogeneous plasma column with temperature \( T = 2 \text{ keV} \). The calculations give a value of 0.23 G/cm for the constant \( A \). At large values of \( A \) the thermonuclear combustion wave was initiated at an earlier stage of sausage instability growth, and at smaller values of \( A \) the wave damped.

We also performed calculations for different ratios between radiative energy losses from the plasma and thermonuclear heat release substantially exceeds radiative energy losses at the initial instant of time, we see a rapid expansion of the plasma column over its whole length; \( \rho r \) decreases, and no thermonuclear combustion wave is formed.

On the other hand, when thermonuclear heat release is -
small by comparison with radiative energy losses, a radiation compression regime is observed which leads to the formation, in the sausage instability region, of a plasma with substantially higher parameters; this in turn increases the parameter $\rho_r$ and thus facilitates the process by which a thermonuclear combustion wave is produced and propagated along the pinch.

The inhomogeneity which arises through the growth of the $m=0$ mode of MHD instability in the pinch is not only no obstacle to the emergence and propagation of a thermonuclear combustion wave but in fact leads to the initiation of such a wave. The emergence and propagation of a thermonuclear combustion wave is actually facilitated by large radiative energy losses from the plasma.

REFERENCES.

CHARACTERISTICS OF FAST PARTICLES AND AN ANALYSIS OF X-RAY SPECTRA IN THE PLASMA FOCUS DISCHARGE

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The dense high-temperature plasma produced under non-cylindrical compression of a pulsed pinched discharge is studied. It has been shown previously [1] that fast Al-ions with the energies in the MeV-range are generated in such a plasma. The multiply-ionized atoms of Ta, W, Mo are studied in a given paper. For determining the parameters of fast particles, as well as electron temperature $T_e$ and density $N_e$, the X-ray spectroscopy techniques are used. A technique for singling the electron distribution function out from the produced Bremsstrahlung spectrum is proposed. The application of the developed by us technique of filters operating in the vicinity to their K-edge has allowed us to exclude an effect of the resuperposition of reflection orders, fog, parasitic exposure.

INSTRUMENTATION

The experiments were done in the Plasma Focus Test Stand (PFTS) - facility in the geometry shown in Fig.1.

The discharge chamber was filled with argon up to the pressure $P = 30$ Pa, anode was made of copper with Ta, Mo, Fe - insertions at the centre. The diagnostic instrumentation shown in Fig.1 had the following main characteristics:

1. Pinhole camera with $d_{orif} = 100$ μm, the distance from the plasma to the orifice in the camera is 320 mm, magnification is $1:8$. 

Fig.1.

1. Pinhole
2. Convex mica crystal device
3. Iogansson's device
4. Caushua's device
2. The X-ray spectrograph with a convex mica crystal [3]. The plasma–crystal distance is \( L \approx 490 \text{ mm} \), magnification 1:2.

3. The X-ray spectrograph used the Iogansson's circuit-diagram. Roland's radius is \( R = 250 \text{ mm} \); quartz-crystal is \( 2d = 6.68 \, \text{Å} \), \( \Delta \lambda/\lambda = 10^{-4} \).

4. The Caushua spectrograph is for the energy range 10 - 150 keV. The plasma–crystal distance is about 600 mm, \( R_{er} = 600 \text{ mm} \).

A film with amplifying screen was used as a detector. The measurements were done under two operating conditions of the facility:

a) ion beam mode of operation, the anode has a small, conical deepening;

b) electron beam mode of operation is realized in case of a flat anode.

EXPERIMENTAL RESULTS

1. ION BEAM MODE OF OPERATION. The He-like Ar–spectrum and its satellites, registered in the third order of reflection with the spectrograph 2 (See, Fig.1), are given in Fig.2. The Ta-lines in the second order of reflection are observed from the short wavelength side in the vicinity to Ar-lines. Using the data from the papers [4], the most bright Ta-lines were identified as the transition \( 3d^9 - 3d^84f \) in the ion Ta XLVII with the wavelengths \( \lambda_1 = 5.886 \, \text{Å} \) and \( \lambda_2 = 5.916 \, \text{Å} \). The spectra were registered per one discharge with the spatial resolution \( \sim 300 \, \text{µm} \). The plasma images obtained in the wavelength of argon and tantalum allow one to obtain the dimensions of regions radiating these lines. In Table 1 the dimensions of the regions of glow along the pinch radius, \( d_r \), and along the pinch height, \( d_h \), for the He-like Ar and for the Co-like Ta are given. \( T_e \) values are given for the same regions under conditions of corona equilibrium. The electron density is estimated from the ratio of intensities for intercombinational and resonance Ar-lines.
An analysis of the spectra related to different discharges has showed that the pair of Ta-lines is shifted respective to the group of He-like Ar-lines. It can be explained by a spread in the sources of glow for these ions in space. The ions of a metal glow from the near-anode "pillow" formed under anode bombardment by fast electrons, meanwhile the ions of a gas glow from the pinch "body" which is located above. This shift is varied from one discharge to another and can be estimated from the pinhole images. Another reason for the Ta-line shift can be Doppler's shift provided by Ta-ion motion towards the cathode. The due regard for these factors has allowed us to estimate $V_i = (0.9 - 1.3) \cdot 10^9$ cm/s that correspondent to $E_i = (0.8 - 1.6)$ MeV in different discharges.

2. ELECTRON BEAM MODE OF OPERATION. The absence Ar-lines was registered in this geometry, the transitions (3d$^n$-3d$^{n-1}$kf), $k = 5, 6, 7, 9$ in the Ni-like Ta in the range 3 + 5 Å were registered. Along the lines of Ta-ions, the most bright lines of M-series for a Ta−atom − $M_{\alpha, \beta}$, $M_{\text{III}}$ − $N_I$ are present. At the same time $K_{\alpha_1, \alpha_1, \beta_1}$, $L_{\alpha, \beta}$, $M_{\alpha_1, \beta, \gamma_1, \gamma_2}$ Ta were fixed with the Caushua spectrograph. $K_{\alpha_1, \beta}$ for Mo were obtained similarly, and $K_{\alpha, \beta}$ for Fe were obtained with the Iogansson's spectrograph. These measurements confirm the presence of an electron beam in the plasma with the energy not less than 60 keV, bombarding the anode. The following technique for the electron beam distribution function recovery with respect to the energies $I_{\gamma}(E_{\gamma})$ is proposed.

The intensity of Bremsstrahlung, $I_{\gamma}$, emerges as a result of deceleration of monoenergetic electrons with the energy $E_e$ in a thick target is represented by the relationship:

$$I_{\gamma} = \text{Const} \left( E_e - E_{\gamma} \right)$$  \hspace{1cm} (1)

In case of electrons with the distribution function,
\( I_e(E) \), this dependence has the form:

\[
I_\gamma = \text{Const} \int_{E}^{E_{\text{max}}} (E - E_\gamma) I_e(E) dE
\]

(2)

Double differentiation (2) with due regard for the boundary conditions gives:

\[
I_e = \text{const} \frac{d^2 I_\gamma}{dE_\gamma^2}
\]

(3)

From (3) it follows that one can obtain \( I_e(E) \), having measured the Bremsstrahlung spectrum and having approximated it by the corresponding functions. To calculate \( I_e(E) \) the correct measurements of a Bremsstrahlung spectrum, taking account of the resuperposition of orders, fog, background, are necessary. To exclude the contributions of these factors in the Caushua instrument we have developed the technique of filters, operation in the region of their K-edges. At the place of a detector, corresponding to the K-edges of some element, a filter having an affective thickness and containing this element located. Then, the true intensity of the spectrum, \( I_\gamma \), is calculated from the relationship:

\[
I_\gamma = \frac{(I_{1\gamma} + I_{2\gamma})}{\exp(-\mu_1 \delta) - \exp(-\mu_2 \delta)}
\]

(4)

where \( I_{1\gamma}, I_{2\gamma} \) are the Bremsstrahlung intensity beyond the filter before and after the K-edge;

\( \mu_1, \mu_2 \) are the linear coefficients of absorption before and after the K-edge, \( \mu_2 > \mu_1 \).

In this case, there is an optimal filter thickness

\[
d_{\text{opt}} = \frac{(\ln \mu_2 - \ln \mu_1)}{\mu_2 - \mu_1}
\]

(5)

The experimental verification of this technique is now in progress, the results will be published in the next paper.

References.

Z-pin...Turbulent Energy Considerations
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Two plasma jets and typical dense central region have been created in the low energy z-pin...pulse discharge (0.5 kJ, $T \sim 4 \mu s$, 50 kA) in nitrogen and argon (under pressure 1 $\pm$ 50 kPa) [1, 2, 3].

The existence of nondissipative turbulence as a considerable form of energy balance is discussed in this contribution.

Plasma is accelerated due to the Lorentz force in a radial direction towards the discharge symmetry axis. In the vicinity of the jet's boundary the parameters of the plasma are: the electron velocity $(3 \pm 8) \times 10^3 \, \text{ms}^{-1}$, the concentration $(1 \pm 3) \times 10^{24} \, \text{m}^{-3}$ and temperature $(2 \pm 2.5) \, \text{eV}$. The velocity values are higher than those of the Alfvén velocity and so the instability begins to develop. The current sheath disrupts into filaments and simultaneously plasma turbulent motion starts in the central region (characteristic time constant is approximately $0.1 \, \mu s$). The filaments (diameter 0.1 mm, average mutual distance 0.5 mm) are relatively stable. The accelerated turbulent plasma penetrates into the jets through the interfilament's slots and the radial velocity changes in the axial one inside the jets.

The jet's electron density was estimated from schlieren pictures and it is comparable with that of the jet's neighbourhood. The electron temperature was calculated from the relative intensities of visual continuum. Its value for the jets (0.3 eV) is substantially lower than that for the jet's neighbourhood (2.5 eV). This difference indicates that the major part of the heat energy and the kinetic energy of ions penetrating into the jets should be hidden in the turbulence rotational motion.

The turbulences are formed in strong magnetic fields, and high electric current density can be induced under this condition. The feedback of the electromagnetic and inertial forces should probably result in the existence of relatively stable turbulent structures with mechanic and electromagnetic energy density balance, and partial electron and ion paths separation. Therefore the electron-ion collision frequency probably decreases. Consequently, Joule's heating and the rotational motion energy dissipation are reduced.
Further plasma jets development, when the plasma is ejected from the jets in opposite directions, into the discharge central region, supports this hypothesis.

High electron density ($>10^{25} \text{ m}^{-3}$) and low electron temperature (0.3 eV) are typical for this central region structure, but its shape strongly depends on the discharge chamber gas filling (nitrogen, argon).

In nitrogen, the central region structure has spherical shape characterized by increasing cross-section. The surface of this structure consists of little spheres similar to the shape of the cauliflower surface. Outside this surface some spiral filaments like whirlwinds are ejected with high speed (higher than $10^4 \text{ ms}^{-1}$). With increasing distance from the central region structure outer surface, the filament's diameter also increases (from 0.05 mm near the surface to 0.2 mm at the opposite end where the filaments have spherical cloud-like structure). The structure disintegrates during several tens of $\mu$s after first current half period. It is probable, that the jet's turbulence (its life time about $1 \mu$s) ceases in the little spheres of the central region structure.

In argon (under pressure 1±50 kPa), the central region structure shape is characterized with quite a different form and evolution. A compact structure with sharp density gradients on the boundary is created during the 0.5 $\mu$s at the time of maximum current ($1 \mu$s). During further 1 $\mu$s interval this structure keeps its diameter and then very quickly disintegrates.

It seems, that general conditions of the magnetohydrodynamic turbulence interaction and correlation phenomena are realized due to the frozen magnetic lines, high conductivity and plasma selforganisation.

For better understanding there is necessary further study at higher energy equipments.

References:

EXPERIMENTAL STUDIES OF DENSE MAGNETIZED PLASMAS PRODUCED BY PF-TYPE DISCHARGES

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Introduction

Experimental facilities producing dense magnetized plasmas, e.g., PF-type devices have been investigated for many years as pulsed sources of soft- and hard- X-rays, energetic deuterons, fast impurity ions, relativistic electron beams, and fusion products (fast neutrons and protons). Emission characteristics of such facilities depend strongly on their technical parameters and operational conditions. Knowledge of these characteristics is needed in order to explain physical phenomena and to determine possible technological applications.

The main aim of this paper was to report on recent studies concentrated upon time correlations of the X-ray and charged particle emission from PF-devices operated at SINS.

Experimental facilities

The PF-360 device [1-2] was equipped with two coaxial electrodes of 120 mm and 170 mm diameter, respectively. They were made of 300-mm-long copper tubes, and a root of the inner electrode was embraced by an 80-mm-long ceramic insulator. The system was supplied from a 288-μF condenser bank charged up to 30 or 35 kV.

The MAJA-PF device [3] possessed a 72-mm-dia. inner electrode and a 124-mm-dia. outer one made of 16 copper rods. The electrodes were 298 mm long, and the tubular insulator was made of the same material as in the PF-360 facility. The MAJA system was supplied by a 72-μF condenser bank charged up to 35 kV.

The two experimental facilities possessed standard equipment for U(t), dI/dt, Y_x, and Y_n measurements. Additionally, there were applied various scintillation detectors for X-ray measurements, nuclear track detectors for the registration of ions, and special Čerenkov-type radiators [4] for the detection of fast electrons.

Experimental results

Within a framework of previous experiments with the two devices described above, there were performed detailed time-integrated measurements of X-ray and neutron yield

1 Participation within a framework of a fellowship granted by the Polish Ministry of Education
Since, in comparison with other PF experiments of a similar scale, there have been observed some differences in characteristics of the X-ray and charged particle emission, particular attention has been paid to time correlations.

During recent experiments with the PF-360 device there have been registered U(t) signals from the HV-divider and I(t) waveforms from the Rogovski coil. Those traces have been compared with R_{min}(t) signals obtained from a fast photodiode observing a pinch column. Time-resolved X-ray and neutron signals have been measured with scintillation detectors coupled with fast X-2020 photomultipliers. Those monitors were placed side-on at the distance of 250 cm from the electrode axis. In order to measure fast ions and fusion protons, use was made of several miniature scintillation detectors located within the experimental chamber, at different angles to the z-axis, but at the same distance of 15 cm from the pinch region. Typical time-resolved traces have been presented in Fig.1. It can be seen that the first voltage peculiarity appears during the maximum compression, and two successive peaks (corresponding to neutron pulses) appear with some delay. An analysis of ion detector signals shows that those detectors were sensitive to hard X-rays (the first peaks), and in some extent also to (for deuterium discharges), as well as observations of energetic ions and fast electrons. The results of those measurements have just been summarized [5] and they have shown that the optimization of the both experimental facilities remains in agreement with known scaling laws [1].

![Graph](image.png)

**Fig.1.** Time correlations observed for PF-360 device operated at \( p_o = 5 \) torr \( D_2 \), \( U_o = 30 \) kV, \( W_o = 130 \) kJ, \( Y_n = 5.7 \times 10^9 \).

**Fig.2.** Correlation of U(t) and \( Y_e + Y_n \) waveforms with electron signals from Čerenkov-type diamond and plastic-radiators, as observed for PF-360 device at \( p_o = 5 \) torr \( D_2 \), \( U_o = 33 \) kV, \( W_o = 157 \) kJ, \( Y_n = 1.5 \times 10^{10} \).
neutrons (the successive maxima). Nevertheless, a comparison of signals obtained from different detectors suggests that multiple spikes can be identified with very fast deuterons and fusion-produced protons. The two neutron pulses have the FWHM values equal to about 40 ns. The signals obtained from the internal ion detectors are characterized by relatively long tails, which are probably induced by very fast deuterons and fusion protons. The detectors oriented at different angles give often irreproducible signals, what proves that the charged particle emission shows distinct anisotropy and a stochastic character.

In order to study fast electron beams emitted through the inner electrode in the upstream direction, use was made of different Čerenkov-type radiators placed at the distance of 48 cm from the pinch region and coupled with fast photomultipliers. Typical traces have been shown in Fig.2. It can be seen that the electron pulses have $\text{FWHM} = 30-50\,\text{ns}$, and they are delayed in relation to the hard X-ray signals because of differences in times of flight. Taking into account the energy threshold for a diamond radiator (with a 30-μm Cu-filter) and that for a plastic radiator (with the same filter), one can conclude that the population of electrons above 360 keV is about 4 times smaller than that above 215 keV.

Another series of studies has been performed with the MAJA—PF device. Typical correlation traces have been presented in Fig.3. It was observed that the maximum pinch compression ($R_{\text{min}}$ signal) corresponds to the first maximum of $\text{d}I/\text{d}t$ waveform, at the very beginning of a small "plateau". It can be explained if after the pinch formation the discharge current is stabilized for some time. Almost simultaneously, there also appears the emission of soft ($<10\,\text{keV}$) X-rays connected with interactions of electrons with a dense pinch plasma. The second X-ray peak is usually observed when the discharge current starts to change once again after a rise of some instabilities within the pinch column. These instabilities are evidently accompanied by the generation of fast electrons. The maximum of an electron-induced signal is usually delayed in relation to the second X-ray peak because of a time of flight of electrons before their registration.

Recently, to make possible to modify a plasma concentration locally and to probe a pinch region with solid microtargets, there has been developed a special cryogenic system [6]. This facility consists of a liquid-He cryostat equipped with a "cold nose" and an injection
channel. It has been demonstrated that the arrangement can produce a narrow (below 100 μm) stream of a liquid nitrogen or argon. Using an additional piezo-electric modulator one can divide this stream into a sequence of miniature droplets, as shown in Fig.4. Such droplets can in turn be frozen before entering the main experimental chamber. A new injection system equipped with diaphragms enabling the differential pumping and transport of the cryogenic microtargets into the plasma region (at pressure 1-10 torr) is now under final laboratory tests.

![Fig.4. Picture of a cryogenic stream divided into miniature droplets (on the left) and light reflection signals from the microtargets formed with frequency of 13 kHz.](image)

Conclusions

The experimental results presented above can be summarized as follows:

1. Signals induced by very fast deuterons and fusion protons are rather poorly correlated with X-ray and neutron pulses because neutrons are not influenced by a magnetic field, while motions of charged particles depend strongly on it.

2. Fast electron beams emitted along the z-axis in the upstream direction are rather well correlated with the hard X-ray pulses, taking into account differences in a time of flight.

3. A new special equipment has just been developed for experiments with injection of cryogenic micropellets into a dense plasma region.

References

Investigation of Plasma Light Output in a Mather-Type Plasma Focus Device

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Introduction

Among the different types of focus devices, there is the Mather-Type plasma focus (Mather, 1964). The device is useful for the small laboratory experiments of fusion (Jager and Herold, 1987) and as a source of electron and ion beams or X-ray, (Decker and Pross, 1982). In the present work, emission of X-ray is investigated and light output was analyzed using three gases; hydrogen, argon, and air. Plasma lifetime was estimated and the atomic processes were studied. Minimum and maximum pressure limits were defined.

Experimental set up:

A schematic diagram of the Mather-type device is shown in Fig. (1) where: C is a 30 μF capacitor bank, T.P. is a triggering pulse generator, S.G. is a spark gap, P.D. is a potential divider probe, R.C. is a Rogowski coil, V. ch. is a vacuum chamber, A and C are the anode and cathode, G is a glass sleeve, a, b and c are the electrodes dimensions, P are the pin diodes, M.P. is a magnetic probe, G.I. is the gas inlet and CRO to the oscilloscope. The two pin diodes were placed inside the vacuum vessel; one was covered by Al-foil of different thickness. Signals of the Rogowski coil I(t), potential probe V(t), uncovered pin diode L and the covered pin diode (X-ray) are shown in Fig. (2). The current analysis are based upon the equation:

\[ I(t) = C \cdot V \cdot \omega \exp \left( \frac{1}{2} t \right) \sin \omega t \]  

where \( \gamma = r/L \) and \( \omega_0^2 = 1/LC \)

Results and Discussion:

Since focus device is considered as an electromagnetic shock tube, the transmitted power by shock wave must be
sufficient for gas ionization, otherwise plasma becomes transparent to magnetic field and the expected effective magnetic piston lose its efficiency. Thus upper pressure limits are determined at which the transmitted power is so low that focus mechanism does not work. Measurements of the power indicates that only less than 50% was transmitted at higher pressures. Thus, the upper pressure limits were 8 torr for H₂ and 1-0.8 torr for argon and air. These limits were also confirmed by estimating the magnetic Renold numbers.

The transit time $t_a$, of the formed sheath, in the axial phase using the Snow-plow model (Lee, 1983) where

$$t_a = \text{const} \left( \frac{b^2-a^2}{\ln(b/a)} \right)^{1/2} Z \frac{1}{\rho_0^{1/2} I_0^{1/2}}$$

(2)

$Z$ is the axial dimension and $\rho$ is the gas density. The sheath velocity $v_s$ was computed for H₂, Ar and air as a function of gas pressure (see Fig. 3). The minimum gas pressure limits to operate the device efficiently, were defined whereas the sheath velocity exceeds $10^7$ cm/sec. and current spoke formates. The reasonable axial speeds were $(3-10) \times 10^6$ cm which agree with the measured velocities at pressures of 2 torr for H₂ and less than 0.5 torr for argon and air. The maximum and minimum pressure limits were confirmed by the experimental observation of focus formation, within the accuracy of the pressure measurements.

The majority of the produced X-ray, Fig. (2), were observed in the break up phase. Its production can be related to the "runaway' process of the fast electrons in the fully ionized gas (Dreicer, 1959). The required critical electric field $E_C$ is given by (Landau and Lifshitz, 1981):

$$E_C = 4mN_e e^3 \ln A/kT_e$$

(3)

For the present Mather-type device $E_C$ must exceeds 1300 V/cm. The electron-ion collision time $\tau_{e-i}$ has to be longer than the travelling time necessary for leaving the focus zone. $\tau_{e-i}$ was calculated from (Chen, 1974):

$$\tau_{e-i} = \frac{2.3 \sqrt{2}}{2N_i n_e^4}$$

(4)
The two criteria were well satisfied. The formed electric field was about $10^5 \text{ V/cm}$, $\tau_{e-i}$ was about $10^{-10}$ sec. while the travelling time was about $10^{-15}$ sec.

Decaying of the light output is attributed either to the atomic relaxation processes or to particles diffusion. Among the expected atomic processes, the Bremsstrahlung process was the dominant since the criterion $kT_e > \beta z^2$ was satisfied. Recombination and line radiation may only be considered as a result of the slow electron. The experimental data fit closely to the empirical formula (Fig. 2).

$$L = L_p \exp\left(-\frac{t}{\tau}\right)$$  \hspace{1cm} (5)

where $L_p$ is the maximum light output intensity and $\tau$ is the plasma life time. Fig. (4) shows the estimated values of $\tau$ for the different gases and at different pressures. $\tau$ can also be computed using Fick's law and Bohm's formula i.e. that

$$\tau = \frac{N_e R^2}{\Gamma} = \frac{N_e R^2}{2D}$$  \hspace{1cm} (6)

where $N$ is the particles flux, $R$ is the focus radius and $D$ is the diffusion coefficient. Fig. (4) shows fair agreement between the computed values of $\tau$, (solid curves) using equation (6) and the experimental data (equation 5). Decreasing of $\tau$ at high pressures and large atomic numbers is related to plasma expansion into vacuum, which depends on the ion sound speed, and to particles diffusion.

References


Fig. 1

Fig. 2:

Fig. 3:

Fig. 4:
Stability study related to pressure anisotropy in the GAMMA 10 tandem mirror

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1. Introduction
Effects of pressure anisotropy on micro- and macrostabilities are studied in ICRF(Ion Cyclotron Range of Frequency) heated plasmas of the GAMMA 10 tandem mirror. The pressure anisotropy is defined as ratio of pressures perpendicular to parallel to the magnetic field line ($P_L/P_\parallel$). Alfvén Ion Cyclotron(AIC) mode is one of microinstabilities in the range of the ion cyclotron frequency, which is driven by plasma $\beta$ (ratio of plasma pressure to magnetic pressure) and the pressure anisotropy. Flute-interchange mode is one of macroinstabilities which is driven by the pressure weighted on the bad curvature region of the magnetic field line. The flute-interchange and the AIC modes are influenced by an axial profile of the pressure which is related to the pressure anisotropy in term of the magnetohydrodynamic(MHD) equilibrium along the magnetic field line. The purpose of this paper is to clarify the relations between the pressure anisotropy and the stability of these modes in the tandem mirror plasma.

2. Experimental Apparatus
GAMMA 10 is an axisymmetric tandem mirror with a thermal barrier. A typical axial profile of the magnetic field strength is shown in Fig.1. GAMMA 10 consists of five mirror cells, which are a central cell, minimum-B anchor cells and end plug/barrier cells. The central cell is the main confining region. The anchor cell is provided with baseball seam coils for MHD stability of the tandem mirror plasma. Beach heating is adopted to heat the central and anchor cell ions. Alfvén waves with two frequencies are excited at both ends of the central cell. One is 6.3 MHz corresponding to the ion cyclotron frequency near the central cell midplane and the other is 9.9 MHz corresponding to the ion cyclotron frequency near the anchor cell midplane. The ion pressure profile has a peak near the resonance region.

Several diamagnetic loops arranged along the magnetic field line are used to measure the axial profile of the perpendicular pressure. Ratio of the pressures between parallel and perpendicular to the magnetic field line is determined from the pressure balance equation.
experimental point at which the AIC mode are observed and the open circles no AIC mode. When the amplitude of the AIC mode increase, relaxation of the pressure anisotropy has been observed\textsuperscript{79}. In the parameter region including the solid and open circles the plasma is macroscopically stable, which is in good agreement with the predicted stable region. In the region above the broken line, strong fluctuations appear and the plasma can not be sustained.

Conclusion

Depending on the $\beta$ and ion pressure anisotropy, AIC modes with a few discrete line spectra are excited in the frequency range a little below the ion cyclotron frequency. The frequency spectra could be explained by the observed standing wave structure associated with a wave reflection at an axial position dependent on the AIC drive. Relaxation of the pressure anisotropy has been observed with an increase of the AIC mode amplitudes.

The ratio of the central cell $\beta_{LC}$ to the anchor cell $\beta_{LA}$ determines the MHD stability boundary. The stability boundary for the flute-interchange mode is obtained experimentally and in good agreement with the theoretical prediction. Due to the decrement of the pressure weighting on the bad curvature region, the stable region greatly expands on $\beta_{LC}$ vs. $\beta_{LA}$ diagram.

Acknowledgement

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References


along the magnetic field line\(^6\). Fluctuations are detected with magnetic probes, electrostatic probe array and a millimeter wave reflectometer system\(^7\).

3. Experimental results

3.1 AIC mode

The AIC modes are identified from the magnetic probe measurements when both pressure anisotropy and central cell beta value\(\beta_{LC}\) are relatively high. As shown in Fig.2 frequency spectra consist of a few discrete lines the number of which depends on the plasma parameter. Dependence of the amplitude of the AIC modes on drive term \(\eta_{AIC}\) defined by \(\beta_{LC} \left( \frac{P_{LC}}{P_{\|c}} \right)^2\) is shown in Fig.3(a). Open and solid circles are the AIC amplitudes measured at \(z = -1.29\) m and \(z = -1.14\) m in the central cell. Phase difference between two magnetic probe signals strongly depends on the AIC drive term as shown in Fig.3(b). The phase difference becomes nearly zero when the AIC drive term exceeds a critical value \(\eta_c\).

3.2 Flute-Interchange mode

Stability boundary of GAMMA 10 for flute interchange modes is experimentally obtained. Additional gas puffing into the anchor cell or modulation of the ICRF heating power are successfully used to vary beta values at the anchor and the central cells in a wide range. Figure 4 shows temporal evolutions of the central cell and anchor cell pressures with a strong and a weak anisotropic pressure profiles at the central cell. When the additional gas is puffed into the anchor cell, the anchor cell pressure decreases due to an increase in charge exchange energy loss. When the beta ratio reaches a critical value, the central cell pressure dumps abruptly. It is found that the MHD stability is determined by the ratio of the central cell pressure\(\beta_{LC}\) to the anchor cell pressure\(\beta_{LA}\).

![Fig.3](image1.png)

(a) Dependence of amplitude on the AIC drive term. Open and closed circles correspond to amplitudes of AIC mode at \(z=-1.29\) m and \(z=-1.14\) m.

(b) Phase difference between two magnetic probe signals strongly depend on AIC drive term. Phase difference nearly zero when AIC drive term exceed the critical value, \(\eta_c\).

![Fig.4](image2.png)

Fig.4. Typical temporal evolutions of the central and anchor cell beta values in additional gas puffing experiments. Solid line is stronger anisotropic case \((P_{LC}/P_{\|c}=10.2)\) and dotted line is weaker anisotropic case \((P_{LC}/P_{\|c}=4.7)\).
4. Discussion

The magnetic field strength and the normal curvature $\kappa_\varphi$ of the central cell are shown by a dotted and solid lines in Fig.5(a). Hot ion pressure profiles in the direction perpendicular and parallel to the magnetic field line are indicated by a solid and a dotted lines in Fig.5(b). The pressure profiles are determined from the diamagnetic signals in term of the MHD equilibrium equation. The axial profile of the AIC drive term $\eta_{AIC}$ evaluated from the profile of the parallel and perpendicular pressure is also peaked near the midplane of the central cell. In the inner region of the central cell where $\eta_{AIC} \geq \eta_c$, the AIC mode appears as a standing wave which may be caused by some kind of wave reflection at the point where the $\eta_{AIC} = \eta_c$ is satisfied. In the outer region where the $\eta_{AIC} < \eta_c$, the AIC mode propagates towards the anchor cell. The critical value $\eta_c$ corresponds to the value above which the region with the standing wave structure expands over the magnetic probe location. This type of the wave structure has been observed also in experiments of low frequency ballooning mode.

The flute-interchange stability is determined by integrating the product of the total pressure and the normal curvature along the magnetic field line as follows:

$$\Gamma = \int \frac{(P_\perp + P_\parallel)}{B} \kappa_\varphi \, dl \geq 0. \quad (1)$$

When the pressure anisotropy is strong, the pressure profile is more peaked near the midplane of the central cell. It is found that the stable region on $\beta_{LC}$ vs. $\beta_{LA}$ diagram expands due to the decrease of the pressure weighting on the bad curvature region.

The stability region for the flute-interchange mode and the growth rate of the AIC mode are shown on $\beta_{LC}$ vs. $P_\perp/P_\parallel$ diagram in Fig.6. Solid lines show the theoretical contour lines of the growth rate of the AIC mode. Dot line is theoretically predicted as a boundary between absolute and convective unstable region, $\eta_{AIC} = 3.5$. Broken line is the flute stability boundary which is calculated from Eq.(1). The solid circles indicate the
FOKKER-PLANCK SIMULATION OF SLOSHING IONS IN FEF-II AND GDT

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The sloshing ions are produced by oblique injection of a high energy neutral beam into a mirror confined plasma. These sloshing ions play important roles to use mirror confinement scheme for two-compartment neutron source, to produce strong neutron flux by DT fusion reaction between the sloshing ions and the target plasma ions. In the present work, we obtain the distribution function of the sloshing ions by Fokker-Planck simulation and estimate the axial distribution of the sloshing ion density in the cases of FEF-II, which is mirror based neutron source, and GDT. Then, we have compared the numerical simulation results with the experimental results in GDT. We have also estimated the strength of the neutron flux expected in FEF-II by this simulation.

The conceptual design study of mirror based fusion plasma neutron source FEF has been carried out since 1980. The first stage of this study concluded in 1989 [1]. Recently, FEF-II design study for the purpose of early construction of the neutron source has been started [2]. FEF-II is the two-component plasma system, and multi-pole field (Version 1) or RF plugged cusp (Version 2) is used for MHD stability. Configuration of the magnetic field of FEF-II is formed basically by long solenoid with
mirror field at both end.

GDT (gas dynamic trap) is an axisymmetric mirror device with a large mirror ratio [3]. The main advantage of this trap is that it can provide a plasma MHD stability in an axisymmetric geometry. Recently, design study of GDT-based neutron source (GDTNS) has been started. The parameters of FEF-II and GDT are shown in Table 1.

The Fokker-Planck equation for the velocity distribution function $f(v, \zeta, t)$ of the sloshing ions is given by

$$\frac{\partial f}{\partial t} = \frac{1}{v^2} \frac{\partial}{\partial v} \left[v^2 \left(Af + D_{\parallel} \frac{\partial f}{\partial v}\right)\right] + \frac{1}{v^2} D_{\perp} \frac{\partial}{\partial \zeta} \left[(1 - \zeta^2) \frac{\partial f}{\partial \zeta}\right] - \frac{f}{\tau_{\text{ex}}} + S$$

where $S$ is the production term of the sloshing ions due to NBI, $\tau_{\text{ex}}$ is the charge-exchange loss time, $\zeta$ is the cosine of the pitch angle [4]. The Fokker-Planck coefficients $A$, $D_{\parallel}$ and $D_{\perp}$ represent the slowing down, the energy diffusion and the pitch angle scattering, respectively. We assume that the target plasma is Maxwellian to obtain the Fokker-Planck coefficients. We use the production term as

$$S = \frac{I}{\pi e \Delta v \Delta \zeta V} \exp \left[-\left(\frac{v-v_0}{\Delta v}\right)^2 - \left(\frac{\zeta - \zeta_0}{\Delta \zeta}\right)^2\right]$$

where $I$ is the NBI current, $V$ is the volume of plasmas, $v_0$ is the velocity of the sloshing ions produced and $\zeta_0$ is the cosine of injection angle.

We obtain the axial distribution of the sloshing ion density from $f(v, \zeta, t)$. $P$ is the probability of existence of plasma particles at $z = z_A \sim z_A + \Delta z$. The axial distribution of the sloshing ion density is given by

$$n_s(z) = \frac{B(z) \Delta z}{B_0} \frac{\Delta \zeta}{L} \int_0^\infty v^2 dv \int_0^{\zeta_c} P \cdot f(v, \zeta, t) d\zeta$$

In numerical calculation, we set $1/2m_s v_0^2 = 100\text{keV}$ (FEF-II) and $15\text{keV}$ (GDT), $1/2m_s (\Delta v)^2 = 1\text{keV}$ (FEF-II) and $0.2\text{keV}$ (GDT), $\zeta_c = 0.82$ (FEF-II) and $0.98$ (GDT), $\Delta \zeta = 0.01$ (FEF-II) and $0.03$ (GDT) and $\zeta_0 = \pi/4$ (FEF-II,GDT).

The results from the simulation are as follows. In the case of FEF-II, the averaged sloshing ion density in the steady state ($\geq 5\text{msec}$) goes up
to $1.6 \times 10^{13}$ cm$^{-3}$. The axial distribution of the sloshing ion density in the steady state is shown in Fig. 1. The density of sloshing ions at the turning points is about 2 times higher than that at the midplane. Fig. 2 shows the axial distribution of the 14 MeV neutron flux at the first wall ($r = 20$ cm). The neutron flux at the midplane is $1.6 \times 10^{14}$ cm$^{-2}$s$^{-1}$ (3.5 MW/m$^2$) which comes from mainly the fusion reaction between the sloshing ions (D) and tritons in the target plasma. It is shown that the peaking of neutron flux at the turning points does not appear, which is interpreted by averaging over the longitudinal distribution.

In the case of GDT, we include the effect of temporal variation of the target plasma parameters which is obtained experimentally. Fig. 3 shows the time evolution of the sloshing ion averaged density in GDT. The axial distribution of the sloshing ion density is shown in Fig. 4. The results of the simulation indicate that the average density of the sloshing ions in GDT goes up to about $7 \times 10^{12}$ cm$^{-3}$, and that the density of the sloshing ions at the turning points is about 3 times higher than that at the midplane. These values are same order of magnitude with the ones in the GDT experiment. From these results, it could be said that this simulation code would give realistic plasma parameters in the plasma based neutron source which is expected to be built in near future.

References


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Table 1 Parameters of FEF-II and GDT

Fig.1 Axial distribution of the sloshing ion density in FEF-II

Fig.2 Axial distribution of the neutron flux in FEF-II

Fig.3 Time evolution of the sloshing ion averaged density in GDT

Fig.4 Axial distribution of the sloshing ion density in GDT
Topic 4

Plasma Edge Physics
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MEASUREMENTS OF PLASMA CONVECTION IN THE SOL OF JET USING A LANGMUIR/MACH PROBE

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1. Introduction. Measurements of plasma convection in the SOL (scrape-off layer) have been carried out for JET divertor discharges using a Langmuir/Mach reciprocating probe [1, 2] with sensitive elements facing the divertor (ion drift side for normal $B_n$) and the main SOL (electron drift side for normal $B_n$) (see Fig.1).

Plasma convection competes with thermal conduction as a mechanism to carry energy to the divertor target and can diminish temperature gradients along the field. Convective flow also provides the basic force that favours the retention of impurities near the target. Hence, it is necessary to assess experimentally its importance in JET divertor discharges and to compare it with results from models for the SOL plasma [3], which have been validated with experimental measurements [4].

2. Experimental Measurements. Mach probe measurements in a series of upper single null discharges in various regimes have been considered (OH, L-mode, H-mode). Of the available models we use the model described in [2] that accounts for the reduction of ion flux to the side facing the divertor due to the finite size of the probe (the disturbance length for Ohmic and L-mode conditions is of about 4 m for the JET reciprocating probe, safely smaller than the 20 m from the probe to the outer divertor target for the discharges studied) and viscous and non viscous effects.

Measurements obtained for a medium density ($2.5 \times 10^{19} \text{ m}^{-3}$) low additional heating (L-mode) discharge are shown in Fig.2. The plasma flow at the reciprocating probe is directed towards the divertor all across the SOL, the ratio of the fluxes at both sides of the probe being $(2.3 \pm 0.2)$, which corresponds to a Mach number of $(0.5 \pm 0.03)$ if perpendicular viscosity is ignored and to $(0.3 \pm 0.03)$ including it. Similar values are obtained for ohmic discharges in the same density range.

The pattern of the plasma flow at the reciprocating probe is more complicated for H-mode and low density L-mode discharges. In Fig.3 the measurements obtained for a low density L-mode ($9.5 \times 10^{18} \text{ m}^{-3}$) are shown. The plasma flow is directed towards the divertor in the external part of the SOL with values for the ratio of ion fluxes on both sides of the probe similar to those obtained in the medium density case. However, close to the separatrix a region in which the flux on the ion side is larger than on the electron side is found. This is interpreted as a region in which the flow is away from the divertor at the reciprocating probe position. The reversal of the flow close to the separatrix is also found in H-mode discharges, in agreement with previous observations [5]. The ratio of ion fluxes in both sides of the probe for the external part of the SOL in H-mode discharges is similar to that found in L-mode. The radial
location of the flow reversal region is more uncertain due to possible errors in the relative position of the probe to the magnetic separatrix and the fact that the profiles of plasma parameters are very steep close to the separatrix in H-modes [6].

The ion flux measurements with the probe could be affected by thermionic emission and by the presence of large negative currents close to the separatrix in JET [7]. However, these effects would affect predominantly the electron side of the probe (it suffers a higher power deposition and intercepts the electron currents), but in experiment is the ion side which shows the largest changes, indicating that flow reversal must take place.

3. Modelling of the discharges. These discharges have been modelled with EDGE1D [3] for a pure hydrogen plasma using measured plasma parameters as inputs (no shift of the relative position of the probe to the magnetic separatrix has been imposed). Assuming equal power into the SOL via electrons and ions, the calculated total power arriving at the divertor target is in agreement with the power determined from main plasma measurements ($P_{\text{targ}}$) in the medium density case (Fig. 2). At low densities (Fig. 3) the calculations account for only 30% of $P_{\text{targ}}$.

The calculated Mach number (modelled as constant across the SOL) at the probe is in the range (0.15 - 0.20) in reasonable agreement with experimental estimates (viscous case) for the outer regions of the SOL. The results for low densities tend to produce a smaller Mach number despite a larger neutral escape from the divertor to the main plasma (30% at low densities, 20% at high densities) due to the influence of parallel viscosity (larger at higher temperatures [8]). This is also the trend barely detectable in the experiment (compare outer SOL in Fig. 2 and Fig. 3).

The calculated heat flux into the the divertor is shared between conduction (60%) and convection (40%), the Mach number at the divertor entrance (X-point) being 0.5. The conductive heat fluxes are calculated in the collisional approximation, and non-local effects can be estimated by comparison with the ‘free streaming flux’ ($q_{\text{el}} = n \sqrt{K T_{\text{el}}/m_{\text{el}}}$). The ratio between conductive and ‘free streaming flux’ is about 0.07 for electrons and 1.0 for the ions. The electron heat flux is a factor 2.3 higher than that allowed by collisional theory [9] but much smaller than the collisionless flux (0.2 - 0.3 of the ‘free streaming flux’), hence at most a reduction of 25% in the electron conductive flux is expected due to non-local effects. These effects may be stronger for the ion conductive heat flux.

Some 2D aspects of the plasma flow were studied using the ionization source calculated in the model with the NIMBUS Monte-Carlo code [10] for various flux tubes in the SOL. Magnetic geometry effects cause the appearance of the point of maximum plasma flux onto the target at a finite distance from the separatrix [11]. This produces an ionization source with different radial dependence than the plasma density (and flux) profile (exponentially decreasing from the separatrix, in magnetic flux), shown in Fig. 4. This, in turn, leads to a complicated behaviour for the
ionization source along the field which causes a reversal of the flow in some flux tubes and two stagnation points for the flow between the target and the symmetry point (itself an stagnation point), also found in 2D models [12]. Flow reversal can be measured experimentally if the probe is inserted between these stagnation points. An upper estimate for the position of the stagnation point closer to the target is obtained by integrating along the field the particle source for every flux tube (momentum transfer from neighbouring flux tubes will move the stagnation points closer to the target). The distance along the field from the target to the stagnation point \( (\xi') \) is

\[
n_t c_{\xi'} = \int_0^{\xi'} S(\xi') \, d\xi'
\]

where \( n_t c_{\xi'} \) is the flux to the target and \( S(\xi') \) the particle source in the flux tube. For flux tubes close to the separatrix no additional stagnation points associated with flow reversal are obtained. However, away from the separatrix additional stagnation points are found. For the higher density case the stagnation point is approximately at the position of the reciprocating probe for the flux tube number 6 (see Fig.4). For flux tubes 7 and 8 flow reversal occurs but within the divertor channel. At lower densities the results are similar but the distance to the stagnation points is longer. This is consistent with the experimental measurements in which flow reversal is only measured for the low density case. However, to obtain an exact description of this flow behaviour accurate 2D modelling is needed.

4. Conclusions. Plasma convection has been measured in the SOL of JET diverted discharges with a Mach/Langmuir reciprocating probe. The Mach number deduced at the probe position is about 0.3 within the range of densities 0.9 - 2.5 \( 10^{19} \, m^{-3} \) for ohmic and L-mode discharges with low additional heating. A similar Mach number is deduced for H-mode discharges. In the low density L-mode and in H-modes a flow reversal for field lines close to the separatrix is found. 1-D modelling of these discharges produces a proportion of 40% convected heat flux into the divertor and 60% conducted heat flux, although these values may change slightly due to non local effects in the ion conduction and they ignore flow reversal. A qualitative explanation of the dependence of the flow reversal observed experimentally is found in the analysis of the modelled particle sources in the various flux tubes of the SOL which are influenced by magnetic geometry effects.

Acknowledgment: P.J. Harbour is acknowledged for enlightening discussions.

5. References.

Fig. 1. Position of the reciprocating probe and upper single null MHD equilibrium in JET.

Fig. 2. Measured ion fluxes versus distance normal to the flux surfaces at the probe (medium density L-mode).

Fig. 3. Measured ion fluxes versus distance normal to the flux surfaces at the probe (low density L-mode).

Fig. 4. Modelled fluxes at the divertor target ($\Gamma_\perp$, along $\vec{B}$, $\Gamma_{\text{normal}}$ onto the target) and ionization source for various flux tubes. Modelled and observed regions of flow reversal are shown.
THE IMPORTANCE OF THE ION GRAD B DRIFT DIRECTION FOR THE DIVERTOR PLASMA AT JET

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Introduction: Due to a lower power threshold for the H-mode, single null X-point discharges in JET as well as in other divertor Tokamaks are normally operated with the ion grad B drift pointing towards the divertor target (normal ion grad B drift direction). In such discharges a strong asymmetry in conducted power (up to 1 : 3) between the two divertor branches is commonly observed in JET and elsewhere [1,2]. This asymmetry is a well known problem for the power exhaust in a next step machine. However, its consequences for the plasma parameters in the divertor and so for its performance have in general been insufficiently emphasised so far. In this paper we demonstrate that optimum divertor plasma parameters are achieved when a single null divertor is operated with the ion grad B drift pointing away from the target (reversed ion grad B drift direction).

Experiments: During the 91/92 experimental campaign, extensive power and density scans (n_e ~ 2.5 to 6x10^{19} m^{-3}, P_{IN} = 3 to 12 MW) were performed in single null X-point discharges with normal ion grad B drift direction on both the Beryllium and the Carbon targets. Regrettably no such systematic data were obtained with reversed ion grad B drift direction. In the latter case only three different types of single null X-point discharges, namely hot ion H-mode discharges (n_e ~ 2 to 3x10^{19} m^{-3}, P_{IN} = 15 MW, C- and Be-target), radiative divertor discharges (n_e ~ 6 to 7x10^{19} m^{-3}, P_{IN} = 10 and 22 MW, Be-target) and a few medium density H-mode discharges (n_e ~ 4x10^{19} m^{-3}, P_{IN} = 12 MW, Be-target) are available in the JET data base. Due to the detachment of both strike zones in radiative divertor discharges, the plasma parameters in front of the target could not be measured, which excludes these discharges from a systematic comparison. In hot ion H-mode discharges, the strong temporal variation of the power conducted into the scrape off layer / divertor (dW/dt, C-bloom), which has a significant effect on the divertor plasma, complicates the data interpretation. The medium density H-mode discharges were performed late in the experimental campaign. The Be target was already slightly damaged and only a few of the target Langmuir probes were operational. Despite of the consequently bigger error bars, their qualitative behaviour was similar to that seen in hot ion H-mode discharges which is discussed below. This very restricted data set for the reversed ion grad B drift direction therefore does not allow direct comparison of similar discharges (same density, same heating power, similar power into the SOL) with opposite ion grad B drift directions. However, due to the significantly different behaviour of the divertor plasma when operating with reversed ion grad B drift compared to normal ion grad B drift, conclusions can nevertheless be drawn. All the data discussed in this paper were obtained during H-modes due to the larger amount of data available. The L-mode behaviour is very similar.

The plasma parameters (electron density - , D^+-flux, and Te- profiles) in the divertor are measured by Langmuir probes embedded in the target tiles applying the same evaluation method as in [3]. In addition spatially resolved Hα measurements (CCD) are available for the C-target. Combining these Hα data with the Te profiles from the Langmuir probes yields neutral deuterium influx profiles. While quantitatively the Hα based D-flux deviates in some cases up to a factor of two from the D^+-flux obtained by Langmuir probes, the qualitative behaviour (in-out asymmetry) is very similar (Fig. 1). Due to the fact that the spatially resolved
Results and discussion: Fig. 1 shows the $T_e$ and the $D^+$-flux profiles on the C-target during the ELM free H-mode phase of two low density ($n_e \sim 3 \times 10^{19} \text{m}^{-3}$) discharges: i.) a 12 MW NBI heated discharge with normal ion grad B drift direction (Fig. 1a, 1b); ii.) a 15 MW NBI heated hot ion H-mode discharge (V-H-mode) with reversed ion grad B drift direction (Fig. 1c, 1d).

In the case with normal ion grad B drift direction (Fig. 1a, 1b) a relatively high $T_e$ (~40 eV) and a low $D^+$-flux (~$5 \times 10^{21} \text{m}^{-2} \text{s}^{-1}$) can be observed at the outer strike zone, while the inner strike zone displays low $T_e$ (~10 eV) and a relatively high $D^+$-flux (~$2 \times 10^{22} \text{m}^{-2} \text{s}^{-1}$). This behaviour indicates that already at low mid plane density two different divertor regimes (low recycling, high recycling) start to develop at the two strike zones. The hot ion H-mode discharge with reversed ion grad B drift direction (Fig. 1c, 1d) has a more symmetrical distribution of $T_e$ (~60 eV) and the $D^+$-flux (~$2 \times 10^{22} \text{m}^{-2} \text{s}^{-1}$) between the two divertor divertor legs. In this case both divertor branches are in the same divertor regime. The difference in the absolute values of the divertor plasma parameters and in particular of $T_e$ between these two discharges can be explained by their different heating power and confinement regimes. The behaviour of the divertor plasma parameters shown in Fig. 1 is not only representative for the whole duration of these two discharges but also for the effect of the ion grad B drift on the divertor plasma in general. Fig. 2 shows the separatrix density at the inner and outer strike zone (from Langmuir probes) plotted versus the line averaged density in the main plasma for both ion grad B drift directions and many discharges.
different. In addition, different dependencies on the scrape off layer density $n_s$ will occur, resulting in increasingly divergent $n_d$ and $T_d$ between the divertor legs with growing $n_s$. In the Single Null X-point discharges described above, with normal ion grad B drift direction (Fig. 2, solid symbols) the behaviour of the plasma parameters in the two divertor branches is therefore qualitatively in line with expectations from the 2 point model for a strong asymmetry in conducted power between the two divertor branches. In discharges with reversed ion grad B drift direction (open symbols), the measured plasma parameters in the two divertor legs suggest a more even distribution of conducted power. In this case the two divertor legs are expected to display similar plasma parameters over a wide range of densities and powers conducted into the divertor, as seen in the experiment (Fig. 2). This allows for a given SOL power, higher main plasma densities to be achieved before divertor detachment occurs, and makes it possible to obtain a low temperature high density plasma in both inner and outer divertors simultaneously. When assessing the power loading of the Be divertor target from tile temperature measurements (by a CCD camera through a 844.5 nm filter) during discharges with reversed ion grad B drift direction the picture is not so clear anymore. These data show a pronounced effect of the ion grad B drift direction at low mid plane density, which seems to vanish at high densities. From the bigger outer surface of the main plasma one would always expect a higher power loading at the outer divertor strike zone. This geometric effect seems to be counteracted by a force which depends on the direction of B as well as on the density (collisionality). The density dependence can be inferred by comparing hot ion H-mode discharges with radiative divertor discharges [6], representing the two extreme ends in the scanned density (collisionality) range. While the hot ion H-mode discharges show clearly a stronger heating at the inner strike zone, the radiative divertor discharges display burn through and consequent target tile heating predominantly at the outer strike zone. Grad $T_e \times B$ and Grad $T_i \times B$ forces, which would vanish if $T_e = T_i$ and which would be strong if $T_e \ll T_i$, are suggested as an explanation of the observed effects [7].

In addition to the importance of the ion grad B drift direction described above, the X-point to target distance (connection length, divertor volume) also has some influence on the divertor in-out asymmetry. During an X-point to target distance scan ($\Delta x = 8, 16, 25$ cm) with normal ion grad B drift, a reduced asymmetry in the densities and temperatures measured at the two strike zones was observed at the maximum $\Delta x$ [3].

Conclusions: The achievement of a high density low temperature divertor plasma (high recycling regime) is essential for good impurity retention as well as for low target power load and low target erosion (impurity production). In a next step machine stable divertor regimes beyond the high recycling regime (radiative divertor, gas target) have to be obtained in order to solve the power exhaust problem. Such a divertor can only work if the same divertor regime is obtained simultaneously in both strike zones over a wide parameter range. In order to achieve this, the conducted power to outer and inner strike zone has to be approximately equal. The only way found so far in JET which accomplishes this is to have the ion grad B drift pointing away from the divertor target.

[7] A. Herrmann et. al. this conference
In the case of 12 MW NBI heated ELM free H-mode discharges with normal ion grad B drift direction (solid symbols) the inner strike zone density (solid points) starts by a factor of 2 higher than the outer strike zone density (solid triangles) and detaches when the main plasma density is increased (high recycling regime -> gas target). The outer strike zone density on the other hand remains relatively low (low recycling regime) over the whole density range. When the main plasma density is increased to the point where the outer strike zone should enter the high recycling regime, the inner strike zone is starved of power, detaches and ultimately a density limit disruption (Marfe) occurs. During the 91/92 campaign, it was therefore not possible to achieve a high recycling outer divertor strike zone with normal ion grad B drift direction, even for high mid plane densities. As already discussed above, hot ion H-mode discharges with reversed ion grad B drift direction (Fig. 2: open symbols) display more symmetrical plasma parameters in the two divertor channels. In this case both divertor branches are in the same regime for a wide range of densities, allowing the achievement of high recycling conditions simultaneously at the inner (Fig. 2: crosses) and outer divertor leg (Fig. 2: open triangles). Due to the high radiation losses during C-blooms and the consequent reduction in the power conducted into the divertor late in the heating phase of these discharges, very high divertor densities are achieved at moderate main plasma densities (Fig. 2: open symbols).

The above described behaviour of the divertor plasma can be qualitatively understood by a simple one dimensional analytical model for the scrape off plasma such as the two point model [4,5]. In this simple model the divertor regime is mainly determined by the density at the mid plane separatrix and by the power flux into the divertor channel. When solving the equations of this model, the dependence of the divertor density \( n_d \) and temperature \( T_d \) on the mid plane scrape off layer density \( n_s \) and on the power flux parallel to the field lines \( q_B \) is approximately

\[
  n_d \approx \frac{n_s^3 L_{||}}{q_B}, \quad T_d \approx \frac{q_B}{n_s^2 L_{||}} \tag{1}
\]

Therefore for fixed \( n_s \), but different powers \( q_B \) conducted into each divertor branch at a given connection length \( L_{||} \), the corresponding divertor densities \( n_d \) and temperatures \( T_d \) will be
PLASMA FUELLING EXPERIMENTS IN JET AND IMPLICATIONS FOR FUTURE DIVERTOR OPERATIONS

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1) Introduction

The size of the deuterium density reached by fuelling a tokamak discharge with a given external particle source depends, among others, on three factors [1,2,3]: the plasma particle confinement time, the neutral particle screening properties of the plasma edge, and the properties of the material surfaces. These factors affect also the size of the recycling flux which constitutes another but internal particle source for fuelling. Fuelling experiments were performed to study the role of these factors in JET discharges.

2) Experimental

a) Dependence of plasma density on poloidal position of gas valve. In a series of ohmic discharges (3MA, 2.8T, densities 1-3*10¹⁹ m⁻³) in double null X-point configurations plasmas were fuelled by gas valves that were located at the plasma midplane and at the private region of the divertor, respectively. At these two positions the fuelling conditions for neutral deuterium are rather different. At the plasma midplane the plasma scrape off layer falls off radially within one to two centimeters whereas at the divertor this fall-off expands by a factor of about three and the divertor plasma poloidally extends about ten centimeters between the divertor target plates and the X-point. The effect of this difference can be seen experimentally when the plasma changes from limiter (close to plasma midplane) to X-point configuration (figure 1(b)). The density decreases due to a decrease of the fuelling efficiency for recycling [3,5]. However, in the experiments here where the location of the active gas valve was changed from plasma midplane to X-point position nevertheless the same density and recycling fluxes were achieved.

b) Plasma density and flux variation during neutral beam injection. A limiter discharge was heated by a 15MW deuterium beam providing a fuelling rate of 1.5*10²¹ particles/s. At the start of the beam injection there is a slow increase of the plasma density but a rapid increase of the recycling flux at the limiter, see Figs. 2 (b), (d).

3) Analysis

For studying the fuelling situation a particle balance model is employed details of which can be found in reference [3]. The fundamental assumption of the model, corroborated by experiments is that particles that are admitted into the torus mainly reside in the material surfaces rather than in the plasma. However, there is a mutual particle exchange. This is affected by reflection at the material surface (reflection coefficient r) and neutral particle screening at the plasma edge (coefficient f). Particles are assumed to stay for an average particle confinement time τ [1,2] within the confined plasma (enclosed by the last closed magnetic flux
surface). Particles that leave the plasma can be absorbed by material surfaces. Since beryllium has been used the JET first wall in quasi steady state exhibit a characteristic metallic behaviour with regard to hydrogen bombardment [3] and so it is assumed that particles in the material surfaces are subject to diffusion into the bulk material and recombination at the surface forming a desorbing particle flux $\Phi_w$. The full particle balance equation reads:

$$\frac{dN_p}{dt} = -\frac{N_p}{\tau} + \frac{f}{1-r(1-f)} \left[ r \frac{N_p}{\tau} + \Phi_w + r(1-f_{ex})\Phi_{ex} \right] + f_{ex}\Phi_{ex} \quad (1)$$

On the right hand side the first term is the particle flux leaving the confined plasma. Then, within parantheses, there is the recycling flux composed of the reflected plasma particle flux, the particle flux that desorbs from material surfaces (molecules), and the reflected particle flux of external particle sources that do not directly fuel the plasma due to a restricted external fuelling efficiency ($f_{ex}$). These fluxes are assumed to fuel the plasma with an efficiency $f$. The fraction $(1-f)$ is assumed to return to the surface. The forefactor in front of the parantheses in equation (2) takes account of multiple recycling process. The last term of eq.(1) describes the direct fuelling by external particle sources. The model [3] combines eq.(1) with the diffusion equation for particles in the material to calculate selfconsistently the desorbed particle flux $\Phi_w$ and the recycling flux. Calculations are time dependent to follow the entire discharge period.

a) Application to gas fuelling experiments in X-point configuration

It is instructive to simplify equation (1) and to derive an expression for the plasma particle inventory in steady state. The reflected and desorbed particle fluxes can be combined into a recycling flux which is usually assumed to be proportional to the total impacting flux with a proportionality factor $R$ as the recycling coefficient at the material surface. With this and substitution of $r$ by $R$ in eq. (2) the plasma particle inventory reads:

$$N_p = \tau \left[ \frac{fR}{1-R} + f_{ex} \right] \Phi_{ex} = \tau f_{eff} \Phi_{ex} \quad (2)$$

The first term within parantheses is the contribution by recycling and the second term is the contribution by direct fuelling. It is useful to define the effective fuelling efficiency $f_{eff}$ (see eq 2) and to calculate it by computing $f$ and $R$ using the detailed model of above. For $R$ approaching one, $f_{eff}$ can get much larger than $f_{ex}$, and recycling is then dominating the fuelling. The measured plasma particle inventory and the external particle fluxes are used as an input for the calculations. The fuelling efficiency $f$ for recycling is allowed to be time dependent, other parameters are assumed to be constant. By iteration, the set of parameters is found with which the calculated recycling flux fits best the experimental flux derived from H-alpha measurements. The reflection coefficient had to be assumed to be about 0.7, suggesting low energy ($\leq 100$eV) deuterium impacting at shallow angles on carbon/beryllium [5]. A particle confinement time of 0.3s was assumed [6]. However, the important quantity is the product of confinement time with
the recycling fuelling efficiency \( f \). Errors in \( \tau \) can be tolerated as long as \( f \) stays within credible limits (0<\( f \)≤1). Assuming molecular ionisation, dissociation, and charge exchange the fuelling efficiency of neutral molecular hydrogen thermally released from a position that is a scrape off layer width behind the last closed flux surface can be estimated to be about 0.7 or less for plasma edge densities of above \( 10^{18} \text{ m}^{-3} \) and edge electron temperatures of above 20eV. The fuelling efficiency of an externally fuelled hydrogen molecule is likely to be smaller as it has to pass through the entire scrape off layer. For \( f_{ex} \), during X-point phase two values of 0.5 and 0.1 were assumed to simulate plasma midplane and divertor fuelling, respectively. Particle up-take by materials is determined by the parameter (\( A*D/K \)), where \( A \) is the effective material area bombarded by particles, \( D \) is the effective particle diffusion and \( K \) is the effective recombination coefficient. For limiter and divertor configuration a value of \( 2*10^{22} \) gave good results.

Results of the analysis are indicated in figures 1, (f)-(k). During the initial start up limiter phase (\( t<2s \)) the recycling flux is not well simulated suggesting inaccuracies in the coefficients assumed to be constant during plasma current ramp up. During the constant current phase agreement is good and results are almost independent of \( f_{ex} \). As expected \( f \) decreases after transition to the X-point. The recycling coefficient is approaching one. The effective fuelling efficiency reaches a value of 2 in X-point configuration (even 5 in limiter phase). Compared with the values for \( f_{ex} \) of 0.5 or 0.1, \( f_{eff} \) is clearly larger indicating recycling to be the most important fuelling source in both fuelling schemes, hence the similarity of experimental results.

b) Application to neutral beam heated limiter discharges

Additional plasma heating is likely to increase the power deposition on the limiter surface. It is therefore suitable to allow the surface related parameter (\( A*D/K \)) to be time dependent. The external fuelling efficiency for neutral beams was assumed to be 1. The recycling fuelling efficiency was set to a constant value (limiter phase) of 0.5 and a particle confinement time of 0.3s was assumed. For \( r \) a value of 0.5 proved to be necessary.

Results are indicated in fig.2. Important is the section where \( t>9s \). For good simulation of the recycling fluxes during neutral beam phase a decrease of the parameter (\( A*D/K \)), is necessary, probably caused by enhanced power deposition, as well as a decrease of the particle confinement time from 0.3s to 0.2 s (contrary to the initial assumption). This is consistent with an observed 30% decrease of the energy confinement during transition from ohmnic to L-mode.

4. Conclusions

It has been shown that plasma density and recycling fluxes are affected by the fuelling efficiency controlled by plasma edge conditions, by the recycling coefficient at the material surface controlled by material properties as well as particle fluxes, and by the plasma particle confinement time controlled by transport. For recycling coefficients close to one the recycling fluxes dominate plasma fuelling and the efficiency of external fuelling becomes less important. Using eq. (3) \textit{the recycling flux normalized to the plasma particle losses} can be calculated:
\( \Phi_{r, \text{norm}} = \Phi_r \frac{1}{\tau R} = \frac{1}{\frac{R}{\tau} + \frac{1}{\tau (1-R)}} \)  

It is clear that if \( f_{ex} = f \), for instance by puffing gas into the divertor region, the normalized recycling flux does not depend on the recycling coefficient and hence not on the pumping. In order to maximize the recycling flux the fuelling efficiency should then be made as low as possible, for example by introducing neutral particle baffling and operating at high densities. If \( f_{ex} = 1 \), for instance by using neutral beams or pellets as the only external particle source, then \( f \) should be again minimized but \( R \) should be maximized, hence little pumping would be desired.

5. References

6. Figures
Figure 1. Almost identical fuelling situations in plasma discharges with gas fuelling either from the divertor (dashed in fig 1(g)) or from the plasma midplane. \( I_p := \text{plasma current} \), \( N_p := \text{plasma particle inventory} \), \( \Phi_L := \text{limiter flux} \), \( \Phi_{xp} := X\text{-point flux} \), other parameters see text.

Figure 2. Fuelling situation of a neutral beam fuelled limiter discharge. Meaning of parameters as in figure 1) and text. During the NB-phase a decrease of \( \tau \) (not shown) must be additionally assumed to simulate experimental fluxes.
Asymmetric energy flux to the ASDEX-Upgrade divertor plates determined by thermography and calorimetry

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Introduction

The high power loads on the divertor plates in future fusion reactors like ITER require a power distribution on the plates as uniform as possible. Experiments made on ASDEX-Upgrade show, in agreement with observations on other tokamaks [1,2], that in discharges where the ion grad-B drift direction points away from the X-Point the energy and power distribution on the inner and outer divertor plates is more symmetric than in discharges with reversed magnetic field.

Diagnostics

Thermographic measurements on ASDEX-Upgrade are performed with a new high time resolution infrared camera consisting of a linear focal plane array of 256 indium antimonide diodes (pixels) working at liquid nitrogen temperature. The detector line is orientated along the major radius viewing both the inner and outer target plates. The spatial resolution is 2.72 mm/pixel at the outer divertor plate and 3.4 mm/pixel at the inner plate. The camera is calibrated and so the absolute surface temperatures of the divertor plates, which are made of carbon tiles, are determined. The highest time resolution is 128 μs per profile. From the time-dependent surface temperature distribution the corresponding energy flux onto the divertor plates is calculated by solving numerically the 2D nonlinear heat conduction equation (temperature dependent heat diffusion coefficient).

The cooling water calorimetry system monitors the energy deposited on the inner and outer target plates and on the heat shield in each of the 16 toroidal sections of ASDEX-Upgrade. The cooling water calorimetry system consists of resistor thermometers and flow meters located in the cooling water system downstream and upstream of the individual cooling units, respectively.

Experiments

The energy flux to the ASDEX-Upgrade divertor plates is routinely measured. The shots under investigation are single null divertor discharges with a density flat top. The inner vessel was boronized and the discharge parameters were as follows: \( I_p = (600–800) \text{kA} \), \( B_t = (1.35–2) \text{T} \) and \( B_s = + 2 \text{T} \) giving \( q = (2.8 – 5) \). ICRH up to 1.8 MW was sometimes used.

Results

Investigations of the behaviour of the power and energy deposition were made for various types of discharges.

The toroidal distribution of the energy load to the divertor plates as measured by cooling water calorimetry can be taken as toroidally uniform within the limits of the experimental errors (±15 kJ) (fig.1).

The power density calculated from the thermography measurement spatially integrated over the inner and outer divertor plates, respectively, and time integrated over a stable discharge phase (1.4 s — 1.9 s) is used to calculate the ratio between the energy flux to the
outer and to the inner divertor plates (fig. 2). The ratio of outboard to inboard energy load depends on the orientation of \( B_t \). If the ion grad-B drift direction is towards the X-point the mean value of the asymmetry ratio is of the order of two, otherwise close to one (1.3). The time and spatially resolved power density profiles as obtained from thermography exhibit a clear dependence on the confinement regime and on the direction of the magnetic field.

In ohmic L-Mode discharges with the ion grad-B drift direction towards the X-point we find a ratio of the total power deposited onto the inner divertor plate to the total power load onto the outer plate of 1.7. When the plasma changes to ohmic H-Mode [3] this asymmetry is found to rise by 20% (fig. 3). The power density profiles are significantly different for both modes. The profiles are sharper during the H-Mode (fig. 4 a,b). Also the asymmetry ratio of the spatial maxima of the power density raised during the H-Mode indicating that the power distribution on both the inner and outer plates is essentially the same. H-Mode experiments with reversed ion grad-B drift have not yet been performed.

Discussion
The observed variation of the poloidal power deposition behaviour with the direction of the ion grad-B drift suggests a strong dependence of the poloidal fluxes in the scrape-off layer on the magnetic field. Gradients along field lines that might drive these fluxes are recently discussed in the literature [4]. A strong contribution can be due to classical electron and ion diamagnetic heat fluxes

\[
\vec{q}_{e/i} = \pm \frac{5}{2} \frac{P_{e/i}}{Z_{e/i} e B} \left[ \vec{B} \times \nabla T_{e/i} \right].
\]

The radial temperature gradients drive poloidal heat fluxes with different directions for electrons and ions. In a pure hydrogen plasma both terms exactly cancel, if electrons and ions have identical temperature profiles. A net contribution requires different electron and ion temperatures and/or different temperature gradients.

Assuming the ion temperature and temperature gradient at the separatrix to be higher by a factor of 5 than the corresponding values for electrons, one can account for the change of the asymmetry in the energy deposition when the magnetic field is reversed. An analysis with data from the cooling-water-calorimetry was made for a series of ICR-heated discharges with \( B = -2T \) (ion grad-B drift towards the X-point) and \( B = +2T \) (ion grad-B drift away from the X-point), which showed, that, with the ion contribution prevailing, the diamagnetic heat flux could yield a quantitatively suitable reduction in the asymmetry when the magnetic field is reversed. First detailed calculations with the 2D multi-species fluid BRAAMS-code [5] coupled to the neutral gas Monte-Carlo code EIRENE [6] taking into account the diamagnetic heat fluxes confirm this result.

As during H-Mode a steepening of the ion temperature profile has to be expected [7], the observation, that during the transition from L-Mode to H-Mode the asymmetry increases for discharges where the ion grad-B drift direction is towards the X-point, can be taken as additional evidence for the importance of the contribution of the ion diamagnetic heat flux.

Conclusions
Measurements of the thermography and the cooling water calorimetry show that the power load to the inner and outer divertor plates is more symmetric in discharges where the ion grad-B direction points away from the X-point. This holds not only for time integrated but also time resolved measurements. If the ion grad-B drift direction is towards the X-point an increase of the asymmetry ratio during the H-Mode is observed. The findings are qualitatively explained by poloidal heat fluxes driven by different ion and electron temperature and profiles. From this it is expected that the asymmetry decreases in the H-
Mode if the ion grad-B direction is away from the X-point. More experiments are needed to prove this hypothesis. BRAAMS-EIRENE-code modelling is also envisaged.

References
7. R.Schneider et al, 18th EPS Conference on Controlled Fusion and plasma heating, Berlin 199, Europhysics Conference Abstracts Vol.15C, part III, p.117

Figure captions
Fig. 1 Toroidal distribution of the energy flux to ASDEX-Upgrade divertor plates measured with the cooling water calorimetry (#2969, I_p=600 kA, B_t=-2T)
Fig. 2 Distribution of the asymmetry ratio — power on outer/inner plates — for discharges with ion-grad-B-drift direction towards and away the X-point.
Fig. 3 Time-dependence of the asymmetry ratio for an discharge with ohmic H-Mode (#2393, I_p=800 kA, B_t=-1.35T, ohmic H-Mode: (1.32—1.98) s)
Fig. 4 3D plot of the power density deposited onto the target plates (a) and power density profiles (b) for different confinement modes (t=1.53 s — ohmic H-Mode (solid), t=2.13 s — H-L transition (dashed), t=2.23 s — L-Mode (dash-dot)).
Distribution of asymmetry ratio

Fig. 2

ASDEX Upgrade

Fig. 3

Fig. 4a

Fig. 4b
THE INFLUENCE OF SPATIALLY AND TEMPORALLY VARYING EDGE CONDITIONS ON THE INTERPRETATION OF SPECTROSCOPIC PARTICLE FLUX MEASUREMENTS

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1 Introduction  Photon fluxes emerging from the plasma boundary in the vicinity of material surfaces are often directly interpreted in terms of particle fluxes. The factor of proportionality (inverse photon efficiency) is the ratio of the ionization and excitation probabilities, $S/X_B$, including the branching ratio $B$. Generally, $S/X_B$ is taken to be constant and represented by an average value which is typical for the plasma parameters where ionization and excitation is essential. This approach is convenient and justified, as long as $S/X_B$ is nearly independent of the density and a weak function of temperature. If these conditions do not apply, however, or very fast changes in the plasma conditions occur, a more sophisticated analysis is necessary.

We have developed a simple numerical model to interpret the measured photon fluxes under such general conditions. The penetration, ionization and transport of impurity atoms is modelled using measured density and temperature profiles and estimated values for the energy of the sputtered particles. The photon efficiencies obtained exhibit considerable variations with the plasma density and diffusion coefficient.

2 A simple particle transport code for the interpretation of photon fluxes. Impurity particle transport is treated in the following, simplified 1-d geometry: Neutral impurity atoms start at the wall monoenergetically with a spatial $cos \theta$ distribution, leading to a radial velocity distribution in the 1-d model which determines the spatial profile of the ionization source. The ions are then subsequently ionized to higher states. The only transport taken into account is radial diffusion with an impurity diffusion coefficient $D$. This picture is appropriate, as long as measurements in front of a large, homogeneous surface are concerned, such as the ASDEX-Upgrade heatshield considered here. All charge states are assumed to be subject to the same value of $D$. Recombination is also included in the model, but is usually negligible for the plasma parameters under consideration.

Rate coefficients for ionization and excitation are calculated from standard formulas (Lotz, Van Regemorter) neglecting multistep collisions. For cases, where more sophisticated values for $S$ and $X$ are available in the literature (e.g., Behringer et al.), these values are reproduced by adjustment of parameters. As an example, the $S/X_B$ is plotted as a function of the electron temperature in Fig. 1 for the transition $C^{III}(464.7nm)$, which is further investigated in the following.

The particle transport calculations predicts the photon flux measured by an observer for a given impurity influx. The photon flux is used to determine the effective inverse
 photon efficiency, $S/XB_{eff} = \Gamma_{in, particle}/\Gamma_{photon}$.

Special emphasis is placed in the simulation of the measured ion temperature, $T_I$, in order to obtain information about the deuteron temperature. $T_I$ is determined by the deuteron temperature and the ion-ion collision rate during the ionization time. To reconstruct the measurements, Gaussian profiles are summed up along the line-of-sight, whose line widths and heights are given by the local impurity ion temperature and line emissivity. Apparent $C^{2+}(2s2p^3P)$ ion temperatures are determined experimentally from the Doppler broadening of the spectral line, whose $\sigma$-components are suppressed by a polarizer. The comparison of simulated and measured linewidths can be used to determine the edge ion temperature. The procedure described above is applied to a number of discharges in the ASDEX-Upgrade tokamak.

3 Experimental results

Transport calculations were carried out to interprete a series of ohmically heated ASDEX-Upgrade discharges with different line-averaged densities.
CIII photon fluxes and line profiles were measured using radial viewing chords against the inner heat shield. Fig. 2 shows corresponding radial profiles on the high-field-side for a discharge with $\bar{n}_e = 3.25 \times 10^{19} \text{m}^{-3}$. The electron density and temperature profiles were constructed from measurements with a moveable Langmuir probe in the divertor region which were smoothly linked to Thomson scattering and interferometric data in the main plasma. Profile mapping into the region of interest was done assuming $T_e$ and $n_e$ to be constant on flux surfaces. This assumption was found to be acceptable also for the separatrix and SOL plasma for the low-recycling divertor conditions. The overall uncertainty of the local parameters is estimated to be within a factor of 2. The ion temperature profile was obtained by multiplying $T_e(x)$ with a factor 1.6 in order to match the simulated and measured $T_{\text{C}^{2+}}$. The energy of the neutral C-atoms and diffusion coefficient were set to 3 eV and 1 m$^2$/s, respectively. Width and position of the simulated emission shell have been confirmed within the uncertainty of ±1 cm by the deconvolution of corresponding measurements on the low field side using tangential viewing chords.

In order to demonstrate its dependence on plasma parameters, we have calculated $S/XB_{\text{eff}}$ for various line-averaged densities. For the transport simulations, we started with the profiles shown in Fig. 2 and scaled the local density with $E$ and the electron temperature with $W_e/5$. Fig. 3 shows the dependence of $S/XB_{\text{eff}}$ on $\bar{n}_e$ and the diffusion coefficient $D$, which is used to model the carbon transport in general.

![Figure 3](image)

**Figure 3** $S/XB_{\text{eff}}$ versus density and diffusion coefficient. The density-dependence originates from the reduction of $T_e$ at the position of ionization. The increase of $S/XB_{\text{eff}}$ with $D$ results from the loss of $\text{C}^{2+}$ ions before photon emission. $n_e-T_e$-variation with $D=1 \text{ m}^2/\text{s}$, $D$-variation with $\bar{n}_e = 4 \times 10^{19} \text{ m}^{-3}$.

The density-dependent, effective photon efficiencies shown in Fig. 3 are used to evaluate measured particle fluxes for discharges with different line-averaged densities in Fig. 4. As a result, a linear dependence of the particle influx on $\bar{n}_e$ is found. A comparison of measured and simulated ion temperatures is shown in Fig. 5. While simulated and measured temperatures match at high $\bar{n}_e$, $T_i$ has to be increased at low densities to obtain agreement. For $\bar{n}_e = 2 \times 10^{19} \text{ m}^{-3}$, the enhancement factor for the ion temperature profile is 2. The enhancement factor $T_i/T_e$ is not simply proportional to the ratio $T_z/T_i$ owing to the competition of ionization and impurity-ion collisions ($\tau_{iz}$...
increases with $T_1$). The predicted $C^{2+}$ ion temperature depends also on the choice of the value of $D$: When $D$ is changed from 0.3 to 3 m$^2$/s, the calculated ($n_e = 4 \times 10^{18}$ m$^{-3}$) $T_{C^{2+}}$ increases from 20 to 25 eV owing to an inward shift of the emission radius.

The analysis gets more complicated, if fast changes of plasma parameters (e.g., owing to density fluctuations) have to be taken into account. In this case, the ionization time of the observed species, $\tau_{ion}$, is found to be the critical parameter: For parameter changes which are faster than $\tau_{ion}$, the proportionality between photon- and particle flux is lost, and the photon flux directly follows the density variations.

4 Discussion

The analysis of the spectral line emission produced by carbon ions entering the plasma edge has shown that transport calculations are essential to obtain a quantitative interpretation in terms of particle flux. The simulation of the measured, apparent impurity ion temperature yields information about the temperature of the plasma ions, including the spatial localisation.

The application of the transport calculations to photon fluxes emitted by $C^{2+}$ entering the plasma in front of the inner heat shield of ASDEX-Upgrade indicates that the carbon influx increases in proportion to the plasma density $n_e$. The ion temperature at the separatrix equals $T_e$ at high densities but exceeds $T_e$ considerably for low values of $n_e$.

References

/3/ S. Pitcher et al., this conference.
/4/ H.-S. Bosch et al., this conference.
Spectroscopic investigation of ELM phenomena in the ASDEX-Upgrade divertor with high time resolution

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Introduction: Improved tokamak H-mode confinement is associated with the formation of an insulating zone just within the separatrix. At a critical pressure gradient a sudden burst of MHD activity (an ELM) degrades edge confinement, releasing particles and energy into the scrape-off layer (SOL) which is subsequently transported to the divertor. Here, these phenomena are studied using spectroscopic diagnostics and target plate thermography of high spatial and temporal resolution.

Spectroscopic diagnostics: The SOL/divertor region of ASDEX-Upgrade (AUG) is viewed by multi-chord UV/visible spectrometer systems. One set of 16 chords views tangentially to the outer target plate covering the region from 1.5mm above the plate to above the X-point with good spatial (≈2mm) and temporal (τ<1μs) resolution [1](see Fig. 1). The outer target plate is also observed from just above the mid-plane by a grating spectrometer fitted with a photo-diode array giving ≈600 chords across the plate but with less time resolution (τ≈25ms).

Results: The results discussed below are from an ICRH heated (30MHz, P_{ICRH}=1.47MW) single-null discharge in deuterium (shot #2309) with I_p = 600kA, B_t = -2T, n_e = 4 x 10^{19} m^{-3}.

ELM Power and Particle Balances: Another detailed study [2] presents the following scenario for a type III ELM on ASDEX-Upgrade. During the first phase of 10–100 μs MHD turbulence ergodises the last closed flux surfaces releasing a pulse of energy by electron conduction into the SOL. Particles are subsequently lost on a slower timescale of 0.1–1.0 ms. The total energy loss per ELM is about 5–10 kJ (3–6% of the total plasma energy), most being deposited on the outer target plate (see Fig. 3(b)). The energy convected by particles to the plates is much smaller (∼10%) than that conducted through the electrons. Investigating the ELMs H-mode (shot #2309, t=1.83 –1.9 s), in the quiescent phase between ELMs the total D-flux to the target plates as determined from Dα measurements is ≈1.6x10^{21} s^{-1} with an out/in asymmetry of ∼1.3 (see Fig. 3(a)). The mean total particle loss per ELM to the target plates (also determined from Dα measurements) is ≈ 9.6x10^{18}, the particle losses...
to the inner and outer plates being about equal. Because of the short duration of the ELMs (≈1 ms) only about 25% of the mean D-flux to the plates is attributable to them. Note that the inverse emission efficiency of Dα (S/XB) increases moderately with increasing Te and ne (by a factor 2 from Te=5–100 eV at ne=10¹⁹ m⁻³ and by a factor 2.7 from ne=0.5–5.0x10¹⁹ m⁻³ at Te=20 eV). Hence, the fluxes quoted above have an estimated uncertainty of a factor ≈ 2. Furthermore, because Te and ne increase during the ELM, the actual D-flux during the ELM may be higher relative to that between ELMs than stated above by up to a factor ≈ 2.

**Fig. 3 (a)** Deuterium flux from target plate; **(b)** deposited power on target plate; and **(c)** CIII (464.9 nm) intensity 1 cm above target plate (Shot #2309).

*Temporal behaviour:* Fig. 2 shows, at the commencement of magnetic activity (a), enhanced heat and particle fluxes into the SOL as seen by a fast bolometer channel viewing outer SOL (c). This is associated with an increase of Dα (b) and CIII(464.9 nm) line radiation (d & e) above the outer target plate. Channels viewing 1 cm above the plate (e) and through the x-point (d) show a simultaneous increase in CIII intensity. This behaviour is shown particularly clearly in Fig. 5 which shows Dα (a) and CIII (b) intensity profiles above the outer target (y = chord elevation above plate centre) as a function of time. The CIII intensity increases within 11 μs throughout the profile which spans from the plate to above the X-point. This is much faster than for transport of deuterons from the mid-plane to the plate (τ_e = 0.2 ms at Te=50 eV, L_e = 15 m) and implies that the intensity increase is due to an increase in Te caused by electron conduction (τ_e = 3.6 μs). The CIII(464.9 nm) [3p 3P–3s 3S] line is excited from the metastable 3s 3S state (excitation energy 2.67 eV). Due to the large energy gap of E_m = 29.5 eV between the 3s 1S ground state and the metastable 3s 3S state the excitation coefficient has a strong Te dependence, increasing by a factor of 25 from 5 eV to 50 eV. The initial rise in CIII intensity by a factor 2 to 3 can thus easily be explained by increased excitation due to an increase in Te. An increase of Te would, however, also cause an immediate increase in the sputtered C-flux due to the increased sheath potential and hence incident particle energies. Over the energy range of interest the sputtering yield is, however, not a strong function of D+ energy (E_E = 60–300 eV, Y = 1–3%). A proposed explanation [2] for the factor ≈ 3 increase in intensity in the channels intersecting the X-point region is that C²⁺ ions (x_i = 47.9 eV) in the cold ‘private’ region (Te = 5–10 eV) are excited to the metastable state after mixing with hotter plasma by ELM-generated magnetic turbulence around the X-point. This radiation decays after ≈ 0.3 ms implying ionisation of the C²⁺ ions (τ_i = 0.14 ms at Te = 20 eV and ne = 10¹⁹ m⁻³) or a reduction of ergodisation. After this initial 0.3 ms the Dα and CIII intensities remain above the pre-ELM level (factor ≈ 2 for Dα), subsequently decaying over the following ≈ 2 ms. This can
be explained by the arrival of the background ions exhausted by the ELM and the associated increase in the sputtering of C from the plate. Some of the increased C-flux may, however, be attributable to a temporary increase of $T_i$ and hence sputtering yield.

**Fig. 4** Profiles of (a) $D\alpha$ and (b) CIII line intensity as function of elevation above the outer target plate during L- and H-modes before an ELM (* - 1.855s) and during the ELM (—- - 1.855s and o - 1.855s).

![Graph](image1)

**Fig. 5** Line intensity profiles of (a) $D\alpha$ (656.1 nm) and (b) CIII (464.9 nm) above outer target plate during an ELM (Shot # 2309)

![Graph](image2)

**Intensity Profiles:** When comparing profiles of $D\alpha$ and CIII intensity as a function of elevation above the target before and during an ELM (see Figs. 4 (a,b)) large increases in the $D\alpha$ intensity in front of the plate and CIII intensity in the vicinity of the X-point are observed. Otherwise, little change in the form of the profiles is observed. Figs. 5 (a, b) show the temporal development of these profiles in more detail. Figs. 6 (a,b) show time averaged ($\tau_a=25$ms) intensity profiles of $D\alpha$ and CII(656.1nm) lines across the outer plate under L- and H-mode conditions ($I_p=800$ kA, $B_l=-2$ T, $n_e=4x10^{19}$ m$^{-3}$, $P_{ICRH}=1.3$ MW). Such data have been used to study erosion of the plate under Ohmic conditions. Profiles of $T_e$ at the plate are needed for their interpretation in terms of particle fluxes. These are available for some Ohmic shots from a swept Langmuir probe [3]. Erosion yields in agreement with those expected for physical sputtering are found ($Y_C \approx 3-5\%$). The broad wings within the private region, where $T_e$ and $n_e$ are low, can be explained by small particle fluxes but low values of the inverse emission efficiency ($S/XB$). Outside the separatrix the half widths of the profiles are 2.6cm (L-mode) and 3.2cm (H-mode) for $D\alpha$ and 4.5cm (L-mode) and 5.8cm (H-mode) for CII (the H-mode profiles are integrated over several ELMs).
Fig. 6 Profiles of (a) Dα and (b) CII line intensities measured across the outer target plate during ICRH heated L- and H-modes.

Measurements similar to those of Fig. 4 of CII(658.1nm) line intensity show the C⁺ ions to be localised within 2cm of the plate. Preliminary modelling shows that the approximate form of this profile can be explained by ionisation of sputtered C⁰ atoms and parallel diffusion of C⁺ before further ionisation. Similarly the CIII profile may be explained by a source of ionised C⁺ ions, parallel diffusion of C²⁺ (with some backstreaming due to friction with background ions) and further ionisation. Although recombination of C³⁺ does not contribute significantly to the C²⁺ population, explanation of the peak of the CIII intensity near the plate (Fig. 4(b)) may require a contribution from charge exchange recombination with thermal D⁰ ($r_{cx}=0.3$ms and $\tau_i=0.13$ms) with $n_0=10^{12}$m⁻³, $T_e=20$eV and $n_e=10^{13}$m⁻³.

As previously mentioned, about 25% of the D⁺ flux arrives during the ELMs. Thus, under the assumption that $T_e$ and $T_i$ remain constant throughout the ELMs giving constant incident D⁺ energy ($E_0 = 3eT_e + 2eT_i$) and hence sputtering yield, 25% of the sputtered C-flux could be attributed to the ELMs. A detailed investigation, however, awaits $T_e$ measurements of sufficient temporal resolution without which it is not possible to calculate reliable C-fluxes from the spectroscopic measurements.

Conclusions: The behaviour of D and C line emission profiles in the divertor region during ELMs can be understood in conjunction with supporting studies [2]. An initial conducted heat pulse from the ELM (10–100µs) causes an increase in $T_e$ and hence CIII emission throughout this region, particularly near the X-point where C²⁺ ions are mixed with hotter plasma by ergodisation. An immediate increase of sputtered C-flux may be expected due to the increased $T_e$ at the plate. Later (0.1–1.0ms) a particle pulse reaches the target plates causing enhanced sputtering and hence CIII emission. More detailed interpretation in terms of erosion and particle transport will require fast $T_e$ and $n_e$ measurements. More quantitative modelling of the profiles is under development.

References:
INTERPRETATION OF LOW IONIZED IMPURITY DISTRIBUTIONS IN THE ASDEX UPGRADE DIVERTOR


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Introduction
Design studies for reactor-like devices, like ITER, have particularly emphasized the importance of erosion and transport of material from the divertor target plates. In this context experimental measurements which can lead to a better understanding of the underlying physics are highly desirable.

We discuss the spatial profiles of line emission from impurities measured in the divertor of ASDEX Upgrade with a recently developed multi-chord divertor spectrometer system. These profiles are obtained from observations in the ultra-violet/visible spectral range. The divertor spectrometer system was developed particularly to measure the erosion of the divertor plates and to study transport of the impurities and the ionization and recombination processes in the divertor region.

Experimental
As shown in Fig. 1, the divertor spectrometer views along 16 lines of sight tangential to the lower outer divertor plate. Its spatial resolution is better than 3 mm. The chord integrated plasma emission is collected by quartz lenses via a faceted mirror and led to the detector systems by quartz optical fibers. Two detection systems are simultaneously available: A 1 m Czerny-Turner spectrometer fitted with a freely programmable and gateable 2D CCD-camera giving a spectral and temporal resolution of up to 0.03 nm and 20 ms respectively. This system is used for measurements over a limited spectral range in order to select appropriate emission lines and to monitor these under the various operating conditions. The second system consists of two 16 channel photomultiplier arrays which are equipped with different interference filters and is used to measure the intensity of two isolated spectral lines simultaneously with a temporal resolution as high as 11 μs. Interference filters with a spectral bandwidth (FWHM) of about 1 nm are used. The fast detection system has been designed especially to investigate transient phenomena such as MARFEs 1,3 and ELMs 2.

The complete spectrometer system was calibrated between experimental campaigns using an Ulbricht sphere inside the ASDEX Upgrade vessel. The total errors in the relative calibration between the chords and the absolute intensity measurements are estimated as ± 20 % and ± 30 %, respectively. The measurements described below refer to ASDEX Upgrade lower single null operation with the following discharge conditions: deuterium filling gas, plasma current $I_p = 800$ kA, toroidal field $B_t = -2$ T, pulse duration (flat top) $t_d = 4$ s, line-averaged density $n_e = 3 \cdot 10^{19} \text{ m}^{-3}$ and boronized plasma facing components.

Results and discussion
Measurements of intensity profiles in the divertor have been performed under various operational conditions. In order to identify the impurities present in He-discharges and during D-discharges under boronized conditions survey measurements have been made over the whole accessible spectral region. The following impurity species could be identified in He-discharges: H, O, C, CH and as minorities Cr, Cu, Ca, Cl and N and in D-discharges: B, BD, C, CD, O and the minorities He, Ca and Cr.
Intensity distributions have been measured of neutral (BI, CI, Cal, CII, OI) and ionized (BII-BIV, CIII-CV, OII-OIV) atoms. Fig. 2 shows as an example the spatial profile of Cal a trace impurity in the graphite tiles, plotted versus the chord elevation in front of the plate (see Fig. 1). The emission is localized within 2 cm above the plate and peaks sharply at the plate. This intensity profile of the Cal (422.7 nm) line can be explained in terms of ionization according to \( I(x) \propto \exp \left( -\frac{x}{\lambda_{\text{ion}}} \right) \) with the ionization length \( \lambda_{\text{ion}} = \frac{\nu_0}{S} \) where \( S \) is the ionization rate coefficient and \( \nu_0 \) is the initial velocity of the eroded atoms. The measured decay length of 1 cm is in good agreement with estimates based upon plasma parameters measured by Langmuir probes. We can use measured values for \( \lambda_{\text{ion}} \) and \( I(0) \), the intensity at the plate surface, to estimate the integrated particle flux from the plate. Assuming the line emission to be localized within a region of width \( d \) along the lines of sight, this quantity is given by \( \Gamma = \frac{4\pi S}{X_B} \frac{I(0)}{d} \lambda_{\text{ion}} \), where \( X \) and \( B \) are the excitation rate coefficient and the branching ratio, respectively.

Turning now to the ion line emission profiles, we present in Fig. 3 profiles for BII, BIII and BIV. The BI profile is also shown for comparison. As an aid to the discussion of the profiles we give in Table 2 some typical times and lengths characterizing the various collision processes in the divertor. Boron ions entering the scrape-off layer (SOL) from the main plasma will mainly be in the BV (He-like) and BVI (fully ionized) stages. At typical SOL-temperatures

<table>
<thead>
<tr>
<th>Species</th>
<th>( \tau_{\text{eq}} ) [( \mu s )]</th>
<th>( \tau_{\text{dwell}} ) [( \mu s )]</th>
<th>( \tau_{\text{ion}} ) [( \mu s )]</th>
<th>( \tau_{\text{rec}} ) [ms]</th>
<th>( \lambda_{\text{mfp}} ) [m]</th>
<th>( \lambda_{\text{ion}} ) [m]</th>
</tr>
</thead>
<tbody>
<tr>
<td>BI</td>
<td>---</td>
<td>2.9</td>
<td>---</td>
<td>---</td>
<td>---</td>
<td>0.027</td>
</tr>
<tr>
<td>BII</td>
<td>34</td>
<td>180</td>
<td>33</td>
<td>56</td>
<td>0.54</td>
<td>0.53</td>
</tr>
<tr>
<td>BIII</td>
<td>9.1</td>
<td>180</td>
<td>330</td>
<td>77</td>
<td>0.15</td>
<td>5.3</td>
</tr>
<tr>
<td>BIV</td>
<td>4.2</td>
<td>180</td>
<td>&gt; 10^6</td>
<td>29</td>
<td>0.067</td>
<td>1e4</td>
</tr>
<tr>
<td>BV</td>
<td>2.4 (33)</td>
<td>180 (72)</td>
<td>&gt;&gt; 10^6 (1.5e4)</td>
<td>50 (150)</td>
<td>0.038 (1.38)</td>
<td>--- (630)</td>
</tr>
<tr>
<td>BVI</td>
<td>1.6 (22)</td>
<td>180 (72)</td>
<td>--- (---)</td>
<td>45 (200)</td>
<td>0.026 (0.92)</td>
<td>--- (---)</td>
</tr>
</tbody>
</table>

Tab. 2 Estimations of characteristic times \( \tau \) and lengths \( \lambda \) in the divertor based on \( T_e = 15 \) eV above the plate and \( T_e = 100 \) eV in the SOL and \( n_e = 1 \cdot 10^{16} \) m\(^{-3}\). An X-point to target distance of 3 m along the field lines is assumed. The times are the thermal equilibration time \( \tau_{\text{eq}} \) with the background ions, the collision-free dwell time \( \tau_{\text{dwell}} \) in the divertor, the ionization time \( \tau_{\text{ion}} \) and the recombination time \( \tau_{\text{rec}} \). The lengths are the collision mean free path \( \lambda_{\text{mfp}} \) and the ionization length \( \lambda_{\text{ion}} \). Values in brackets are for ions in the SOL with \( T_i = 100 \) eV.

Tab. 1: Integrated particle flux \( \Gamma \) of neutral atoms emanating from the divertor plate.

<table>
<thead>
<tr>
<th>Species</th>
<th>( E_{\text{ion}} ) [eV]</th>
<th>( \lambda ) [nm]</th>
<th>( E_n ) [eV]</th>
<th>Shot # (filling gas)</th>
<th>( \Gamma ) [10^20 m(^{-2}) s(^{-1})]</th>
</tr>
</thead>
<tbody>
<tr>
<td>D</td>
<td>13.6</td>
<td>656.1</td>
<td>12.1</td>
<td>2635 (D)</td>
<td>50</td>
</tr>
<tr>
<td>H</td>
<td>13.6</td>
<td>656.3</td>
<td>12.1</td>
<td>2635 (D)</td>
<td>20</td>
</tr>
<tr>
<td>Cal</td>
<td>6.1</td>
<td>422.7</td>
<td>2.9</td>
<td>2389 (D)</td>
<td>0.003</td>
</tr>
<tr>
<td>HeI</td>
<td>24.5</td>
<td>667.8</td>
<td>23.1</td>
<td>1406 (He)</td>
<td>65</td>
</tr>
</tbody>
</table>

(Integration energy \( E_{\text{ion}} \), excitation energy \( E_n \) and spectral line wavelength \( \lambda \)).
$T_e = T_i = 100 \text{ eV}$ the equilibration times with the background ions $\tau_{\text{equ}}$ are much shorter than the collision-free dwell time in the SOL ($\tau_{\text{dwell}}$) estimated assuming $v_{\text{par}} = v_{\text{th}}$. These ions will therefore thermalize to the deuteron temperature in the SOL and divertor region and be entrained in their flow. They will, nevertheless, flow rapidly to the target in a time much shorter than the recombination time $\tau_{\text{rec}} \approx 1 \text{ s}$ (radiation and dielectronic recombination). The observed lower ionization stages are therefore not populated by recombination but are produced by ionization of neutrals produced at the target.

Sputtered B atoms have a typical energy of $5 \text{ eV}$ which at $T_e = 15 \text{ eV}$ and $n_e = 1 \cdot 10^{19} \text{ m}^{-3}$ gives $\lambda_{\text{ion}} = 2.7 \text{ cm}$. The observed BI profile exhibits an exponential decay of similar decay length but peaks a few mm from the plate. This may be explained by dissociation of e.g. BD molecules providing an additional source of B atoms a few mm from the plate.

As a further point of interest we learn from Table 2 that, for the singly ionized BII ions, the ionization and equilibration times are about equal and much shorter than the dwell time in the divertor. These particles can therefore only marginally be described by a fluid approximation. The relatively large mean free path length of 1.5 m causes a broadening of the emission region so that these particles are also visible close to the plate. In contrast, BIII and the higher ionized species are collision dominated and will flow together with the background plasma but with a superimposed parallel and perpendicular diffusion. According to Table 2 the sputtered boron atoms will be rapidly ionized to the He-like state (BIV). Excitation of the BIV-levels or ionization to BV by electron collisions, however, is practically impossible for energetic reasons ($T_e = 15 \text{ eV}, E_n > 200 \text{ eV}, E_{\text{ion}} = 259 \text{ eV}$). The observed emission of BIV shown in Fig. 3 is therefore most remarkable. As a tentative interpretation we assume population of the corresponding $2p\;^3\text{P}$-level by charge exchange recombination of BV with D$^0$ atoms. This process has a rather high cross-section ($\sigma \approx 1 - 3 \cdot 10^{-15} \text{ cm}^2$) with preferential capture into the $n = 3$ level. Further support for this explanation is provided by the increase of the emission near to the plate which probably reflects the D$^0$ density. We should note that the existence of such BV (and BVI) ions in the divertor region is compatible with Table 2. These particles are produced in the main plasma and — after diffusing into the SOL — will reach the target plate without recombining. Generally, confirmation of the presence of H-like and He-like ions in the divertor is problematical since at the typically low values of $T_e$ these ions are not excited by electron collisions.

**Summary**

We presented first measurements of impurity emission profiles in a divertor. These measurements will be used to check the advanced divertor modelling codes. One important fact is the observation of He-like boron in the whole divertor plasma which confirms the existence of such highly ionized particles due to low recombination probability.

**References**

1. Mertens V., et al., "Experimental investigation and interpretation of MARFE's and density limit in ASDEX Upgrade", same conference
5. Herrmann A., et al., "Asymmetric energy flux to the ASDEX Upgrade divertor plates determined by thermography and calorimetry", same conference
6. Schweinzer J., private communication
Figures

Fig. 1 Schematic diagram of the divertor spectrometer showing lines of sight and elevation.

Fig. 2 Profile of Ca I intensity versus elevation above the outer divertor plate

Fig. 3 Profiles of BI, BII, BIII and BIV intensities versus elevation above the outer divertor plate
Experimental and theoretical investigation of density and potential fluctuations in the scrape-off layer of ASDEX

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Introduction
Electrostatic fluctuations (i.e. the magnetic field is assumed constant) are candidates for the explanation of the anomalous transport of particles and energy in both tokamaks and stellarators. While most theoretical effort has been directed to an explanation of the anomalous transport in the bulk plasma, it is now widely being realized that the anomalous radial transport in the scrape-off layer, determining the width of the power flow channel at limiter or divertor plates, may be equally important to a future reactor experiment.

In the divertor tokamak ASDEX density and potential fluctuations in the scrape-off layer were investigated with high temporal and spatial resolution by Langmuir probes and an HDI diagnostic. Many results of these measurements were reported in [1] and are summarized below. Several of these properties of the fluctuations have also been reported from other experiments.

Basic properties of the observed fluctuations

- The fluctuations can be described by independent individual “events” superimposed randomly rather than by modes following a dispersion relation.
- In poloidal direction, these events show a periodicity with a wavelength of typically 4–6 cm, but the poloidal correlation decays within 2–3 wavelengths (the poloidal half width is about as large as the wavelength).
- Parallel to the magnetic field the fluctuations are in phase and highly correlated over a distance of at least 10 m.
- The typical lifetime of the largest fluctuation events is of the order of a few 10 μs.
- On the high field side, such fluctuations are only observed in “single null” discharges when there is a connection between inboard scrape-off layer and outboard scrape-off layer, whereas on the low field side the fluctuations are present also in “double null” discharges.
- Outside the separatrix the fluctuations propagate into ion diamagnetic drift direction with velocities of up to a few 100 m s⁻¹, but at low temperatures the velocity can also be very close to 0 or even reverse its sign; inside the separatrix the fluctuations always propagate into electron diamagnetic drift direction.
- The phase between the fluctuations of ion saturation current and of floating potential yields maximum transport radially outward and suggests the picture of eddy-like events exchanging plasma between the inner and outer parts of the scrape-off layer; the lifetime of an individual “event” just fits to allow half a turn with the radial E×B drift velocity and radial width of the scrape-off layer measured.
The radial transport derived from the fluctuation measurements is of the right order of magnitude to explain the global particle confinement.

Modelling the scrape-off layer

In addition to the gradients of density and temperature and the curvature of the magnetic field, in the scrape-off layer the boundary conditions at the intersection of the magnetic flux tubes with the limiter or divertor plates have to be taken into account. At these intersection points electrostatic sheaths are formed and the flows of ions, electrons and energy to the target plate are connected to the density and temperature in front of the sheath and to the potential drop in the sheath. In [2] it was shown for a fluid model that under these conditions the scrape-off layer will be unstable in the region of unfavourable magnetic curvature, if potential fluctuations are regarded in connection with either density or temperature fluctuations. Our analysis takes into account simultaneous fluctuations of all three quantities in a simplified cylindrical geometry. We assume that fluctuations are in phase along each magnetic flux tube, which is supported by the correlation measurements mentioned above and is consistent with the physical picture to be presented below. Thus, the time-dependent (fluctuating) part of the equations can be treated two-dimensionally, and a non-linear numerical analysis would be desirable, but has not been performed yet.

We performed a linear stability analysis and arrived at the following results within the limits of our model:

- The scrape-off layer is stable against perturbations
  - of sufficiently large poloidal wavelength \( \lambda_{\text{pol}} \)
  - or for slab geometry (no magnetic curvature) and homogeneous magnetic field.
- For unfavourable magnetic curvature (low field side in a tokamak) a stability limit \( \lambda_0 \) for \( \lambda_{\text{pol}} \) exists, and perturbations with wavelengths \( < \lambda_0 \) will grow with a rate \( \propto 1/\lambda_{\text{pol}}^{2} \) for small \( \lambda_{\text{pol}} \).
- The instability mechanism in this case can be illustrated by assuming a fluctuation with poloidal pressure gradients: The resulting radial diamagnetic currents have a non-vanishing divergence due to the cylindrical geometry. They have to be balanced by parallel currents through the sheath forcing changes in the electric potential of the magnetic flux tube and thus leading to poloidal electric fields. The resulting radial \( E \times B \) drift leads to a pressure perturbation in phase with the original perturbation, and thus instability results.
- In the region of favourable magnetic curvature the scrape-off layer is stable in the limit of small and of large wavelengths, but if the density gradient is sufficiently large compared to the pressure gradient, an intermediate range of unstable wavelengths can exist. This result is due to the inclusion of both density and temperature fluctuations in the model.

If we use mixing length arguments, trying to describe the convective transport due to the fluctuations by diffusion, a diffusion coefficient results which itself is proportional
to the pressure gradient. We use this diffusion coefficient to add an anomalous radial transport term to the stationary version of our model equations and thus get estimates for a diffusion coefficient and decay lengths for density and pressure. The diffusion coefficient is \( \sim 1 \text{ m}^2\text{s}^{-1} \cdot \left( \frac{T_0}{10^3 \text{ eV}} \right)^{7/6} \left( \frac{L_m}{1 \text{ m}} \right)^{1/3} \left( \frac{B}{1 \text{ T}} \right)^{-4/3} \left( \frac{R}{1 \text{ m}} \right)^{-2/3} \) where \( T_0 \) is the mean electron temperature, \( L_m \) is the connection length between the target plates, \( B \) is the magnetic field and \( R \) is the major radius. The decay lengths for density \( L_n \) and pressure \( L_p \) are in the order of 1–3 cm and scale as \( T_0^{1/3} L_m^{2/3} B^{-2/3} R^{-1/3} \).

These scalings are only valid, however, if the secondary electron emission coefficient at the sheath and the mean net current through the sheath are constant. But there are indications from Langmuir probe measurements in the divertor that the net current through the sheath can vary strongly with different plasma conditions, and for this case our model predicts opposite behaviour of \( L_n \) and \( L_p \) which in the divertor scrape-off layer indeed has been observed.

### Comparison between model and experiments

For density ramp discharges at ASDEX, several model results are plotted versus parameters of the correlation functions in figs. 1–3. As the model is based on several crude simplifications up to now, the qualitative agreement is quite satisfactory. The results of the model are not explicitly dependent on density but on temperature, and the temperature in the scrape-off layer scales inversely with density. Fig. 4 indicates that the typical wavelengths of the fluctuations are closely correlated to the radial width of the energy-carrying layer.

[3] gives a dependence of \( L_n \propto T_0^{-0.23} L_m^{0.45} \) from lithium beam experiments for a set of comparable discharges, which is not so far away from \( L_n \propto T_0^{1/3} L_m^{2/3} \) given above (without taking into account variations of the secondary electron emission coefficient and the net current through the sheath), but no dependence of \( L_n \) on the magnetic field could be detected in those experiments.


FIG. 1  Typical growth time from our model versus correlation time from the temporal and spatial correlation functions of the fluctuations in the scrape-off layer as measured by the H_a diagnostic. The data are from three ASDEX discharges with density ramps (different symbols for each discharge).

FIG. 2  Poloidal phase velocities of the fluctuations from our model versus the measured values from the correlation functions. Same data as in fig. 1.

FIG. 3  Density decay lengths from mixing lengths estimates in our model versus the values given by the scalings in [3]. The straight line is fitted through the origin. The qualitative behaviour and the order of magnitude agree. Details may be affected by the secondary electron emission coefficient and the mean net current to the target plates.

FIG. 4  Experimental wavelengths of the fluctuations as determined from the temporal and spatial correlation functions versus equivalent width of the power-carrying layer in the divertor as given by the scaling in [3]. The straight line is fitted through the origin. The data are from the same discharges as in fig. 1.


Particle and Energy Transport in the ASDEX Scrape-Off Layer

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Introduction: The increasing maturity in edge/SOL modeling approaches within the IPP, and the crystallization of clear trends within edge data bases for ASDEX now begin to permit a more systematic exploration of the underlying physics behind edge phenomena observed in the past. This paper first documents the parametrical behavior of a number of edge/SOL/divertor quantities for ohmically- and NI-heated plasmas. Then, the results of parametric calculations with the B2 code /1/ in conjunction with a simple gas recycling model are compared to experiment. Hereby, the goal is to explore the systematic trends predicted for experimentally-measured quantities with the aims of (a) discerning the transport model most appropriate to the data, (b) elucidating the interdependencies of salient code output quantities, and (c) delineation of the direction of future code work of this nature.

Experimental Scallings: In the edge plasma of the main chamber, temperature (T₉39.4) and density (n₉39.4) values are available from the YAG Thomson scattering system at a point approximately one cm inside the magnetically-determined separatrix. Regressions involving Tₑ and nₑ are based on this point in order to minimize scatter of the data. Extrapolations of n₉39.4 and T₉39.4 to the approximate separatrix position yields values about 2/3 of those at the point of measurement. The lithium beam probe delivers relative nₑ(r) profiles in the outer midplane. A reciprocating Langmuir probe in the divertor yields nₑ(r), Tₑ(r) and equivalent power flux qₓ profiles widths Δₚ. Interferometric line densities nₑL near the divertor throat as well as the divertor neutral gas density n₀div and H₀ intensity are studied.

Table I presents regression fits for the quantities above from an (OH,Nl) experimental data base in terms of those parameters expected to be relevant for edge/SOL physics. The "machine parameter" B₉ must also be used to obtain reasonable results. Of note is: (a) qₐ and B₉ often are of major importance in the fit, and (b) among edge quantities there is no isotope effect.

In the outer midplane λₑ is derived by fitting an exponential to the density profile over the first two cm outside the separatrix position Rₛ. λₑ is typically given by λₑ=qₑ0.45-0.8 T₉39.40.23-0.41, λₑ≈1.6-2.6cm /2,3/. The smaller qa exponent is more representative. However for Tₑ₉39.4<40eV, corresponding to high-qa ohmic discharges, λₑ abruptly changes its behavior: λₑ(Tₑ₉39.4<40eV)=qₑ0.79 T₉39.4₀.51. Phenomenologically, the density shoulder seen /2,3,4/ for R-Rₛ≈2cm in normal regions of operation becomes characteristic of the entire SOL for low Tₑ and high qa. Δₑ in the divertor scales as Δₑ=1.05±0.013 qₑ0.53±0.025 (R=0.9) for an ohmic deuterium data base (Δₑ=1-4 cm). The general trend with decreasing Tₑₖ as the density limit is approached is that the Tₑₖ profile becomes very flat in the range 5-8eV, whereas nₑ(r) maintains a somewhat peaked profile near the separatrix. A shoulder in the qₓ profiles is also seen, which dominates Δₑ for low Tₑₖ and high qa /5/.

In previous work /2,3/ an initial effort to deduce the perpendicular diffusion coefficient D in the SOL from the formulation "D"=λₑ²/τₓ, under the assumption of a constant mach number led to the result "D" = T₉₃9.4 for a wide variety of conditions for ohmically- and NI-heated plasmas. This result is summarized for an OH data base in fig.2.

Braams Code Results: The "D" = T₉₃9.4 relation of above is predicated on a primitive model and should be checked against complete B2-EIRENE /1/ calculations. However, the approach discussed here employs a simplified impurity and neutral gas model, with a single-fluid B2 version, in order to gain a feeling for leading terms. Heat transport along field lines is taken to be classical with a flux limit of 0.2evTe/λₑ. For perpendicular transport, three models are examined: (A) D=χₑ=χ₉ₑ/3=0.5,1,1.5m²/s; constant over the SOL, (B) D=χₑ=χ₉ₑ/3=
$D_0 10^{13}$ cm$^{-3}$/n$_e$ (i.e. $D = 1/n_e$); $D_0 = 0.5, 1$ m$^2$/s, and (C) $D = 1$ m$^2$/s, $\chi_1 = \chi_2 = 0.01$ (i.e. energy transport mainly by convection). The code input parameters are varied over typical operational ranges; the results are summarized in Table II in the form of regressions for the quantities listed in Table I, whereby the B2 input quantities $n_{es}$, $P_{Sol}$, $q_a$, impurity concentration $n_{imp}/n_e$ and $D_0$ are used as the regression parameters. This statistical approach has the advantage that the validity of individual transport models, as well as inadequacies of other model assumptions, can readily be examined.

Classical drifts or an inward pinch are not considered. Only one equilibrium grid is used with a spatial extent of 1.8 cm inside $R_S$ and 2.6 cm outside $R_S$ at the outer midplane. The code operates with feedback loops so that $n_{es}$ and $P_{Sol}$ can be specified; power is divided between electrons in the ratio 2:1. $Z_{eff} = 3$ is taken for all calculations. The neutral gas model assumes the $n_{odiv}$ profile - motivated by Degas/Eirene calculations - is invariant in the divertor (reasonable for the $n_{es}$ values encountered in the plasma fan, see Table I). The absolute value of $n_{odiv}$ is calculated from the divertor particle balance. $n_{es}$ is regulated by balancing all ion losses from the SOL and divertor by a poloidally homogeneous gas puff in the main chamber.$6$. Other parameters in the code are selected to give close agreement with complete B2-EIRENE runs for a generic situation.

Fig. 1 illustrates B2-generated $n_{es}$, $T_e$ profiles for $n_{es} = 1 \times 10^{13}$ cm$^{-3}$ and $P_{Sol} = 0.6$ MW. For $R - R_S = 2$ cm the differences in the three diffusive models are minimal. The $1/n_e$ variation has the pleasing aspect that a shoulder is produced for $R - R_S = 2$ cm, as seen in experiment. But, the shoulder is always there, not scaling in the proper way. For $T_e(r)$ the conv. case is characterized by rather high $T_e$ at the separatrix with a very steep falloff in the SOL.

Instead of a point-by-point comparison between experimental and code regressions, the overall trends will be reviewed. The inclusion of $D$ in Table II as regression parameter demonstrates the sensitivity of changes in D to the quantity under question. Quite generally, the experimental $q_a$-scalings are not duplicated in the calculations. There is good agreement for $H_a$ in all cases. The code-predicted $P_{Sol}$ dependence of $n_{odiv}$ is absolutely not found in experiment. The density dependence for $A_D$ and $\lambda_n$ predicted by case B is in strong contrast to experiment. Case C gives an extremely strong relationship between $T_{es}$ and $n_{es}$, $P_{Sol}$ - again in strong contrast to experiment. Case C does have the correct $T_{ed}$ exponent for $A_D$. However, as mentioned above, the $q_{\|}$ profile for low $T_{ed}$ develops a prominant shoulder which surely strongly contributes to the $T_{ed}$-$A_D$ dependence seen in experiment, and none of the models produce a shoulder in the correct sense. From these considerations, case A is the preferred form of diffusion coefficient simply from the standpoint that it nowhere strongly violates experimental behavior as do B and C for certain quantities.

For case A, the fit for $\lambda_n$ is obtained only if values $n_{es} > 0.8$ are excluded from the data base. For $n_{es}$ below roughly this limit (i.e. in the direction of higher $T_{es}$, $T_{ed}$) $\lambda_n$ begins to decrease with temperature. This behavior is correlated with the Mach number: centered around $T_{es} = 35$ eV, $M$ augments rapidly with increasing or decreasing $T_{es}$ such that, heuristically, $\lambda_n \approx (D_0 L/MC_S)^{0.5} \approx (D_0 L/MT_e)^{0.5}$ would be expected to diminish, which is indeed observed in the calculation. Thus, the experimental scaling of $\lambda_n^{-T_{esg,4}}$ comes about not to the diffusion coefficient scaling as $T_{esg,4}$, but due to the variation in $M$ with $T_{es}$. Fig. 2 illustrates this point, where $D = \lambda_n^{2/r_{\|}}$ is used to derive a "diffusion coefficient" for case A (with $D_0 = 0.5, 1.5$ m$^2$/s) in the same way as the experimental data base. For each $D_0$ a set of points arises which behave very much as the experimental data set in both a relative and absolute sense. Going further, keeping in mind the very rough $T_{esg,4}^{-1.5}T_{es}$ relationship, the experimental set is compatible with $D_0 = 0.5-1$ at high $T_{es}$, and around $T_{esg,4} = 40$ eV $D_0 = 1.5$ is more appropriate.

Note that in fig. 3, the clean variation of "$D$" with $T_{es}$ comes about only because $q_a$ has been held constant, and the data set restricted to $n_{es} > 0.8$. The point scatter arising when these limitations are lifted is not compatible with experiment, where with $q_a = 2.5-5.9$ and high temperatures, the scaling for "$D" continues. Further code work is necessary here.

Overall, in terms of absolute numbers, the B2 results are satisfactorily in agreement with experiment with respect to $T_{es}$ and $n_{es} T_{es}$, keeping in mind that $T_{esg,4}$, $n_{esg,4}$ when extrapolated to the separatrix are roughly 2/3 of the values given in table I, and that no
attempt was made to model any particular shot in detail. $T_{ed}$ is predicted to be too high, as $T_{ed}>35$eV is seldom found in experiment. On the other hand, for such high temperatures the impurity content of the plasma will increase, which would exacerbate excessive $T_{ed}$ values if included in the model. $n_{eL}$, $n_{ediv}$ and $n_{ed}$ are all about a factor of two too low; increase of recycling in the model will help to alleviate this inconsistency.

Conclusions and Discussion: Case A, $D=\chi_1=\chi_2/3$, is most often closer to experiment, including the fit to $\lambda_n$. A $D^{-1/n_e}$ scaling is of interest, since indications on the stellarator W7-AS are that such behavior prevails/7/. Nonetheless, the $n_{es}$ scalings for $\lambda_n$ and $\Delta_p$ implicit for the model are simply not present. For example, during a density ramp, $n_{es13}$ might vary from 0.5 to 2.0; the $D(1/n_{es})$scaling $\lambda_n-n_{es}^{-0.51}$ requires that $\lambda_n$ decrease by a factor of two, i.e. from two to one cm. This is far outside measurement error. The tendency of the convective model is for $T_{es}$ to respond too sharply to changes in $n_{es}$ and $P_{sol}$.

The strong $q_\theta$, $B_1$, $I_0$ variations of some experimental observables might be an indication that transport is related to these global parameters. It is interesting, for instance, to find for cases where a pure $q_\theta$-scaling is present (i.e. not in combination with $B_1$), that the assumption $D-q_\theta$ leads immediately to the correct $q_\theta$ dependence for $\Delta_p$ and $\lambda_n$. Such a scaling will also presumably allow the code to more correctly duplicate the $q_\theta$-scaling of the density limit - without this caveat the density limit is largely unaffected by $q_\theta$ variations.

The code results can identify combinations of experimental parameters suitable to extract $D$. For case A, $D=\lambda_n^{1.99\pm0.076}T_{es}^{-0.61\pm0.35}$; that is, $\lambda_n^{2}$ is directly proportional to $D$ as expected from simple 1D models, but the exponent of $T_{es}$ is exactly opposite to that normally assumed. This arises from the variation of mach number, both with $T_{es}$ and $q_\theta$.

Finally, the presence of a shoulder in $n_{e}(r)$ and $q_\theta (r)$ cannot be explained within the code, other than to assume a diffusion coefficient in excess of Bohm in the low-temperature region where the shoulder occurs. If transport is effected by ExB-driven turbulent eddies, of the size of $\lambda_n$ and with correlation times of the order of parallel equilibration times-for which some evidence exists/8/, then during one rotation the plasma may drift outwards such that the effective value of the anomalous diffusion coefficient may exceed Bohm. One expects, then, $B_1$ to be correlated with quantities associated with the SOL. The regressions of table I, where $q_\theta$ and $B_1$ have been used as principle components, indicate core/edge quantities are roughly proportional to $B_1$ whereas divertor parameters are inversely proportional.

<table>
<thead>
<tr>
<th>const. $n_{e39.4}$</th>
<th>$P_{sol}$</th>
<th>$q_\theta$</th>
<th>$B_1$</th>
<th>$Z_{eff^{-1}}$</th>
<th>$A$</th>
<th>$R$</th>
<th>RANGE</th>
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<tr>
<td>$T_{e 39.4}$</td>
<td>92</td>
<td>-0.52±0.03</td>
<td>0.48±0.01</td>
<td>-0.61±0.02</td>
<td>0.67±0.05</td>
<td>0</td>
<td>0</td>
</tr>
<tr>
<td>$n_{e39.4}$</td>
<td>0.29</td>
<td>0.94±0.03*</td>
<td>0.24±0.01</td>
<td>0.14±0.01</td>
<td>0</td>
<td>-0.15±0.01</td>
<td>-0.14±0.01</td>
</tr>
<tr>
<td>$(n_{e T e 39.4})$</td>
<td>95</td>
<td>-</td>
<td>0.70±0.01</td>
<td>-0.51±0.03</td>
<td>0.78±0.06</td>
<td>-0.25±0.02</td>
<td>0</td>
</tr>
<tr>
<td>$n_{ediv}$</td>
<td>0.96</td>
<td>1.81±0.08</td>
<td>0.24±0.04</td>
<td>1.64±0.05</td>
<td>-0.85±0.08</td>
<td>0.16±0.05</td>
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</tr>
<tr>
<td>$H_{ediv}$</td>
<td>a.u.</td>
<td>0.78±0.04</td>
<td>0.77±0.02</td>
<td>0.21±0.04</td>
<td>-0.79±0.07</td>
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<td>0</td>
</tr>
<tr>
<td>$q_{ediv}$</td>
<td>14</td>
<td>1.53±0.05</td>
<td>0</td>
<td>0.8±0.06</td>
<td>-1.08±0.10</td>
<td>0.23±0.03</td>
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<tr>
<td>$n_{ed}$</td>
<td>0.70</td>
<td>1.67±0.05</td>
<td>0</td>
<td>0.8±0.06</td>
<td>-1.08±0.10</td>
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<td>$T_{ed}$</td>
<td>17</td>
<td>-0.84±0.06</td>
<td>0.58±0.04</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>0.97</td>
</tr>
<tr>
<td>$n_{e T e 39.4}$</td>
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<td>0.82±0.08</td>
<td>0.56±0.06</td>
<td>0</td>
<td>0</td>
<td>0</td>
<td>0.97</td>
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Table I: Regression fits for experimental quantities. The star indicates the fit with $n_{e}$/, not $n_{e39.4}$. See the subtitle of table II for more details. The div. data contains only an $n_{e^{-1}}$, $P_{NI}$-scan.

References:

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Fig. 1 Above: $n_e$ and $T_\theta$ profiles for the three diffusion models. $D_\perp =$ (A) const., (B) $1/n$, (C) convection only. $n_{e813}=1.0$, $P_{\text{Sol}}=0.86$ MW.

Fig. 2 Right: $\lambda_n^{2/\tau_{\parallel}}$ (for mach no. $=0.3$) vs. $T_\theta$ for an experimental OH data base (top), and the $D_\parallel=$const. part of the B2 data base (bottom).

| Table II: Fits for the B2 data base, $T_\theta = T_{\text{ed}}$ for the $\Delta p$ regression, and $T_\theta = T_{\text{es}}$ for the $\lambda_n$ regression. In both tables, exponents are given for fits of the form: const. $\times n_e^a \times P_{\text{Sol}}^b \times \lambda_n^c \times n_{e839.4}[10^{13} \text{ cm}^{-3}] \times P_{\text{Sol}}^{[\text{MW}]} \times B_\parallel[T] \times A[\text{amu}] \times D[\text{m}^2/\text{s}] \times T_\theta[\text{eV}] \times n_{\text{es}}^{[\text{div}]}[10^{13} \text{ cm}^{-3}]$; $\Delta p$, $\lambda_n[\text{cm}]$. R is the regression coefficient. Range = the min. and max values of the parameter range. |
Gaseous Impurity Production in ASDEX UPGRADE Discharges.

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Introduction.
Similar to the arrangement on former ASDEX, a Quadrupole mass analyzer (QMA) is attached downstream near to one of the fourteen turbomolecular pumps on ASDEX UPGRADE. Owing to about twice the distance between the QMA and the torus volume as compared to ASDEX, the response time (1) has increased to 120 ms for deuterium, scaling with the square root of mass. As it turned out, too, the magnetic shielding which was completely sufficient for ASDEX had to be improved by an additional encasement of 16 mm Armco steel. Calibration of the QMA is done under regular pumping conditions against the pressure in the torus vessel as measured by a MKS Baratron. Thus, given partial pressures always refer to an equivalent upstream vessel pressure and not to the actual pressure at the QMA which is lower by a factor of three or more because of the pressure drop in the pumping port.

Early Results.
The QMA was in operation from the very first discharges, starting in April 91. At this time all experiments were performed in He, since experience at ASDEX and other plasma machines has shown that problems with density control and reproducibility as well as with low density limits due to high gaseous impurity recycling are mitigated. Figure 1 shows the partial pressures of M=4, M=28 and M=44, during and after one of these early discharges. At this time still circular and non diverted plasmas at Ip=350 kA were run, mostly with the inner heat shield acting as limiter. Here the high levels of CO and CO2 were typical, especially well seen in the outgassing after the discharge. The partial pressure of CO2 reaches its maximum much later than CO, which cannot be explained by the differences in pumping speed. Obviously, outgassing characteristics are different. This behaviour together with the high contribution of CO2 seems to be rather typical for limiter tokamak discharges. In our experience it is associated with strong local heat-up of limiting components. CO, and even more so, CO2 are diminished as soon as discharges are better position controlled. CO2 is usually insignificant in good, diverted discharges, which is also plausible in view of the now generally accepted mechanism of main oxygen recycling via CO due to a chemical reaction between carbon and an oxygen atom stemming from the discharge zone. The formation of CO2 on the same line is obviously less likely, at least at low temperature wall conditions. On the other side, outgassing of CO and CO2 from hot wall components is well known and described in terms of thermodynamics. In contrast to the behaviour of CO in Fig. 1, strong wall pumping of CO is usually observed after divertor discharges, especially if parts of the wall are metallic (2,3).

The contribution of water vapour in the impurity recycling in ASDEX UPGRADE was and is less significant, though it might well be part of the initial source of oxygen at the start of the discharge. This latter observation is rather typical for plasma machines with mostly carbon covered walls. The situation was definitely different in former ASDEX with stainless steel walls (3).
The Effect of Boronization.
Unfortunately, the data base for a comparison of gaseous impurity production before and after boronization of ASDEX UPGRADE is somewhat scarce. This comes about as most of the discharges before boronization, including divertor discharges, were performed in He and only very few in D$_2$ are available. On the other hand, the effect of boronization, as documented on other machines as well, is so significant that better statistics is not necessarily the point. In He discharges hydrocarbons are only produced to the extent as hydrogen from the walls is outgassing into the discharge, mostly originating from previous glow discharge conditioning or boronization with B$_2$H$_6$. Thus, interference problems in the mass spectra, as explained in the next chapter, are minor and the evaluation for H$_2$O, CO and CO$_2$ is straightforward. He discharges before boronization show about half the amount of CO as those in D$_2$. This and the absence of the hydrocarbons might be partly responsible for the better performance and higher density limit achieved in He. Boronization reduces CO by about an order of magnitude with CO$_2$ and H$_2$O almost disappearing in the noise level. A similar finding applies to D$_2$ discharges as far as all the O-compound molecules are concerned. However, the reduction of hydrocarbons by boronization is much less pronounced. Here the data base is yet somewhat too insufficient, but the tendency corresponds well with the observations on ASDEX (4).

The Role of Hydrocarbons.
Most of what remains as gaseous impurity fluxes under boronized wall conditions are hydrocarbons. One problem in the interpretation of the mass spectra is that we usually have to deal with a mixture of deuterium and hydrogen in D$_2$ fuelled discharges, as already mentioned. For the hydrocarbons this leads to rather complicated spectra with some peaks strongly interfering with the mass peaks of water-vapour, CO and CO$_2$. In an earlier paper (5) pertaining to this problem on ASDEX, a model was used to describe the spectra for methane as a function of the ratio of H/D simultaneously found from the H$_2$, HD and D$_2$ mixture in the divertor exhaust. The model assumes nearly stationary conditions of atomic fluxes leaving the plasma and being adsorbed by the wall, and molecular fluxes leaving the wall by recombination and being again pumped by the plasma. This in effect affords conditions similar to thermodynamic equilibrium. The model which rendered very good agreement with the observed spectra of methane was recently extended to the C$_2$H$_n$D$_m$ group of hydrocarbons with their higher degree of complexity. This was especially indicated in view of a better assessment of the residual CO component after boronization.

Figure 2 shows a mass spectrum of the mass range 12 to 50 as taken during the plateau phase of a typical ohmic divertor discharge with I$_p$: 600 kA and <n$_e$>: $3 \times 10^{19}$ m$^{-3}$. The mass peaks pertaining to the C, C$_2$ and C$_3$ hydrocarbon groups are well discernible. A comparison with the model is shown in Figures 3 and 4 for the C and C$_2$ group respectively. The problem which arises with the latter is that ethane as well as ethylene and acetylene can be formed. An analysis of many spectra shows, however, that acetylene may be neglected and that best agreement between data and model is usually obtained with the assumption of 70% ethane and 30% ethylene. The peak patterns of Figures 3 and 4 are given for a concentration H/H+D of 0.1, 0.15, and 0.2. A comparison with Figure 2 shows best agreement around 0.15, which also is the ratio directly obtained from the measured partial pressures of H$_2$, HD and D$_2$. The only disagreement is seen on M=28 which is due to the contribution of CO which can be assessed now with fair accuracy to about 2/3 of the peak in Figure 2. Equally, 10% of water vapour of same isotope composition were
added to methane to improve the fit for M=18 to M=20. But this amount of water also constitutes a limit of what seems likely.

The characteristic pattern of the hydrocarbon peaks shows a very fast and pronounced change with the concentration. By comparing the measured pattern with those of the model, a quick and rather accurate estimate of the H/H+D ratio is possible which, for instance, is of interest for H-minority ICRH.

On the basis of Figure 2, one can derive the fluxes of the various hydrocarbons versus the flux of CO (CO₂ being obviously negligible). Adding up the peaks, one finds for methane, ethylene and ethane a contribution of 54%, 33% and 13% respectively, to the total hydrocarbon pressure of 6 *10⁻⁵ mbar which makes up about 20% of the hydrogen isotope pressure in the divertor exhaust. Concerning the recycling of these gases we have to convert these figures to the corresponding fluxes of carbon atoms which are in the ratio of 41%, 40% and 19% for CH₃Dₙ, C₂H₄Dₙ and C₃H₆Dₙ respectively. Compared with these total carbon fluxes, the reflux of oxygen, carried by CO, amounts to less than 10%. With the closed divertor of former ASDEX we had a reasonable model and a good data base to convert such data to absolute global fluxes entering the main plasma, which were found in good agreement with the radiated power and spectroscopical data. In ASDEX UPGRADE a direct connection of the divertor exhaust exists only to the zone of the outer divertor fan, and no data about effective conductances of this volume back to the main plasma are yet available. Still, the similarities of the findings make us believe that at least the relations between the fluxes are typical in a global sense. In view of the difficulties encountered in optical spectroscopy to quantify carbon versus oxygen, this constitutes an important statement in tying the main source of edge radiation then to carbon.

For a better understanding of the origin of hydrocarbons, a few more annotations should be made. Considering the conditions in the plasma fan in front of the collector plates it is not very likely that hydrocarbons there produced leave the plasma fan. Assumptions about redeposition go in the same direction. On the other hand, the high neutral gas pressure in the divertor volume leads to a strong parasitic recycling which results in a high flux of Franck-Condon and CX atoms hitting the surfaces in the vicinity of the divertor plasma. The reactions on these surfaces constitute the main source of the molecular hydrogen and hydrocarbons actually seen in the divertor exhaust. This interpretation is well in line with laboratory experiments which show that the observed high production rate of heavier hydrocarbons is typical for low energy bombardment of carbon surfaces as is obviously the case here (6).

One rather noteworthy aspect of the model for hydrocarbon generation is that the yields for the actual carbon erosion (chemical sputtering) in the machine can directly be derived from the measured fluxes of hydrocarbons versus the molecular hydrogen fluxes. The only approximation is that the molecular fluxes leaving the wall are balanced by the impinging atomic flux which is verified by the observed high recycling rate (low fuelling). Also, the fluxes have to originate from the same surfaces. The erosion yield obtained in this way is between 5% and 10% which is much higher than room temperature lab. data would suggest(6). Even less chemical erosion should be expected after boronization. Such discrepancies might be explained by actual wall conditions rather different from the well defined ones of a laboratory experiment and, last not least, also by the higher hydrogen fluxes encountered in plasma machines.
Figure 1. Partial pressures during and after limiter discharge #465.

Figure 2. Mass spectrum M: 12 to 50 during plateau phase of divertor discharge #2424.

Figure 3. Model calculation of the $\text{CH}_m\text{D}_n + 10\% \text{H}_m\text{D}_n\text{O}$ cracking pattern.

Figure 4. Model calculation of the cracking pattern $\text{C}_2\text{H}_m\text{D}_n$, 70%: $n+m=6$, 30%: $n+m=4$.

References:

(4) U. Schneider et al., J. Nucl. Mater. 176/177 (1990) 350
(5) W. Poschenrieder et al. J. Nucl. Mat. 176/177 (1990) 381
Tore Supra recycling properties with boronized walls

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I) Introduction

In a previous paper, the main improvements of the Tore Supra plasma operation after boronization were described: an oxygen plasma content decrease by a factor of 5 and strong effects on wall recycling properties [1]. These latter are the subject of this study.

In the first part of this paper, the hydrogen depletion after an helium glow discharge from boronized and carbonized walls is compared. The very slow decay of the hydrogen production can be explained by a much larger hydrognoïd atoms diffusion coefficient for boronized walls (BW) than for carbonized walls (CW).

The wall recycling properties during plasma shots for H and D atoms in BW and CW are then compared. The observations are in agreement with the assumed diffusion increase in BW.

II) Helium glow discharge analysis

After a boronization or a carbonization process, the Tore Supra inner walls are conditioned by means of an helium glow discharge (HeGD) (helium pressure: 0.3 Pa, and GD current: 3 A). The H and D production is recorded all along the conditioning procedure.

In the case of CW, the total amount of released atoms is $5.5 \times 10^{22}$ after 10 hours HeGD (wall temperature: 170 °C).

In the BW case, $7 \times 10^{22}$ atoms are detrapped which is about 20% more in spite of a lower wall temperature (160 °C).

Moreover, the time evolution of the H$_2$ production ($Q_{H_2}$) is recorded during all the HeGD. In the CW case, $Q_{H_2}$ follows $t^{0.5}$ evolution which is the signature of a limited diffusion process [2]. In BW conditions, $Q_{H_2}$ evolution is rather different since it follows a $t^{-0.15}$ law (see figure 1). In contrast with the carbonized case, this slow evolution seems to prove that, in BW, desorption is not diffusion limited. This implies that the diffusion coefficient is very high in the boronized layers, allowing a continuous hydrogen supply of the desorption layer from the bulk. It is worthwhile noting that penetration during HeGD is at the most 5 nm. $7*10^{22}$ atoms detrapped from a fully saturated wall correspond to about 90 fully desaturated monolayers which is equal to 22 nm of desaturated depth [3]. Consequently, this hydrognoïd production requires that the particles diffuse during HeGD in BW. Accordingly, in the frame of a diffusion process, a $t^{0.5}$ evolution should be observed for $Q_{H_2}$.

This slow evolution of $Q_{H_2}$ could be attributed to a catalytic effect of boron on the creation of H$_2$ molecules which can then diffuse very rapidly in an amorphous layer. However, when the HeGD is turned off, the H$_2$ production stops which is contradictory with the catalytic hypothesis.

In conclusion, to explain the observations, the particles must have a wall diffusion coefficient much greater in BW than in CW ($D_{D:B} > D_{D:C}$).
III) Plasma shot properties

For the forthcoming analysis, the wall temperature is kept constant at 170°C and the plasma shots which are compared, have the same time history with Ip=1.5MA without auxiliary heating.

a) Fuelling efficiency

The fuelling efficiency (FE) is defined as \( \frac{dNe}{dt} \), where Ne is the total number of electron in the plasma and \( \phi_{\text{inj}} \) is the injected gas influx. In figure 2, FE is represented in BW and in CW as a function of time.

In the CW conditions, FE decreases continuously from a value much greater than one to about 20% at the current plateau. At the beginning of the ramp up, the prefill gas after being trapped in the wall is rapidly recovered and FE exhibits a high value. The ionisation length being proportional to \( 1/ne \) where \( ne \) is the plasma density, FE appears to decrease in the same way as the ionisation length.

In the BW conditions, FE is always smaller than 25% even in the ramp up. It should be noted that before entering the plasma bulk, particles once ionized in the S.O.L may recycle many times on the walls. If, as supposed in part one, \( DD:B > > DD:C \), the D particles can diffuse more deeply in the bulk in BW. So, they are more easily trapped and FE is thus smaller in boronized conditions.

b) Vessel pressure evolution (after a shot)

In figure 3, time evolution of \( D_2 \) partial pressure after a shot is presented for BW and CW conditions. This different evolutions are always observed after a shot : in BW, \( P_{D2} = P_0t^{-0.8} \) and in CW \( P_{D2} = P_0t^{-1} \). This difference represents a specific characteristic of the wall conditioning status. At the end of the shot, the particle density profile in the walls is out of equilibrium and the particles can diffuse from the layer (corresponding to the maximum concentration of deposition) to the wall surface. Under the assumption that \( DD:B > > DD:C \), the desorption is larger in BW than in CW.

The total number of particles recovered in BW after a shot is much greater than in CW. In figure 4, the ratio of the total amount of pumped particles (Np) to the total amount of injected particles (Ni) is plotted versus the maximum plasma electron content (Ne_m) reached during a shot at the end of the gas puff. In BW, Np/Ni is always higher than in CW. For Ne_m smaller than \( 4*10^{20} \), Np/Ni is greater than 1 and the walls are desaturated even by deuterium shots. Moreover, high density are reached only in BW and the ratio Np/Ni is then equal to 25%. In this case, \( 6.4*10^{21} \) atoms are injected and \( 1.6*10^{21} \) atoms are pumped. It is noticeable that for all the shots presented in this figure, Np is always of the same order.

c) Particle balance

A new technique to study particle balance and to estimate the wall saturation status has been developped for Tore Supra shots.
This technique consists [4] of making a series of shots keeping the same procedure all along the series. First, the plasma is initiated on the outboard pump limiter (area/0.5m²) and during this phase a large and variable quantity of gas is puffed into the machine. Then, 3 seconds after the end of the gas puff, the plasma is pushed on the inner wall (more than 15 m² graphite). This displacement causes a decrease of the plasma density characterised by τ₀ which is the time scale of the density evolution. We have shown previously that, in CW conditions, τ₀ increases continuously until the total saturation of the wall is reached [4].

For every shot, Qdep (=Np-Ni) which is the total amount of particles deposited in the wall is measured. On figure 5, τ₀ (for BW and CW conditions) is plotted as a function of Qtot which is the integrated value of Qdep over the entire series of shots.

At the beginning of each serie, τ₀ is the same in BW and in CW indicating that the wall saturation status is equal in both conditions.

In the carbonized case, τ₀ increases rapidly from 60ms to 200ms : the wall is reaching its saturation. On the contrary, in the boronized case, τ₀ remains almost constant : the wall interacting with the plasma stays desaturated.

This is confirmed by (Hα, Dα) measurements recorded during the gas puff which are presented in figure 6 where shots labeled as I (CW) and J (BW) in figure 5 are compared. For similar Qtot, τHα which is the time scale of (Hα, Dα) evolution is longer in CW than in BW. This confirms that the wall is more saturated (for the same total D₂ load) in CW conditions than in BW ones.

IV) Conclusions

We have shown that the particle diffusion properties are very different for BW and CW conditions. To explain this discrepancy, we assume that the particle diffusion coefficient is much higher for BW than for CW. Moreover, the wall recycling properties are compared in BW and CW conditions. We have shown that : i) the fuelling efficiency is lower in BW than in CW, ii) the total amount of particles recovered after a shot is greater in BW than in CW, for any plasma density, iii) using a particle balance analysis, no net wall saturation is observed in BW compared with CW for similar deuterium loads. These observations confirm the presumed diffusion increase. This particular behaviour of the BW allows an efficient recycling control for very long Tore Supra discharges up to 67s where a wall saturation is not reached in spite of a very high injected gas puff (4.24*10²⁰ atoms/s) [5].

REFERENCES:
Figure 1: Time evolution of H2 partial pressure during HeGD.

Figure 2: Time evolution of fuelling efficiency (dashed line: CW, solid line: BW).

Figure 3: Time evolution of D2 partial pressure after a shot (dashed line: CW, solid line: BW).

Figure 4: Np/Ni = f(Ne max).

Figure 5: Time scale density decrease evolution versus the total amount of particle deposited in Tore Supra.

Figure 6: Time evolution of H4D4 measurements during gas puff (dashed line: CW, solid line: BW).
1. Introduction. The generation of impurities due to plasma interaction with graphite surfaces has been the focus of a great deal of study. We present results from an experiment in which the plasma-surface conditions can be relatively well characterized. Impurity spectroscopic observations have been made in the throat region of the Outboard Pump Limiter in Tore Supra. The temperature of the neutralizer plate can be estimated from measurements of the temperature of the front face of this limiter and extensive thermal modeling using the actual limiter geometry. The plasma conditions are characterized by Langmuir probe measurements in the pump limiter throat, approximately 10 cm from the emitting surface. The full visible region of impurity emission is observed spectroscopically. These observations lead to the identification of hydrocarbon generation from chemical sputtering processes. Modeling with a 3D impurity transport code (BBQ) allows a comparison with expected emission due to chemical sputtering. The experimental set-up is described in Section 2, the spectroscopic observations and identification is given in Section 3, the transport modeling and comparison is discussed in Section 4, and conclusions are given in Section 5.

2. Experimental Setup. The neutralizer plate (deflector) of the Tore Supra outboard pump limiter is viewed by means of a telescope located at the end of the pumping chamber and which is fiber-optically linked to a visible spectrometer with a 1-D Optical Multichannel Analyser (OMA). The linear dispersion of the system is ~ 0.1 Å/mm. The viewing chord intercepts with the deflector are ~6 cm in diameter, roughly spanning the toroidal extent of the plate. Absolute intensity calibration of the system is done by placing the telescope at the aperture of a blackbody radiation source. The pump limiter is also equipped with a large number of other diagnostics, including a set of Langmuir probes that measure the ion flux at the entrance of the pump limiter throat, as well as the T_e and n_e in the throat region (Fig. 1).

3. Experimental Results. Spectra containing both atomic lines and molecular bands are obtained in the throat region. Fig. 2(a) shows a spectrum containing the CD band at
4315Å together with the CII line at 4267Å. (Note that some of the rotational structure of the molecule can be resolved). Fig. 2(b) shows the molecular band of C₂ at 5161Å (from the Swan series) together with a series of CII lines. The CIII line at 4647Å has also been observed in this region. The presence of the molecules in the throat suggests that chemical sputtering is a dominant carbon impurity generation mechanism. This is further supported by estimates of the deflector surface temperature in the range of 450-550°C (near the maximum of chemical sputtering for 100-200eV ions), as obtained from the front face IR measurements and a heat conduction model [1]. The presence of the Langmuir probes, which can measure directly the deuteron flux into the throat, allows us to estimate the chemical sputtering yield.

For example, for a high density (nₑₚ=6.10¹⁹ m⁻²) case (TS 10447 and 10453), the throat Langmuir probe data give an average ion flux to the deflector of jD⁺ = 2•10¹⁹ ions/cm²/s, while the OMA measures an integrated intensity of I=2.5x10⁶ counts/s (cps) for the CD band. Using the calibration of the optical system for 4300Å, which is C = 1.8x10⁻⁸ cps/(photons/cm²/sr/s), we get a photon flux of (4π sr)/C = 1.7x10¹⁵ photons/cm²-s. Then using molecular data from [2], for the approximate plasma conditions in the throat (25eV) the ratio of CD₄ dissociations per CD band photon emitted is 100. Therefore, the photon flux above corresponds to a "molecular loss event flux" for CD₄ of 1.7·10¹⁷ losses/cm²/s. If every CD₄ chemically sputtered from the deflector is also lost near the deflector, then there are 1.7·10¹⁷ CD₄ molecules/cm²/s produced near the surface. From the impinging ion flux and the CD₄ emerging flux, we compute a sputtering yield Y = 0.009. This is consistent with extrapolations of existing studies of chemical sputtering on graphite (most studies are limited to ion fluxes of 10¹³ to 10¹⁷ ions/cm²/s) [3].

A similar calculation can be carried out for C₂Dₓ molecules. Taking for example another high density shot (TS 9132), where the C₂ molecular band and the C⁺ line were measured simultaneously, we have for the C₂ band intensity a plateau value of 9.5·10⁶ cps. For this wavelength region C=3.6·10⁻⁸ cps/(photons/cm²/sr). Therefore, the photon flux at the surface was 3.3·10¹⁴ photons/cm²/s. Data from [2] give 500±200 loss-events/photon, i.e. losses of C₂D₂ or C₂D₄ per C₂ photon emitted. This leads to a dissociation rate of 1.65·10¹⁷ molecules/cm²/s. Using an average ion flux of 1.6·10¹⁹ ions/cm²-s from the Langmuir probes, we get a sputtering yield of Y ≈ 0.01 molecules/ion or 0.02 C-atoms/ion, which is a bit higher than for CD₄ and consistent with the scaling (for ~200eV ions) in [3].

Finally, using data from [2], we check the consistency of these calculations by apply-
ing it to CII as well. For this same shot 9132, from the plateau intensity of the C+ line we compute, with the same procedure as above, a photon flux of $1.7 \times 10^{14}$ photons/cm$^2$/s. If we assumed that all the C+ comes from dissociation of CD$_4$, then [2] gives $7000 \pm 200$ losses/photon and we compute $2.4 \times 10^7$ losses/cm$^2$/s. Assuming that all the C+ comes from C$_2$D$_2$, [2] would give $6000 \pm 200$ losses/photon and we compute $2.0 \times 10^{17}$ losses/cm$^2$/s. This is to be compared with the $1.7 \times 10^{16}$ losses of C$_2$D$_2$/cm$^2$/s as computed from the C$_2$ band. Therefore the C+ radiation is within about 20% consistent with this chemical sputtering mechanism. This is also supported by the time histories shown in Fig. 3, where the increasing CII flux replaces the C$_2$ flux as $T_e$ increases after the D2 gas feed is stopped. Some discrepancy may be expected to the presence of other sputtering mechanisms or due to the recombination of C++ backstreaming into the throat from the main plasma.

In both examples above, the pumps were active. In such cases, modeling with the DEGAS neutrals code has shown that flux amplification is ~1. As a result, the ion flux has been assumed to be the same at the probes and at the deflector (i.e. the sputtering yield estimates are conservative). We note, however, that these estimates have been based on global considerations, without consideration of profile effects or the details of the geometry. An examination of the latter effects has been made with the BBQ impurity transport code.

4. Impurity transport calculations. In order to relate the spectroscopic observations to fundamental sources of impurities we have modeled the expected spatial distributions of hydrocarbon and carbon populations in the pump limiter throat to compute the expected CD band emission for comparison with the OMA measurements. The ORNL BBQ code, an impurity Monte Carlo code patterned after the LIM code of P. Stangeby et al., has been used. The 3D geometry in the vicinity of the neutralizer plate has been treated and profiles of SOL plasma parameters have been used which are based on Langmuir probe estimates of the SOL $n_e$ and $T_e$ decay lengths. The semi-empirical fits for physical and chemical sputtering given in [3] are used to calculate impurity generation rates at the deflector. The neutral C (or C$_n$D$_m$) fluxes predicted by these rates are emitted from the neutralizer plate with Thomson energy and cosine angular distributions, and are tracked as neutrals or ions until they enter the pump, exit the throat region to the main SOL, strike the roof or floor of the pump limiter throat or strike the side walls. Figure 3 shows the expected distributions of the breakup products from chemical sputtering of CD$_4$, for the case described above. The calculations show that ~40% of the emitted C flux (starting as CD$_4$) goes to the pump, ~50% to the floor,
3% to the roof, 1% to the side walls, and 6% escapes the throat region to the main SOL. The BBQ results for this case predict 3.06 \times 10^6 cps for the OMA system, while the measured value is 2.5 \times 10^6 cps. The average sputtering rate calculated by BBQ is 0.009. For this case the fraction of chemically sputtered CD$_4$ relative to the total flux is assumed to be 0.46, in accordance with the results in [3]. However, recent calculations with TRIM suggest that the ratio of CD$_3$/CD$_4$ may be different than assumed here[4]. In addition, the inferred yields are very sensitive to the assumed T$_e$; varying T$_e$ from 20 - 35 eV in BBQ produces estimated yields from 0.005 to 0.012. Calculations in which the sputtering due to the deuteron flux to the deflector is augmented by C self-sputtering, (T$_C$/T$_D$ = 15%, Z$_C$=4) indicate that the self-sputtered C source may be roughly comparable to the deuteron-sputtered source.

5. Discussion. Evidence has been found for chemical sputtering from the neutralizer plate of the Tore Supra Outboard Pump Limiter, when the plate has a temperature \sim 700K, which should enhance such sputtering. The chemical sputtering yield has been found to be approximately consistent with extrapolations to high ion flux conditions (T$_D$\sim 10^{19}\text{ }/\text{cm}^2/\text{s}) in Ref. [3]. Impurity transport simulations based on the 3D geometry of the neutralizer plate region, and using the semi-empirical fits of [3], are also consistent with this estimate.

6. References.
[2]. A. Pospieszczky, et.al., UCLA-PPG-1251 (December 1989) Institute of Plasma and Fusion Research, University of California, Los Angeles, California, U.S.A.
[3]. J. Roth, E. Vietzke, A.A. Haasz, in ATOMIC AND PLASMA-MATERIAL INTERACTION DATA FOR FUSION (Supplement to the journal Nuclear Fusion), 1991, pp. 63-78
[4]. E. Gauthier. Private Communication.

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THE USE of LARGE SURFACE AREA for PARTICLE and POWER DEPOSITION

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I) INTRODUCTION

Since the parallel heat flux passing through the LCFS has increased dramatically with the size of machines one has to cope with very large particle and power fluxes on the limiters. Thus the size of the limiters has been increased by the use of inner bumper limiters (for example in JET, TFTR, TORE-SUPRA and JT60). The 'exponential-sine' model is widely used to estimate the heat flux (Q) to a wall for a plasma flux surface with incident angle θ. The model predict $Q = q_\parallel(0) \sin \theta \ e^{-p/\lambda_q} + q_\perp(0) \cos \theta \ e^{-p/\lambda_q}$, where $\theta = 0^\circ$ when the flux surface is exactly tangential to the limiting surface, $p$ is the minor radius measured from the last closed flux surface (LCFS), $\lambda_q$ is the SOL decay length of the heat flux density and $q(0)$ is the heat flux density at the last closed surface. If we approximate the heat flux as $Q = q_\parallel(0) \ e^{-p/\lambda_q} \sin (\theta + \alpha)$, with $\alpha = \tan^{-1}[q_\perp(0) / q_\parallel(0)]$, then $\alpha$ can be interpreted as an effective 'minimum angle of incidence'. Under conditions where the geometric angle $\theta$ has been made almost grazing (below 5°) the predictions of the simplest model (with $\alpha = 0^\circ$) is not adequate to represent the observation made in TORE-SUPRA; a similar result is found in TFTR [1].

Experimental observations of heat and particle deposition on the large area limiter on the inner wall of TORE-SUPRA are presented. These results have been analyzed with a Monte Carlo code (THOR) describing the diffusion of hydrogenic particles across the LCFS to the limiting objects in the Scrape Off Layer (SOL), and by impurity generation calculations using the full 'exponential-sine' model ($\alpha=0$) used as input to an impurity (carbon) Monte Carlo code (BBQ).

II) EXPERIMENT

A sequence of shots has been dedicated to the understanding of the behavior of power and particle deposition and its consequences when the plasma is limited by the large inner bumper limiter. The TORE-SUPRA plasma cross section has a circular shape ($a < 0.78$ m) with a main radius of $R = 2.34$ m. The inner bumper limiter is made of more than 7500 tiles brazed to a mechanical stainless steel structure actively cooled by pressurized high velocity water ($P = 3.5$ MPa, $v = 7$ m/s, $T < 230^\circ$C) so that very long pulses (up to one minute to date) can be handled with steady state plasma conditions. The inner wall represents a surface of 24 m² but the total surface of the tiles represents 12 m². The ohmic shots presented here have the following characteristics:

- $I_p = 1$ MA, $Z_{eff} = 1.5$, $<n_0> = 2 \times 10^{19}$ m⁻³, $B_T = 3.85$ T, $a = 0.72, 0.75$ and $0.78$ m.
III) MEASUREMENTS

TORE-SUPRA plasmas are diagnosed by a set of standard diagnostics which will not be described in detail here. The major diagnostics and their main findings are: a CCD camera equipped with a narrow 0.89 μm filter, which views the hot surface of the inner wall tiles. The camera was not absolutely calibrated but relative profiles were obtained. When averaging in the toroidal direction one find a poloidal profile which is extremely broad (Fig. 1); a fast reciprocating Langmuir probe located at the top of the machine measures the plasma edge parameters. A complete profile was taken at three times during a shot. The electron temperature and density profiles were unfolded taking into account the poloidal variation of the edge magnetic geometry due to the Shafranov shift ($\Delta = \beta_p + \lambda/2 = 0.8$). The $e$-folding lengths we find are very large and are in good agreement with those found for the same situation on TFTR [2]. Other Langmuir probes are placed at the front face of each of the two lower hybrid grills and give information on the SOL. They also indicate a broadening of the SOL when the plasma is leaning on the inner bumper limiter [3]. Magnetic probes indicate a relatively small deformation (2% from an ideal circular shape) of the LCFS. A visible spectrometer is used to look at CII and $H_\alpha$ emission locally at the inboard bumper. The general shape of the profiles can be explained if one takes into account both the detailed geometry of the TORE-SUPRA inner first wall (see the modeling below) and the actual viewing geometry of the telescope system.

This information is summarised for the largest and smallest plasmas in Table 1. These are relatively low power, low density cases, thus $n_e$ and $T_e$ in the scrape-off layer are rather low. The results are presented at the inboard and outboard equatorial plane, when the plasma is leaning on the inner bumper.

Table 1: Characteristics of the plasma scrape off layer at the inner (outer) wall position measured by a Langmuir probe positioned at the top of the machine (the Shafranov shift has been taken into account) and code results when one assumes a constant cross field diffusion coefficient of 1 or 10 m$^2$/s. Here $a$ is the minor radius, $z$ is the vertical displacement, $\lambda_{n,T,I}$ are the SOL decay lengths for $n_e,T_e$, $l_{sat}$ and $\lambda_{1,10}$ are the predicted SOL decay lengths for $D_I=1,10$ m$^2$/s

<table>
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<th>$a$ (cm)</th>
<th>$z$ (mm)</th>
<th>$\lambda_n$ (cm)</th>
<th>$n_0$ ($10^{12}$ m$^{-3}$)</th>
<th>$\lambda_T$ (cm)</th>
<th>$T_0$ (eV)</th>
<th>$\lambda_i$ (cm)</th>
<th>$\lambda(1$m$^2$/s)</th>
<th>$\lambda(10$m$^2$/s)</th>
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<td>0</td>
<td>(5.7)96</td>
<td>2.5</td>
<td>(9.2)15.6</td>
<td>22</td>
<td>(4.8)8.1</td>
<td>1.8</td>
<td>----</td>
</tr>
</tbody>
</table>

IV) INTERPRETATION

A) Monte Carlo calculation of deposition on Inner wall. The "THOR" code has been developed to calculate the particle deposition on the objects located in the Scrape Off Layer. The results are the impact positions of test particles on the objects. The geometry introduced so far in the code is simplified. The magnetic equi-
librium gives circular cross sections (which is very close to the experiment) where the Shafranov shift due to plasma pressure is taken into account. The ripple due to the discrete coil configuration is ignored since it is small at the inner first wall (0.6 < δ < 2 mm). Up to 100,000 particles are launched randomly inside the LCFS. They are assumed to flow along field lines at a fraction of the ion sound speed and diffuse randomly in the poloidal and radial direction (1 m²/s < D⊥ < 10 m²/s is assumed poloidally uniform). The result of the code is quite different from the simple "exponential-sine model" with α=0 (Figure 1). At the tangency point (θ=0) where the simple model with α=0 stipulates that there is no particle or power deposition one finds a hole (whose depth depends mainly on the cross field diffusion) or a maximum. The deposition pattern has a bell shape for the TORE-SUPRA configuration and plasma parameters. When the cross field diffusion is turned off only near the inner bumper limiter (half of the machine at the high field side) one finds a double peak structure and a null at the mid-plane as one can expect from the simple "sine model" with α=0. We find a good agreement between the experimental and the calculated profiles when the cross field is D⊥ = 10 m²/s. This value seems large (DBohm=Te/(16eB)=0.5 m²/s) but we have to bear in mind that the code does not account for the ionisation of neutrals within the SOL. The e-folding length for density when a large ionisation source is present near the LCFS can be expressed as: λn = [D⊥ Lc/(Vth n₀ <σv> L)]¹/² where n₀ is the density of the neutrals and <σv> the total rate coefficient for ionisation. The only way to match the large λn seen experimentally is to assume a large ionization source in the SOL. This means that the particle recycling occurs in the high heat flux region at the inner wall, a situation that we can call "closed recycling" [4]. This is comparable to the TFTR situation where an e-folding length for power deposition λq is found to be 3.5cm, which is to be compared to the λq = 1cm found (as on TORE-SUPRA [5]) when the plasma is leaning on the outboard limiter.
B) Impurity distribution in the SOL

The measurement of the poloidal distribution of impurity sources gives another indication of the flux deposition. The BBQ impurity transport code has been used to calculate the expected distribution of C as a function of poloidal angle. This code incorporates the physical processes treated by the LIM code [6], but it is specialized to the detailed Tore Supra geometry with discrete tiles. The calculated poloidal profile of impurity generation along with Thomson energy and cosine angular emission distributions, is used as a starting condition for a Monte Carlo transport calculation. The code tracks the wall-emitted C neutrals until they return to the wall or penetrate (by random cross-field diffusion) to a 1 cm depth beyond the last closed flux surface. Figure 2 shows the poloidal distribution of particle deposition calculated with the full exponential-sine model, for \( q_\| (0)/q_\perp (0) = 2.5 \) [the value expected for TORE-SUPRA according to the Stangeby et al model in which \( \Phi_\parallel/\Phi_\perp = (L_\parallel/L_\perp)^{0.5} \); \( 2L_\parallel = \) field line length along the inner wall]. The predicted poloidal distributions of C density from the BBQ code are shown in Figure 3, for \( q_\| (0)/q_\perp (0) = 2.5 \) and \( D_\perp = 1 \) m²/s. The CII poloidal profile matches the main characteristic features of spectroscopic distributions on the inner wall, when account is taken of the spatial position and resolution of the present TORE-SUPRA spectroscopic system. These profiles are very sensitive to the magnetic configuration; the predicted peak on the equator, for example, is strongly reduced if the elongation is changed from 0.98 to 1.05, thus suggesting the possibility of influencing the heat flux deposition in nominally circular tokamaks using small changes in plasma shape.

V) CONCLUSIONS

The 'exponential-sine' model for the deposition of heat and particles on a large area limiter can give a good representation of the TORE-SUPRA observations only if several conditions are fulfilled: 1) the perpendicular heat flux must be included (\( \alpha \neq 0 \)); 2) the detailed geometry must be taken into account, and 3) the actual SOL decay length (which includes the effect of sources due to localized recycling at the inner wall) must be used.

REFERENCES
PARTICLE CONTROL WITH A PUMP LIMITER ON TORE SUPRA

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Introduction

Tore Supra has been designed for long pulse operation which requires continuous and simultaneous control of plasma density and particle removal. A set of seven pump limiters (6 bottom and 1 outboard) has been installed and the outboard pump limiter (OPL) has been equipped [1, 2] with diagnostics to study the physics inside the pump limiter as well as to assess the efficiency of such a system to control the plasma density. Previous ohmic experiments have shown that the exhaust efficiency depends strongly on the volume average plasma density [3].

In this paper we describe a series of successive ohmic shots of different plasma densities with boronized wall. It is shown that without active pumping by the limiter the wall content becomes sufficient to maintain the requested density. Activation of the pumps in the OPL plenum allows both the plasma and the wall contents to be pumped out. In the first part of this paper, we describe the experimental sequence. Quantitative analysis of the balance between externally injected particles versus pumping by the wall and by the pumps of the OPL are addressed in the second part.

Experiment

The reported experiments concern a series of 15 consecutive ohmic shots (I_p=1.6 MA) with a deuterium plasma leaning on the OPL (surface interacting with the plasma 0.3 m²). The total number N of particles in the discharge can be described by the conventional balance dN/dt = Γ - N/τ_p*, where Γ is the injected particle flow, and τ_p* the apparent confinement time τ_p* = τ_p/(1 - R + e), where τ_p is the particle confinement time, R the global recycling coefficient, and e the exhaust efficiency of the OPL.

The features of a single shot of the sequence included a gas puff with feedback on the density from t=0 to 5s; at t=5s the gas puff was cut off. The first eight shots (10442 to 10449) were “non-pumped” and made prior to activating the titanium getter in the pump. The six following shots (10450 to 10455) were pumped and finally, the last shot (10456) of the series was non pumped. The average plasma density was increased from \langle n_e \rangle = 2.4 \times 10^{19} \text{m}^{-3} for the first shot to \langle n_e \rangle = 4.4 \times 10^{19} \text{m}^{-3} for the last non-pumped shot (#10449). For t>5s, in the absence of a source (Γ=0) τ_p* is deduced from τ_p* = -N/(dN/dt). In these experiments the
part of the wall in contact with the plasma had been initially saturated; the global recycling coefficient \( R \) was close to 1 (\( \tau_p^* > 10-15 \text{s} \)), even for the largest densities (\( \langle n_e \rangle = 4.4 \times 10^{19} \text{m}^{-3} \)), in contrast to a depleted wall where \( \tau_p^* \) has been found to decrease with density [2]. Figure 1 displays the density evolution versus time for two consecutive shots, non-pumped (#10449) and pumped (#10450) respectively.

The effect of pumping by the OPL on the density is very clear when the valve of the gas injection is shut off (t=5s) and it can be observed consistently that the total quantity of gas injected is about twice as much for the pumped shot as compared with the non-pumped shot. As t approaches 5s, the injected flux is close to zero for the non pumped shot (\( \tau_p^* > 10 \text{s} \)), whereas for the pumped shot the injection approaches a value close to 3 \( \text{Pa m}^3 \text{s}^{-1} \) (\( \tau_p^* = 2 \text{s} \)).

The exhaust efficiency is defined as the ratio of the particle flux exhausted by the pump limiter, divided by the plasma efflux. With \( p \) and \( S \) being pressure and pumping speed in the pump limiter respectively, the exhausted flux is \( \Gamma_{\text{exh}} = pS \). To determine the plasma efflux \( N/\tau_p \), the particle confinement time has to be estimated. Assuming a typical value around 100 ms, exhaust efficiencies of 10% result. In these experiments the exhaust efficiency also increases with the average plasma density [3], but it should be noted that it does not depend crucially on the saturation state of the wall.

**Particle balance**

The particle balance has been realised for the overall series of the 15 shots and each shot has been divided in three phases. The first phase concerns the plasma build-up between 0 and 3.25s. The second part describes the phase of stationary plasma density between 3.25 and 5s. And the third phase deals with the density decay between 5 and 8s after the gas puff is turned off. The particle balance for each of these three phases is shown as a function of shot number in figure 2.

During the first phase, the total number of particles injected is four times higher than the total number in the discharge. This behaviour is the same for non-pumped and for pumped discharges, and is also independent of the density. This is explained by the fact that the plasma is built up on the outboard pump limiter but has sufficient connection to the inboard wall. This way, the charge exchange area around the outboard limiter is depleted while the parti-
cles are absorbed by the inboard wall. Under such conditions, the part of the wall responsible for the plasma build-up releases particles during the first three seconds of the discharge while the inboard wall due to its large capacity now absorbs the majority of the particles.

The second part of the discharge concerns the steady state conditions for the density, between 3.25 and 5s. The gas puff is essentially balanced by the pump limiter exhaust for the pumped shots. For the non-pumped shots, the gas puff is going to zero while the density remained roughly constant showing an equilibrium between the outflux from the plasma and the recycling flux.

The third phase describes the final part of the discharge between 5 and 8s, when the pump limiter exhaust exceeds the density decay of the plasma. Since the plasma is at the outboard limiter, this means that the charge exchange footprint around the limiter is depleted and that the fuelling ratio [4] (defined as the ratio of the total number of particles in the plasma to the total number of particles injected $F = N/\phi_{gas}$) of the following discharge is lowered. The fuelling ratio decreases in subsequent discharges whereas a saturation of the wall would have shown an increase. Figure 3 displays the fuelling ratio $F$ and the...
averaged plasma density at \( t=5s \) as a function of the shot number. \( F \) decreases from shot to shot even when the pumps are not activated, and it drops to less than typically 20% when the pumps are activated. Indeed, the pumps of the OPL were shut off prior to shot 10456 (i.e. the OPL pumps were not active) and both the fuelling ratio and the total number of particles injected for that shot show that the wall has built up some pumping capacity. Figure 4 displays the density evolution and the injected gas flux \( \Gamma \) versus time for this shot.

The total amount of injected gas has been multiplied by 2 for shot 10456 compared to shot 10449 (which is a similar shot without active pumping) to obtain a density of \( \langle n_e \rangle = 5 \times 10^{19} \text{ m}^{-3} \) i.e. 15% higher than for shot 10449. Furthermore, in contrast to the "non-pumped" shot 10449 where the gas puff decays to zero at \( t=5s \), the gas puff for the "non-pumped" shot 10456 remains high during the plateau phase (\( t=3.25-5s \)). Since the pump limiter is "off", this indicates that the wall in contact with the plasma has developed some pumping capacity during the pumping sequence.

Summary

This series of ohmic discharges demonstrates the potential of the pump limiter for particle control in Tore Supra. While the plasma build-up phase between \( t = 0 \) and 3.25s shows similar characteristics for "non-pumped" and pumped discharges, the stationary and decay phases of the discharge show very clearly the effect of the pump limiter. During the stationary phase, the density is maintained constant with a gas puff which is balanced by the pump limiter exhaust. After the gas puff is turned off, the density decays rapidly. During this phase, the particle exhaust, measured in the pump limiter, exceeds the density decay of the plasma, indicating that part of the exhausted gas is due to depletion of the wall inventory. Consequently, the overall effect of the pump limiters is not only plasma density control, but also control of the wall hydrogen inventory, with boronized wall.

References

EXTENSION OF THE OPERATIONAL DOMAIN OF THE ERGODIC DIVERTOR IN TORE SUPRA


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INTRODUCTION

On TORE SUPRA over the past years, the Ergodic Divertor (a set of 6 octopolar coils producing the stochasticity of magnetic field lines outside the q=2 surface) has exhibited some global macroscopic properties which may be used to improve plasma discharges in at least two fields: 1) the MHD stability of edge kink and tearing modes and 2) the production of stable edge radiating layers.

The serious problem of the fuelling with gas injection of ED discharges has been recently partly overcome with the boronization of the vacuum vessel. This has allowed to extend the operational domain of the ED to 3 times larger plasma densities.

In parallel, the stabilizing effect of the ED on the edge MHD activity has been recently used to stabilize the kink mode triggered at q=3 during current rise as well as the saturated m=2,n=1 tearing mode which is always observed without ED in the low q domain (2<q<2.7) i.e. far below the ED resonance (qa=3.2). In this sense the operational domain of the ED has been extended to low q discharges.

1) HIGH DENSITY DISCHARGES WITH THE ED

It has been stressed that the obtainment of high density deuterium plasmas in ED experiments is a difficult problem [1]. Without gas injection, the deconfinement induced at the edge by the ED produces a drastic density decrease (up to a factor of 2 at resonance) [2,3].

In previous TORE SUPRA experiments with carbonized walls any attempt to compensate the density decrease by gas injection has led to disruptions. The disruption is the result of the combination of several factors: 1) The temperature in the ergodic layer is low (below 100 eV); 2) The fuelling efficiency is very low (the screening of the ED acts also on Deuterium), consequently very hard gas puffing is needed (few Pa m^3/s) which increases the recycling flux and further cools down the plasma edge; 3) this influences the oxygen production rate, the brightnesses of the low ionization state of oxygen increasing dramatically in agreement with the behaviour of the total radiated power. The disruption occurs as soon as the ratio Prad/Pohm reaches unity (as in standard density limit disruption). Consequently, the density operating regime in ED experiments was restricted to volume averaged densities <ne> lower than 1.5 \times 10^{19} m^{-3}.

Recent boronization of the vacuum vessel has allowed much higher stationary densities to be obtained, in spite of the fact that the fuelling efficiency remains very low (1-2%). It has been possible to increase and sustain densities up to <ne> = 4.5 \times 10^{19} m^{-3} without disrupting. This was achieved by strong gas injections of the order of 1 to 5 Pa m^3/s (1 Pa m^3 = 5.33 \times 10^{20} D^0) over the entire pulse length (Fig.1&Fig.2). The main effect of the boronization has been to reduce the oxygen contamination of the plasma at the highest densities by at least a factor of 5 [4]. In a boronized vessel the intrinsic impurity which increases with density (and recycling) and which determines by its radiation rate the density limit is chlorine /5/. Possibly due to its shorter ionization length the chlorine is screened better than oxygen; consequently in boronized vessel the temperature inside the ergodic layer can be lowered more with hard gas puff before chlorine enters the discharge and causes a radiative collapse. This allows the obtainment of much higher plasma densities.

One cannot note that the effect of the boronization is more important in ED experiments (an increase of a factor 3 of the working density) than in limiter experiments (an increase of 20-30% of the density limit), thus illustrating the interest of an effective screening of impurities.

Two examples of high density discharges are shown on Fig.1&2 (R=2.4 m, a=0.75 m, Bt=3 T, Ip=1.4 MA, qa=3.2).

On Fig.1 shot 9610 is the highest plasma density (<ne>=3.8 \times 10^{19} m^{-3}) achieved in ohmically heated ED discharges. The screening of impurities can be observed by the strong reduction of the impurity lines brightness (CVI OVI) as the plasma current reaches the resonant
FIG.1 High density ohmic discharge with the ED. In such discharges the temperature measured in the ergodic layer by reciprocal Langmuir probes or by Doppler broadening of a CVI charge exchange line is of the order of few tens eV /s. On shot 9610 the ratio Prad/Pohm increases with density up to 60% while Zeff is kept to about 1.2. The increase of the radiation rate is usually correlated with an increase of the carbon line brightness, but carbon is probably not the only radiating impurity in the ergodic layer; low ionization states of chlorine may also play a role. This will be studied in future experiments. In such ED discharges the brightness profile measured with the bolometer array exhibits the unusual shape of a stationary low field side "Marfe" (Fig.3).

FIG.2 Ohmic+LHCD ED discharges.

On Fig.2 shot 10077 is the highest plasma density achieved yet (\(<n_e>=4.510^{19}\)m\(^{-3}\)) with the addition of 1.5MW of Lower Hybrid wave. Zeff is of order of 1.3 and the ratio Prad/Ptot reaches about 80%. The bolometer profiles exhibit the shape of a stationary high field side Marfe (Fig.4). The local emissivity inside the Marfe is of the order of 3MW/m\(^3\). The Lower Hybrid power increases the carbon sources which are necessary to create and sustain the Marfe.
It has to be noted that the neutral pressure in the pump limiter plenum (Fig.1&2) reaches 1 to .2 Pa in spite of a very low ion influx in the pump limiter throat.

II) LOW q DISCHARGES, CURRENT PROFILE AND STABILITY

It has been already reported that ergodization of magnetic field lines on a thin layer at the plasma boundary has a strong stabilizing effect on both the surface kink modes and the m=2 n=1 tearing mode [7,8]. Although the ED resonance is centered at q=3.2 these effects have been observed over a wide range of edge safety factor: 2.2<q(a)<4.5. In previous experiments these effects have been used (1) to stabilize the m=2 mode which develops close to the density limit; this, combined with the particles deconfinement induced by the ED, has allowed to define a strategy for preventing density limit disruptions [7]; (2) to stabilize the surface kink mode which develops during the current rise at q(a)=4; the destabilizing effect of the ED allows to increase by about a factor of 2 the critical current rise rate at which the mode locks and triggers disruptions [8].

In a new series of experiments the ED has been used to stabilize the kink mode which appears at q(a)=3. As reported in [8] for q(a)=4, two stabilizing effects are observed: 1) the ED tends to increase the internal inductance (Fig.5) and with the ED the kink stability limit is shifted towards lower r_i. With the ED it has been possible to cross q(a)=3 at a rate of 1MA/s without disrupting.

Moreover in the low q domain (2.2<q(a)<2.7) it has been observed that the ED stabilizes the saturated m=2 tearing mode which always perturbs the plasma equilibrium in limiter discharges. It has been found that the ED magnetic perturbation amplitude needed to fully stabilize this mode decreases with q(a). On Fig.6 only one fifth (0.5kA) of the maximum available ED current is needed to stabilize the m=2 mode. As pointed out in [11] this indicates that the stability of the m=2 mode depends on the relative extension of the ergodic layer with respect to the radius of the the q=2 surface.

![Graph](https://example.com/graph.png)

SAFETY FACTOR q

FIG. 5 li versus q(a). These curves illustrate the effect of the ED on the current density profile and illustrate also the resonant response of the plasma to the ED perturbation.

In such low q discharges with the ED the only edge MHD activity which is observed is correlated with the sawtooth crash and corresponds to the “gong” mode [9] (Fig.6).

These stabilizing effects on kink and tearing modes have been interpreted as resulting from a strong modification of the equilibrium current density profile. In the ergodic layer the connection of magnetic field lines with the wall induces a large increase of the plasma resistivity in this region. This leads to an increase of the current diffusion toward the center and to an increase of the peaking factor of the current density profile. This interpretation was supported by the characteristic behaviour of the ohmic power and of the internal inductance li. For example Fig.5 shows the behaviour of li versus q(a) for two shots with and without ED. The li curve again illustrates the resonant response of the plasma to the ED perturbation.

The current density profile identified with the IDENT-D equilibrium code fed with magnetic and polarimetric measurements [10] has allowed to confirm this theoretical interpretation developed in [11]. Current density profiles with and without ED are shown in Fig.7 at the ED resonance q(a)=3 and Fig.8 at q(a)=2.5. It must be stressed that the current density profiles with ED shown on Fig.7&8 are obtained from IDENTD without any additional assumption taking into account the effect of the ED on the magnetic structure of the plasma edge to cancel the current in the ergodic layer.

![Graph](https://example.com/graph.png)

FIG. 6 SHOT 9661 the switching off and on of the m=2 oscillations by the ED at q=2.5. The spike t=5.11s during the quiet phase is linked to a sawtooth crash and corresponds to a “gong mode”.
The measurements are coherent with the fact that the stabilizing effect on the surface kink mode comes from a decrease of the current density gradient at the edge.

III) CONCLUSIONS

1) Boronization of the vacuum vessel has allowed to extend the working density with the ED by a factor of 3. The reduction of the oxygen contamination of the plasma allows to tolerate much higher recycling fluxes between the wall and the ergodic layer and to approach the formation of a "cold mantle". The extremely low edge temperature allows to sustain significant radiation rates. It is observed that the radiation rate of carbon increases. The contribution of the low ionization states of chlorine which is expected also to play a role has not yet been quantified.

2) The stabilizing effect of the ED on the edge MHD activity has been observed on the kink mode when crossing \( q(a) = 3 \) and on the \( m=2 \) tearing mode down to \( q(a) = 2.2 \).

3) Current density profiles identified with IDENTD prove that in fact the stabilizing effect results from a strong modification of the equilibrium current density as expected earlier from theoretical considerations.

4) Next steps will be: a) to use the ED to produce and control edge radiating layers in plasma heated with a significant power level; b) to produce low \( q \) discharges (\( q(a) = 2.5 \)) fully free of MHD activity by combining the stabilizing effect of the ED on the \( m=2 \) mode with the non inductive lower hybrid current drive which is known to have the capability to stabilize the sawtooth activity and thus the "gong mode".


1/C.BRETON & al. NF 1991 Vol.31 p.1774
2/Ph. GLÉNDE & P. Schmitz & al. 15th EPS Vol.I p.253
3/Ph. GLÉNDE & P. Schmitz & al. 15th EPS Vol.II p.847
4/G. FUCHS & al. 15th EPS Vol.II p.2662
7/J. C. VALLET & al. 19th EPS Vol.II p.253
8/E. JOFFRIN & al. 19th EPS Vol.III p.855
9/P. A. DUPEREK & al. 19th EPS Vol.II p.1421
10/P. A. DUPEREK & al. 19th EPS Vol.II p.1421
11/D. EDERY & al. 19th EPS Vol.III p.1421
12/D. EDERY & al. 19th EPS Vol.III p.1421

FIG.7 current density profiles with/without ED at the ED resonance (\( qa = 3.2 \))

FIG.8 current density profiles at \( qa = 2.5 \)
EFFECTS OF THE ERGODIC DIVERTOR
DURING LHCD EXPERIMENTS
IN TORE SUPRA

and Lower Hybrid group

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The compatibility of lower hybrid (LH) current drive in the presence of a stochastic layer created by the ergodic divertor has been investigated in Tore Supra with the plasma leaning on the outboard limiter in contrast with preliminary results, at low density, with the plasma leaning on the inner wall /1/. Various important issues have been addressed such as: LH waves coupling, power deposition on plasma facing components, screening of incoming impurities and finally effect of the ergodic divertor operation on fast electrons and on the LH current drive efficiency.

1. LH waves coupling

Coupling of LH waves, which is sensitive to the electron density of the plasma layer facing the antenna (grill), is well understood and can be accurately modeled /2/. When a stochastic edge is produced by the ergodic divertor configuration /3/, field line connections are created between the grills and the ergodic divertor coils; consequently the density at the grills aperture is no longer homogeneous all across the radiating surface but is modulated according to the flux tubes length which can be as low as 1 m /4/. This typical poloidal size of the modulation is about 20 cm i.e. the RF module (4x2 waveguides) size for which incident and reflected RF power measurements are performed. Modification of the coupling is indeed observed, when the ergodic divertor is applied, with different behaviours of grill 1 and grill 2 (which have a different location with respect to the magnetic perturbation). For the 16 modules composing each grill, different trends generally occur as shown in figure 1 where the reflection coefficient R averaged on a quarter of a grill (4 modules) is plotted versus time. In fact zones of higher reflection are those for which the particles and heat fluxes are decreased as a consequence of the connection length decrease. A Langmuir probe, located in the midplane of the grill close to the electron drift side of the guard limiter, indicates a strong decrease of the density from 1.5 x 10^{18} cm^{-3} to 3 x 10^{17} cm^{-3}. This is consistent with a local increase of R between 2 and 5 %. In most cases the global reflection coefficient is maintained at a low value and the maximum values are still manageable for high power transmission with no significant distortion of the N_{//} spectrum.
II-616 4-16

2. Heat deposition

The local connection of field lines induces a strong anisotropy in heat deposition as observed on the outboard pumped limiter (OPL) /5/. However the deposition pattern in LH heated plasmas is the same as during ohmic plasmas. With an injected LH power of 3 MW, the maximum surface temperature of the carbon tiles remains below 1000°C. It is worthwhile to note that infrared images of the guard limiter of grill 2, which has the same position as the OPL with respect to the magnetic configuration, exhibit the same location of the hotter zones and are consistent with experimental coupling modification given by fig.1. The heat load on the limiter (and on the grills) can be strongly reduced by switching to a highly radiating regime /6/. However, in all cases, when the plasma is leaning on the OPL, calorimetric measurements of the integrated heat flux indicate a strong asymmetry for the neutralizer plates of the 6 divertor modules. Figure 2 indicates that module 6, located close to the grills, extracts 6 times more than module 3, when the grills are at the same radial position that the divertor modules. This is confirmed by IR images which show overheating of neutralizers magnetically connected to the grills particularly on the electronic drift side. However this additional power loading (60 kW), focussed on small areas, is only 2% of the injected power. This could be interpreted by a diffusion of electrons along the magnetic field lines connected to the neutralizer plate which appears possible due to the ED induced radial deflection of about 3 cm in front of each module /4/. This assumption is supported by the different heat load patterns when the radial distance between the grills and the divertor modules is changed.

3. Impurity screening

The main beneficial effect of the ergodic divertor is a reduction of the impurity content in the plasma core. In ohmic discharges, a reduction of the carbon concentration by a factor of 2.5-3 is deduced from UV spectroscopic measurements of the CVI line /7/. In LH current drive experiments, such a reduction is also measured for carbon and copper (Cu XIX). With 2 MW of LH power, at a volumic density \( \langle n_e \rangle = 2.3 \times 10^{19} \text{m}^{-3} \) in helium, carbon and copper concentrations, at \( r/a = 0.7 \), are reduced by 40 and 25% respectively (Fig.3). On the other hand, the CII signal shows an exponential increase after 1 s, as a possible consequence of enhanced carbon local source or \( T_e \) lowering at the very edge. At higher LH power (3 MW), good screening of carbon is still obtained but localized copper bursts have been noticed when the plasma is leaning on the OPL. These bursts are well correlated to the surface temperature of the divertor neutralizer plates. For the different experimental conditions (\( P_{\text{LH}} = 1-3.5 \text{ MW}, \langle n_e \rangle = 1.5-5 \times 10^{19} \text{m}^{-3} \)), the carbon concentration is reduced by a factor = 2, significantly smaller than in ohmic discharges (Fig.4).
4. Current drive efficiency and effect on fast electrons

The high diffusion coefficient in the ergodic zone /8/ might affect the confinement of fast electrons and therefore the current drive efficiency. Analysis of hard X-ray emission shows no significant change for the central chord but a 2-fold increase, at $E = 75$ keV, for the external one ($r/a=0.71$). After Abel inversion, radial profiles confirm a significant change in the edge: X-ray emission profile is flattened in the outer region ($0.6 < r/a < 0.75$) during the entire LH + ergodic divertor phase (Fig.5). The increase of impurities at the edge, between $t = 7.9$ s and $t = 8.9$ s does not affect the profile and $Z_{eff}$ effect has to be discarded. In contrast, the surface loop voltage is constant up to $t = 8.5$s and then slightly increases due to an increase of $Z_{eff}$ at the edge (Fig. 6). Thomson scattering $T_e$ profiles indicate no change in the central region but the usual cooling of the very edge when the ergodic divertor is activated. However second harmonic ECE emission (viewing also the down-shifted 3rd harmonic emission of fast electrons) is the same for all the experimental chords after density restoration giving confidence in the conservation of fast electrons content at least in the core of the discharge.

Conclusion

During LHCD experiments with the ergodic divertor, the new connection of edge field lines leads to inhomogeneity in LH coupling and heat loading of plasma facing components. However satisfactory LH coupling conditions can be obtained and heat loads, up to 3 MW of LH power, have been managed. However copper bursts from highly thermally loaded neutralizers are observed at high power. Good impurity screening is maintained. With 2 MW of LH, at medium density, current drive efficiency is not changed by the activation of ergodic divertor. Accordingly, the fast electrons content is not affected in the core of the discharge while hard X-ray emission indicates modification at the edge.

References:
/1/ T. Evans et al., J. Nucl. Mat. 196-198 (1992) 421
/2/ X. Litauodon et al. Nucl. Fusion 32(1992) 1883
/6/ J. C. Vallet et al., this conference
/7/ C. De Michelis et al., J. Nucl. Mat. 196-198 (1992) 985
Fig.1-Reflection coefficients (Grill 1)

Fig.2-Integrated heat flux of the 6 neutralizer plates

Fig.3-Carbon and copper content (Shot 10005)

Fig.4 Carbon concentration vs $P_{100}/<n_e>$

Fig.5-Hard X-ray emission vs normalized radius

Fig.6-LHCD experiments (Shot 10005)
CONDITIONING OF FTU CRYOGENIC VACUUM CHAMBER

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INTRODUCTION

On a high density tokamak as FTU the content of low Z impurities is crucial to assess the plasma operational limits and to determine the ability to recover stable plasma discharges after major disruptions.

The optimization of the surface conditioning is very important and particularly relevant for FTU because of the complexity of the inner wall. The vacuum chamber is completely covered by thermal shields that originate hidden regions that can trap neutral gas and impurities. The only way to clean, at the same time, both the exposed and not exposed surfaces is the combined operation of baking plus glow discharge in hydrogen [1,2].

In the paper, after a brief description of the experimental apparatus and the operative procedures, we compare the plasma results obtained before and after glow discharge was applied to clean the vacuum wall, accidentally contaminated by a carbon and oxygen layer. As a consequence of the reduction of low Z impurities, an improvement of the density limit was obtained. In addition after the conditioning, also the recovery from plasma disruptions was easier.

EXPERIMENTAL APPARATUS

A simple electrical arrangement is used to produce the glow discharge [3]. Two cylindrical stainless-steel electrodes are inserted into the vacuum vessel at the midplane through two vertical ports located 180° toroidally from each other. The anodes are supported by Cu conductors which are insulated with alumina sleeves to prevent breakdown along the ports. The electrical circuit has been designed to supply the electrodes with high voltage (1200 V) at low currents (≤1 A) to start the discharge and a lower voltage (800 V) at higher currents (3-10 A) in the quiescent phase of the discharge. The negative pole of the power supply and the vacuum vessel are grounded.

During the conditioning the production of gases is monitored by a differentially pumped quadrupole mass analyzer (QMA) which is calibrated by means of an ion gauge with four different gases (H₂, CH₄, N₂, Ar).

OPERATIVE PROCEDURE

Glow discharge in hydrogen is routinely utilized on FTU before cooling the machine to liquid nitrogen temperature for an experimental campaign. The cleaning procedure begins with the baking of the machine at 120 °C performed by inducing a toroidal current on the vessel using the vertical field conductors and the transformer as primary coil; generally the machine is maintained at this temperature for five days and every day several hours of glow discharge are applied. No reconditioning of the wall, during the operation of the machine, is required because no degradation of the plasma performance is observed.

Under the typical working conditions the hydrogen pressure during the glow discharge is around $5 \times 10^{-3}$ mbar while the voltage drop across the discharge is of
350 V at the total current of 3.25 A. This value corresponds to 25 µA/cm² on the FTU vacuum chamber.

In these conditions, according to our previous experiments [2], the glow fills the whole inner volume and very good cleaning efficiency is obtained.

EXPERIMENTAL RESULTS

In the following the experimental results obtained during the conditioning campaign of June 1992 are shown.

Glow discharge was applied to remove a carbon-oxygen layer deposited on the wall for the accidental exposure of a plastic insulator to the plasma.

In Fig. 1 is shown the temporal evolution of CO, the highest peak monitored on the QMA, during four glow discharges performed in subsequent days. A considerable increase of the signal was observed starting the first glow discharge rapidly decreasing in about 1 hour to an asymptotic value higher than the QMA background level (black squares in the Fig. 1). This behaviour is typical of a contaminated metallic wall dominated initially by the interaction of hydrogen with the carbon and oxygen deposited on the surface and at longer times with the carbon and oxygen present in the material as metal carbides and oxides.

It is interesting to observe the increase of the signal amplitude at the beginning of every glow discharges performed in subsequent days. This behaviour is a sign of a contamination of the plasma chamber probably due to the migration of gaseous impurities from the regions not directly exposed to the discharge. It is not surprising since the ports area (~36 m²) and the internal structure of the thermal shields (~25 m³) are larger than the inner vacuum chamber (~13 m³) and so dominate the outgassing.

In Table 1 are summarized the results of cleaning efficiency expressed in terms of number of monolayers desorbed during the glow discharge in the first 2 hours and after other 12 hours.

The following formula was applied:

\[ v_i = \frac{S_{pi} \times \Delta P_i}{KT \times C \times A} \]

where \( S_{pi} \) is the effective pumping speed in the plasma chamber for the impurity of mass i, \( \Delta P_i \) is the corresponding partial pressure increase, C is the number of molecules which forms a monolayer and A is the vacuum chamber area. Details of the evaluation are given in [3]. About 0.6 gr of carbon (CO, CH₄, CO₂) were removed in 14 hours of glow discharge from the FTU vacuum chamber.

<table>
<thead>
<tr>
<th>M</th>
<th>In the first 2 h</th>
<th>After other 12 h</th>
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<tr>
<td>CO</td>
<td>90</td>
<td>7</td>
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<tr>
<td>CH₄</td>
<td>28</td>
<td>3</td>
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<tr>
<td>CO₂</td>
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Fig. 1 Temporal evolution of CO quadrupole signal during glow discharge
To evaluate the effect of the glow discharge on plasma performances we compare the experimental results obtained before and after glow discharge (G.D.) cleaning, where in the first case only the baking of the machine was performed.

As a consequence of the better wall conditioning a significant reduction of low Z impurities from the wall was observed by visible spectroscopy. The ratio of fluxes of CIII or OII to the fluxes of the working gas were 6.4% and 3.0% respectively before the conditioning. After the conditioning this ratio decreased to 2.5% for carbon and 1.9% for oxygen. These results were obtained as mean values on the following experimental conditions: plasma current between 300 and 700 kA, electron density $3\times10^{19}\text{m}^{-3}$.

These data are consistent with the evolution of $Z_{\text{eff}}$ from resistivity measurements, as is shown in Fig. 2. Here the $Z_{\text{eff}}$ values are plotted as a function of the electron density at plasma currents $I_p<400$ kA. The crosses represent the values obtained in the period before G.D. whereas the other symbols in the subsequent period. It is easy to recognize that for all the density range, the $Z_{\text{eff}}$ values obtained after G.D. are systematically lower than the values obtained after G.D.

The first immediate indication of the low Z impurity reduction during the machine operations was the higher gas fuelling required for the breakdown of the plasma discharge; this means that it was easier to exceed the radiation barrier during the increase of plasma current.

The most important effect was the extension of the operational plasma regime with respect of the density limit (Fig. 3). The vertical dotted line represents the operation limit obtained before glow discharge cleaning. In the subsequent phase of operations, it was possible to extend the density limit even beyond the Greenwald limit (by 1.7 times) [4].

An other very important effect was the recovery from plasma disruptions. This result is shown in terms of number of shots required to reproduce standard plasma discharges. Before G.D. (Fig. 4), after a disruption (curve 1), two other plasma discharges were necessary (curves 2,3) before obtaining a normal discharge (curve 4).

Fig. 2 $Z_{\text{eff}}$ values as a function of electron density; the crosses represent the values obtained before G.D. while the other symbols the values obtained after G.D.

Fig. 3 Operational plasma regime: the inverse of the safety factor is plotted as a function of the Murakami parameter.
CONCLUSIONS

The glow discharge at high temperature (120 °C) is very efficient for the conditioning of FTU vacuum chamber characterized by a very complex geometry. The typical working conditions are: discharge current of 3.25 A (25 μA/cm²) and hydrogen pressure of about 5×10⁻³ mbar.

Low Z_eff values (<2), high density (>2×10²⁰ m⁻³), better recovery from plasma disruptions have been obtained in the experimental campaign following the first application of the glow discharge.

FOOTNOTE AND REFERENCES

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Impurity sources and impurity concentrations in FTU

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Introduction

The analysis of impurities in FTU has given special emphasis to the different behaviour of metal and light impurities in the discharge, in fact due to the poloidal metal limiter and the absence of deliberate low Z deposition on the wall, most of the operation is influenced by the limiter constituents.

FTU discharges are characterised by a $Z_{\text{eff}}$ value decreasing with the line averaged density and increasing with the toroidal plasma current. Previous work [1] has shown that metal impurities, produced by the interaction of the plasma with the limiter, dominate the low density plasmas while at high density only Oxygen and Carbon lines are detected in the VUV emission of the discharge.

In this paper spectroscopic observations in the visible, VUV and Soft X-ray spectral regions are exploited to gain a better understanding of impurity sources and to clarify the leading production mechanism.

Impurity diagnostics

The FTU poloidal limiter is composed of two sections bearing respectively 7 and 15 Inconel mushrooms which constitute the contact surface with the plasma, their total area represents 1/200th of the vessel surface. Spectroscopic comparisons of the emissions from equal areas of the limiter and the wall has shown [2] that its contribution to the global recycling is of order 30% for hydrogen and 10 times less for oxygen. We cannot assess any estimate for the production of the metals near the wall since their neutral emission is below the detection limit; it seems nevertheless reasonable (considering also the observed surface damage) that the overall metal flux in the plasma is coming from the limiter surfaces.

For determining the whole flux of particles exiting the limiter we developed a collection optics viewing half of the limiter mushrooms through two endoscopes pointed to its inner and outer sections. A visible OMA spectrometer is used to analyse the emitted light.

The spectra of the two halves are similar and the total emission is derived by the sum of the two contributions averaged on the number of mushrooms in the field of view; nevertheless a deeper analysis will be necessary for some of the lines (e.g. H$_\alpha$) which show a clear asymmetry in some plasma conditions.

Fig. 1 shows spectra taken on the two halves of the limiter, it must also be considered that in the external view some contribution arises from molecular emission inside the port.

Line emission of metal ions such as NiXVIII and CrXXII, measured by a SPRED spectrometer in the VUV, were selected to describe the plasma emission at intermediate radii.
Experimental results

Experimental results

Z_{eff} determinations in FTU show that for high density operation, when the main impurity production is due to the release of light molecules weakly bound to the wall, plasma purity is very high (Z_{eff} = 1.2 in hydrogen plasmas). In these conditions Z_{eff} is approximately constant with density and the flux-dependent OI brightness suggests saturation of its recycling at the limiter. A similar behaviour is shown by the OVI line measured in the VUV (see Figs. 2, 3). Oxygen concentrations, deduced from the bremsstrahlung enhancement factor measured by a soft X-ray Pulse Height Analyser, are below 1% in this density range in accordance with the values of Z_{eff} deduced from the plasma resistivity. Both OVI and OI decrease to very small levels for decreasing densities.

On the contrary, Ni and Cr lines in the VUV and in the visible spectra, decrease both steadily with density (Figs. 4, 5) and fall below the detection limit when Z_{eff} approaches its lower value. This corresponds to a reduction in the Ni
concentration derived from Kα line intensity, from 1% at \( n_e = 2 \times 10^{19} \text{m}^{-3} \) to 0.01% at \( n_e = 1.2 \times 10^{20} \text{m}^{-3} \).

Comparing the data in Figs 4 and 5 with \( Z_{\text{eff}} \) plotted in Fig. 6, it appears that the large variations of \( Z_{\text{eff}} \) in the density range explored are accompanied by a difference in the metal input fluxes of about one order of magnitude. These observations confirm that \( Z_{\text{eff}} \) at low density is entirely determined by metal impurities.

**Code simulations**

From the data presented above it appears that the brightness of an inner shell ion e.g. CrXXII is an increasing function of the neutral flux of the same element from the limiter. To understand if this implies equal penetrations of the metals inside the discharge at different densities, we compared the experimental results with the predictions of a numerical code describing the transport of impurities with an anomalous flux \( \Gamma = -D \frac{\partial n_z}{\partial r} + v n_z \).
The numerical simulation of CrXXII brightness takes into account the experimental temperature and density profiles as well as the above reported input fluxes of neutral Chromium.

Typical anomalous diffusion coefficients $D = 0.4 \text{ m}^2/\text{s}$ and inward pinch velocities $v = -2 \text{ m/s}$ are initially assumed at all densities and the results are compared with the values of the CrXXII deduced by the VUV spectrometer. In this approximation it is found that the predicted brightnesses result compatible with the experimental values at low density. On the other side going toward high densities CrXXII decreases faster than expected by the code.

The experimental CrXXII dependence on $n_e$ can be reproduced by assuming a diffusion coefficient increasing with density (keeping $v/D$ ratio constant, it is necessary to vary $D$ by a factor $= 8$ when the density varies from $3.5 \times 10^{19}$ to $1 \times 10^{20} \text{ m}^{-3}$). The analysis above is not sufficient to distinguish between a real variation of the impurity transport coefficients and a screening of the metals entering the plasma by the scrape off layer. The experimental determination of the real impurity transport coefficients in FTU, using impurity injection, is presently under development.

The possibility that a significant contribution to the metal fluxes is coming from the wall does not seem to be compatible with the experimental results.

References

*ENEA Guest
HEAT FLUXES AND ENERGY BALANCE IN THE FTU MACHINE

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ABSTRACT

Thermal loads on the FTU limiter are routinely measured and energy losses via conduction/convection are inferred. A quite small fraction of the input power (4 to 8%) has been measured from mushrooms temperature increase. Numerical evaluation and comparison with thermocouples located at different radial positions in the S.O.L. suggest a long energy decay length $\lambda_E$. The power loads inferred from the estimated $\lambda_E$ in the actual geometry of the limiter and first wall lead to a global energy balance close to be satisfied.

I. INTRODUCTION

The FTU is a medium size, high field machine ($B_T$ up to 8T) presently operating with a main segmented poloidal limiter [1]. Figure 1 shows a schematic view of the limiter itself: it covers the poloidal section for 160° in the inner part and 80° in the outer one. The plasma facing part of the limiter structure is composed of 22 mushrooms. It is instrumented with 8 thermocouples poloidally distributed as shown in Fig. 1. These thermocouples are embedded in the limiter mushrooms being the hot junction at a distance of 0.4 cm from the tip of the mushroom head. Other 4 thermocouples are located on the limiter support directly facing the plasma (2 on the electron side, 2 on the ion). The distance of these 4 thermocouples from the limiter's head is 2 cm and the distance between the head of the limiter and the first wall is 3.3 cm.

Fig. 1 - FTU limiter
Measurements of energy losses via conduction/convection channel are now routinely carried out on the FTU machine. The evaluation is made by measuring the thermocouple increase as described in the next section. A quite small fraction of the input power ranging from 4 to 8% is usually detected suggesting the presence of high radiation losses and/or a long energy decay length that let the plasma particles reach the limiter support.

Since last plasma campaign also bolometric measurements have been available confirming high radiation losses in FTU, ranging from 50 to 75% of the input power; nevertheless the missing power in the energy balance is still too large.

All this indications pushed us in evaluating the energy decay length at the limiter for a set of plasma wall interaction dedicated shots.

II. RESULTS AND DISCUSSION

Magnetic surfaces in FTU are usually slightly elongated (typically k=1.05) in order to allow for a better magnetic diagnostic. For minimizing the asymmetries when characterizing SOL parameters, during three shot days dedicated to the plasma-wall interaction, great attention was devoted to get circular magnetic surfaces wetting a limiter surface as large as possible.

The energy collected by the limiter mushrooms was evaluated in an adiabatic way by measuring the temperature rise after 25 s from the shot start (it lasts 1.4 s), being the mushroom heads nearly thermalized after this time. Each mushroom temperature was assumed to be representative of the temperature of the neighbouring mushrooms.

The so calculated fraction of the total input energy that reaches the mushrooms is plotted in Fig. 2 as a function of the line average density measured along the chord passing through the magnetic axis of the machine. Its value is very low (<10%) and increases with density for all the currents considered; a sudden decrease of this energy fraction can be observed only at the lowest current in correspondence with the attainment of the critical $n_u/I_p$ value needed for the onset of a mafes, as confirmed by the enhancement of the $H_u$ signals. The maximum power

![Fig. 2 - Energy fraction lost to limiter mushrooms vs. line average density ($\Delta I_p = 400$ kA; $\circ I_p = 600$ kA; + $I_p = 800$ kA)](image)

![Fig. 3 - Energy decay length vs. line average density ($\Delta I_p = 400$ kA; $\circ I_p = 600$ kA; + $I_p = 800$ kA)](image)
load in these discharges amounts to 3-5 MW/m², which rules out the material evaporation as an impurity source.

Given the very low value of the measured energy flux to the mushroom heads, a careful evaluation of the energy decay length seemed to be suitable especially considering that the thermocouples directly facing plasma at 2 cm from the LCMS recorded a temperature increase much higher than the expected one on the basis of the \( \lambda_E \) values extrapolated from the FT results [2,3]. Moreover the large \( \lambda_n \) and \( \lambda_T \) that can be inferred from the Langmuir probe measurements [4] suggested that in FTU the energy decay length is rather large. In order to get \( \lambda_E \) values in a self-consistent way starting from the measured heat fluxes, numerical simulations were made by using a 2D finite elements heat transfer code (PDE—PROTRAN). A power density \( P(r) = P_0 \times \exp(-r/\lambda_E) \), where \( r \) is the radial distance from the LCMS, was allowed to flow on a toroidal cross-section of the mushroom for a time step equal to the shot duration. Results showed that the rate of the temperature increase at the site where the thermocouples are embedded depends only on the spatial profile of the power flux, namely on \( \lambda_E \). In practice an analytical relation between the time elapsed before the temperature increment reached an half of its maximum value and the quantity \( \lambda_E \) was established. A careful analysis was carried out only for the two smallest mushrooms in Fig. 1, being the others not completely thermalized within the acquisition time. The corresponding \( \lambda_E \) values are reported in Fig. 3. Nevertheless from the \( \Delta T \) vs time curves of all the other mushrooms at least as long \( \lambda_E \) values are to be expected. The \( \lambda_E \)'s in Fig. 3 were therefore assumed as poloidal average values.

The validity of this computer simulation was checked by comparing the energy flux at 2 cm from the LCMS, as resulting from the calculated \( \lambda_E \) and the experimental heat load on the mushroom heads, and the flux measured by the thermocouples facing the plasma. A satisfactory agreement was found between the calculated and the measured values.

By using the \( \lambda_E \) values of Fig. 3, the energy deposited on the mushroom heads (thickness \( d_m = 1 \) cm) was corrected by the factor \( 1/(1-\exp(-d_m/\lambda_E)) \) in order to get the actual energy lost by conduction-convection in the scrape-off layer.

In this way the fraction \( P_{cc}/P_{oh} \) of the total input power lost by non radiative mechanisms amounts to the 20-35%, a value close to the one inferred for FT [5]. On the basis of these corrected values of \( P_{cc}/P_{oh} \) a global energy balance for the examined discharges was temptatively outlined by evaluating the amount of radiated energy. Since January 1993 bolometric data are routinely available in FTU [6]. The bolometer consists of a vertical array of 16 golden foils looking at the lower half of the plasma cross-section at the same toroidal position (60° from the limiter). The ratio \( P_{rad}/P_{oh} \) between the total radiated power and the total input power as a function of the line average density is plotted in Fig. 4. A trend with density complementary to the one of the conduction-convection losses can be observed. The decrease of radiation losses starts at higher density higher is the plasma current, in agreement with the trend of \( Z_{eff} \) [7]. This decrease at high density could be explained with an enhanced screening efficiency of the impurities by the SOL [8] or with a decrease of the edge ion temperature, and therefore of the metallic impurity production by sputtering, as predicted by a 2D multifluid code for impurity production and retention in FTU SOL [9].

Finally the ratios \( P_{cc}/P_{oh} \), \( P_{rad}/P_{oh} \) and \( (P_{cc}+P_{rad})/P_{oh} \) were plotted together in Fig. 5. The "missing" power in the energy balance ranges from 5 to 20%.

III. CONCLUSION

The energy balance for a set of FTU discharges was tentatively outlined, based on the measurement of the limiter temperature increase by thermocouples and of the radiated power by a bolometric array.
The missing power in the energy balance was found to be 5-20% of the total ohmic input power.

Energy decay lengths as long as 3-4.5 cm were inferred from the temperature increase rate of the mushroom heads by using a finite elements heat conduction code and confirmed by the thermocouple measurements at 2 cm from the LCMS.

Given the amount of the energy deposited on the limiter, material evaporation can be ruled out as an impurity source during well centered steady state discharges.

**FOOTNOTE AND REFERENCES**

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SCALING LAWS AND POLOIDAL ASYMMETRIES
IN THE SCRAPE OFF LAYER OF FTU

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The scrape-off layer (SOL) plasma of FTU is routinely studied by means of 5 fast reciprocating probes, located in the poloidal plane as sketched in Fig. 1. Each of them bears 4 electrodes acting as 4 single Langmuir probes, and is able to span the entire SOL in less than 160 msec. The whole system (electrical + mechanical part) is remotely controlled via computer.

The first aim of this diagnostics is to build a solid data base in order to study the scaling of the main quantities relevant for the SOL, with the core plasma parameters and to characterise possible poloidal asymmetries.

The high line averaged densities $n_e \leq 3 \times 10^{20} \text{m}^{-3}$, and the reactor relevant power fluxes across the last closed magnetic surface LCMS (0.1 MW/m² in ohmic phase and up to 0.4 in the future lower hybrid phase) that are achievable in FTU, strongly support this task.

Local values of electron density $n_e$, temperature $T_e$, and plasma potential $V_p$ are calculated from a full non linear fit of the current-voltage (I-V) characteristics for each electrode. Up to 4 radial profiles are so obtained in a single shot, that overlap very satisfactorily in case of stationary plasmas, due to the particular care spent in the arrangement of both the mechanical and electrical part of the equipment.

The most important quantities, which are normally recovered from such profiles are the decay lengths of density and temperature $\lambda_n, \lambda_T$, and their values at LCMS, $n_{e,LCMS}, T_{e,LCMS}$. However, because of the actual shape of the LCMS and of the complex geometry of the material obstacles within the vacuum vessel, it is not straightforward to obtain the above quantities. A code is used that matches the experimental profile $n_e(x)$ with the following formula:

$$ n_e(x) = n_{e,LCMS} e^{-\int d\xi/\lambda(x)} $$

The integration is performed along the actual path of the electrode, starting from LCMS up to the generic coordinate $x$, and $\lambda(x)$ is the local value of the e-folding decay length, given as usual by $\lambda(x) = \frac{1}{f} \left[ \frac{D_\perp \cdot L_{\parallel}(x)}{(2 \cdot c_s) (\lambda(x)/2 \cdot D_\perp)} \right]^{1/2}$. Here $f$ is a factor to take into account the poloidally varying radial separation between flux surfaces [1]. $D_\perp$ is the perpendicular particle diffusion coefficient, $c_s$ is the ion acoustic speed at LCMS. $L_{\parallel}(x)$, the length of the local flux tube, is computed through a full tracing of the actual magnetic field lines up until they strike the first material obstacle. All such detailed calculations have proved many times to be necessary in order to avoid substantial inaccuracies. The outputs of the code are therefore the two free parameters $n_{e,LCMS}$ and $D_\perp$.

A similar procedure is applied to $T_e(x)$. 
RESULTS

Our database is limited at present to ohmic operation and stationary plasmas, within the following ranges: $0.3 \leq \bar{n}_e \leq 1.8 \times 10^{20} \text{ m}^{-3}$; $0.4 \leq I_p \leq 0.6 \text{ MA}$; $B_T=4$ and $6 \text{ T}$; $3.15 \leq q_\psi \leq 7.7$ (with $q_\psi$ the safety factor).

D. Scaling

Scaling laws have been searched for $n_{e,LCMS}$, $T_{e,LCMS}$, $D_L$ (or $\lambda_n$). The last two quantities have, however, been averaged over the poloidal angle $\theta$, since they result strongly varying functions of $\theta$, while density is almost constant over the sampled LCMS, as described further on in the paper.

A regression analysis has been applied to all the data, taking as independent variables $\overline{n}_e$, $q_\psi$, and the density peaking factor $p=n_0/\bar{n}_e$, being $n_0$ the peak plasma density. A high degree of correlation has been found only for the LCMS density. This is shown in fig. 2, where $n_{e,LCMS}$ is plotted versus the regression variable $x$, together with the old data of FT [2]. For FTU it is $n_{e,LCMS} \propto (\bar{n}_e/p)^{1.46} \times q_\psi^{0.98}$. The relevance of $p$, suggested by simple models [3], has been tested directly on experimental data with $1.26 \leq p \leq 1.7$, but the same $q_\psi$.

The scaling is very similar to that of FT, whose data therefore present a good correlation also in fig. 2 resulting a factor $s=1.92$ greater than those of FTU. All the major features of the scaling are thus confirmed: the almost linear increase of $n_{e,LCMS}$ with $q_\psi$ (or with $1/I_p$), even now that the inward pinch velocity is explicitly considered through $p$ [3], and the power dependence $\bar{n}_e^\alpha$. We however would expect $\alpha >1.46$ and closer to 2, which is more appropriate when ionization is unimportant within the SOL [3], and when $\lambda_n$ does not show appreciable variations with $\bar{n}_e$ [3]. Both these requirements are satisfied for FTU, as already for FT [2], over more than one order of magnitude of LCMS density variation. The clearer constancy now of the decay lengths constitutes a further support to neglecting SOL ionization. A possible explanation of this value of $\alpha$ could come from the fact that the velocity of the refuelling neutral atoms may increase with $\bar{n}_e$, as found for JET with the neutral code NIMBUS [4].

The factor $s=1.92$, quoted above, within the errors is the same as given by the law $n_{e,LCMS} \propto 1/(R^2-a)$ (R,a=plasma major and minor radius), which approximates several accepted scalings of the energy confinement time $\tau_E$.

The variation of $T_{e,LCMS}$ is, instead, not very clear, as noted before, except for a very mild increase with the total input power into the SOL $P_{SOL}$. It is striking the negligible effect that the LCMS density has on $T_{e,LCMS}$, differently from what predicted and observed in low density tokamaks as JET. Similar considerations hold also for FT and Alcator C. We are at present investigating whether the ion temperature $T_i$, usually assumed $\approx T_e$, can instead play an important role.

Preliminary computer simulations with a bidimensional multifluid model show indeed that $T_i$ can vary [5] from $T_i=5T_e$ at $\bar{n}_e=4 \times 10^{19} \text{ m}^{-3}$ to $T_i=2T_e$ at $\bar{n}_e=1.5 \times 10^{20} \text{ m}^{-3}$. A poor coupling between ions and electrons can be expected also from the simple criterion for the energy equiparition: $L_{she}/n_e/T \geq 10^{17} \text{ m}^{-2}/\text{ev}$, which would give FTU mostly in a marginal equiparition regime. The analysis of the poloidal asymmetries of $T_{e,LCMS}$ discussed later further supports this view.

The transport properties, after averaging, are almost constant for the whole present data base, as pointed out before, but show considerable poloidal asymmetries which are described in the next paragraph.

II. Poloidal asymmetries

As already said, no asymmetry is evident in the LCMS density.
In fig. 3a $T_{e,LCMS}$ is plotted versus $\theta$ for $B_T=6$ T, $n_{e,LCMS}=10^{19}$ m$^{-3}$, but for different values of $I_p$ as there reported. In fig 3b instead, $D_\perp$ is shown together with the estimated Bohm value multiplied by 10, for $\bar{n}_e=0.7 \times 10^{20}$ m$^{-3}$, $I_p=600$ kA, $B_T=6$ T.

The reason and the structure of such asymmetries are quite unclear at present. The explanations offered by fluid models cannot account at the same time for the temperature variation and for the density uniformity. Pure geometrical effects [2,3], linked to the complex structure of FTU SOL flux tubes, fail in describe both the magnitude and the pattern of the measured poloidal variation of $T_{e,LCMS}$. This is true also in the hypothesis of asymmetric power input into SOL, since the flux tubes connected to the probes should average it over large poloidal angles, typically 70-200 degrees. Nevertheless, the only appreciable correlation found up to now is a fair direct link of $T_{e,LCMS}$ with $L_{\parallel}$, as shown in fig. 4, where the same data of fig 3a are plotted against the relative $L_{\parallel}$ values, calculated as described before. This behaviour can be qualitatively understood, reminding that FTU SOL is marginally collisional. In this case the length of the flux tube can play a crucial role in determining the amount of power transferred from ions to electrons, and hence in establishing $T_e$.

Turning now to $D_\perp$, fig. 3b, we must note how large it is, compared to Bohm's, especially in the inner side of the poloidal plane. This means very long decay lengths, $\lambda_{\parallel}>3$ cm have been actually measured, which in turn imply high plasma density and possible power deposition directly on the vessel walls. Such high values for $\lambda_{\parallel}$ agree with those inferred from thermocouple measurements of thermal load on the FTU limiter [6]. Besides, they can become a very important element in the puzzling question of the missing power, which so often appears when the balance between the input and output power is attempted in a tokamak.

High values for $D_\perp$ are not new for high field tokamaks as well as the asymmetry along the poloidal plane, as the measurements made in the past on Alcator C confirm, even if these show a different poloidal pattern.

**CONCLUSIONS**

The Langmuir probe system in FTU has shown that density at LCMS can be considered uniform on the poloidal plane and scales approximately as $(\bar{n}_e/p)^{1.46} \times r^{0.98}/(R^2 \times a)$ (with the data of FT included), over more than one order of magnitude for $n_{e,LCMS}$.

Temperature at LCMS, instead, is almost independent of many macroscopic plasma parameters, particularly density, but shows a quite complex poloidal structure. Model analysis and some experimental observations, as the effect of the length of the flux tubes on $T_{e,LCMS}$, suggest that the role of the ion temperature may be essential in this issue.

The particle diffusion coefficient results much higher than expected, up to 30 times the Bohm's value, implying very long SOL decay lengths. Its poloidal variation, however, is not yet understood.

**REFERENCES**

[6] M. Ciotti et al 'Thermal load on the FTU limiter', This Conference
Fig 1 - Poloidal cross section of FTU with the location of the 5 fast Langmuir probes.

Fig 2 - Density at LCMS plotted versus the indicated the regression variable, FT and FTU, several $I_p$ and $B_T$ (see text).

Fig 3a - Poloidal variation of $T_e$ at LCMS for different $I_p$ values, $B_T=6T$, (i.e. different $q_\psi$), but same $n_e$ at LCMS.

Fig 3b - Poloidal variation of the particle diffusion coefficient for $\bar{n}_e=0.7*10^{20} \text{ m}^{-3}$, $I_p=600 \text{ kA}$, $B_T=6 \text{T}$.

Fig 4 - Temperature at LCMS plotted vs. the length of the flux tube connected to the probe. Same data as fig. 3a. The differences in L// come from the different poloidal angles and $q$'s.
ACTIVE PROBING OF PLASMA EDGE TURBULENCE AND FEEDBACK STUDIES ON THE TEXAS EXPERIMENTAL TOKAMAK (TEXT)

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Introduction

The edge fluctuations play a critical role in the overall tokamak confinement.¹ Experiments on TEXT show that electrostatic fluctuations in the edge plasma are the dominant mechanism for energy and particle transport.² The basic mechanisms responsible for the edge turbulence are the subject of ongoing research in fusion devices. To understand the driving forces responsible for edge fluctuations, a novel experiment is underway on TEXT to actively modify the turbulence at the plasma edge by launching waves using electrostatic probes in the shadow of the limiter. This technique permits active probing of the spectral properties of the edge turbulence. This new approach to the study of edge fluctuations can provide more insight into the basic dynamics of the turbulence and may, in turn, enable detailed comparison with the theory. These experiments, which rely on the use of oscillating electric fields at the plasma edge, complement edge fluctuation control studies that are presently limited to the use of applied dc biasing to influence the edge electric field profile.³ These experiments have been extended to control of the edge plasma fluctuation level, using feedback to explore its effects on the edge turbulence characteristics as well as on confinement.

Experimental Arrangement and Diagnostics

The experiments are carried out with a wave launching system consisting of two Langmuir probes (L₁, L₂) which are separated by \( d = \lambda / 2 \sim 1.8 \text{ cm} \), where \( \lambda \) is the wavelength of the electrostatic edge fluctuations, in the poloidal direction with respect to the toroidal magnetic field, \( B \). The L₁, L₂ are operated in the electron side of the (I, V) characteristic. Each probe tip is fed separately by independent ac power supplies capable of providing up to 1.5 kW of power in the frequency range of 9 to 250 kHz. The power sources are driven by a signal generator through a phase shifter, which allows control of the ac phase difference \( \Delta \phi \) between the L₁ and L₂, and a band pass filter (BPF), as shown schematically in Fig. 1.

![Fig. 1. Schematic of the fast reciprocating active probe (FRAP).](image)

The L₁, L₂ are designed to handle an ac probe current of up to \( I_{ac} \sim 15 \text{ A} \), which corresponds to \( \sim 30\% \) of the estimated total fluctuation current within the correlation volume of edge plasma that has relative density fluctuations of \( \tilde{n}/n \sim 20\% \) at typical averaged frequency of \( f = \omega / 2 \pi \sim 50 \text{ kHz} \).² Besides these wave-launching (or exciting) tips there are two small probe tips (S₁, S₂) separated by \( d / 2 \) placed on the same probe head (see Fig. 1) to measure the local plasma floating potential, \( \Phi_0 \). One of these tips, S₁, is utilized for feedback experiments to provide the input signal for driving the L₁ and L₂. This launcher system is called the fast reciprocating active probe (FRAP) because of its fast plunging action into plasma, which takes about 50 ms for a 5-cm stroke, to reduce the heat load on the probes during the discharge. The specific edge fluctuation diagnostics used for these experiments are a fast reciprocating Langmuir probe.
(FRLP) array\textsuperscript{4}, used as sensing probe, located halfway around the torus from FRAP, separated by $\sim 157^\circ$ toroidally, and two sets of fast H\textsubscript{\alpha} radiation measuring arrays.\textsuperscript{5}

**Experimental Observations**

**a. Active probing of edge turbulence**

The series of active probing experiments are carried out by launching waves from FRAP in ohmically heated plasmas with a flat top of $\sim 300$ ms in hydrogenic discharges. The toroidal magnetic field is $B \sim 2.1$ T; the average plasma density is $n_0 \sim 3 \times 10^{13}$ cm$^{-3}$. The rail limiters (top, bottom, and outside) are located at $r_a = 27$ cm, while FRAP is at $r = 27.5$ cm on the machine top. The following experiments are performed with FRAP ac current $I_{ac} \sim 5$--8 A in the frequency range of 15 to 50 kHz with broadband BPF settings. Measurements of potential fluctuations $\phi_f$ from FRLP indicate that the excited waves are received by FRLP, which is located $r = 27.5$ cm at the bottom of the torus. For example, in Fig. 2, the FFT of $f = 30$ kHz signal launched from FRAP is shown together with the fast Fourier transform (FFT) of the received signal $\overline{\phi}_r$, which is about 25 dB above the background fluctuations level. This experiment is performed with a plasma current of $I_p = 180$ kA corresponding to an edge safety factor of $q = 4.3$ at $r = 27.5$ cm. Earlier experiments\textsuperscript{6} indicated that for $q = 4.3$ FRAP and FRLP have measured the highest turbulence coherence, and at the same time the magnetic field line (FL) plots show that these probes and one of the H\textsubscript{\alpha} arrays (located at P1) are all on the same magnetic flux tube, while the second H\textsubscript{\alpha} array (located at P14) is not. For this experiment the measured traveling time of the wave is $\tau_d \sim 0.2$ ms from FRAP to FRLP. Given the distance along the field line, $L \sim 12$ m, it is estimated that these waves have a speed of $v_2 = L_2/\tau_d \sim 5 \times 10^6$ cm/s which is about the ion sound speed $c_s$ for $T_e \sim 25$ eV edge plasma. The detected signal strength of $\overline{\phi}_r$ weakly depends on the frequency of the wave, the plasma current, and the phasing of the applied ac signal between $L_1$ and $L_2$. It is observed that for $f = 15$-25 kHz the amplitude of $\overline{\phi}_r$ is slightly higher, by about factor of 1.5, than the rest of the frequencies used (15 to 50 kHz) during these experiments. The effect of $I_p$ on $\overline{\phi}_r$ measured at the launching frequency of $\sim 30$ kHz is shown in Fig. 3(a), and the intensity of the fluctuating H\textsubscript{\alpha} radiation, $I_{H\alpha}$, from the two arrays (port P1 is on the same FL with FRAP for $q \sim 4.3$, but port P14 is not) is plotted.

![Fig. 2](image-url)  
**Fig. 2.** (a) The 30 kHz injected signal from FRAP, $r = 27.5$ cm, (b) received by the sensing probe FRLP, $r = 27.5$ cm, located halfway around the torus from FRAP.

![Fig. 3](image-url)  
**Fig. 3.** (a) The effect of plasma current $I_p$ on $\overline{\phi}_r$ measured at the launching frequency of $\sim 30$ kHz is shown. (b) The intensity of the fluctuating H\textsubscript{\alpha} radiation, $I_{H\alpha}$, from the two arrays (port P1 is on the same FL with FRAP for $q \sim 4.3$, but port P14 is not) is plotted.
from the two arrays is plotted in Fig. 3(b). The $I_\phi$ from both arrays has similar dependence on the plasma current. This observation may suggest excitations of these waves on the flux surface rather than on the flux tube. The radial extent of the wave is also measured by scanning FRLP. The signal strength of $\phi_r$ from FRLP is shown in Fig. 4 as a function of the radial position of FRLP. It is observed that the excited wave penetrates into the core plasma even though it is being launched from $\sim$0.5 cm behind the limiter, and $\phi_r$ has $\sim$2 cm of radial width. This result indicates that by actively perturbing the edge plasma fluctuations from the limiter shadow, it may be possible to influence the core plasma characteristics inside the limiter. The ac phase shift $\Delta \phi_{12}$ between $L_1$ and $L_2$ is also varied, and the result is given in Fig. 5. The $\phi_l$ measured with FRLP is about a factor of two higher when $\Delta \phi_{12} = 0$. At the same time, modifications to the frequency spectrum of the poloidal wavenumber $k_\theta$, Fig. 6(a), the fluctuations of the edge density, Fig. 6(b), and the plasma potential, Fig. 6(c), are observed at the launching frequency of $f \sim 30$ kHz. For example, at $f \sim 30$ kHz, $\bar{n}(ac\ on)/\bar{n}(ac\ off) \sim 2$, while $\bar{\phi}(ac\ on)/\bar{\phi}(ac\ off) \sim k_\theta(\phi\ (ac\ off)/k_\theta(\phi\ (ac\ on)) \sim 5$, which is consistent with the predictions of the simple mixing length theory.7

b. Feedback experiments

When the launcher is driven by the floating potential fluctuations, sensed by $S_1$, and used as an input to the system through the BPF = 10–30 kHz at the location of FRAP, the edge fluctuations can be suppressed ($\leq 40$ kHz), without enhancing other modes, or excited ($\sim 0$ kHz), depending on $\Delta \phi_{12}$, both locally and at the downstream sensing probe, FRLP. This feedback arrangement is similar to the ones used on earlier experiments.8 The results are shown in Fig. 7(a), obtained from the FRAP sensing tip $S_2$, and Fig. 7(b), obtained from FRLP. These measurements indicate that the feedback affects the fluctuations not only locally but also halfway around the torus. The estimated fluctuation-induced radial particle flux $\Gamma$ also varies with $\Delta \phi_{12}$. For example, $\Gamma$ is $\sim$20% higher without the feedback when $\Delta \phi_{12} = 0$, but it becomes $\sim$20% lower when $\Delta \phi_{12} = \pi/2$. The global core plasma parameters have not been affected by the feedback except for slight variations on the edge $n_e$ and $T_e$.

Discussion

These preliminary observations have successfully demonstrated the feasibility of exciting waves at the plasma edge to actively probe the spectral properties of the edge turbulence. The initial feedback trials are also encouraging for controlling edge turbulence. Using these initial observations detailed experiments are planned at various feedback gain and phase shift settings for the upcoming TEXT-U with additional sensing probes located at various locations around the torus. The poloidal extent of the feedback excitations will also be investigated. Meanwhile, detailed data interpretation and modelling studies are underway.

Acknowledgements

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References

8. For example, see D. P. Dixon et al., Plasma Phys. 20, 225 (1978).
Fig. 4. The signal strength of $\Phi_f$ from FRLP is shown as a function of the radial position of FRLP.

Fig. 5. The effects ac phase shift $\Delta \phi_{12}$ between the $L_1$ and $L_2$ on the signal strength of $\Phi_f$ from FRLP is shown.

Fig. 6. (a) Spectrum of the wavenumber; (b) spectrum of the normalized density fluctuations; (c) spectrum of the plasma potential fluctuations normalized to $T_e$ are shown for zero phase shift, $\Delta \phi_{12} = 0$.

Fig. 7. Effects of feedback on the potential fluctuations spectrum measured at (a) FRAP and also at (b) FRLP for the phase shift settings of $\Delta \phi_{12} = 0$ and $\pi$ compared to no feedback case (BG) are shown.
Plasma Conditions Under Various Divertor Biasing Configurations on TdeV*

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Introduction
The closed divertors of TdeV feature eight independent neutralization plates electrically insulated from the vacuum vessel assembly. The fact that the inner and outer plates can be biased separately in a single-null or double-null geometry makes possible the study of a wide variety of biasing schemes. Current injection, in which the potential difference is applied on the same flux surface, and plasma biasing, in which the potential difference appears between the separatrix and the wall, are pure biasing modes of particular interest. Either of these modes is selected by floating or referencing the power supply with respect to the vacuum chamber. Current injection is shown to drive parallel electric fields with no significant change in the plasma potential or fundamental properties. The plasma biasing mode is observed to set up a radial electric field and current with major effects on both the edge and central plasma.

Current injection.
The current injection mode is characterized by a current flowing in the SOL from one divertor plate to another in a direction parallel to the magnetic field lines with no perpendicular components. By floating the external voltage source with respect to the vacuum chamber no radial current or E_r field can be sustained. In TdeV, current injection can be achieved in a single-null or a double-null magnetic configuration. The latter configuration is schematized in Fig.1. The external power supply is connected to the upper and lower outboard divertor plates with no referencing to the grounded vessel wall. In general, the injection of a current of either polarity, up to the saturation limit, has almost no influence on the SOL and central plasma basic parameters. Particularly, the plasma potential in the SOL is practically left unaffected and the plate potentials adjust themselves consequently. This is explained by the fact that the SOL plasma is connected transversally to the vessel wall at ground potential and that no perpendicular current can be drawn by the floating voltage source to pull the separatrix potential up or down. In keeping with the properties of sheaths, a greater fraction of the voltage drop is found at the negative plate as it is forced to ion saturation. The I-V characteristic of the upper outer plates is presented in Fig.2 and the I-V curve of the combined upper and lower plates is given in Fig.3a. (I_p=210 kA and n_e=2.5 \times 10^{19} m^{-3}). This latter curve resembles the characteristic curve of a double electrostatic probe, the two divertor plates being considered as the plasma electrodes. It features a central linear section whose slope is given by the electron temperature, by a current offset at zero voltage due to thermoelectricity and by a current saturation at large potential differences limited by the collection of ions at the negative electrode. To analyze our data we use a simple model based on linear probe theory and constant pressure fluid.
transport. The resulting fit to our data is presented in Fig. 2 as a dotted line. Although the model recovers most of the characteristics of the measurements, it is however obvious that some of the data are not well simulated particularly in the knee region. Since the shape of the curve before saturation is a signature of the energy distribution of the electrons repelled by the plate potential, we postulate that the disagreement between the model and the data could be caused by the presence of hotter electrons. In fact, the mean free path of 60-eV electrons in the SOL is almost equal to the connection length, so that a fraction of high energy electrons leaving the central plasma can reach the sheath area unthermalized (Fig. 3a). A small amount of these uncoupled electrons can change substantially the collected electron current characteristics. A modification of the model to include two Maxwellian distributions, a cold population and a hot one, leads to the solid lines in Fig. 2 and Fig. 3. The new fit shows that a hot population representing \( n = 3.5\% \) of the total electrons is sufficient to explain the whole set of data.

A similar experiment conducted in a more collisional plasma (\( n_0 = 4.5 \times 10^{19} \text{ m}^{-3} \)), shows that the fraction necessary to fit the data is reduced to 0.08% (see Fig. 3b).

**Plasma biasing**

The other fundamental mode, called plasma biasing, consists in biasing the separatrix with respect to the wall. Figure 4 illustrates this biasing configuration for a double-null geometry. The external voltage source is applied between the outboard (upper and lower) divertor plates and the vessel. In this way, a radial current is forced to flow from the separatrix in close contact with the plates and the outer tenuous SOL plasma connected to the wall. We have evidence from instrumented tiles located at different positions inside and outside the divertor, that the return current does not come from the interior wall of the divertor but from the main chamber wall near the throat entrance. In addition, it is observed that the floating tiles insulating the field lines lying between the plates and the throat entrance reach a potential of...
approximately half the applied voltage, indicating that a perpendicular electric field is in fact produced near the plates. The presence of a radial electric field is also confirmed in the equatorial plane by Langmuir probe measurements.

This bias induced radial field and its associated current are the driving source of all the biasing effects observed in TdeV. Induced SOL plasma rotation, edge density profile modification, change in gas fuelling and recycling and pressure build up or decrease in the divertor chamber are examples of plasma biasing phenomena\cite{21}. Effects are also observed inside the separatrix. In fact, the poloidal and toroidal plasma rotation and the impurity level in the core are seen to be affected by the application of a biasing voltage.

The I-V characteristic of the plasma biasing mode is plotted in Fig.5 \((I_p=210 \text{ kA and } n_p=2.5 \times 10^{19} \text{ m}^{-3})\). It features a linear impedance of \(-0.6 \Omega\) with a rolloff at positive voltages above 100 V, presumably due to ion saturation at the wall electrode. Over the available range of applied negative biasing voltages, the measurements show no sign of saturation. The I-V curve also shows a positive current offset at \(V=0\), corresponding to the expected thermoelectric current flowing from the hot plate to the cold wall. To interpret our data, we have developed a simple 1-D model including non ambipolar radial current and sheath controlled parallel current. Presented in Fig.6, the geometry consists of a straitened SOL with a parallel coordinate, \(z\), along the magnetic field line and a perpendicular coordinate, \(r\), in the radial direction. The SOL flux lines are bounded by the divertor plates near the separatrix \((0<r<r_s)\), insulated by floating tiles between \(r_a\) and \(r_b\), and finally connected to the wall at a larger radius. The model is simply based on current continuity and assumes that the edge density, plasma potential and perpendicular current are toroidally symmetrical \((\partial (n_e, \mu, \phi, T_e)/\partial z=0)\). For the computation, we use exponential profiles for the density, the temperature and the mobility. The calculated plate and the radial current density profiles are presented in Fig.6 for \(V_{bias}=150\ V\). It should be noticed that, as confirmed by the flush mounted current probes installed in the divertor plates, the plate current is not centred about the separatrix as in the case of current injection but extends to the plate edge. The calculated plasma potential profiles are presented in Fig.7 for three biasing voltages and

![Figure 5 I-V characteristics in the plasma biasing mode.](image)

![Figure 6 Model for the generation of the E_r field.](image)
compared to the Langmuir probe measurements in the equatorial plane. For these calculations, the SOL temperatures are adjusted to the measured values. Although the model prediction and the probe data differ slightly at $V_{bias}=+150$ V, the general agreement is satisfactory and indicates that the essential physics is already included. It is believed that the 18-mm floating region between the plate edge and the wall is a critical element in the model since it allows more efficient E-field generation by not shorting the field lines as do the divertor plates and the walls. Calculations are also compared to the experimental I-V characteristic in Fig. 5. Using the measured edge density of $4 \times 10^{16}$ m$^{-3}$, the best fit to the data is obtained with a mobility of $0.02$ m$^2$/Vs. Using realistic plasma parameters, the fit reproduces remarkably well the major features of the data: the current offset, the linear response for negative biasing voltages and the saturation observed at large positive voltages. In addition, the mobility used in the fit is in very good agreement with the value of $0.03$ m$^2$/Vs deduced from modelling the edge density profile modification during biasing[3].

**Discussion**

We have identified and studied two basic modes of divertor plate biasing in TdC: first, a current injection mode consisting essentially of parallel current and electric fields with no radial components, ensured by floating the voltage source, and, second, a plasma biasing mode in which primarily a radial current and electric field are set up by electrically biasing the divertor plates with respect to the vessel wall. We believe that all the other biasing modes can be decomposed into these two modes. As an example, we have simulated the DIII-D biasing configuration by grounding both the inboard plates and the vessel wall and biasing the outer plates. When a negative biasing voltage is applied, the outboard plates, being negative, develop an electrostatic sheath with a voltage drop almost equal to the applied voltage. In this situation, the SOL potential remains unaffected as in the case of current injection. In the case of positive biasing, the situation is reversed and the plasma potential follows the outboard plate potential, the major potential drop appearing across the sheath at the more negative grounded electrode. The latter case resembles the plasma biasing mode, since the separatrix is biased with respect to the wall, and develops a radial electric field with all the associated effects. During these experiments we have used the spectroscopically determined toroidal plasma rotation as a test diagnostic. For negative biasing, no rotation is observed while, for positive biasing, a toroidal rotation is set up with the amplitude and direction usually seen in the plasma biasing mode. We conclude that the DIII-D biasing configuration leads to a hybrid mode, the negative polarity resembling the current injection mode and the positive polarity giving rise to a mode similar to plasma biasing.

Understanding the physics of divertor plate biasing will allow the study of a multitude of interesting hybrid modes by decomposing them into fundamental biasing modes with well identified characteristics. Effort in improving modelling of the two basic modes will continue.


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VARIATION OF DIVERTOR PLASMA PARAMETERS WITH DIVERTOR DEPTH FOR H–MODE DISCHARGES IN DIII–D*


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We report here the results of experiments aimed at quantifying the advantages of increasing the X–point to target-plate distance in a divertor tokamak operating with H–mode confinement. Larger distances should lower the peak electron temperature at the target plates, thereby reducing sputtering and lowering the impurity concentration in the core plasma. When gas puffing is used to reduce the divertor heat flux, extra field-line length may increase the volume available for radiation and increase gas isolation between the core and divertor regions.

These experiments were carried out using a lower single-null open divertor configuration ($L_p = 1.4$ MA, $B_T = 2.1$ T) with neutral beam heating ($P_{NBH} = 4.8$ and 6.8 MW) to produce ELMing H–mode discharges lasting about 3 s. The X–point height ($z_g$) was varied from 1.5–32 cm above the target plates by changing the plasma elongation on a shot by shot basis; the X–point radius was also varied in order to keep the outer strike point aligned with divertor Langmuir probe tips, as in Fig. 1. Though there was no gas fueling during the H–mode phase of the discharge, the plasma density remained constant for all $z_g$ obtained. Additional $D_2$ gas puffing for radiative divertor experiments was applied for the last 1.5 s of the H–mode period.

The poloidal distance ($L_{pol}$) from the X–point to target plates along the flux surfaces was varied over the range 4 to 40 cm (corresponding to $L_{pol}/a \leq 0.7$). This changed the parallel connection length between the target plate and the X–point from 2 to 8 m for a field line in the SOL 0.5 cm from the separatrix. The magnetic expansion of this flux tube going from the midplane to the target plates ranged from 13:1 to 4:1 for the lowest and highest X–point heights, respectively.

Core plasma parameters were affected very little by changing $z_g$. The energy confinement time remained nearly constant, as shown in Fig. 2(a), even though there was some variation in the ELM frequency, Fig. 2(b). More significantly, no systematic variation in $Z_{eff}$ was seen, Fig. 2(c). This result is supported by measurements of the 182 Å CVI line brightness produced by charge-exchange recombination with beam neutrals in the core of the plasma; the intensity does not change noticeably except for the very lowest X–point, where the plasma has nearly a limiter configuration anyway. This may simply mean that most of the carbon in the core plasma results from wall sputtering, as reported† for JET, and not from the divertor targets.

Upstream scrape-off layer conditions also did not change measurably, consistent with the fact that the total connection length in the SOL varied by only about 15% as the X–point was lowered ($z_{sep}$ went from 65 to 55 m). The density and temperature on the separatrix, as determined from Thomson scattering measurements, remained at $n_e,sep \approx 3–4 \times 10^{18}$ m$^{-3}$ and $T_e,sep \approx 200$ eV. The characteristic lengths also stayed about the same: $\lambda_n \approx 0.7$ cm and $\lambda_T \approx 0.6$ cm inside of 1 cm from the separatrix. In the far SOL (2 cm from the separatrix at the midplane) these lengths were considerably longer (several cm).

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§ Oak Ridge National Laboratory.
¶ University of California at Los Angeles.
Fig. 1. Cross section of lower half of DIII–D showing range of X–point locations studied. 1.0 cm and 0.5 cm midplane flux surfaces shown.

The divertor plasma conditions near the separatrix changed with X–point height as expected. Peak electron temperature at the outboard strike point decreased from about 35 to 15 eV as $L_{pol}$ increased from 3 to 40 cm, while the peak heat flux increased from 0.7 to 2 MW/m², as shown in Fig. 3. At the very lowest X–point heights (2 cm and below) the electron temperature increased sharply to about 90 eV; however, the geometry in this case was essentially that of a limiter plasma, where recycled neutrals are reionized on closed field lines in the core plasma. The radiated power fraction and the fraction of power reaching the target plates remained nearly constant except for the lowest $z_e$.

The measured change in peak divertor parameters agreed quite well with predictions based on the standard theory of the high recycling divertor as contained in the LEDGE 2-D fluid code. In these simulations we assumed standard ITER CDA values for the radial transport coefficients; $\chi_x = 2 m^2/s$ and $\chi_i = D_\perp = 1/3 \chi_x$, $n_{sep} = 0.3 n_e$ for the upstream density, and 40% for the fraction of heating power lost by radiative processes (half in the core and half in the SOL). The midplane temperature was adjusted to produce the correct power flow across the separatrix given the measured input power. This procedure is the same as was used for predicting conditions in the ITER divertor.

The model results for variation with power and X–point height are plotted alongside the experimental data in Fig. 3. The drop in temperature with increasing $z_e$ results from increased radial diffusion (mainly into the private region) and not from increased radiation or recycling (divertor density remained constant in the simulation). The change in peak heat flux largely follows the magnetic flux expansion. Though not shown, the biggest difference between model and experiment is that the code is predicting much broader SOL profiles than observed (both in the divertor and near the midplane) so that the total power reaching the target plates is higher than in the experiment. A better match to the profile width could be obtained by using smaller values for $\chi_{x,i}$ ($\sim 0.5 m^2/s$) as reported elsewhere, but then the calculated peak values would be much higher than measured. The narrow profiles in the experiment may be due to scrape-off on the inner wall of the tokamak.

In these experiments we added D₂ gas puffing late in the discharge to see how radiative divertor operation was affected by X–point position. Figure 4 shows the main results: maximum heat flux reduction was obtained over a broad range of $z_e$ values (7 to ≥22 cm), but the gas isolation steadily improved as the divertor length was increased, as inferred from the density rise produced by 1 s of gas puffing at 100 T-L/s. Very little heat flux reduction was observed at the lowest X–point position ($z_e \leq 2$ cm), though the radial profile did broaden. As reported in earlier work on radiative divertors on DIII–D, the largest reduction in divertor heat flux is obtained after a MARFE forms near the X–point. At either extreme of X–point...
height, MARFE formation was delayed until very near the end of the discharge. In all cases, gas injection lowered $Z_{\text{eff}}$ in the core plasma, especially after the MARFE was present.

These experiments have shown that, in an ELMing H-mode plasma, increased divertor depth does produce lower plasma temperatures at the divertor targets. However, it does not reduce the impurity content of the core plasma and comes at the expense of significantly higher peak divertor heat flux due to compression of the magnetic flux surfaces ($2.5\times$ increase in $\dot{q}_d$ for a $2\times$ reduction in $T_{\text{e,div}}$). Successful radiative divertor operation ($\dot{q}_d$ reduced $\sim 2\times$ with little change in $\tau_E$) was obtained over a broad range of X-point heights, though core fueling from the gas puff steadily increased as the divertor depth was reduced.

Fig. 4. Radiative divertor results. (a) peak divertor heat flux with and without gas puffing. (b) density rise resulting from gas puffing.
ACTIVE DENSITY CONTROL IN DIII-D H-MODE PLASMAS*


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Particle control in H-mode plasmas is of great importance to the viability of a tokamak fusion reactor. Particle control is essential for helium ash exhaust. ECH access, as well as increasing $T_e$ to enhance fast wave current drive efficiency which for $\frac{\alpha}{\rho_i}$ $\approx 1$ increases with $T_e$. In addition, density control in present devices would allow separate determination of the scaling of energy confinement with $I_p$ and $n_e$ in H-mode plasmas.

The enhanced energy confinement of the H-mode is invariably coupled to improved particle confinement and uncontrolled density rise. In DIII-D, following the H-mode transition, the line average plasma density rises to $n_e (10^{19} \text{m}^{-3}) \approx 6 \times I_p$ (MA). Previous attempts to control the H-mode density have not been successful. In DIII-D, H-mode density can be increased by no more than 20% by gas puffing before degradation of confinement and, ultimately, a transition back into the L-mode. Divertor biasing strongly compresses particles under the divertor baffle, however, in the absence of pumping, biasing reduces the line average H-mode density by no more than $\approx 10%$.

The previously accessible density range in the DIII-D H-mode plasmas has been too narrow to allow a determination of the scaling of confinement with $I_p$ and $n_e$ separately. Furthermore, the desired densities for planned rf heating and current drive experiments are mostly below this range. Therefore, we have installed a unique divertor system for density control as well as a number of other divertor studies. The system, shown in Fig. 1, and referred to as the Advanced Divertor, consists of a divertor baffle, an in-vessel cryopump ($D_2$ pumping speed $= 35000$ l/s, at a pressure of 2 m-Torr) and a bias electrode. Detailed descriptions of the Advanced Divertor subsystems, and the results of biasing experiments are described in detail elsewhere.1–3 In this paper we present the first results of divertor cryopumping experiments in the absence of biasing on DIII-D H-mode plasmas.

All plasmas described here were started with the Outer Divertor Strike Point (ODSP) initially away from the pump entrance aperture. After the plasma current ramp, the ODSP was moved to the vicinity of the aperture at $t \approx 1–2$ s, where divertor pumping was effective. Figure 2 shows time histories of the number of particles injected by gas puffing, stored in the plasma, injected by neutral beams, and pumped away by the divertor pump. The number of particles pumped is determined from the measured pumping speed and measured neutral pressure under the baffle. The difference between the source and sink terms is attributed to changes in the wall particle inventory.

During the plasma current ramp more than 100 Torr-l ($\approx 6.7 \times 10^{21}$ deuterons) is injected to achieve a plasma particle content of only $\approx 10$ Torr-l. Most of the particles are absorbed by

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the well conditioned graphite first wall. In contrast, the initial rate of density rise after the L–H transition is 3–4 times greater than the particle input rate from the external sources. In this phase, as a result of improved plasma particle confinement, an oversaturated wall relaxes to a new equilibrium with the plasma.

With the onset of ELMs, divertor pumping arrests the rate of density rise and ultimately lowers the particle content of the plasma as well as the vessel walls. By the end of the shot, the number of particles pumped away exceeds that due to the external sources by \( \approx 100 \text{ Torr}^{-1} \).

For the first divertor pumping shot, the excess number of particles pumped away is even a factor of two higher than this quantity, but within the following several shots it settles at this value. This observation indicates a gradual depletion of the wall particle inventory by divertor pumping. It should be noted that the particle removal rate during H–mode is \( \approx 80 \text{ Torr}^{-1}/s \), which is of the same magnitude as desired for a burning plasma.

The plasma parameters of two otherwise similar shots (\( I_p = 1 \text{ MA}, B_T = 2 \text{ T}, P_{\text{beam}} = 4.5 \text{ MW} \)), one with and the other without divertor pumping, are shown in Fig. 3. Divertor pumping does not significantly reduce the rate of density rise for the first 50 ms after the L–H transition. However, within two seconds from the onset of ELMs the line average electron density decays to a minimum of \( 3.3 \times 10^{13} \text{ cm}^{-3} \), 40% below that of the reference shot. As shown in Fig. 4, typically, late in the H–mode, the density profiles of pumped plasmas were slightly peaked. On the other hand, depending on the wall conditions, the density profiles of standard DIII–D H–mode plasmas are flat to slightly peaked. Concomitant with the density drop, there is a roughly proportionate rise in electron temperature; therefore, the plasma collisionality becomes a factor of two lower. Because of the reduced ion-electron energy exchange at lower collisionality and better ion energy confinement, the ion temperature rises more than the electron temperature. (Central \( T_i = 6.5 \text{ keV} \) vs. \( T_e = 3.5 \text{ keV} \)). This behavior is similar to the hot ion mode observed on JET and DIII–D.

In general the pumped shots were reproducible and reached quasi-steady state within 3 seconds after the L–H transition. However, when the NBI power was increased to 6.7 MW, locked modes grew when \( n_e \) fell to \( \approx 3.0 \times 10^{13} \text{ cm}^{-3} \), and confinement deteriorated. A similar locked mode density limit of \( 3.8 \times 10^{13} \text{ cm}^{-3} \) was encountered at 1.5 MA. These locked mode density limits are roughly twice the values reported earlier\(^5\) and are attributed to the higher \( \beta_N [\beta_N \equiv \beta \times B_T / (I_p \times a)] \) of the low density H–mode plasmas in the present experiments. A density feedback scheme, using a combination of position of the separatrix relative to the
entrance aperture and gas puffing was used successfully to regulate the plasma density above the locked mode boundary.

We have succeeded in reducing the electron density of H-mode plasmas by up to 40%, which results in a roughly proportional increase in the plasma temperature and a factor of two reduction in collisionality. This demonstration of density control in H-mode plasmas has several far reaching implications: (1) it greatly enhances the prospects of a steady-state tokamak based on rf current drive, in particular ECH and fast wave; (2) the accessible density range, at constant $I_p$, for DIII-D H-mode plasmas is now sufficiently large to allow a separate determination of the scaling of confinement with $n_e$ and $I_p$ with moderate accuracy; (3) support feasibility of helium ash exhaust in ITER with a modest pumping speed.

Fig. 4. Density profiles of two shots with and without divertor pumping 2.5 s after the L-H transition. Typically the pumped plasmas have a slightly more peaked profile late in the H-mode phase.
EROSION AND REDEPOSITION ON CARBON PROBES IN TEXTOR

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1. INTRODUCTION

Plasma induced erosion at the limiters and divertor plates will limit the lifetime of these plasma exposed wall tiles and it represents a critical source of impurities for a future fusion reactor. However, there are hardly any quantitative measurements about the plasma induced erosion at limiters and the possible redeposition of the eroded material. In order to obtain such quantitative data a limiter like carbon probe has been exposed in the tokamak TEXTOR and the measured erosion at different areas is compared with calculated values obtained with the computer simulation program ERO [1].

2. EXPERIMENTAL

The special designed limiter like carbon probe with dimensions 50*60*90 mm$^3$ is shown in fig.1. Different markers with a diameter of 3 mm and thickness 220 Å (Mo), 200 Å (V) and 500 Å (Al) have been exposed in the boundary plasma of TEXTOR during six ohmic D discharges (from #53486 to #53491) with a total exposure time of $t_{\text{exp}} = 15$ s. The probe has been inserted from the bottom into TEXTOR beyond the last closed flux surface (LCFS) given by the main ALT II toroidal limiter system. The positions have been $z_0 = -7$ mm for discharge number #53486, $z_0 = -12$ mm for #53487 and $z_0 = -16$ mm for #53488 to #53491. The profiles of electron temperature $T_e(r)$ and density $n_e(r)$ (r is the radius) in the plasma edge (SOL) were measured during the
3. MODEL

The erosion/redeposition has been calculated with the program ERO [1]. Atoms eroded from the probe by sputtering enter the plasma, they are ionized and subsequently guided along the magnetic field lines. They may be pushed back toward the probe surface by the streaming plasma flow. In the version of ERO used here the velocity of the sputtered ions after ionization is given by [2]

$$\frac{dv}{dt} = \frac{q^*e}{m_I} \left( E + [v \times B] \right) + \frac{\delta v}{\delta t},$$

(1)

where \( q = q(t) \) is the charge state and \( m_I \) denotes the mass of the impurity ion. The second term on the right hand side of eqn. (1) describes the changing of the ion velocity by collisions with the plasma ions, e.g. on account of parallel friction, parallel and perpendicular diffusion. Equation (1) is solved using the leap-frog method for the integration and the Boris method for the \([v \times B]\) rotation [3]. The electric field near the probe surface is calculated following [2].

The marker erosion of element 'I' is given by

$$d^I = (n_I)^{-1} \left( \Gamma^e - \Gamma^n \right) \left( \gamma_{D^+I} f^D + \gamma_{C^+I} f^C + \gamma_{O^+I} f^O \right) \exp,$$

(2)

where \( n_I \) is the atomic density of element 'I' and \( \gamma \) presents the sputtering yields, respectively for the various ion-target combinations. The flux \( \Gamma^e \)

$$\Gamma^e = 0.5 n_e(r) c_s \sin(\alpha)$$

(3)

with \( \alpha \) being the angle between the surface of the probe and the magnetic field lines is reduced by \( \Gamma^n \)

$$\Gamma^n = \frac{\Gamma^e f^C (1 - R^{C\to I})}{y_{D\to C} f^D + y_{C\to C} f^C + y_{O\to C} f^O}.$$

(4)

The concentration of the ion species 'i' (e.g. \( D^+, O^{++}, C^{++} \)) are denoted by \( f^i \) and satisfy with charge states \( q^i \) the condition \( \sum f^i q^i = 1 \). The ion acoustic speed is
\[ c_s = \left(2\pi k T_e/m_D\right)^{1/2} \]

with \( m_D \) being the atomic mass of the \( D^+ \) ions. The ion flux \( \Gamma_{\text{e}} f_C \) is needed to sputter the amount of carbon ions \( \Gamma_{\text{e}} f_C \) coming from the plasma and sticking on the surface of the marker with the probability of \( (1-R^{C\rightarrow I}) \), where \( R^{C\rightarrow I} \) is the particle reflection coefficient of carbon ions scattered from a surface of element 'I'.

4. RESULTS AND DISCUSSION

4.1. Marker erosion

After exposure of the probe to the plasma at the sides 1-8 (indicated in fig.1) samples were cut out for surface layer analysis with MeV ion beam techniques (RBS, PIXE and nuclear reactions). For the Mo marker an erosion of about 120 Å, for V of about 90 Å, for Al\(^{(1)}\) of about 90 Å, for Al\(^{(2)}\) of about 120 Å, and for Al\(^{(3)}\) of about 40 Å was measured. Though, the sputtering yields for these materials at low energy D bombardment are very different, the measured erosion is nearly equal. The variation of electron density and temperature with decreasing distance from the LCFS is not large enough to explain this behavior. For understanding these observations we must take into account that besides D the probe will also be bombarded with C and O ions. The sticking probability for the C ions on low-Z elements V and Al is larger than on the higher-Z element Mo.

Fig.2 shows the calculated erosion using eqn. (2) for C, Mo, V, and Al\(^{(1)}\) in dependence on the carbon concentration in the plasma \( f_C^{\text{LCFS}}(n_e^{\text{LCFS}} = 1.6 \times 10^{12} \text{ cm}^{-3}, T_e^{\text{LCFS}} = 40 \text{ eV}, \text{ ratio } O/D = 0.01, \alpha = 25^\circ) \). The sputtering yields were calculated by using the approach proposed in [6], where the modifications to
the energy distribution function of incident ions due to the acceleration in the sheath electric field are included. An agreement with the measured value of about 100 Å erosion for all elements can only be achieved by using particle reflection coefficient \( R^{C_{11}} = 0.11 \, \text{Å}, \ R^{C_{21}} = 0.18 \, \text{Å}, \ R^{C_{31}} = 0.26 \, \text{Å}, \ R^{C_{30}} = 0.31 \, \text{Å} \) which are half of those calculated with TRIM [8] and an assumed value of about 0.06 for \( f^{C} \). Surface roughness can be the cause for the decrease of the carbon reflection from the marker surfaces [7].

4.2. Redeposition of the marker materials

The slit construction (see fig.1) was designed in order to prevent a further erosion of the redeposited there marker material. Within the error level of the RBS technique only on side 4 and only for Mo a redeposition could be observed (fig.3). This redeposition is clearly caused by direct deposition of the neutrals sputtered from the Mo marker with a likely cosine distribution and impinge directly on side 4.

![Fig.3: Distribution of Mo atoms on side 4](image)

REFERENCES

[1] D. Naujoks, R. Behrisch, J.P. Coad, L. de Kock, JET-P(92)64, to be published in Nuclear Fusion, in press
Control of Poloidal Asymmetry in TEXTOR Edge Plasma
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Under conditions of poloidal asymmetry of plasma in tokamaks, longitudinal equilibrium currents (Pfirsch-Schlüter currents) are closed in a radial direction. In so doing, the force moments $\mathbf{I} \times \mathbf{B}_1$ for positive and negative values of the radial current $\mathbf{I}$ are not equal by the absolute value because of the dependence of the toroidal magnetic field $\mathbf{B}_1$ on the poloidal angle. As a result, plasma rotation must occur. This effect may be especially significant in a scrape-off layer (SOL) where the radial current flowing in the plasma closes on the limiter [1]. Poloidal rotation, in its turn, may change the level of wall turbulence.

The rate of poloidal rotation can be controlled by affecting the symmetry, say, by varying the position of gas puffing [2]. This paper describes the first experiments in varying the poloidal asymmetry of the edge plasma in the TEXTOR tokamak by applying a bias current to the limiter.

**Experiment**

The current was flowing between the ALT-II segmented toroidal limiter and a special graphite test limiter located in the bottom of the vessel. The voltage between the latter limiter and earth was varied within $\pm 300 \text{ V}$. The position of the limiters is illustrated in Fig. 1. For both limiters, $r_L = 0.46 \text{ m}$.

Measurements were taken of the current to each one of the eight segments of the ALT-II limiter, earthed via 1-Ohm resistor. The current distribution on the ALT-II limiter was approximately uniform. The ratio of total ALT-II current to total test limiter current is given in Fig. 2.

The radial profile of the plasma density at the plasma edge was measured in three different cross-sections at three poloidal angles, Fig. 3 namely, at the top, in the equatorial plane at the outer board (in the same cross-section as that of the test limiter) and at the bottom, using atomic lithium beams [3]. The bias voltage was applied during the time interval of 0.8-1.8 s after the beginning of the discharge ($I_p = 340 \text{ kA}$, $B_t = 2.2 \text{ T}$, $n_e = (2.5 - 3.8) \times 10^{19} \text{ m}^{-3}$). The upper limit of density measured by Li-beams was $5 \times 10^{18} \text{ m}^{-3}$.

Measurements were also taken of the velocity distribution of neutral deuterium atoms near the test limiter surface, analogous to the measurements described in [4].
Discussion of the Results

1. Poloidal asymmetry of the plasma in the SOL and on closed magnetic surfaces near the limiter is observed in the absence of biasing as well (Fig. 3a). Its causes include the toroidal drift and the asymmetric position of the ALT-II limiter [5].

2. Under conditions of positive voltage, when the test limiter serves as anode, a considerable decrease is observed of the edge plasma density at all three poloidal angles at which measurements were performed (Fig. 4). In so doing, the total number of particles in TEXTOR remains practically the same. From this it follows that the density on the inner board increases. Therefore, applying a bias voltage to generate an electric current flowing along the magnetic field changes the poloidal distribution of the plasma density in the SOL. Also a corresponding decrease in $D_e$ intensity at ALT-II is observed.

In the case of Fig. 3b a decrease of the density on the outer board to values below $5 \cdot 10^{18} \text{cm}^{-3}$ permitted the measurement of $n_e$ on closed magnetic surfaces at distances up to 1cm from the limiter edge (Fig. 3b). Pronounced poloidal asymmetry inside the separatrix is indicative of the existence of poloidal flows in the edge plasma. It can also be controlled by biasing. The data obtained suggest further investigations of the effect of asymmetry on the poloidal rotation of the edge plasma [6].

Acknowledgement:

We should like to thank Drs. Bogen, Pospieszczyn and Samm for their help and advice.

Figure captions

Fig. 1:

Fig. 2: The ratio between the current to the ALT-II limiter and the current from the test limiter: 1, test limiter as anode ($V > 0$); 2, test limiter as cathode ($V < 0$).

Fig. 3: Projection of Li-beam diagnostics into one cross-section showing their location in poloidal direction.

Fig. 4: Plasma density profiles: 1, on the outer board; 2, at the top, at the bottom. $a) V_{bias} = 0$; $b) V_{bias} = +300 \text{ V}$, $t = 1.5s$.

Fig. 5: The density at $r = 46 \text{ cm}$ as a function of current to the ALT-II limiter. Designations are the same as in Fig. 3.
References:

5. U. Samm et al., Nucl. Fusion 31, 1386 (1991)
6. H. Gerhauser, H.A. Claassen (this conference)
Fig. 3

Fig. 4

Fig. 5
Investigations on Particle Removal Processes in the ALT-II Toroidal Belt Limiter Scoop


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Introduction

In a first approximation a pump limiter removes the flux of particles entering the throat with a probability which is equal to the conductance between neutralizer plate and pump divided by the sum of this value and the throat conductance. This picture is altered if the recycling of particles at the limiter front and processes inside the throat become important. The latter is the topic of this study. The relevant processes are ionization by electron impact and charge exchange, dissociation and finally at highest fluxes change of the flow pattern inside the scoop. The result of these processes is in the following written as a change of the throat conductance.

Experimental Conditions

The ALT–II pump limiter is positioned on the outboard side of the TEXTOR torus at an angle of 45° below the midplane /1,2,3/. Various diagnostic devices are included, such as Langmuir probes in the scoop and two different types of ion gauges. The total pressure is measured by fast ionization gauges /4/ and a Penning ion gauge set-up /5/; the latter is capable of measuring simultaneously the partial pressure of helium in the scoop as well as the partial deuterium or hydrogen pressure with a time-resolution of 26 ms during our experiments. Additionally, the helium concentration in the plasma center (r=26 cm) and the He and H/D fluxes recycling at the ALT–II limiter blade #5 are measured as described in /6/.

Figure 1. Schematic description of the throat of ALT–II blade #6. The helium is injected into the scoop (ΦHe gas-flow) with a width δ=27 mm and a length L=18 cm; the high (h) is 15 cm. The effective pumping speed at the plenum Seff is 650 l/s.
For analyzing the throat conductance under the influence of plasma, helium has been injected into the scoop of blade #6 which is open to the ion drift side only. The single sided scoop opening enables a simpler interpretation of the data than in the double-sided case. Blade #6 is covered presently with 12 mm thick carbon tiles. A schematic set up is shown in figure 1. The effective pumping speed $S_{\text{eff}}$ of the turbopump is 650 l/s in the plenum and 190 l/s at the entrance of the scoop. This module can remove 1.2 mbar l/s of gas from the discharge without degradation of the pumping speed /7/.

**Throat Conductance Variations**

Previous measurements on the plasma influence on throat conductance of the TEXTOR pump limiter ALT-I /5/ show a decrease of the throat conductance. There was no active helium pumping possible since cryopumps were used. These experiments as well as first experiments on ALT-II /1,2,8/ have been carried out by injecting a short neutral helium gas puff either into the plenum of the pump limiter /5/ or into the discharge /8/. For the throat conductance variations reported here, neutral helium was injected into the scoop of the ALT-II pump limiter during the discharge so that steady gas flow conditions during the flat top phase of the TEXTOR discharge (total length 4 s, $B_{\perp}$=2.25 T, $I=340$ kA) are reached. In comparison to the other experiments the investigated electron temperature and electron density range was extended to higher values. The gas flow injected into the scoop was varied between 0.5 mbar l/s and 0.9 mbar l/s resulting in a partial helium pressure between 2.4 - 10$^{-1}$ mbar and 4.5 · 10$^{-4}$ mbar depending on the discharge. Experimentally determined is the injected gas flow $\Phi_{\text{He}}$ and the partial pressure of helium, which is proportional to the neutral particle density in the scoop $n_{\text{He}}$, from these values follows

$$\frac{\Phi_{\text{He}}}{n_{\text{He}} \cdot \delta} = \frac{S_{\text{eff}}}{C_{\text{He}}},$$

where $C_{\text{He}}$ is the throat conductance.

In order to have a wide variation of the local electron density and the electron temperature in the scoop of blade #6, ohmic as well as neutral beam heated discharges (1.5 MW of neutral beam power) have been studied under two different conditions. Firstly, all eight ALT-II limiter blades were positioned at a minor plasma radius of 46 cm; secondly, to obtain higher local electron temperature and density only blade #6 was inserted to 44 cm while all the other blades were retracted to 47.5 cm. All discharges have been L-mode discharges in a boronized machine. The central line averaged density has been varied between 1.6 · 10$^{13}$ cm$^{-3}$ up to 7.9 · 10$^{13}$ cm$^{-3}$ widely extending the range of previous measurements /6/.

**Results**

As it is not possible to vary the electron temperature in the scoop independently from the electron density our measurements are always including both effects. As the neutral helium flow into the scoop normalized to the partial helium pressure ($\Phi_{\text{He}}/n_{\text{He}}$) is measured a comparison with theoretical deductions of the throat conductance following equation (1) is easily possible. The value of the neutral helium flow into the scoop normalized to the partial helium pressure in the scoop, increases with increasing electron density in the scoop for ohmic and neutral beam heated plasmas (fig. 2 a.). The normalized gas flow seems to be constant after a rapid drop for electron temperatures higher than 15 eV (fig. 2 b.). Spectroscopical measurements at the front surface of blade #5 show an increasing helium flux recycling on the ALT-II limiter as a result of the increasing amount of helium escaping from the scoop. The helium concentration in the center during neutral beam heated discharges at the flat top phase of the He flow measurements at 1.5 - 1.8 s does not follow the helium recycling signal at the limiter but is nearly constant for all discharges, with the exception of the lowest density discharge. The influence of the electron density and temperature variation together have to be taken into account to interpret the
experimental findings. Part of helium may be stored transient in the walls /6/ which has been noticed in our experiment too. To another part it can be attributed to the observation that the fuelling efficiency for helium decreases with the line averaged electron density /6/.

![Figure 2a) and b)](image.png)

**Figure 2 a) and b).** Dependence of the helium flow into the scoop normalized to the partial He-pressure from the electron density in the scoop (fig. 2a.) and the electron temperature (fig. 2 b.). x: neutral beam heated discharges with all eight ALT-II blades at 46 cm minor radius; o: ohmic and Δ neutral beam heated discharges with blade # 6 inserted to 44 cm and the other blades retracted to r=47.4 cm.

An analytical hydrodynamic model /9/ as well as the Monte Carlo EIRENE code /10/ was applied to study the transport of neutral helium in the throat. Both approaches give in agreement with previous measurements a decrease of the conductance for low plasma density ( < 5 \cdot 10^{12} \text{ cm}^3) in comparison to the conductance off an empty throat. For this range of n_e, the results of the calculations agree with the plasma free case and underestimate the experimental data which show only a small decrease of \( \Phi_{He}/n_{He0} \). If n_e exceeds a value of 5 \cdot 10^{12} \text{ cm}^3 the increase of the throat conductance as defined in (1) cannot be explained by transport of neutrals. Reverse flow of plasma /11/ was treated as a possible cause of this phenomenon. Such a flow is generated by ionization of neutral gas entering the throat plasma from the plenum if the gas pressure is high enough. Estimates for blade #6 show that this critical level is reached at a scoop plasma density of 5 \cdot 10^{12} \text{ cm}^2. The contribution of the reverse flow to the conductance was estimated under the assumption that all helium ions born in the flow region escaped into the scrape off layer. The results including reverse flow of helium ions are in agreement with the experimental data. The dependency from n_e as shown in figure 3 a) is nearly the same as in figure 2 a). Calculations for the highest densities could not be made due to missing electron temperature values. For the T_e dependency the analytical model (figure 3 b)) shows the same behavior as the experimental data.

Helium exhaust and particle removal studies have been carried out by use of the toroidal belt limiter ALT-II. The measured values of the conductance for low electron temperature and density agree within the error with previous results. The Monte Carlo calculations as well as the analytical model describe this situation very well. For higher scoop electron density (>5 \cdot 10^{12} \text{ cm}^3) an increase of \( \Phi_{He}/n_{He0} \) with increasing density was measured. This change of the conductance can be explained by reversal flow of He^+ in the scoop of the belt limiter. This results in an escape of pumped particles and hence in a reduction of the exhaust efficiency.

The authors would like to thank the TEXTOR and NBI groups, especially A. Pospieszczyk and B. Schweer for providing data and fruitful discussions and A. Hiller for technical assistance.
Figure 3a) and b) Theoretical deduction of $\frac{\Phi_{He}}{n_{He}}$ following the hydrodynamic model /9/. (Symbols as in figure 2.)

References

Impurity production and plasma edge parameters under various wall conditions in TEXTOR

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1. Introduction

Ever since tokamaks have been operated oxygen is one of the dominating impurity problems. Major progress has been achieved by conditioning methods reducing oxygen via gettering. In particular, the low-Z wall coating methods pioneered at Jülich, i.e. carbonization, boronization and siliconization [1][2], improved the general performance of the machines significantly. Furthermore, wall coatings have influence on the radiation level, on plasma edge parameters like electron density $n_e$ and temperature $T_e$ and on the release of wall materials. The relevant processes at the plasma boundary compose a highly non-linear system: $T_e$ affects the energy of the impinging ions and thus the sputtering yields via the sheath potential. The impurities released from the limiter penetrate into the plasma boundary, are ionized and affect the heat flux balance by their line radiation causing a reduction of $T_e$.

Owing to the different wall coatings applied on TEXTOR and the comprehensive edge diagnostics available a rather large database on particle fluxes and edge parameters $n_e$, $T_e$ has been obtained. This is used to study the complex processes at the edge including radiation effects. It is shown how with different wall coatings quite distinct edge parameters and radiation levels develop, due to their specific radiation characteristics and sputtering yields. Special emphasis is put on synergistic effects of combined impurities like C, B, O and Si. The analysis of the experimental data is made with the help of a model developed for this special purpose and described in chapter 3.

2. Experiment

The tokamak TEXTOR has been operated under standard conditions: $R=1.75$ m, $a=46$ cm, $I_p=340$ kA, $B_t=2.25$ T and $D_2$ filling. The plasma heating used was ohmic (0.35 MW) or with auxiliary heating employing neutral beam co-injection (1.5 MW).

A variety of spectroscopic measurements (filter systems, grating spectrometer) at the different limiters provide integral measurements of the line radiation from neutral and singly charged particles like D, C, O, B, Si. The edge parameters $n_e$ and $T_e$ are determined with active beam diagnostics (Li, C, He) [3]. From the spectroscopic data and the $T_e$ measurements the particle flux released from the limiter can be deduced [4]. A bolometer provides measurements of the total radiation.

3. Model

The model developed for this analysis is based on a 0-d particle and heat balance at the limiter edge. Similar models have been developed earlier [5]. The synergistic effects of different impurities and materials and the consideration of radiation from the plasma boundary based on new data are the main improvements in our model. The heating power and the central density are external parameters, all other parameters like $n_e$, $T_e$ and impurity fluxes (except oxygen) can be deduced. The particle transport is calculated based on a simple diffusion model [6] considered to be adequate for this purpose. Thus the radial particle flux is given by $\Gamma=D_\perp \frac{dn_e}{dr}$. The radial diffusion coefficient $D_\perp$ varies with $T_e$ according to a Bohm-like diffusion. The electron density gradient is approximated by the ionization length of deuterium and the central density to $\frac{dn_e}{dr}=\frac{n_e}{\lambda_i}$. The total radiation $P_{rad}$ is related to the particle fluxes via the radiation potential $E_{rad,i}$ [7][8][9], i.e. the energy radiated per impurity particle during its dwell time in the plasma. Thus $P_{rad}$ is the sum of contributions from the various impurity species $i$

$$P_{rad} = \Gamma_D \sum c_i E_{rad,i}$$
The impurity level $\alpha_i = \Gamma_i/\Gamma_0$ represents the "effective" sputtering yield and is calculated with a sputtering model [10][11] as a function of the average energy of the impinging ions (assuming $T_i = T_e$ and normal incidence). This energy is deduced from the sheath potential and an average charge state $q$ for the ions. We used $q=4$ for C and $q=6$ for Si. The yield for oxygen on carbon is treated as a special case, based on the knowledge that O impinging on a graphite surface produces CO with a yield of about 1 [12]. Because the primary source of oxygen and its mechanisms are not well known, we use the measured oxygen level $\alpha_o = \Gamma_o/\Gamma_0$ as an external parameter in the model. The effective sputtering yield (impurity level) for the element $i$ is

$$\alpha_i = (Y_{Di} + \sum_j \alpha_j Y_{Dj}) / (1 - Y_{Di}).$$

This term represents the sum over all sputtering yields with the projectiles of D and the impurity j on the element i, including the self-sputtering $Y_{Di}$.

4. Results and discussion

With boronized and siliconized walls the oxygen level ($\alpha_o = \Gamma_o/\Gamma_0$) is significantly reduced compared to carbonized walls (Fig.1a); e.g. at a line averaged central electron density $n_e = 4 \times 10^{13}$ cm$^{-3}$ the oxygen level $\alpha_o$ is 1.3% with carbonized, 0.4% with boronized and 0.3% with siliconized walls, clearly showing the enhanced oxygen gettering capability of the different coatings.

It is evident from the comparison of Fig.1a and b that the carbon level $\alpha_c$ is closely linked to the oxygen level (solid lines represent model calculations). This behaviour can be understood following our assumption about the formation of CO and comparing its yield with the sputtering of C by deuterium. Within the model calculations we can distinguish the different contributions for carbon release from CO formation, sputtering by D and self sputtering as shown in the insert of Fig.1b. CO formation can give a significant contribution, provided the oxygen level is as high as the deuterium sputtering yield (about 1-2%), it is even dominating in a cold edge plasma (ohmic heating, high $n_e$, $T_e = 10$ eV) when sputtering is low [13][14]. If the oxygen level is low ($< 1\%$) and sputtering is reduced due to low $T_e$, the relative contributions from other carbon release channels (like hydro- carbon formation) may become important.

In a carbonized and a boronized machine the radiation level $\gamma = P_{\text{rad}}/P_{\text{host}}$ (Fig.1d) is mainly determined by oxygen because of its relatively high radiation potential $E_{\text{rad}}$. In a boronized machine $\gamma$ is lower due to the reduced oxygen level. To show the importance of oxygen radiation it is also indicated in Fig.1d what would be expected for $\gamma$ in an idealized case without any oxygen. Boron behaves similar to carbon with respect to sputtering and radiation. Therefore the measured and modelled plasma edge parameters $n_e$ and $T_e$ do not differ between carbonized and boronized walls (Fig.1e/1f).

The edge parameters change after siliconization. $T_e$ is generally lower (Fig.1e). This is mainly due to the higher radiation potential of silicon (about a factor 4 higher than O at $T_e = 30$ eV), as reflected in the elevation of $\gamma$ (Fig.1d). This figure also shows a decrease of $\gamma$ with increasing central density, in contrast to the carbonized/boronized cases. This drop of radiation with increasing density can be explained by the fact that sputtering is the dominant release channel for silicon and that the variation of the silicon level $\alpha_{si}$ is very strong (Fig.1c) because $T_e$ at the boundary decreases to values at which the energy of the impinging projectiles on the limiter comes close to the thresholds for the sputtering processes. Table 1 summarizes the variation of relevant parameters which determine the impurity radiation level $\gamma = E_{\text{rad}} \alpha_i \Gamma_i / P_{\text{host}}$ when going from low to high central densities for a siliconized (silicon dominant) and a carbonized wall (oxygen dominant). Shown are the ratios between high and low density of the measured impurity levels $\alpha_i(high n_e)/\alpha_i(low n_e)$, of the measured deuterium fluxes $\Gamma_0(high n_e)/\Gamma_0(low n_e)$, of the radiation potential $E_{\text{rad}}(high n_e)/E_{\text{rad}}(low n_e)$ and of the resulting ratios for the radiation level $\gamma(high n_e)/\gamma(low n_e)$. 

Table 1
The increase of the radiation potential and of the deuterium flux is similar in both cases, but the decrease of $\alpha_i$ is significantly stronger for silicon than for oxygen. Therefore, with silicon the decrease in $\alpha_i$ overcompensates the increase of the product $E_{\text{rad}}\Gamma_D$ leading to a reduction of radiation, whereas with oxygen the variation is too small to inhibit the usual increase of $\gamma$. On the other hand the silicon system is rather sensitive to oxygen added to it. Already 1% of O more at high $n_e$ may reverse the behaviour such that the radiation again increases with density.

5. Conclusions

The data presented in this paper demonstrate that the application of different wall coatings and thus the influence of different impurities allows to establish quite distinct ranges of plasma edge parameters. In particular, the key role of oxygen for plasma edge behaviour under different wall coatings is shown. However, the origin of oxygen and its release mechanism from the walls are still open questions and have to be the subject of further studies.

Silicon wall coating represents a system in which impurity release is dominated by physical sputtering of silicon resulting in a decrease of radiation with increasing central density. The comparison of the experimental data with the model calculations show that it is possible to describe some important elements out of the complex relations of plasma edge processes even with a simple 0-d edge model. To give a more detailed description of the plasma edge and to achieve a better agreement with experimental data other models are necessary which include 1-d or 2-d structure of the plasma, improved modelling of transport, oblique incidence of sputtering ions and a more accurate treatment of $T_i$.

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References

[9] U. Samm et al., invited paper, this conference
Fig. 1: Variation of measured edge and global parameters as a function of $n_e$ for different coatings (carbonization - C, boronization - B, siliconization - Si), auxiliary heating with neutral beam injection ($P_{\text{beam}} = 1.5$ MW), in comparison with results from edge modelling (solid lines):

a) oxygen level $\alpha_0 = \Gamma_0/\Gamma_D$, b) carbon level $\alpha_C = \Gamma_C/\Gamma_D$, insert: modelled contributions from CO-formation and physical sputtering to the total carbon release (carbonized conditions, neutral beam and pure ohmic heating), c) silicon level $\alpha_{Si} = \Gamma_{Si}/\Gamma_D$, d) radiation level $\gamma = P_{\text{rad}}/P_{\text{heat}}$, e) electron temperature $T_e$ at $r=46\,\text{cm}$, f) electron density $n_e$ at $r=46\,\text{cm}$.
Investigation of MARFEs in TEXTOR


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Introduction

A MARFE (Multifaceted Asymmetric Radiation From the Edge) is a toroidally symmetric but poloidally asymmetric highly radiating zone located near the inner wall of a tokamak [1]. The MARFE, like the solar protuberance [2], is the result of a radiative condensation process.

On the tokamak TEXTOR the formation of MARFEs as well as their suppression has been studied since some years [3]. MARFEs are obtained in plasmas close to the density limit and with a fairly high power density. MARFE formation can be triggered by increasing the average density, injecting additional impurities, reducing the heating power or increasing the local impurity sources on the high field side. The latter can be obtained by shifting the plasma closer to the inner bumper limiter. A common feature of all these methods is an increase of the ratio of local radiation over heating power. The possibility of forming reproducible MARFEs on TEXTOR, in particular MARFEs having a life time of several 100 ms, together with special diagnostics based on optical methods offered the opportunity to study for the first time the radial electron density $n_e$ and temperature $T_e$ profiles across the MARFE region. Short living MARFEs (ca. 20 ms), as they often occur with a disruption, are not considered here.

Experimental arrangement

TEXTOR was operated with $B_T = 2.25$ T, $I_p = 350$ kA and $R = 175$ cm. The direction of $B_T$ was clockwise and $B_T \perp I_p$. The toroidal belt-limiter ALT-II and the poloidal limiter segments (upper, lower and outer segments) were located at $r_{\perp} = 46$ cm. The position of the inner bumper limiter was $r = 49.8$ cm. Co-injection of a neutral beam with 1.3 MW was used as auxiliary heating. Wall coating methods applied on TEXTOR determine the level and type of impurities in the plasma. With siliconization [4] silicon is the dominating impurity with respect to radiation. With boronization [5] the radiated power is mainly coming from oxygen. In this case additional injection of neon helps to increase the radiation level and thus to obtain long living MARFEs.

A CCD-camera equipped with an interference filter was used to observe a full poloidal cross section tangentially at the location of the poloidal limiters in the light of H$_e$. Spectral measurements in the MARFE region were performed by means of a grating spectrometer. The viewing optics collected visible and near infrared light along a tangential chord passing through the region near the bumper limiter (inner wall) and directed it onto the entrance slit of the spectrometer. The entrance slit has been oriented radially along the horizontal midplane of the torus. A CCD-camera with an image amplifier was used to detect the spectra. The width of the spectral interval covered by the CCD-camera was about 25 nm. The time resolution was 20 ms. The detection system was absolutely calibrated against a standard tungsten ribbon lamp. The radial resolution along the slit was about 0.4 cm and the spectral resolution was about 0.06 nm. With the help of this spectrometer the intensities of two He-lines - a singlet line at $\lambda = 728$ nm and a triplet line at $\lambda = 706$ nm - have been measured simultaneously (Fig.1). A small amount of helium was introduced into the plasma to provide a flux of neutral He atoms emitted from the bumper limiter due to recycling. This source of He can be used similar to the thermal helium beam employed routinely at the outboard side of TEXTOR [6]. According to this method the measured ratio of the He-line intensities.
allows to deduce the radial $T_e$-profiles. Because the density in the MARFE is high enough we can use the continuum radiation, also measured by the spectrometer, to deduce $n_e$ [7]. In addition to the optical methods also probes located inside the bumper limiter near the equatorial plane have been used. Several electrode pins have been aligned vertically with a distance of 7 cm between each. The electrodes are operated as double probes by applying an alternating voltage with a frequency of 200 Hz.

Measurements and results

Apart from the different methods to trigger a MARFE, the general evolution of a MARFE is similar in all cases. As a typical example the growth and the motion of a MARFE in a discharge with a siliconized wall (shot #53433) will be discussed in detail. In Fig.3 the evolution of the MARFE is exhibited as contour plots of H$_\alpha$-light as seen in the tangential view. Fig.3a represents a typical distribution of H$_\alpha$ radiation without a MARFE, where the only intense radiation comes from recycling particles at the limiters. In this discharge the central line averaged electron density was about $n_e=5.3 \times 10^{13}$ cm$^{-3}$. The MARFE occurs during the ramp down of the plasma current. When the current dropped from 350 kA to about 300 kA and the global radiation level $P_{\text{rad}}/P_{\text{heat}}$ increased at the same time from 30% to 40% a MARFE started to grow.

About 100 ms before the MARFE appears a zone of slightly enhanced H$_\alpha$ emission (factor 2-3 over background) can clearly be seen. This "pre-MARFE" is located at the horizontal midplane close to the bumper limiter and extends in poloidal direction by about 5 cm (Fig.3a). It indicates the location where the MARFE will be born. When the MARFE starts to grow it first moves downwards (Fig.3b) to a position about 15° below the horizontal midplane. There it reaches an H$_\alpha$-radiation level comparable to the intensity at the limiters (Fig.3c). At this phase the inner edge of the MARFE is at $r=40$ cm. From this position the MARFE begins to move upwards steadily in poloidal direction (direction opposite to $B_x\times B_y$) and continues to grow as can be seen in the Figs.3c–f. When it reaches the midplane position its poloidal extent becomes about 50° and its inner edge is located at $r=36.5$ cm. When the MARFE reaches the top of the plasma column (upper limiter) its poloidal motion stops (Fig.3f). The MARFE attempts to move towards the plasma centre in radial direction but then disappears in less than 20 ms (Fig.3g). At the same time a new MARFE starts to grow at the same location where the first MARFE began its upward motion and the process is repeated (Fig.3h). In our example a MARFE was formed four times the same way as described above lasting about 700 ms in total before the discharge ended with a major disruption. The average poloidal velocity $<v>$ of the MARFE motion changed. The first MARFE had $<v>=0.77$ m/s and the third one $<v>=0.95$ m/s. $<v>$ appears to increase linearly with time.

The measured He line intensities and the continuum radiation across the MARFE zone provide information about $n_e$ and $T_e$. An Abel inversion was used to deduce the radial distribution of the volume emission inside the MARFE zone. The resulting electron density and temperature profiles are given in Fig.2. These profiles correspond to the MARFE phase shown in Fig.3d. The main part of the MARFE is located inside the limiter radius, i.e. $r<46$ cm. The electron density reaches a peak value of $2 \times 10^{14}$ cm$^{-3}$ at a radius of $r=42$ cm. The electron temperature is about 5 eV at this location. Data from other experiments obtained under boronized conditions with an addition of neon gave similar peak values: $n_e=1.2 \times 10^{14}$ cm$^{-3}$ and $T_e=5$ eV. The drop of $T_e$ measured by the probes at the bumper limiter during MARFE formation is as pronounced as seen by spectroscopy. After the MARFE disappears the electron temperature grows to the old value.

Discussion and conclusions

The measured electron density and $T_e$ in the MARFE have been measured for the first time. The peak values in density are remarkably high, even exceeding those in the centre of the core plasma.
Existing theoretical models [8][9][10][11] allow partly to understand the MARFE's behaviour described above and help to distinguish the role of different physical processes. The existence of a "pre-MARFE" stage with enhanced Hα emission indicates that localized recycling at the bumper limiter, which leads to a local plasma cooling, plays probably a trigger role in the appearance of the first MARFE [9]. The characteristic time of MARFE development agrees with the results of our modelling of plasma detachment arising as a result of a radiative instability with impurities. The plasma parameters in a MARFE at the stage of slow evolution are consistent with estimations of Ref.[10], where the MARFE is treated as a non-linear wave whose location at the inner edge is linked to the effect of the Shafranov shift of magnetic surfaces on energy losses. According to these calculations the temperature inside the MARFE is expected to drop to $T_e \approx 3$ eV. Owing to approximate constancy of the plasma pressure on flux surfaces a reduced plasma temperature in the MARFE region consequently must lead to an increased plasma density. During the MARFE growth the pressure inside the MARFE must be somewhat smaller than at the surrounding plasma. Generally a layer with reduced pressure is displaced towards higher magnetic field. This mechanism together with the effect of the Shafranov shift may explain why a MARFE has never been observed at the outside and that the MARFE disappears when it reaches the top of the torus. An interpretation of other details of the MARFE behaviour, such as direction and velocity of the poloidal motion, needs further development of theoretical approaches.

Acknowledgements

The authors thank Prof. E. Hintz and the TEXTOR team for help.

References


Fig. 1 Emission spectra obtained from a MARFE showing the two HeI lines used for deducing $T_e$.

Fig. 2 The radial distribution of $n_e$ and $T_e$ across a MARFE.
Fig. 3  Evolution of a MARFE in shot #53433 in the light of $H_\alpha$ shown as contour plots (lines of constant intensity). The view shows the poloidal cross-section where the main limiter segements (poloidal) are located.
EFFECTS OF ICRH ON THE TEXTOR SCRAPE-OFF LAYER WITH SILICON COATED WALLS

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1. Introduction

The interaction between waves in the ion cyclotron range of frequencies and the plasma boundary plays - in view of the relation between the edge and the overall plasma properties and performance - an important role in the success of ICRF to heat plasmas to ignition conditions\textsuperscript{1}. Its use as a current drive method may be even more affected by interactions with the edge. The interaction between ICRH and the plasma edge is a vast and interlinked domain\textsuperscript{1}: the edge affects the RF heating and the RF heating affects the edge. Methods to avoid RF-induced edge effects or to mitigate their consequences have been empirically developed over the years. In recent years important progress has been achieved in the understanding of the interaction between ICRF waves and the plasma boundary, in particular due to the identification of RF sheath formation as the main underlying physical mechanism for RF-induced impurity generation and edge effects. This understanding has led to improvements in the empirical methods and to the development of new concepts. It has been shown on several machines that it is possible to routinely apply high power ICRH without large edge effects.

The ICRH-induced edge effects can be influenced either by trying to avoid the underlying mechanisms or by trying to mitigate their consequences\textsuperscript{1}. Sometimes the mechanisms can be suppressed by choosing the appropriate plasma- and heating conditions and by reducing the sheath rectification effects. The consequences of the mechanisms can be reduced, e.g. by keeping the plasma edge temperature low, by antenna conditioning and appropriate antenna materials, and by using low-Z coating techniques.

An RF sheath is formed wherever parallel RF-fields are present in the plasma near the wall\textsuperscript{1}. In large machines, where the waves are well focused and the absorption is good, the RF-fields and sheath effects are essentially concentrated in front of the antenna (Faraday screen). In machines like TEXTOR the sheath effect can also occur on walls that are further away from the antenna. On TEXTOR it has been shown that the DC sheath-rectified current\textsuperscript{2}, generated at a powered antenna, is mostly drawn by the antenna protection limiters, the contribution of the Faraday screen (FS) blades (where the plasma density is much lower) being negligible\textsuperscript{3}. This is confirmed by the fact that no significant difference in DC current drawn by an antenna with or without FS has observed. A toroidal RF current at the generator frequency and its harmonics accompanies this DC current\textsuperscript{2}.

The sheath rectification effects can also be reduced by supporting the voltage drop along the field lines not in the plasma sheath but in an insulator\textsuperscript{4}. Therefore, one of the TEXTOR antenna pairs was equipped with insulated protection tiles.
Prior to carbonization on TEXTOR, all ICRF-heated discharges switched into strong detachment at low power levels and rapidly evolved to disruption. Changing from an "all-metal" to an "all-carbon" surrounding resulted in stationary conditions without high-Z impurity problems. The use of boronization resulted in an erosion rate lower than for carbonized walls, a significantly reduced carbon and especially oxygen influx, and a recycling coefficient below one, allowing density feedback control. Recently, TEXTOR has been operated with walls and limiters fully coated with silicon using a plasma-assisted chemical vapor deposition from silane gases. The concentration of low-Z impurities B, C and O dropped well below the already low levels with boronization, and Si becomes the dominant plasma impurity.

2. Experimental set-up and results

The experiments were performed for minority hydrogen heating in deuterium under standard TEXTOR conditions (ALT-II belt limiter at 46 cm, $B_t = 2.25$ T, $I_p = 350$ kA, RF frequency 32.5 MHz, TT-phased antennas, ohmic plasma: $n_{eo} = 3 \times 10^{13}$ cm$^{-3}$, $T_{eo} = 1.2$ keV. Both antenna pairs (A1 and A2) were without FS and had feeder protection. Pair A2, separated by 180° from A1 had silicon nitride insulation between the graphite protection limiters and the steel frame. In agreement with our expectations, the DC current drawn by A2 was totally suppressed. Heating performance as well as boundary plasma temperature and density profiles were found to be similar for operation with A1 or A2. However, no advantage w.r.t. to the operation with A1 was found. Unfortunately, in a later phase arcing occurred at the insulation and the silicon nitride tiles broke due thermal stresses, and were removed. Afterwards experiments at RF power levels up to 2.5 MW (i.e. $7 \times P_{OH}$) have been conducted. Some special measurements were performed with a view to the study of plasma surface interactions under siliconized wall conditions.

a) The surface collector, Stockholm-TEXTOR probe system has been used in order to trace deuterium and plasma impurity fluxes (Si, B, and Inconel components Ni + Cr + Fe) 15 - 35 mm deep in the SOL. The probe was facing the electron drift direction. Several graphite samples were exposed in ohmically- and ICRF-heated discharges; 1 MW RF power was coupled for 1.6 $s$ from antennas A1 or A2. The exposures were performed during 3.5 $s$ shots with $n_{eo} = 3.0 \times 10^{13}$ cm$^{-3}$ making them available for direct comparison with other exposures done under siliconized conditions. Following the exposure, the samples were transported to the surface analysis station and then analyzed by means of high energy ion beam techniques such as Rutherford backscattering spectroscopy (RBS) and nuclear reaction analysis for the detection of deuterium and boron. The deposited deuterium was determined with a $^3$He$^+$ analyzing beam [$^3$He(d, p)$^4$He] at an energy of 760 keV and a total dose of $6.25 \times 10^{12}$ cm$^{-2}$ to avoid detrapping effects, whereas boron was determined with a proton beam [$^1$B(p, a)$^4$He] of 640 keV.

The exposures of the probes resulted in the co-deposition of deuterium together with plasma impurity atoms. Table 1 shows the fluxes of different species measured 17 mm deep in the SOL during pulses made in the freshly siliconized machine (# 52193-6) and several hundred discharges later.
Table 1
Fluxes of silicon, boron and metals (Ni + Cr + Fe)

<table>
<thead>
<tr>
<th>Discharge no.</th>
<th>Flux* $F_D$ $10^{15}$cm$^{-2}$s$^{-1}$</th>
<th>Heating</th>
<th>Fluxes (deposition rates)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td></td>
<td>Silicon $10^{15}$cm$^{-2}$s$^{-1}$</td>
</tr>
<tr>
<td>52193-6</td>
<td></td>
<td></td>
<td>11.2</td>
</tr>
<tr>
<td>freshly siliconized</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>52713</td>
<td>2.6</td>
<td>$\Omega$</td>
<td>1.8</td>
</tr>
<tr>
<td>52725 (A1)</td>
<td>5.5</td>
<td>ICR</td>
<td>5.1</td>
</tr>
<tr>
<td>52729 (A2)</td>
<td>5</td>
<td>$\Omega$</td>
<td>4.3</td>
</tr>
</tbody>
</table>

*Deposition rates of D are determined by the co-deposition process$^{10}$.

One can notice distinct variations in the fluxes depending on the wall condition and plasma heating. As expected, the silicon fluxes are the highest just after the siliconization and decrease afterwards, whereas the metal fluxes - suppressed strongly by the fresh siliconization - increase with time. Moreover, ICRF-heated discharges lead to the increased release of atoms from the walls, and influence the impurity fluxes in the SOL. These fluxes are 2.5 - 3 times higher when discharges before ICRH (52713, 6) and with ICRH are compared. However, even these increased fluxes are still much lower than the metal fluxes under carbonized$^9$ or boronized$^{10}$ conditions in TEXTOR.

b) The density of atomic deuterium was measured in the vicinity of the wall, at a given time within the discharge, using laser-induced fluorescence at Lyman-alpha$^{11, 12}$. The near-UV radiation of an excimer-pumped, pulsed dye laser was frequency-tripled down to the wavelength $\lambda_{LD}(D^0) = 121.534$ nm. The fluorescence volume was situated 4 cm in front of the liner, at $r = 51 \pm 0.6$ cm, and observed perpendicularly to the exciting beam with mirror optics and a solar blind photomultiplier. The fast pressure tuning of the laser wavelength also allowed the recording of spectral profiles. The knowledge of density and spectral profiles of atomic deuterium is very important to understand the physical processes underlying recycling phenomena. ICRH (1 MW) does not appreciably change the velocity distribution of the observed neutrals: these are slow atoms ($< 2$ eV), presumably of molecular origin (Franck-Condon atoms). On the other hand, if we consider the density of atomic deuterium the situation seems to differ from that with carbonization observed earlier ($I_p = 340$ kA, $n_e = 2.5 \times 10^{13}$ cm$^{-3}$, $T_e = 0.8 - 1$ keV, $P_{ICRH} = 630$ kW, liner temperature 250 °C), where density increases of up to a factor of 3 were found. The measured increase from ohmric to ICRF-heated discharges was lying below 50 % in the siliconized machine. Since we are dealing with slightly different conditions the exact antenna design, and in the material of some plasma-facing components, further measurements are desirable, including the case of boronization, which could not be studied yet with this method.

c) A sniffer probe$^{14}$ that acts in principle as a small pump limiter system in the SOL in which plasma particles entering the aperture (at $r = 47.5$ cm) interact with a thermally isolated carbonized stainless steel plate at floating potential, was used to measure the
hydrogenic particle fluxes in the SOL and to determine simultaneously the deposited energy by means of calorimetry. From these data the average (ionic) impact energies were evaluated for a variety of ohmic, NBI- and ICRH-heated discharges. This is the energy after the ions have passed the sheath potential. Assuming \( T_e = T_i \) this energy is about 7.2 \( T_e \).

For most of the ICRH-heated discharges the deposited energy and the average energy per incoming hydrogenic particle are similar to those observed in comparable NBI-heated discharges. Typical values range between 80 and 500 eV (see Fig.). A rather large scatter has been observed for different plasma and heating conditions. Within this scatter the results are similar to those under boronized wall conditions. The actual value of the energy per particle seems to depend strongly on the edge density which is determined by the wall recycling condition.

3. Conclusion

The results indicate that siliconized wall conditioning are suitable for high power ICRH operation\(^{15}\), in spite of the relatively high atomic number of silicon (\( Z = 14 \)).

[1] Recent review: J.M. Noterdaeme and G. Van Oost, to be published in PPCF
[6] J. Winter et al., these proceedings
[8] M. Rubel et al., these proceedings
[15] J. Ongena et al., these proceedings
TOROIDAL TRANSPORT MEASUREMENT OF LASER INJECTED LITHIUM IN TEXTOR

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Introduction

Impurities play an important role in the determination of the energy balance of tokamak plasmas through radiation losses, therefore the improvement of our knowledge about the impurity transport in a tokamak plasma is a great challenge in plasma physics. Since the sources of intrinsic impurities are not well known, these particles are not ideal for transport investigations. The injection of non intrinsic impurities overcomes this problem allowing a systematic transport investigation.

Impurity transport phenomena have received considerable attention the last few years, but up to this time the focus of the investigations was mainly directed to the radial transport examinations of highly ionized impurities supposing toroidal symmetry of the tokamak plasma [1-3]. Our experiments are concentrated on the study of the toroidal transport of non-intrinsic impurities. In order to study the transport properties of lithium atoms and Li⁺ ions, lithium was injected into the TEXTOR tokamak using the laser blow-off method. At different toroidal positions of the tokamak, spectroscopic diagnostics with high temporal and spatial resolution were employed in the detection of line radiation from the injected particles.

Experiments

The experimental arrangement at the cross section of the injection is shown in Fig.1. The toroidally collimated lithium beam was radially introduced into the plasma from the top of the TEXTOR tokamak. A 1 μm thick LiF target on a 500Å Cr backing was used for the injection and the emitted line radiation of the Li atom and Li⁺ ion was detected at 6708Å and 5485Å respectively. The reason for the choice to inject lithium was that this atom has a low ionization potential (5.4 eV) and will therefore be ionized in the plasma in a very short timescale. On the contrary, the Li⁺ ion has a high ionization potential (75.6 eV), therefore this ion has a long lifetime and can move far in the toroidal direction in the plasma before it will be further ionized. Spectroscopic detectors were used perpendicular to the axis of the injected beam. In order to obtain time resolved and two dimensional pictures a gated image intensified CCD camera was mounted. This camera covered a 15cm (in radial direction) × 20cm (in toroidal direction) area of the tokamak edge allowing us to get a good spatial resolution. The integration time of this camera was selected to 50μs. The time evolution of the injected beam was monitored by a photomultiplier. For the wavelength selection interference filters were used in front of both detectors.

During lithium injection the radial and toroidal distribution of lithium atoms changed in time (Fig.2.). At the first 100–200μs of the injection the lithium atoms are restricted in toroidal direction to the location determined by the image of the toroidal collimator slit in the plasma (Fig.2a.).
At later times of the injection the lithium atoms can also be detected further away from the location of the injected beam (Fig. 2b.), predominantly at larger plasma radii. At smaller radii the atoms were always concentrated at the position of the original beam only. Atoms which appear in the plasma some hundred microseconds after the beginning of the injection can penetrate obviously deeper into the plasma than the faster atoms which arrive earlier.

The toroidal distribution of the Li$^+$ ions is broader than the distribution of the atoms at every radial position (Fig.3.) and displays the same radial structure as the atomic distribution.

In another series of experiments the time history of the Li$^+$ line radiation was measured along the torus. The detectors were set at three toroidally different cross sections: at the place of the injection, at 45° and 247.5° toroidally away from the injection. Because of the high ionisation potential of Li$^+$, these particles have enough time to move far away from the location of the injection along the magnetic field lines; therefore we could detect them still at a distance of 7m away from the place of the injection. Fig.4. shows the time history of the LiI line radiation measured at the cross section of the injection and that of the Li$^+$ line radiation observed 45° toroidally away from the injection (1.5m) for ohmic and for neutral beam heated plasmas. From the time delay of the maximum of the line radiation, a characteristic velocity of about 1cm/μs was obtained during neutral beam heating for lithium ions. For ohmic plasma this characteristic velocity was a factor of two less (0.4cm/μs).
Fig. 2. Distribution of the atomic line radiation (6708Å) 100µs (a.) and 250µs (b.) after the laser shooting.

Fig. 3. Distribution of the LiII line radiation (5485Å) 150µs (a.) and 250µs (b.) after the laser shooting. The dashed curve indicates the restriction by the port in the break vacuum vessel.
Fig. 4. The time history of the LiII line radiation detected at the cross section of the injection and the time history of the LiII line radiation observed 45° toroidally from the injection (at 1.5m from the injection) for an ohmic (dashed line) and for a neutral beam heated (solid line) plasma.

**Discussion**

At the first 100–200μs of the injection the measured penetration depth of the atoms coincides with the penetration depth which was calculated using the measured edge plasma density profile, the ionization rate coefficient for lithium atoms, and their radial velocity. From this one can conclude that during this time the injected beam consists mainly of atoms. These results can be explained by a simple transport model. Because of the high characteristic toroidal velocity of the Li+ ambipolar diffusion or expansion of the lithium ions seem to be a good explanation.

At later times of the injection the picture is more complicated. A part of the injected atoms penetrates deeper into the plasma than the calculated penetration depth. Two candidates for the explanation of this phenomenon have emerged: the injected beam contains microscopic particles (pellets), which naturally penetrate deeper into the plasma than the atoms, or the injected atomic beam perturbs and cools the plasma locally.

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**References**


Measurement of Dα Sources for Particle Confinement Time Determination in TEXTOR

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1. Introduction

An important quantity in the study of tokamak discharges is the global particle confinement time, defined for each ionic species \(i\) by the equation below, where \(N_i\) is the total population of the species in the plasma and \(S_i\) is the source rate (ionization rate) of the species. Of particular significance is the confinement time of the main plasma component, deuterium; here, in most cases of interest, the time derivative is negligible and the confinement time is given by \(N/S\). The deuterium content \(N\) can be estimated from the electron content, measured by interferometry, if \(Z_{\text{eff}}\) is known. A common method of estimating the fueling rate \(S\) is to measure the emission of Dα light from recycling neutrals in the plasma boundary, since collisional-radiative modeling has shown [1] that, for plasma conditions typical in the tokamak edge, the rate of ionization of D atoms and the rate of emission of Dα photons are related by a factor that varies only weakly with electron density and temperature.

In most fusion devices, the complex spatial distribution of Dα light in the edge plasma makes the measurement of the total emission difficult in practice. Commonly spot measurements are made on the main sources of recycling, using photodiodes or photomultipliers; disadvantages of this approach are difficulty in accounting for the complex distribution of light from surfaces in contact with the plasma, uncertainty of the exact view of the detectors, and the need for a separate calibration for each detector. This paper describes the use of a CCD video camera at TEXTOR [2] for the purpose of spatially resolving the Dα light in order to measure more accurately the total emission so that \(\tau_p\) can be determined reliably.

2. Experimental setup

TEXTOR is a medium-size (major radius=1.75 m) limiter tokamak of circular cross-section, typically with plasma minor radius=460 mm, toroidal magnetic field=2.25 T, plasma current=350 kA, and a pulse length of about 4 s. The main source of recycling is ordinarily the pump limiter ALT-II [3], a toroidal belt divided into eight segments, or blades, each covered by an array of graphite tiles (14 toroidally by 2 poloidally); other components in the boundary that may give rise to significant recycling are a poloidal array of three radially movable graphite limiters (at the top, outboard, and bottom of the machine, all located in one toroidal sector), a graphite inner bumper limiter, and two ICRH antenna pairs. The results shown here are from ohmic discharges in which the plasma-facing surfaces were boronized [4]. Some of the results are from discharges in which the ALT belt was set to a minor radius of 450 mm in order to reduce the amount of recycling from other structures such as the bumper limiter (\(r=498\) mm). In these experiments, the poloidal limiters are retracted to a minor radius of 490 mm.

The camera has a tangential view, as shown in fig. 1. The view includes an entire ALT-II blade, half of one ICRH antenna pair, the top and bottom poloidal limiters, and two of the three gas feeds used for external fueling (in the same sector as the poloidal limiters). The outboard poloidal limiter and the inner bumper limiter cannot be seen simultaneously, but the view can be adjusted to include one or the other. Included in the view is a large opening in the liner. The data here are used to determine the brightness of the near edge of the plasma as well as to estimate the poloidal dependence of emission near, but not in front of, the limiter belt. An interference filter is used to select
the $H_\alpha/D_\alpha$ wavelength. The signal from the camera is stored on videotape for later analysis. The images are read from tape using an 8-bit frame grabber. A typical frame is shown in fig. 2.

In addition to the camera, the ALT-II blade is also viewed from above by a calibrated photodiode. The diode measures emission from a rectangular region 400 mm long poloidally (the blade width is 280 mm) and 57 mm wide toroidally, centered on the center of the blade. This measurement is used for calibration of the camera data. Since the diode view integrates poloidally across the blade, the measurement provides a brightness per unit toroidal length.

3. Analysis
In order to determine the total $D_\alpha$ emission from the discharge for a given time, the appropriate frame is read from tape into computer memory and the data are transformed to correct for the nonlinear response of the camera. Regions of the image corresponding to important sources of recycling are then selected, as in fig. 2, and the data are integrated over each region with the contribution from the edge on the near side of the plasma subtracted. Each integral is then multiplied by a factor proportional to the square of the distance from the lens in order to account for the smaller image size of objects further away.

The most important region is the ALT-II blade. To account properly for distance from the lens, it is necessary to divide the blade into several segments since the near end of the blade is much closer than the far end. The region is divided toroidally into seven segments of equal dimensions in $\phi$–$\theta$ space (see fig. 2), each including four tiles. A region equivalent to the view of the photodiode is generated based on the image coordinates of the central segment, and the integral inside this region is compared to the diode signal in order to calibrate the camera measurement. The emission from the one visible blade is multiplied by eight to determine the emission from the entire limiter belt. Similarly, the emission from the ICRH antennas can be estimated by multiplying the emission from one protection limiter by eight, to include both sides of both boxes of both antenna pairs.

Not all neutrals born at the limiter are ionized in front of the limiter; some can first travel a significant distance parallel to the last closed magnetic surface. It was observed in PDX, for example, that the $H_\alpha$ emission from a rail limiter decayed toroidally with a length of 0.5 m [5]. In TEXTOR, a halo exists around the limiter belt, extending over a substantial fraction of the plasma surface. On the assumption that no reflected light is seen in front of the long opening in the liner, the emission from the halo can be evaluated by investigating the poloidal profile of brightness in front of this opening. For the purpose of this evaluation, it is important to correct the profiles for contributions coming from the near edge of the plasma. To estimate the emission from the halo source, an exponential fit is made to the profile in front of the opening and is extrapolated poloidally to the top edge of the limiter. It is assumed that the halo on the bottom side of the limiter has the same decay length and that the emission there is stronger by the bottom/top brightness ratio seen in front of the opening. Integrating the two profiles poloidally and toroidally (assuming toroidal symmetry) gives an estimate of the total emission from this source. While the source of emission seen in front of the belt is almost entirely inside the last closed magnetic surface, a significant part of the halo emission may come from the scrape-off layer. This should be taken into account in determining $\tau_p$ if the halo source is large.

All emissions determined by this method are dependent on the calibration of the photodiode; thus a coating on the window for the diode affects the absolute measurement, but not the relative strength of various sources. The clarity of the window for the camera is not important. A significant source of uncertainty is errors in limiter blade alignment. The uncertainty of the limiter position for each blade is 1 mm and the scrape-off length for particle flux is typically 9 mm, so this introduces an uncertainty of about 10% in the belt emission. It is assumed that the reflection coefficient of the graphite tiles is zero, which introduces some error. Emission from the bumper limiter is not included in the results shown here; since the limiter is radially 38 mm or more behind ALT-II, however, this is not likely to cause a large error in the discharges under discussion. Two sources of uncertainty in the halo emission determination are (1) the assumption of toroidal symmetry and (2) signal-to-noise ratio, since the brightness in front of the opening is relatively low.

4. Results
The variation with toroidal angle of the brightness seen by the camera in front of the limiter blade is shown for a typical discharge in fig. 3. The brightness is substantially higher near the bottom side of the blade than at the top. This is consistent qualitatively with poloidal asymmetry in the scrape-off layer, although the difference is also sensitive to the positioning of the plasma. The near end of the blade appears brighter because of the varying angle of view along the blade; at the near end, the
line of sight passes through the plasma in front of the limiter at a relatively shallow angle, thus with a longer path length, so one pixel views a larger volume of plasma than at the far end. The effect of ripple in the toroidal magnetic field is clearly visible; the effect is stronger along the middle of the blade than near the edges, meaning the brightness cannot be expressed as a function of toroidal angle times a function of poloidal angle.

Poloidal profiles in the liner opening, corrected for contributions from the near edge of the plasma, are shown for various densities fig. 4. The decay length increases slightly with density and is typically about .25 m. If this is interpreted as a mean free path for electron impact ionization of deuterium atoms, the decay length is consistent (using an estimated $n_e < \sigma_{ion} = 1 \times 10^5 s^{-1}$) with a neutral energy on the order of a few electron volts.

The emission in front of the limiter, in the halo, and from other sources (gas feed plus top poloidal limiter plus ICRH antennas) are shown in fig. 5 versus line-averaged density for two different days of operation. The sum of "other" sources is on the order of 1% of the total emission. An important result of this work is that the halo source is not small compared to the limiter source. Aside from the fact that the sources are of comparable size, it is notable that the ratio is not the same for the two days.

When the calibrated photodiode signal is multiplied by the 12.0 m belt length in order to estimate the total $D_\alpha$ emission from the belt, the result is a factor of 1.3-1.6 larger than that found in front of the belt using the camera, although the diode signal is used to calibrate the camera data. There are two reasons for this. Firstly, the diode views a part of the blade that is relatively bright due to the ripple in the toroidal magnetic field. Secondly, the poloidal extent of the view is larger than that of the limiter, so the halo source as well as reflections from the liner may contribute significantly to the signal.

Estimation of the ionization rate in the plasma (and hence the particle confinement time) from $D_\alpha$ emission can be performed by using a constant ionizations-per-photon ratio; in order to obtain reliable values, however, computer modeling needs to be performed so that plasma profiles and atomic processes are properly taken into account. In particular, a substantial part of the halo source may come from the scrape-off layer. This work is in progress at TEXTOR. Preliminary confinement times based on a constant factor of 14 ionizations per photon, and not taking $Z_{eff}$ into account, are about 30-40 ms with little dependence on density, in contrast to the linear dependence of the energy confinement time $\tau_E$ on density seen under ohmic conditions.

5. Conclusions
A CCD video camera has been used to perform quantitative measurements of $D_\alpha$ emission in TEXTOR. This method offers advantages over measurements with spot detectors. It is found that the emission from a halo around the limiter belt is of the same order as that in front of the limiter and must be taken into account in determining the total $D_\alpha$ emission from the discharge. The emission increases with the particle content of the plasma, which is inconsistent with a confinement time proportional to density.


Fig. 1: Top view of TEXTOR showing region visible to CCD camera.
Fig. 2: Reproduction of a video frame from an ohmic discharge. Shown below are typical regions selected for analysis.

Fig. 3: Toroidal profiles along the ALT-II blade of brightness seen by the camera, shown for three poloidal positions on the blade.

Fig. 4: Poloidal profiles of brightness in the long liner hole, corrected for contribution from the near edge of the plasma.

Fig. 5: Dα emission in front of the limiter belt, from the halo around the belt, and from other sources (sum of gas feed, ICRH antennas, and top poloidal limiter) for two days of ohmic operation.
INFLUENCE OF RESONANT HELICAL WINDINGS ON PLASMA EDGE TURBULENCE

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1. Introduction

In order to study the possibility of externally controlling edge turbulence, and plasma edge properties, experiments using resonant helical windings (rhw) /1/ were performed in TBR-1 Tokamak. We have been investigating the modification of plasma density, potential, and magnetic field fluctuations, and temperature, during high MHD activity discharges. Here we present an analysis of the spectral modifications induced by resonant fields.

The turbulence was modified by $m = 4/n = 1$ resonant helical windings which as reported before /2/ altered the Mirnov oscillation amplitudes and slowed down the connected frequencies. These effects were also observed, although not in a so accentuated form, for density and potential fluctuations. However, electrostatic fluctuation induced particle flux spectrum was significantly changed by rhw.

2. Experimental

The work was carried out in the TBR-1, an ohmically heated tokamak /3/ (minor radius $a = 0.08$ m and major radius $R = 0.30$ m). The machine was operated for this experience with $T_e(0) = 200$ eV, $n(0) = 6 \cdot 10^{18}$ m$^{-3}$, plasma current $I_p = 8.5$ kA, toroidal field $B = 0.4$ T (corresponding to a safety factor $q(a) = 5$). The probe diagnostic system was designed to measure simultaneously, and within a short distance (few millimeters), electrostatic and magnetic fluctuations, besides relevant plasma parameters as density, potential, and temperature. The complex probe consists of a triple probe, four single probes, and two double magnetic probes /4,5,6/. Data were recorded at 1 MHz sampling rate and all spectral analyses were calculated from ensemble averaging over 62 blocks of 128 $\mu$s.

The external magnetic field perturbation was created by $m/n = 4/1$ electric current circulating in a set of helical windings located externally around the torus. The current circulating in those windings was adjusted in $I_h = 180$ A, and it was switched on after the plasma current has reached steady values, during time intervals greater than two milliseconds. All results were obtained with hydrogen working gas.
3. Discussion and Conclusions

Potential and density fluctuations show turbulent spectra with and without the application of \( rhw \). The fluctuation power is confined basically below 150 kHz. Density, \( \bar{n}/\langle n \rangle \), and potential, \( e\tilde{\varphi}/kT_e \), fluctuation levels have not shown significant change with the application of \( rhw \).

The plasma potential with \( rhw \) shows a small reduction for positions inside the plasma.

The relative dispersion \( \sigma_{k_\theta}/k_\theta \) for potential fluctuations inside the plasma changes significantly (almost double its value) with the \( rhw \) application. This seems to indicate a modification in the character of the potential fluctuations. However, this effect is not clearly verified for the density fluctuations.

The correlation times for potential and density fluctuations as well as the edge electron temperature did not change with the field perturbations.

The phase velocity was calculated from the signals of two probes separated by 2 mm. The results obtained with and without the application of helical fields show that the direction of the phase velocity is unaffected by the resonant fields; only a decrease in the value of the phase velocity was noted inside the plasma.

Fig. 1 shows the spectra for density fluctuations inside the plasma with and without the effect of \( rhw \). An accentuated increase of low frequency components was detected and this effect almost disappear in the limiter shadow. For potential fluctuations the observed effect is a small decrease of low frequency components of the spectrum.

Fig. 2 shows the spectra for the magnetic poloidal fluctuations. As reported before, there is a clear decrease in amplitude in the part of the spectrum that corresponds to the Mirnov oscillations and a downshift of the correspondent part of the spectrum was also verified.

Fig. 3 shows the coherency spectra between density and magnetic poloidal fluctuations, for Langmuir probe at \( r/a = 0.87 \) and magnetic probe \( r/a = 1.16 \); with \( rhw \) the coherence decreases for all positions investigated.

The radial particle flux due to electrostatic fluctuations was computed from simultaneous measurements of density and potential fluctuations and from measurements of the wave-number of potential fluctuations \( /6/ \). Fig. 4 is a typical example of the behaviour of particle flux spectra: \( rhw \) raised the values of particle flux inside the plasma. The total particle flux is plotted in Fig. 5, where the enhancement of transport by the application of external helical perturbation is observed. The observed sensitivity of the turbulence-driven cross-field particle transport to changes in the plasma edge properties indicates its relevance in the study of the properties of edge turbulence.

A full stabilization or at least a noticeable amplitude reduction of the MHD activity with the application of resonant helical fields has been recorded in all experiments \( /2/ \), as well as alterations on the turbulence and plasma edge profiles \( /7/ \).
References


Fig. 1: Power spectra of density fluctuations with and without rhw at r/a = 0.87.
Fig. 2: Power spectra of magnetic poloidal fluctuations with and without rhw, r/a = 1.16.
Fig. 3: Coherency spectra between density ($r/a = 0.87$) and magnetic poloidal fluctuations ($r/a = 1.16$) with and without $rhw$.

Fig. 4: Particle flux spectra with and without $rhw$ at $r/a = 0.87$.

Fig. 5: Profile of particle flux with $rhw$ (×) and without $rhw$ (○).
SELF-ORGANIZATION OF A HIGH $\beta$ VERY LOW q TOKAMAK PLASMA

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1. Introduction
In this paper, the existence of inherent self-organized magnetic configurations in a high $\beta$ very low q tokamak plasma ($\beta_\perp$$\geq$1, $q_0$<$2$) of the TPE-2 is described. It means that the plasmas necessarily relax to the inherent configuration if it is deviated by the electron thermal transport loss and the deposition of ohmic input. This might cause the release of excess energy and make the particle and energy flows, then degrade the confinement. These phenomena may occur not always at the core plasma ($q(0)$=1) but also at the edge plasma. This mechanism is proposed as ELMs in tokamak plasmas.

TPE-2 is a toroidal screw pinch device with an elongated cross section (b/a=21cm/13cm, $\kappa$=1.6, A=3.0) and a conducting shell /1,2,3/. High $\beta$ very low q tokamak plasmas ($\beta_\perp$$\geq$1, 1.2<$q_0$<$3$) are produced by ramping up of the toroidal and poloidal fields simultaneously in 2.8 $\mu$sec (fast screw pinch) and/or in 200–500 $\mu$sec (slow screw pinch). The plasmas are stably confined ($\tau_\psi$=1msec, $\tau_\tau$>1msec, here, the duration of $B_\perp$<1msec) in the parameter ranges of $T$=50–400 eV and $n_e$=0.5–5$x$10^{20}$m^{-3}$. It has the features that the initial plasmas with various $q_0$ values, $q$ profiles and various $\beta_\perp$ are able to be produced in a few $\mu$sec without energy and particle losses by a fast programming control of the poloidal and toroidal fields. Then, a fast relaxation process to the state is purely observed.

2. Self-organized $q$ profile
The low $q$ ($q_0$=1.2–1.8) or the high $q$ ($q_0$>$2$) stable plasmas with a nearly constant pitch profile are produced by keeping $q_0$ constant at the wall, as predicted before /4/. The initially implosion-heated plasmas transit to a stable state in 10–50 $\mu$sec, through the relaxation process accompanied by the large one turn voltage spikes, that is, the $q$ profile is self-organized through the relaxation process, if the initially programmed $q$ profiles deviate from the inherent one. Essentially, the profile does not depend on $q_0$, density and discharge modes /2/. The $q$ profile remains constant until the end of the discharge, except for the slight deviation. The $q$ profile of the relaxed stable state is close to the lowest free energy state that is calculated with the constraint of $\beta_\perp$=1 and with the constant plasma current /2/. Here, the pressure profile is assumed as $p=p_0[(\psi_s-\psi)/\psi_s]^{12}$, as the current density decreases toward the edge and is 0 at the wall and the pressure has the sharp boundary near the wall (5–10mm) in the experiment. The calculated $q$ profiles are nearly flat in the cylindrical sense, the current density is uniform over the plasma cross section except the edge region, and the $q(0)$ depends on $q_0$. The current density and $q$ profiles of the relaxed plasma are shown by broken lines in Fig.1. They are fairly flat profiles.

The density profiles also relax to broad profiles in 5–50 $\mu$sec simultaneously with the $q$ profile relaxation. The density profiles broaden with decreasing of $q_0$, The electron temperature profile does not depend on $q_0$, which suggests the profile consistency.

The relaxation transition time is slightly longer than the Alfvén transit time (=1–10$\mu$sec) and significantly shorter than the energy confinement time (=1msec) and the growth time of the resistive kink instability (=1msec).

3. Relaxation behaviour and $\beta_\perp$ effect
The low $q$ plasmas can be produced by decreasing $q_0$ from high $q_0$ plasmas ($I_p$ ramp up mode) as well. The same sort of the relaxation process appears drastically at the transition of $q$~2 and 3. When $q_0$ passes through near the rational number by decreasing the plasma current
of the relaxed plasma, which has the inherent $q$ profile, or by increasing $B_p$, the $q$ profile is distorted from the inherent one so as to keep $q_i$ constant (locked $q_i$ mode). A step appears near the rational number in the $q$ profile and it spreads from the middle region to the edge so that $q_i$ remains constant against the change of the current or $B_p$. In the case of decreasing the plasma current by the external circuit, the plasma internal energy remains constant and consequently $\beta_p$ increases. The $q$ profiles at $q_i \sim 2$ are shown by solid lines in Fig. 1. During the locked $q_i$ mode, the low shear stable surfaces near the rational numbers of $q$ exist and the plasma current has the tendency to flow especially on the surface. This change in $q$ profile may be understood by the conception that the plasma current flows along the magnetic line $/4/$. The edge fluctuations decrease and the energy confinement is improved. The strained state of $q_i$ sustains more than 100 $\mu$s, then, $q_i$ transits to more than 2 and the $q$ profile relaxes to the inherent one with a little energy and particle losses. After the transition, $q_i$ changes gradually, while the profile stays constant. The difference between the ramp up mode and the locked $q_i$ mode is that $\beta_p$ is larger and the plasma rotates in the latter case. At the high $\beta$ plasma the plasma current flows on the plasma edge region where $q$ is near the rational number.

The transition to the low $q_i$ plasma can be hardly attained in the case of an excessively elongated plasma and/or an inward shifted plasma that is attained by an excess vertical field. At the transition, the $q$ profile is not monotonic one with a pitch minimum. The necessary condition of the initial $\beta_p$ to achieve the transition to the low $q_i$ regime is that $\beta_p$ must be higher than the critical value $/2/$. At the transition, the high $\beta$ plasma causes the outward shift of the plasma column itself, which results in the reduction of $q_i((0)$ and deepening of the magnetic well near the axis of the column. It means that the energy loss at the transition must be small. High $\beta$ plasma has an effect of keeping the monotonic $q$ profile without any pitch minimum. The broad current density profile of the high $\beta$ low $q$ plasma seems to match that of the energy minimum state. Then, the high $\beta$ plasma plays an important role in the process of self-organization.

![Fig.1](a) Time variations of $q_i$, $1/\tau_i$, and edge plasma flow $j_e$ of the locked $q$ mode (solid line) and the relaxed mode (broken line). (b),(c) $q$ profile and current density profile at the transition of $q_i \sim 2$ of the locked mode (solid lines) and at the relaxed state (broken line).

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**Fig.1** (a) Time variations of $q_i$, $1/\tau_i$, and edge plasma flow $j_e$ of the locked $q$ mode (solid line) and the relaxed mode (broken line). (b),(c) $q$ profile and current density profile at the transition of $q_i \sim 2$ of the locked mode (solid lines) and at the relaxed state (broken line).
4. Rotation of edge fluctuations

The edge magnetic fluctuations $\beta'_p$ are measured at 37 poloidal and toroidal positions. The magnetic deformations from the relaxed state ($\{B_p - \overline{B}_p(\text{relax})\}/\overline{B}_p(\text{relax})$) are shown in Fig. 2. Just after the transition, $\beta'_p$ increases and rotating fluctuations appear. $\beta'_p$ in $|\theta| \leq \pi/2$ is large, $\beta'_p$ in $|\theta| > \pi/2$ is small and $\beta'_p$ at $\theta = \pi$ nearly equals to 0. It means that the fluctuations are localized near the outside of the torus. Then, the mode seems to be a ballooning mode ($m=2$, $n=2$). The deformation rotates roughly in the magnetic line of force and the direction is opposite to the plasma current. Here, we suppose that the magnetic wave packet does not propagate, but that the magnetic deformed plasma rotates to the direction, as same as usual tokamaks. The mode affects scarcely on the confinement, because the particle and the energy flows by the mode are local. The conducting shell may balance the dynamic motion and suppress the glow of the deformation by the plasma rotation. In the case that the relaxed profile is not produced, $m=2/n=1$ mode is dominant and the deformation occurs over the whole torus. The mode does not rotate and the ballooning mode may not appear until the high $\beta$ plasma is produced nearly at the end of the discharge. Then, the necessary condition for the high $\beta$ plasma to be stable is that the mode is rotating.

The rotating speed of the mode decreases with increasing of $q_t$ (10–25 km/sec). The rotation occurs at the initial phase of the discharge. The initial plasma is sometimes not in the complete relaxed state and has not any toroidal symmetry. Then it relaxes so as to be a toroidal symmetric plasma in 5–50 $\mu$sec.

Fig. 2 (a) Time variations of $I_p$, $q_t$, $\beta'_p$, $\tau_E$, (b),(c) The magnetic deformation from the relaxed state $(B_p - \overline{B}_p(\text{relax})/\overline{B}_p(\text{relax}))$ against $\phi$, $\theta$. Just after $\beta'_p$ increases, ballooning like modes rotate to the poloidal and toroidal directions.
One of causes of the edge ballooning deformation may be the nonuniformity of the magnetic surface which is produced by the shell cut.

5. Relaxation oscillation

The actual profile does not remain at a purely steady relaxed state, because the depositions of the ohmic input and the energy loss are different in space. The configuration of the dissipative plasma might be different from the energy minimum state. It may transfer to that of the dissipation minimum state in given magnetic energy /5/. It has rather peaked profiles and a higher free energy than the relaxed state. The real $q$ profile has a step near the rational number of the plasma edge, while $q(0)$ is more than 1. The current profile deviates from uniform profile to a double peaked one. This strained $q$ profile lasts for 100–500 μsec and relaxes to the inherent state in a few μsec, which is similar to that in Fig.1. As the phenomenon repeats and appears to be a small sawtooth, we call it a relaxation oscillation. The period of it decreases with decreasing $q$, and is 50–200 μsec in low $q$ plasma or more than 1 msec in high $q$ plasma. Usually, the change occurs over the whole torus simultaneously. The rapid macroscopic plasma motion, accompanied by the energy and particle losses to the outward edge of the torus, is observed at the relaxation. The relaxed profile is reconstructed by releasing the excess energy. These phenomena may be recognized to be a kind of ELMs. In the case that $q$ profile is carefully controlled to be the relaxed state, the relaxation oscillation disappears, and consequently, the plasma without particle loss and with a slight energy loss is obtained and then good confinement is established, in spite of fact that the rotating ballooning mode still exists and the density fluctuations appear near the edge region.

At first, the $q$ relaxation, that is, the magnetic flux reconnection, occurs at a certain point of the toroidal direction and the flux change propagates along the magnetic field line in Alfvén transit time (~1 μsec). The relaxation occurs over the 1/3 of full torus. When the deviation is not so large, the magnetic flux change may not be necessary to occur over the full torus. In the short time of the transition, the plasma returns to the relaxed state.

6. Conclusion

The experimental results are summarized as follows: The relaxation and self-organization of $q$ profile have been observed. The configuration with the self-organized $q$ and density profiles might be in a minimum energy state of a high β and very low $q$ plasma. The necessary condition for relaxation to the high β low $q$ stable state is that βp before and after the relaxation is high. The ballooning like deformation ($m=2/n=2$) occurs outside the torus in the relaxed plasma and rotates with the rotation of the plasma nearly to the magnetic line. The plasma rotation may have a stabilization effect on the mode and the deformation may not degrade the confinement. As the profile is deviated by the electron thermal transport loss and the deposition of ohmic input, it has a step near the rational number of $q$ and the plasma current has a tendency of flowing on near the rational surface. The strained state relaxes to the inherent profile. This might cause the release of excess energy and make the particle and energy flows, which then degrades the confinement. These phenomena may be recognized to be a kind of ELMs and might be understood as the relaxation oscillation which might be the repetitive transition from the dissipation minimum state to the energy minimum state.

REFERENCE

Analysis of edge fluctuations on the CASTOR tokamak

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Introduction
Edge fluctuations play a dominant role in the anomalous particle losses from tokamak plasma. In particular, recent lower hybrid current drive (LHCD) experiments in low density plasmas on ASDEX [1] and CASTOR [2] tokamaks demonstrated an improvement of the particle confinement at moderate lower hybrid powers, which is linked closely to a reduction of edge electrostatic fluctuations. The mechanism of the fluctuation reduction is discussed elsewhere [3]. This contribution is devoted to a more detailed characterization of electrostatic and magnetic fluctuations in this regime.

CASTOR tokamak
Experiments were carried out on the CASTOR tokamak ($R = 0.4 \, m$, $a = 0.085 \, m$) at $B_t = 1 \, T$, $I_p = 12 \, kA$ and densities $n_e = 2-6 \times 10^{18} \, m^{-3}$. For LHCD, the lower hybrid wave ($f = 1.25 \, GHz$, $P_{LH} \leq 40 \, kW$) was launched into the plasma via the three-waveguide multijunction grill [2] during the quasistationary phase of discharge. A brief survey of fluctuation measurements follows:

A) Oscillatory technique for $T_e$ - measurements (OH)
A sinusoidal voltage $V = V_0 \sin \omega t$ is applied to a Langmuir probe as shown in Fig. 1. The floating potential of the probe drops due to the rectification of electron current as [4]:

$$\Delta V_{fl} = T_e \ln \mathcal{J}(V_0/T_e) \quad (\approx V_0^2/4T_e \quad \text{for} \quad T_e > V_0)$$

where $\mathcal{J}(V_0/T_e)$ is the modified Bessel function and $T_e \equiv kT_e/e$.

Results of measurements are shown in Fig. 2. The frequency of sinusoidal modulation is 0.5 MHz, the amplitude $V_0$ is gated to determine the drop $\Delta V_{fl}$ by means of a single tip (the top trace). The middle trace shows the raw signal. The marks indicate the interpolated values of the signal with and without modulation. The resulting temperature (the bottom trace) is compared with evolution of $T_e$ measured by a triple probe located in the vicinity of the oscillatory tip. The satisfactory agreement suggests that the oscillatory technique can be used if strictly local measurements of the electron temperature are needed.

Further, we test this technique to determine the $T_e$- fluctuations. For this, three tips spaced poloidally by 2.5 mm were used, the sinusoidal voltage being applied to the central one. The floating potential in this spatial point, necessary for determination of $\Delta V_{fl}$, was estimated by the interpolation of data from the neighbouring floating tips. The relative level of temperature fluctuations $\bar{T}_e/T_e \approx 0.12$ obtained seems to be quite reasonable.

B) Mapping of the floating potential (OH)
The linear array of tips spaced poloidally by 2.5 mm was used to investigate poloidal asymmetries in the plasma edge and to determine the correlation length of electrostatic fluctuations. The probe array was movable on a shot to shot basis.

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1124 only), see Fig. 6b. It is interesting to note that their form is linked to the radial profile of autocorrelation time, see Fig. 6c. When the dimensionality in LHCD is higher than in the OH regime, the autocorrelation time is lower and vice versa. Therefore, we investigated the relation between these two quantities in more detail. Fig. 6d is a plot of the dimensionality of data from a shot versus the time delay d. Note that the typical autocorrelation time for our data is $2 - 7 \mu s$ ($d = 10 - 35$ samples). In this range of d, the dimensionality is practically independent on the delay and the difference between LHCD and OH regimes remains nearly constant. It seems, therefore, that the observed effects are not caused by an improper choice of the delay.

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References

Fig. 3 depicts evolution of the raw signals of floating potential (sampled by 1µs) from eight adjacent tips in a 2D plot for two radial positions of the array. It is evident that the perturbations propagate mostly from the tip No. 8 to the tip No.1 (which corresponds to the direction of electron diamagnetic drift) when the probe is within the limiter radius ($a > r = 83mm$), while a reverse tendency is apparent for a more outward position of the array.

The propagation velocity in poloidal direction was deduced by the correlation analysis. Fig. 4 manifests the variation of the form of crosscorrelation function with the radial position of the probe (2D plot, two tips are spaced by 5 mm). The direction of poloidal velocity corresponds to the cross-field drift velocity $(-\mathbf{E} \times \mathbf{B})/B^2$. The velocity shear layer is located close to the region with $E_r \sim -\partial V_{tf}/\partial r = 0$ (see Fig. 5).

C) Correlation analysis of magnetic fluctuations (OH + LHCD)
Poloidal magnetic fluctuations are monitored by two magnetic probes fixed inside the liner, their poloidal distance is 37 mm. The correlation functions of the probe signals are shown in Fig. 7. Typically, the cross-correlation function has two maxima. The maximum with a higher correlation ($C_{ij} > 0.5$) characterizes magnetic fluctuations propagating in the electron diamagnetic drift direction. We suppose, in concordance with Langmuir probe measurements, that these fluctuations have the origin within the plasma ($r < a$). The fluctuations related to the positive time delay of cross-correlation function are interpreted as a consequence of perturbations located within the scrape-off layer and propagating in the ion diamagnetic drift direction.

The time delay of both the maxima decreases with LHCD, which indicates an enhancement of the poloidal velocity of magnetic fluctuations which is also manifested by narrowing of the autocorrelation function. The similar effects are observed for the electrostatic fluctuations as well [6].

D) Dimensional and correlation analysis of density fluctuations (OH + LHCD)
To characterize whether the fluctuations are stochastic or chaotic (with less degrees of freedom than a noise), we calculated the correlation dimension of density fluctuations using the Grassberger-Procaccia algorithm [5]. The ion saturation current of a Langmuir probe is sampled by 0.2µs and 4096 samples (denoted as $x_i$) can be stored per a shot.

The dimensional analysis consists in construction of $n$-dimensional vectors from the data, $n$ is the embedding dimension:

$$r_i = (x_i + x_{i+d} + x_{i+2d} + \ldots + x_{i+(n-1)d}), \quad i = 1, 2, \ldots, 4096 - (n - 1)d$$

The correlation between the individual vector components should be negligible, therefore, we take the delay $d$ as the autocorrelation time $\tau$, defined as a halfwidth of autocorrelation function. Result for OH and LHCD regimes are summarized in Fig. 6.

Fig. 6a compares the dimensionality in LHCD and OH regimes for the probe position close to the limiter radius. The dimensionality in both the regimes seems to be lower than the same quantity for the computer-generated Gaussian noise in all embedding dimensions. Further, the LHCD data are closer to the noise than the OH ones. This suggests that the density fluctuations in the LHCD regime (with reduced fluctuations and improved confinement [2]) are more stochastic. Our limited set of data do not allow to reach the saturation of dimensionality (even if exists).

Further, we compute the radial profiles of dimensionality (in the embedding dimension
Fig. 3. Evolution of floating potential monitored by 8 tips (2D plot)

Fig. 4. Radial profile of $V_{\text{float}}$

Fig. 5. Crosscorrelation of tips 1-3

Fig. 6. Dimensional and correlation analysis of density fluctuations
CHARACTERIZATION OF THE RFX EDGE PLASMA

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Introduction The edge region of the Reversed Field Pinch (RFP) experiment RFX (R = 2 m, a = 0.457 m at the inner wall) /1/ has been investigated by an array of 6 Langmuir probes and 6 heat sensors mounted on a graphite limiter. The limiter is mushroom shaped and the probes are flush with the limiter surface. Each probe consists of a 3 mm diameter cylindrical graphite tip insulated by a machinable ceramic mount /2/. In this campaign two Langmuir probes, 4 mm beyond the tip of the limiter, have been operated in single probe configuration (50+200 Hz, ±150 V sinusoidal voltage sweep) and the other 4 were floating. The limiter has been protruded into the plasma up to 4 mm without any significant change in main plasma parameters. At the deepest insertion the surface temperature of the graphite rose up to 2000 degrees, with an incident energy flux of the order of 100 MW/m² as derived by the energy sensors 1 mm beyond the tip of the limiter. Taking into account the surface exposed to the plasma, the power collected by the limiter results <<1% of the input ohmic power during the current flat-top. Typical waveforms of the plasma current and line averaged density compared with the limiter floating potential are shown in fig.1.

Measurements The data refer to a range of toroidal current I and line averaged electron density nₑ of 500 < I < 700 kA and 2.5 < nₑ < 5×10¹⁹ m⁻³ respectively. As previously found in other RFP experiments /3/ the energy flux in the outer region of the plasma is strongly directional and exhibits a maximum when the collecting surface is exposed with the normal parallel
to the magnetic field and oriented towards the electron drift side. In fig.2-a
the asymmetry between energy sensors exposed to opposite directions is
reported as a function of the angle $\alpha$. The maximum is at $\alpha = -10^\circ$, where
the probes are aligned with the local magnetic field in the outer region, as
derived by magnetic measurements. In fig.2-b the energy flux asymmetry
is plotted as a function of $X$ (where $X$ is the difference between the inner
wall radius and the location of the tip of the limiter) for a fixed $\alpha = -10^\circ$.
During the limiter insertion the temperature and ion saturation current
have been monitored. The asymmetry shows an increase when the limiter
enters into the plasma and its floating potential changes sign, as shown in
fig.3-a.

The electron temperature on the ion drift side is $T_e \sim 10$ eV almost
uniform, with a slight tendency to decrease at deeper insertions (fig.3-b). On the other hand the ion saturation current, which with uniform
temperature is proportional to the electron density, exhibits an exponential
decay with two different decay lengths inside the port hole pipe and into
the plasma $\lambda = 15$ mm and $\lambda = 3$ mm respectively (fig.3-c). To derive the
particle flux the collection surface has been assumed to be the geometrical
surface of the probe, since the ion Larmor radius in the outer region is
comparable to the radius of the tip. The parallel flux at $X = 0$ mm results
$\Gamma_\parallel \sim 10^{22}$ m$^{-2}$ s$^{-1}$. Applying a particle balance equation in the Scrape Off
Layer originated right inside the pipe, with the decay length inside the port
hole pipe $\lambda = 15$ mm the perpendicular flux results comparable to $\Gamma_\parallel$, in
agreement with spectroscopic measurements of hydrogen influx /4/.

The ion saturation current has been measured at different angles
and the results are shown in fig.4, where the ratio between the current
collected by two probes in opposite directions and located at $X = -4$ mm is
reported. An angular dependence is observed with a maximum at $\alpha \sim -60^\circ$.

At fixed insertion $X = 0$ the electron density $n_e(\alpha)$ and temperature
$T_e(\alpha)$ at the edge have different behaviour with the line averaged density
$n_e$. Indeed $n_e(\alpha)$ tends to increase more than linearly whereas $T_e(\alpha)$ is
almost constant, as shown in fig.5.

Discussion The angular dependence of the energy flux at the edge
confirms that the energy transport in RFP is anisotropic and according to
the Kinetic Dynamo Theory (KDT) /5/ the asymmetry can be related to the presence of fast electrons at the edge. The power lost by these fast electrons is \( \sim 2/3 \) of the total power lost by transport, confirming that they are the main loss channel in RFP, as found in smaller experiments such as ETA BETA II /6/. These fast electrons can also account for the difference observed on the parameters between the ion and the electron drift side. Indeed on the electron drift side the current vs. voltage characteristic of the Langmuir probes changes sharply when the limiter enters the plasma. This behaviour can be explained as a distortion due to the current carried by fast electrons which in turn is affected by electron secondary emission processes. On the other hand the asymmetry between the ion saturation currents can be fitted by a simple model by which \( I_s = 0.5A_n(e_{ion} + v_{\perp}\cos\alpha + v_{\parallel}\sin\alpha) \) where \( A \) is the collection surface and \( e_{ion} \) the ion sound velocity. The perpendicular drift velocity \( v_{\perp} \) results \( \sim 25\% \) of \( c_{ion} \) and the parallel drift velocity \( v_{\parallel} \) \( \sim 10\% \) in the electron drift direction. In particular the perpendicular velocity is consistent with ExB drift due to an outward electrical field, inside the pipe, of the order of 1 kV/m.

**Conclusions** Energy and particle fluxes at the edge of the RFX experiment reveal anisotropic behaviour due to the presence of fast electrons and ion drift related to radial electrical field. In the range of plasma current and density explored so far, the electron temperature at the edge is almost constant at a value \( \sim 10 \) eV, whereas the electron density tends to increase more than linearly with the plasma density.

**References**

I. Introduction

Much effort is being devoted to try to evidence the role of the two radiation related mechanisms likely contributing to the edge turbulent transport [1-3]. Condensation instabilities, for which a significant coupling between density and temperature fluctuations is expected, are under study in a number of devices. In our previous work in the TJ-I tokamak [4], it has been reported that in the bulk side region of the velocity shear layer, temperature and density fluctuate at about the same level and in phase close to opposition. On the other hand, thermal instabilities, extremely sensitive to the detailed radiation cooling profile, are difficult to be studied separately because the plasma region where \[ \frac{\partial I_z(T_e)}{\partial T_e} < 0 \], being \( I_z \) the impurity radiation cooling, is typically well inside the magnetic separatrix and hardly accessible to probe measurements. Nevertheless, under some conditions, the thermal instability contribution to the radiatively driven turbulence can be significantly more relevant than the one due to condensation related phenomena [5].

II. Experimental

Density and temperature fluctuations have been measured in the plasma edge region of the TJ-I tokamak (R=30 cm, a=10 cm) by means of Langmuir probes. Measurements were performed in ohmically heated discharges with \( B_t=1 \) T, \( n_e=1.0 - 1.5 \times 10^{13} \) cm\(^{-3} \) and \( I_p=40 \) kA. The probe system consists of a square array of four tips (2 mm x 2 mm); pins are 2 mm long and 0.4 mm in diameter.

In order to measure electron temperature fluctuations one of the probes is working as a fast swept probe. A 500 W broadband amplifier was used to supply the swept voltage (~200 V) at 400 kHz to that single probe. The electron temperature, the ion saturation current and the deduced electron density (\( n_e \propto I_s T_e^{-1/2} \)) can be determined on a time scale of about 1\( \mu \)s. Further details of the measurement procedure and analysis have been described elsewhere [6].

Radial profiles of the total radiation losses have been measured using a 10-channels bolometer array with a radial resolution of the viewing chords at the plasma centre of about 2 cm. Detectors (germanium bolometers [7]) are arranged in a parallel geometry and look at the same plasma poloidal section than the probes, as displayed in figure 1. A movable stainless steel poloidal limiter was placed in the same port of the vacuum chamber (rectangular section.
vertically elongated) than the probes. The resulting set-up designed for this experiment enables to obtain a strong plasma-limiter interaction at the cross-talk volume of the uppermost detector in the bolometer array and the Langmuir probes. Consequently, different radiation profiles due to different impurity concentrations can be obtained.

III.- Results and Discussion

For the present experiment we have run series of repetitive discharges and studied plasmas were limited with the stainless steel limiter located at \( r_{\text{lim}} = 10.3 \) and at \( r_{\text{lim}} = 10.8 \) cm with respect to the equatorial plane of the vacuum vessel. The edge radial profiles of both, electron temperature and density, practically do not change with the limiter position.

Fourier analysis of the temperature and density fluctuations deduced from the fast swept probe technique show that fluctuations are dominated by frequencies below 100 kHz. The crossphase between \( T_e \) and \( \bar{n}_e \) seems to change from zero (for \( r/a_{\text{lim}} \leq 1 \)) to close to opposition (for \( r/a_{\text{lim}} \geq 1 \)) in the frequency range in which the coherence is well above the noise level.

Figures 2a and 2b show the typical time evolution of the chord integral measurements of the total power losses for discharges with the limiter placed at 10.8 cm (a) and at 10.3 cm (b). As can be seen, power losses profiles remain basically unchanged with exception of the uppermost detector signal. In both cases, losses coming from the edge were increasing all along the discharge, due to the strong interaction of plasmas with the limiter, but in fig. 2b, a more pronounced increase is observed. Contributions to these signals due to charged particles reaching the detectors through their long collimators can be neglected [8], and the flux of charge-exchange neutrals can be considered rather constant [9]. Therefore, the observed increase in bolometer signals is due to the enhancement of impurity radiation near the limiter. In the present case there is no need for an Abel inversion to realize that the emissivity profile is heavily peaked at the upper plasma edge.

In figure 2c, we have plotted the normalized root mean square (rms) values of the ion saturation current, \( \overline{I}_{\text{sat}}/I_{\text{sat}} \), electron density, \( \overline{n}_e/n_e \), and electron temperature fluctuation levels, \( \overline{T}_e/T_e \), versus the edge radiation for shots belonging to the above mentioned series in which probes were located inside the viewing chord of the uppermost bolometer and at \( r_{\text{lim}}-r_{\text{p}} = 0.2 \) cm. Analyses of fluctuation levels were performed during 6 ms in each discharge (the time intervals are delimited by the arrows drawn in figures 3a and 3b). For shots with profiles like the shown in figure 3a, at the lower levels of edge radiation, the rms values of the three magnitudes are very similar. As edge radiation increases, temperature fluctuations slightly increase and density and ion saturation current fluctuations slightly decrease, resulting in a ratio \( (\overline{T}_e/T_e) / (\overline{n}_e/n_e) \) of about 2 (see open symbols). However, and although edge radiation is notably higher and further increasing (up to a factor 3), as is the case shown in figure 3b, the rms values of \( \overline{T}_e/T_e \), \( \overline{n}_e/n_e \) and \( \overline{I}_{\text{sat}}/I_{\text{sat}} \) remain constant (see full symbols).
IV.- Conclusions

Considering that impurity radiation cooling drives electron temperature fluctuations, the high values of the rms electron temperature fluctuations observed in the plasma edge region of the TJ-I tokamak point out the possible role of radiative instabilities as a driving mechanism of edge turbulence [3].

As mentioned above, the present experiment shows an increase in the ratio \( \frac{\tilde{T}_e/T_e}{\tilde{n}_e/n_e} \) as the radiation profile tends to peak near the limiter radius (likely when \( \partial I_x(T_e)/\partial T_e \) tends to be negative). Taking into account that temperature fluctuations can be larger than density fluctuations in the plasma region where \( \partial I_x(T_e)/\partial T_e \) is the dominant drive for turbulence (thermal drive), the observed correlation between edge radiation and fluctuation levels reveals that thermal instabilities likely are an important drive for the edge turbulence measured at the upper plasma edge. Once the radiation profile peaks at the upper plasma edge, no relationship between the strength of the radiation (\( R \propto n_x I_x \)) and fluctuation levels is observed.

Finally, it has to be noticed that the fluctuation levels, especially those of electron density, measured in the present experiment in the proximity of the shear layer location at the upper plasma region of TJ-I, are significantly smaller than those previously reported in the outer region of the equatorial plane [4]. Therefore, and taking into account the additional asymmetry introduced by the presence of the limiter, these results demonstrate the existence of strong poloidal asymmetries in the edge turbulence of the TJ-I tokamak.

Acknowledgements

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References

Figure 1.- Experimental setup. D germanium detectors; C collimator tubes.

Figure 2.- Time evolution of the line integrated profiles of plasma emissivity for shots with the limiter at a) 10.8 cm and b) 10.3 cm. c) Normalized rms fluctuations of the ion saturation current, electron density and electron temperature as a function of the edge radiation. Open symbols refer to a) and black ones to b).
INTRODUCTION

By means of visible spectroscopy, particle confinement property, recycling and fluctuation of the edge plasma on the HL-1 tokamak in various discharges (lower hybrid current drive (LHCD) plasma, electron cyclotron heated plasma, pellet injection plasma, biased electrode and limiter) have been studied. The effect of LHCD on the particle confinement and the statistical results have been obtained in a variety of operating conditions. In our experiments, the global parameters of plasma are as below: \( I_p < 110 \text{kA}, \bar{n}_e = (0.5-3.0) \times 10^{19} \text{cm}^{-3}, B_t < 2.8 \text{T}. \)

EXPERIMENTAL RESULTS

1. In the experiment of LHCD discharge on the HL-1 tokamak, LHW power is coupled into plasma by four wave guide phrased array coupler, the frequency of LHW is 2.45 GHz, while the frequency of ECR is 75 GHz. Figure 1 gives evolution of some parameters in LHCD plasma where LHW is transmitted at the counter direction of the electron drift (i.e., anticurrent drive). The intensity of plasma bremsstralung (5360Å) and the line emission of impurity measured by VUV obviously decreases during the period of LHCD. It is found that the effect of LHW in various densities on the plasma is not the same. In the low density (\( \bar{n}_e < 2.0 \times 10^{19} \text{cm}^{-3} \)), whether it is at the electron drift direction or the counter direction of it, the density increases by (40-70)\%, and H\( \alpha \left( \text{D}\alpha \right) \) emission decreases by (20-40)\%. Because the ionization rate is proportional to the emission intensity, the results above during the period of LHCD, have shown that the ionization rates of the neutral hydrogen (or deuterium) and the impurity or ions decrease. It can be concluded that, the increase of plasma density during the period of LHCD is caused by the improvement of particle confinement. When the plasma density exceeds \( 2.0 \times 10^{19} \text{cm}^{-3} \), the plasma density is not increased, but the H\( \alpha \left( \text{D}\alpha \right) \) and impurity emission are increased obviously, suggesting that the particle confinement of the plasma deteriorates. The statistics has been done during LHCD period in a variety of operating conditions. Fig. 2 and 3 give the dependence of particle confinement on the input power of LHW and the line-average density of plasma respectively. In ECRH experiment plasma parameters are shown in figure 4. During ECRH period, the temperature of the core plasma goes up obviously, suggesting the
heating effect, the increase is obvious in Hα (Da) emission and bremsstrahlung and impurity emission, so the ionization rate goes up of all species, the number of particles should increase if the other conditions did not change, the result of the experiment shows that the average density tends to decrease, so the particle confinement deteriorates. In our experiments, the particle confinement time decreased by (20—70)% . Because of the heating effect on the plasma during the ECRH period, the neutral and ionized impurity at the lower stages is ionized, the effective charge of the plasma increases by 50%-80%.

2. In the pellet injection experiment on the HL-1 tokamak, the typical plasma parameters are shown in figure 5. After the injection of the pellet into the plasma, perturbation will be caused in the plasma, the transports of the particle and the energy have been changed, obvious increase of the density is observed (the relative value is about 140%). Hα (Da) emission increases fastly at various toroidal positions, and soon afterwards, it decreases obviously until it's less than that before the injection, that's shown, the influx of neutral particles decreases, and the same the loss of the plasma particles, the particle confinement is improved (by about 70%). Because of the recycling reduction after the injection, the influx of all species is reduced, this effect, in addition to the dilution caused by the added fuel, leads to the reduction of Zeff (by about 40%).

3. In the LHCD and biasing electrode experiments, the improvement of particle confinement takes place with the suppression of the fluctuation of Hα (Da) emission (see figure 6).

CONCLUSION

In the LHCD experiment, the physical mechanism of the improvement of particle confinement is not clear yet, by the results above, it may be related with the drift instabilities, the suppression of the instability leads to the improvement of the particle confinement. By the analysing above, it can be concluded as bellow:

- The particle confinement time depends on the line average density \( n_e \) during the LHCD period. If \( n_e < 2.0 \times 10^{13} \text{ cm}^{-3} \), the particle confinement is improved when the LHW is coupled into the plasma in whether the current drive (\( \Delta \phi = \pi / 2 \)) or the anticurrent drive (\( \Delta \phi = -\pi / 2 \)) direction, and it is (1.5-3.5) times larger than that in the ohmic heating at the same time, the suppression of Hα (Da) emission fluctuation at the boundary is observed, if \( n_e > 2.0 \times 10^{13} \text{ cm}^{-3} \) the particle confinement deteriorates. On the statistical analysis, it is also found that, if the input power of the LHW is less than 250kW the particle confinement time tends to increase with the input power of lower hybrid wave, and it's maximum value is reached near the density \( 1.0 \times 10^{13} \text{ cm}^{-3} \).
- The deterioration of the particle confinement and the increase of the effective charge are always observed during ECRH period when \( n_e < 2.5 \times 10^{13} \text{ cm}^{-3} \).
- In the experiments of pellet injection, the particle confinement time observed
after pellet injection is increased by about 70%, and the effective charge decreased by about 50%.

- The way by biased electrode and limiter with positive voltage can improve the particle confinement time, and the fluctuation of Hα (Dα) emission of the boundary layer is suppressed.

REFERENCES


Fig. 1 Evolution of some parameters of plasma during LHCD period. \( P_{\text{LH}} = 51 \text{kW}, B_t = 2.2 \text{kG} \) (SHOT 11443)

Fig. 2 Dependence of \( \tau^\text{LH}/\tau^\text{OH} \) on the line-average density (\( \tau^\text{LH} \) and \( \tau^\text{OH} \) represent the particle confinement time during LHCD and ohmic discharg period respectively) (a) anticurrent drive (\( \Delta \phi = -\pi / 2 \)) (b) current drive (\( \Delta \phi = \pi / 2 \))
Fig. 3 Dependence of $\tau_p^{\text{LH}} / \tau_p^{\text{SH}}$ on the input power ($P_{\text{LH}}$) of LHW

Fig. 4 Some parameters of plasma during ECRH period

\[ P_{\text{ECRH}} = 220\text{kW}, \quad B_r = 2.72\text{kG} \]

(SHOT 12544)

Fig. 5 Evolution of some parameters of plasma after pellet injection (SHOT 12357)

Fig. 6 Fluctuation suppression of Hα (D α) emission during LHCD together with biasing electrode. BI - biasing current.

(biasing voltage is 465V, \( P_{\text{pl}} = 114\text{kW} \), \( B_t = 2.3\text{kG} \), SHOT 11831)
Study of Nonlinear Structures in Electrostatic Flute Type Fluctuations

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Techniques of analysing the time series with the concept of modes substituted by quantities such as dimension, Lyapunov exponent and the Kolmogrov entropy have developed in recent years. These are few tools of nonlinear dynamical system analysis. They have similar analog in plasma physics. For example, dimension analysis of turbulent fluctuations (continuous power spectrum) in fusion devices (Prado and Fiedler-Ferrari 1991 and references therein) and in laboratory devices (Strühle and Piel 1989) suggest that for the coherent fluctuations the calculated correlation dimension is low (nearly equal to the number of modes excited) while it is large for turbulent state. The dimension calculated is nearly equal to number of competing modes. Similarly in plasma physics we speak of unstable (growing) and damped (decaying) waves. It is analogous to Lyapunov exponent which determines the divergence (for positive Lyapunov exponent) or convergence (for negative Lyapunov exponent) of neighboring trajectories exponentially in a time series. Number of Lyapunov exponent is same as the dimension of phase space. Matric entropy of the system is sum of the positive Lyapunov exponents which determines overall stability of the system.

During the study of low frequency instabilities in toroidal device (Prasad 1993) we have observed that the density and potential fluctuations makes a transition from coherent multimode state (exhibiting a considerable nonlinearity as seen from the bicoherence spectrum) to a turbulent state with increase in magnetic field. We have applied the tools of non-linear dynamical system analysis on density fluctuations obtained at different magnetic field. The results of the analysis are presented.

Given the phase space trajectory of the system, the tools of non-linear dynamical systems analysis can be used to quantify the dynamics of the system. One such measure is the fractal dimension which gives an estimate of number of effective degrees of freedom in the system. For a chaotic system Hurdorff's dimension is always non-integer. Takens (1981) has demonstrated that it is possible to reconstruct an equivalent phase space from the measured time series of a single variable by the method of delays. One creates a set of $D_E$ dimensional vectors whose components are just time delayed values of original time series: $\tilde{z}(i) = [x(t_i), x(t_{i} + \tau), ..., x(t_{i} + (D_E - 1)\tau)]$ where $t_i = i\delta t$, $\delta t$ is the sampling time, $\tau$ is the time delay between successive elements of vector, and $D_E$ is the embedding dimension. Grassberger and Proccacia (1983a) defines the correlation integral as:

$$C_2(r) = \frac{1}{N} \sum_{i=1}^{N} \frac{1}{N-1} \sum_{\substack{j=1 \atop j \neq i}}^{N} \Theta \left( r - \| \tilde{z}_i - \tilde{z}_j \| \right)$$

where $\Theta$ is called Heaviside function and $\| \|$ is the Euclidean norm.

The time delay $\tau$ could be chosen arbitrarily provided the amount of data is infinite and noise free. However, in experiment finite amount of noisy data makes the choice of $\tau$ critical. In practice either autocorrelation e-folding time or first minima of mutual information are used. If autocorrelation time is large compared to sampling time then highly correlated points are included in the Heaviside function resulting in spurious low dimension. Theiler (1986) has proposed a method which can take care of this effect by slightly modifying the correlation integral:

$$C_2(r) = \frac{2}{(N-W)(N-W+1)} \sum_{j=W+1}^{N} \sum_{i=1}^{N-j} \Theta \left( r - \| \tilde{z}_{i+j} - \tilde{z}_j \| \right)$$

where $W$ is the width of the window.
where \( W > \tau \left( \frac{2}{N} \right)^{2/D_E} \)

For \( W = 1 \), one recovers equation 1. Eckmann and Ruelle (1992) have shown that, at least \( N = 10^{D_E/2} \) data points are necessary to reliably estimate fractal dimension \( D \). These calculations of dimension can be cross checked using the method of surrogated data set suggested by Theiler et al., (1991). In this method, one randomises the phases of the Fourier transform of the original time series and then inverts the transform. Using these surrogated time series, the dimension is recalculated (see for details of implementation Theiler et al., 1991). If the results are not significantly different than those of the original time series, the dimension estimate should not be trusted.

The experiment is performed in toroidal device (Prasad et al., 1992) with major and minor radii of 45 and 15 cm respectively. A variable toroidal magnetic field (TF) up to 1 kG can be applied. The plasma is produced at \( 10^{-4} \) Torr of Argon gas pressure by striking a discharge between the cathode (incandescent tungsten filament is placed vertically at major radius of 36 cm) and the anode (vessel wall). Langmuir probes are used as diagnostics during the experiment. The fluctuations in density \( \bar{n} \) and floating potential \( \bar{\phi} \) are picked up by probes at radial location of 6 cm during this study.

The plasma produced has peak density \( \sim 10^{11} \text{cm}^{-3} \), electron temperature 4 eV and ion-temperature 0.2 eV. Density and potential profiles are given in Prasad et al., 1992. The variation of spectral characteristics as a function of TF is shown in figure 1. The fluctuations below 200 Gauss do not exhibit any marked peak and power spectrum is flat indicating the absence of coherent wave phenomena. At 200 Gauss of TF, the fluctuations observed in \( \bar{n} \) and \( \bar{\phi} \) exhibits well defined peaks at frequency \( f \) \( 3.7 \pm 0.3 \) kHz together with peaks at integer multiples of this frequency with mode number \( m = 1, 2 \) and 3. Measurement of radial, azimuthal and toroidal propagation characteristics of these waves suggests that they are of flute nature. Rayleigh–Taylor instability under the influence of velocity shear is identified as mechanism behind the generation of \( m = 1 \) (Prasad 1993).

The data set at 200 Gauss, consisting of 8192 points with a sampling time of 25 \( \mu \text{sec} \), is divided into 128 records of 64 points each and used for bispectral analysis (Kim and Powers 1977). Bispectrum spectrum (shown in figure 2) exhibit strong peaks at interaction frequencies \( (f, 2f) \) and \( (f, 3f) \) suggesting that \( m = 2 \) and 3 arises due to strong nonlinear interaction of \( 1, 1 \) and \( 1, 2 \) modes respectively. Further increase in TF results in gradual transition to turbulence. The moments of distribution function viz., Kurtosis and skewness as a function of TF (shown in figure 3) increase with increase of TF. This indicates the evolution of bursty nature of oscillations. The study of spectral characteristics at 1000 gauss (Prasad et al., 1992) suggests that long wavelength fluctuations are dominated by R–T instability and short scalelengths by drift waves respectively. Data set at 600 Gauss, consisting of 32000 points with a sampling time of 20 \( \mu \text{sec} \), is divided in to 250 records of 128 points each for the analysis purposes. The bispectrum spectrum at 600 Gauss (see figure 4) exhibit finite bispectrum of 0.3 indicating the presence of coherent structures.

Algorithm suggested by Grassberger and Proccica (1983a) by making use of floating point representation of number is used to calculate correlation dimension. We have used two \( \bar{n} \) time series, one at 200 Gauss where fluctuations are coherent and other at 600 Gauss where fluctuations are turbulent for correlation dimension analysis. The autocorrelation function for 200 and 600 Gauss time series is shown in figure 5a and b. Correlation integral calculated for \( \bar{n} \) time series at 200 gauss with \( \tau = 40 \) and \( W = 1 \) is shown in figure 6a. In order to find linear scaling region we have used local slope analysis. In this analysis slope of a straight line of every three points of the \( \log C(r) \) versus \( \log r \) curve was calculated and plotted as a function of middle point. Local slopes for different em-
bedding dimension (shown in figure 6b) converges to 4.7. However, value of W calculated from equation 3 comes out to be 7.6. Thus the low dimension could be spurious due to autocorrelated effect. We recalculated correlation dimension with W = 10. The results of the analysis shown in figure 7 do not show a significant variation hence autocorrelation effect is negligible. In order to crosscheck dimension calculation we used method of surrogated data suggested by Theiler et al., (1991). \( C_2(r) \), calculated using the method of surrogated data set, and shown in figure 8, exhibit typical characteristics of random noise. All the above analysis suggests that the dynamics of system is $\sim 5$. This is slightly higher than the number of waves excited. The higher dimension could be due to several factors affecting the dimensionality calculation. It could be due to the presence of small amplitude noise and/or the effect of filtering.

The matric entropy determines overall stability of the system and can be calculated from the vertical spacings of log \( C(r) \) versus log \( r \). Matric entropy thus calculated (given in figure 9) approaches a saturated value of 0.018/sampling time, implying that the system is unstable at least in one dimension or at least one of the modes is unstable. Number of Lyapunov exponent is same as the dimension of phase space so in principle more than one mode can be unstable. Due to the statistical (Eckmann and Ruelle 1992) limitation it is difficult to estimate the number of unstable modes.

Correlation integrals and their slopes with \( r = 7 \), \( W = 1 \) for 600 Gauss \( n \) time series (W calculated from equation 3 is < 1) are shown in figure 10. Slopes converges to 5.6. By surrogating this data set the dimension calculated did not increase with increase in embedding dimension (shown in figure 11) indicating spurious saturation and dimension of the system could be greater than 9. Hence a large number of modes are competing.

Our observations are consistent with dimensionality calculation of turbulent fluctuations in fusion devices (Prado and Fiedler-Ferrari 1991) and in laboratory plasmas (Ströhlein and Piel 1989). This can be interpreted in the following way: as the magnetic field is increased more and more modes exceed the finite larmor radius threshold and becomes unstable. They interact nonlinearly in a complicated way resulting in turbulence. Hence observation of high dimensionality is not unexpected.

To summarise, a dimensional analysis of flute type electrostatic fluctuations in toroidal low-$\beta$ plasma as a function of TF is reported. The transition from coherent state to turbulent state with increase in TF is accompanied by increase in bursty nature of oscillation. The system exhibit low dimension at coherent state and dimension increases of TF. These results are in agreement with other laboratory measurements.

References

TAYLOR DISCHARGE CLEANING AND GLOW DISCHARGE CONDITIONING IN NOVILLO TOKAMAK


ABSTRACT

Both Taylor and Glow discharge conditioning techniques in Novillo Tokamak are applied. In each case the results are showed. The effectiveness of each one of these two conditioning techniques was monitored by measuring the gas impurities \( \text{CH}_4, \text{C}_2\text{H}_4, \text{CO}, \text{C} \) and \( \text{O} \) by means of mass spectrometry. In every case tokamak discharges were carried out after conditioning and the better plasma parameters obtained were used to determine the conditioning quality. A 20kW, 17.5kHz power oscillator was used for Taylor discharges. The oscillator energized the OHT, synchronized with a 200-700G toroidal magnetic field. For the Glow discharges the same power oscillator DC voltage source was used to bias two SS electrodes with a 0-1500V and 2A maximum discharge current.

INTRODUCTION

In the tokamak main discharge operation, a clean vessel environment is required in order to obtain good plasma parameters. Usually, this preoperational work is mainly focussed to have an important carbon and oxygen elements reduction in the plasma.

The vessel conditioning is usually a very slow process, that takes a long time. For this, it is necessary to find out the more suitable technique and the parameters of the plasma and machine, which make more efficient the conditioning process.

In the Plasma Physics Laboratory at I.N.I.N., a small tokamak has been constructed and its operational stage has been started. Its main parameters are: \( B_T = 0.5T, I_p = 12kA, T_e = 150eV, T_i = 40eV \) [1]. The vacuum vessel of the Novillo Tokamak is made of stainless steel 316L, it has four 90° elbow sections separated by viton O-rings for vacuum
scaling and electric insulation. The major and minor radii $R = 23\text{ cm}$, $a = 6\text{ cm}$ respectively and 0.32 cm wall thickness. The total area on the chamber surface is approximately $7264\text{ cm}^2$ with 28 access ports representing $617.78\text{ cm}^2$ of effective access area [2].

In the Novillo Tokamak, two conditioning techniques have been applied: Taylor Discharge Cleaning (TDC) and Glow Discharge Cleaning (GDC). The TDC process was carried out energizing the OHT with 2pps using a 20kW, and 17.5kHz oscillator synchronized with a 200-700G toroidal magnetic field, produced with a capacitor bank of $C = 600\mu f$, $V = 900\text{ V}$ [3]. The experimental arrangement for GDC involves two electrodes positioned within the vessel chamber in different vertical ports. These anodes in the experiment are two cylindrical stainless steel bars of 0.62 cm of diameter. They were connected to 0-1500V dc power supply through a load resistor to limit the current flowing in the plasma. The negative pole of the power supply and the vacuum vessel were grounded.

**TAYLOR DISCHARGE CLEANING**

After a base pressure of $2.5 \times 10^{-7}\text{ Torr}$ was reached, the hydrogen was injected up to a work pressure between $6.0 \times 10^{-8}$ to $4.6 \times 10^{-4}\text{ Torr}$. Although sometimes the low energy discharge without preionization was obtained, the curves shown correspond to experiments with a tungsten filament used as a preionization source. The toroidal magnetic field was in the range of 200 to 700G, the corresponding voltage of the synchronized toroidal bank varied from 210 to 900 volts. The OHT power applied was in the range of 10 to 20kW, varying the oscillator power in the same interval [3].

When the toroidal magnetic field and the preionization current are lower, greater is the power oscillator required to obtain the breakdown.

In figure 1 we have a plot of a low energy peak to peak plasma current vs. work pressure. The maximum of these curves is in the range of $1.0 \times 10^{-4}$ to $2.0 \times 10^{-4}\text{ Torr}$.

The range of mass spectrometer used was from 1 to 50 amu, where we can find the most important contributions, in the case of the spectrum displaying the atmosphere inside the vessel during TDC process. The purpose of this was the suppression of the biggest hydrogen spike, in such a way that the other contributions can be seen.
The base pressure of $2.5 \times 10^{-7}$ Torr, before starting TDC process; the water vapor is the highest partial pressure, which is a typical pattern of the analysis of the residual gas in the vessel without conditioning. The spike corresponding to oxygen is also present, figure 2a.

Also a spectrum obtained during the TDC discharge process is presented figure 2b. Rises in methane and water vapor spikes are remarkable. These results coincide with the Taylor report [5]. Finally after the TDC discharge process the spectrum obtained shows that all the spikes have been modified, and the one corresponding to oxygen is practically insignificant.

Figure 2a. Base pressure before TDC.
Figure 2b. Base pressure after TDC.
GLOW DISCHARGE CLEANING

The voltage at which the glow discharge can be started is function of hydrogen pressure. It was not possible to start the discharge at $H_2$ pressure lower than $10^{-2}$Torr because a very high breakdown voltage was required. However, the operating $H_2$ pressure, up to $10^{-1}$Torr permitted reduce the breakdown voltage =500V to the desired value.

The current-voltage characteristics measured between $3 \times 10^{-1}$ and $5.2 \times 10^{-3}$Torr with stable conditions can be obtained. At higher pressures, increasing the current density, this situation can be controlled by limiting the flowing current with a higher resistor in series with the external power supply circuit.

This process did not clean the chamber at all, conversely to the expected, the chamber acquired thin film electrodes metal of deposits. In this case the main discharge had a short time of duration, there was no difference between the tokamak discharge, before and after of the GDC.

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REFERENCES

ANALYSIS OF MAGNETIC TURBULENCE DURING PELLET ABLATION AND RESPONSE OF FUELED PARTICLES BY PELLET INJECTION TO THE SOL AND DIVERTOR


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[1.] Introduction

A divertor study is one of the most important subjects in helical systems. The experimental studies have been performed on this area in Heliotron E [1–4]. The purpose of this paper is to clarify the following points: (1) How long is the connection length of the field line, along which the particles lost from the plasma edge are transported to the divertor? (2) How does the ergodic region observed in mapping [5] at 500 Gauss connect to the divertor region? We analysed a "delayed" particle response to pellet injection at the divertor to estimate the connection length. By changing the region of the particle deposition we studied a relationship between the edge ergodic region and the divertor.

[2.] Density pulse propagation and related magnetic fluctuations

Hg/Dg pellets are injected on the midplane along the major radius [6] with a velocity of 400–600 m/sec. The line density nl is measured with a far infrared interferometer separated toroidally by πRo from the pellet port. Hα′el light from a pellet cloud is measured at the opposite side of the pellet port. The increment of nl is delayed by τdelay with respect to the increment of Hα′el. τdelay is in the range of 250 μs ± 100 μs in various plasma conditions, which indicates that the particles fueled at the injection port propagate along the field line of 3 πRo with a velocity of Cs = √Te/M, where Te ∼ 10 eV. A sharp pulse-like density rise is typically observed in ECH plasmas, in which particles are deposited mainly near the edge. The field line tracing, in which lines are launched on the midplane at the pellet port, is done to understand whether this pulse is related to parallel transport of the ablated dense plasma, or not. A limiter is used to change an initial particle deposition region. For α > L = 21 cm the density pulse was observed along #0 and #1 chords viewing the outboard torus side, but for α > L = 15 cm it was measured along #5 and #6 inboard chords, as shown in Fig.1. From comparison of these observations with line tracing results, it is considered that the former pulse corresponds to the propagation of particles ablated at ΔR(= R – Ro) = 23 – 25 cm and the latter does that at ΔR ∼ 18 cm. Both regions are inside by ΔR = 2 – 5 cm from the last closed magnetic surface.
During this ablation phase bursts of coherent magnetic waves are observed \([7, 8]\). These bursts grow rapidly after pellet injection with a delay time of \(50 - 100 \mu s\) and last for several hundreds microseconds. The mode analysis of Mirnov coils shows that coherent waves have mode numbers of \(m/n = 2/1, 3/2, \) and \(1/1\) etc., where \(m(n)\) is the poloidal (toroidal) mode number. The amplitude of \(B_\theta\) seems to be insensitive to the density rise \(\Delta n\) and bursts are usually damped within < 500 \(\mu s\). This means that \(B_\theta\) may not be caused by the change in radial pressure profile and may be related to the relaxation process in which a new equilibrium state is being established through rapid parallel transport. Fluctuations of \(H_{\phi}^{rel}\) are also found to be coherent with \(B_\theta\) at several frequencies (10–80 kHz) as shown in Fig. 2. Since \(H_{\phi}^{rel}\) measurement is done along the major radius and temporal variation of \(H_{\phi}^{rel}\) corresponds to the spatial trajectory of the pellet, \(H_{\phi}^{rel}\) may reflect the spatial modulation of pellet ablation, that is, incident heat flux on the pellet cloud. Low mode numbers of \(B_\theta\) and good correlation between \(H_{\phi}^{rel}\) and \(B_\theta\) support that a resonant interaction between the incident heat flux on rational surfaces and pellet ablation causes some parts of fluctuations.

Thus, these observations of delayed propagation of a density pulse and fluctuations associated with the rational surfaces \(\left( \epsilon_{s} = \frac{a}{m} \right)\) confirm that a dense plasma propagate along the field line \(\left( \sim \pi R_0/\epsilon_{s} \right)\) with the velocity of \(C_s\) at \(T_e \sim 10 eV\).

[3.] Delayed response of \(I_{++}\) at the divertor

In order to estimate the effective connection length \(L_{eff}\) of the field line from the plasma edge to the divertor, we analyse the delayed response of ion saturation current \(I_{++}\) with Langmuir probes at the divertor after pellet injection. We assume that particles deposited near the edge also propagate to the divertor with the velocity of \(C_s\). For this purpose edge fueling is desirable. Thus, \(\tau_{delay}\) measurement was done in ECH plasmas. By shifting the outermost magnetic surface (by \(\pm 1 \text{ cm}\)) particle deposition near the boundary was studied. Line density along the chord outside the boundary (\(+1 \text{ cm}\)) shows no sharp increase, but that along the chord just inside the boundary (\(-1 \text{ cm}\)) increases rapidly. This indicates that particles are deposited inside the outermost surface. Thus it can be expected that a fast and large \(\Delta I_{++}\) is a response of particles fueled inside the outermost surface. Results of typical probe response are shown in Fig. 3. \(\Delta I_{++}\) is delayed typically with a time of \(100 - 500 \mu s\), and variation in \(\tau_{delay}\) for probes at different divertor footprints is small, although the amplitude of \(\Delta I_{++}\) is fairly different between them. From \(\tau_{delay}\) and an assumed propagation velocity of \(C_s\), \(L_{eff}\) is estimated to be an order of \(\pi R_0\). In order to investigate the properties of “ergodic” region, we studied \(\tau_{delay}\) as a function of \(<a>\), that means if the magnetic surfaces are closed \(\tau_{delay}\) should be longer by \(\tau_{\perp}\) than \(\tau_{||} = L_{eff}/C_s\), where \(\tau_{\perp}\) is a characteristic time for particles to
diffuse from the deposition region to the vacuum outermost surface. This was done by inserting a rail limiter into NBI plasmas, in which the particles are deposited mainly near the core. $\tau_{\text{delay}}$ is increased up to 6ms as shown in Fig.4. A short $\tau_{\text{delay}}$ of < 1ms is, however, observed at particular divertor positions, which suggests that some parts of the "ergodic" region connect directly to the particular divertor positions. It is also found that a density pulse and its response on $I_s^+$ are affected by the global particle confinement properties. Density clamping phenomenon is typically observed in low density ($n_e \lesssim 1 \times 10^{13} \text{cm}^{-3}$) ECH plasmas[9], in which enhanced $I_s^+$ and $H_{\text{wall}}^+$ emission depending on rf power are observed. In this case, a response on $I_s^+$ of a pellet shows a sharp pulse-like increment of $I_s^+$. However in afterglow or plasmas heated at second harmonic cyclotron frequency of 106GHz[10], $I_s^+$ gradually increases by very small amount, or does not increase and furthermore the density pulse is also not observed. These results suggest that improvement of global particle confinement affects both parallel transport in the core region and the outflux to the divertor.

[4.] Summary

Application of pellet injection to an estimation of $L_{\text{eff}}$ is tested by analysing particle parallel transport. From $\tau_{\text{delay}}$ measurement of $\Delta n/l$ and analysis of $B_0$ a resonant interaction between pellet ablation and the rational surfaces, and existence of cold dense plasmoid propagating along the field line are found. From the delayed response on $I_s^+$ at the divertor, $L_{\text{eff}}$ is evaluated to be $\sim \pi R_0$.

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References
[5.] F.SANO et al., 19th EPS (1992)489
[6.] S.SUDO et al., in Tech. Committee Meeting, IAEA(Gut Ising, FRG, 1988)
[8.] H.ZUSHI et al., TCM on Pellet Injection, IAEA (1993)
[10.] K.NAGASAKI et al., this conference (1993)
$n_l(\#0), H^\text{pel}_\alpha$ for $\alpha \geq 21 \text{ cm}$

Fig. 1  (a) $n_l(\#0), H^\text{pel}_\alpha$ for $\alpha \geq 21 \text{ cm}$

$H^\text{pel}_\alpha$ for $\alpha \geq 15 \text{ cm}$

Fig. 1  (b) $n_l(\#6), H^\text{pel}_\alpha$ for $\alpha \geq 15 \text{ cm}$

Fig. 2 cross coherence between $\tilde{D}_\alpha$ and $H^\text{pel}_\alpha$ (a) before injection (b) during ablation

$\tau_{\text{delay}}(\text{ms})$ after pellet $H_\alpha$

Fig. 3 $I^+$ at two different divertors, $n_l$, and $H^\text{pel}_\alpha$ are shown.

Fig. 4 $\tau_{\text{delay}}$ vs. $\alpha$; open/closed circles represent two groups of probes
Bias Experiments in Heliotron-E


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Abstract

The effects of the edge electric field on the plasma performance are studied for ECH/NBI currentless plasmas in Heliotron E. The edge electric field was modified by applying a bias voltage between a limiter and the wall. It is found that the behavior of particles in both the core and the edge regions is affected by this modification. It is suggested that the particle confinement is improved by the negative limiter bias.

1. Introduction

Since the edge plasma works as a boundary condition for the core plasma, it is one of the key issues to study how we can control the core plasma through artificially changing the boundary condition in a helical system. As one step in this study, we have tried to control the edge electric potential or electric field in Heliotron E by applying a bias voltage between a wall and a limiter and/or a "small" electrode which does not work as a limiter. Since the parameters of the confinement field such as $\sqrt{2\pi}$, $\Theta$ and $\Theta_n$ are strong functions of the minor radius, the limiter can change these values at the plasma edge. The combination of the small electrode and the limiter allows us to investigate the role of the region just inside the outermost surface and the effect of the field parameters on this kind of edge modification. In this paper, however, we mainly describe the limiter bias experiment.

In the previous experiment, a mushroom-type carbon limiter (10 cm$^3$ x 10 cm$^2$) was biased for ECH plasma. In that experiment, however, a cold and dense plasma still remained in the outside region of the limiter. Moreover the mechanical shaft of the limiter, which worked as a feeder in the bias experiment, was not insulated from the plasma. Therefore, the flow channel of the bias current was not clear. In this experiment, a rail-type carbon limiter (23 cm$^3$ x (4(top) - 8(bottom) cm$^2$) x 4 cm$^2$) was inserted from the bottom of the torus. Only the limiter head was kept biasing through a discharge.

2. Basic characteristics of the limiter bias in Heliotron E

When a bias voltage, $V_b$, is applied to the limiter, a bias current, $I_b$, flows between the limiter head and the wall. When the limiter is inserted deeply enough in the original outermost surface, $I_b$ must flow across the field line. On the other side, in the case of shallow insertion, $I_b$ can flow along the field line since a part of the limiter head is linked to the wall by the divertor field line. Then, the shallow insertion of the limiter will mainly bias a part of the field lines in the divertor layer, and the bias effect on the magnetic surface will be observed when the limiter is inserted deeply enough.

Figure 1 shows an example of the limiter position, $Z_{lim}$, dependence of $I_b$ (ECH plasmas). The line density of the core plasma, $n_d$, and $V_b$ are also plotted. The bias current remarkably decreases for both polarities of $V_b$ when $Z_{lim}$ is less than 24 cm. The similar tendency of the $Z_{lim}$-dependence of $I_b$ is also observed for NBI plasmas. Taking account of the above discussion, the decrease of $I_b$ is explained by the change of the effective circuit resistance from $R_{para}$ to $R_{perp}$ ($> R_{para}$). It is consistent with the observations that the limiter heat load increased with the limiter insertion and becomes to saturate for $Z_{lim} < 24 - 25$ cm. At the same time, the particle flux to the divertor footprint, $\Gamma_{wall}$ decreased. These mean that the limiter configuration is achieved from this limiter position.
How much current can be drawn for an applied voltage between the wall and an electrode inserted into the magnetic surface depends on $R_{\text{perp}}$. When the electrode works as a limiter, however, the reduction of $\Gamma_{\text{wall}}$ causes another limitation on $I_b$ since the maximum current which can flow into the wall is restricted to the ion (electron)-saturation current for the positive (negative) bias case. It depends on the plasma density and temperature near the wall. Actually, the observed $V_b - I_b$ characteristic shows the saturation of $I_b$ at much lower level for $V_b > 0$ than the current for $V_b < 0$.

3 Bias effects in the divertor configuration

In a shallow insertion case, the bias voltage modulation caused an asymmetric response in the divertor particle flux along the torus. Figure 2 shows the time traces of $V_b$, $I_b$, $\Gamma_{\text{wall}}$ at three different positions (#9.5-1, #9.5-6 and #27.5 sections)\(^{(4, 5)}\) and the $H\alpha_{\text{wall}}$ intensity at two different positions (#17.5 and #33.5), respectively. The solid lines represent for the positive bias case and the other is for the non-bias case. As in Fig. 2, $\Gamma_{\text{wall}}$ (#9.5-1) is increased by the positive bias, but $\Gamma_{\text{wall}}$ (#27.5) is decreased.

Since the divertor field line crosses the wall, only one part of the divertor layer can be biased, where the field lines connect the limiter and the wall. Then, the observed asymmetric response of $\Gamma_{\text{wall}}$ suggests that the flow is modified by the $E \times B$ drift due to the local gap of the plasma potential by the limiter biasing. The similar modification of the divertor flux is observed in IMS\(^{(9)}\), where some of the discrete divertor plates were biased. It is interesting to note that the same effect is obtained in the limiter bias (correctly, an electrode bias) as that in the divertor plate bias. This technique is applicable to control the divertor load; to avoid unfavorable concentration of the divertor heat load or to concentrate the particle flux to a preferable position.
4 Bias effects in the limiter configuration

When the limiter was inserted deeply enough, the limiter bias changed the edge potential or the electric field. Figures 3 shows the radial profile of the space potential (with respect to the wall), $V_s$, in the limiter SOL. The plasma potential was estimated from the floating potential, $V_f$, and electron temperature measured with a fast reciprocating triple probe ($V_s - V_f + 3x T_e$). For $V_b > 0$, $V_s$ is high and its radial profile is similar to that in the non-bias ($V_b = 0$) case. For $V_b < 0$, $V_s$ is almost the same with the non-bias case in the region far from the limiter radius. However, the potential rapidly decreased near the limiter radius. In the limiter SOL, the density fluctuation was remarkably reduced by the negative bias. On the other hand, the positive bias decreased the density in the limiter SOL, but the fluctuation level, $\tilde{n}/n$, was almost the same with that for the non-bias case.

Fig.4 Time traces of $n_\text{p}$, $H\alpha_{\text{wall}}$ (2 positions), $H\alpha_{\text{lim}}$, O V line radiation (62.97 nm), soft X-ray and ECE signals with negative bias (solid line) and without bias (dashed line). (NBI plasma)

Recently, ATF group reported the improvement of the particle confinement by the positive bias of the limiter\(^7\). In Heliotron-E, "good" signs for the particle confinement were also observed in the negative bias case. Figure 4 shows the time traces of $n_\text{p}$, $H\alpha_{\text{wall}}$'s at the two positions, $H\alpha_{\text{lim}}$, an impurity line intensity (OV 62.97 nm), soft X-ray and ECE signals. The limiter position is $Z_{\text{lim}} = 18$ cm. The solid lines are for the negative bias case ($V_b = -200$ V) and the dotted lines are for the non-bias case. The NB injection power were kept constant. Both the $H\alpha_{\text{wall}}$ and the $H\alpha_{\text{lim}}$ was decreased by the negative bias while the core plasma density remained almost the same. This suggests the improvement of $\tau_p$. The observed $\tau_p$ improvement depends on the limiter position and the bias voltage. For the limiter position of 22 cm case, the $H\alpha_{\text{lim}}$ has a minimum value at -100 V. On the other hand
the deeper insertion case shows monotonous decrease with increasing negative bias.

The bolometer signal or the impurity line intensity signal increases with increasing the bias voltage especially for $|V_b| > 100\,\text{V}$. The preliminary analysis shows that $\tau_e$ seems to be almost the same for the negative bias case and slightly decreased for the positive bias case (assuming the same heating power). The increase of the radiation loss might mask the bias effects on the particle or energy transport. Since the radiation from all charge-state of the impurity ions (O, C, Fe, Ni) was enhanced by the limiter bias for NBI plasma, it is considered that the influx of impurities into the core region is increased. As the reasons of this increase of the influx, we must consider the increase of spattering at the wall (and/or the limiter head) and reduction of the shielding effect of the limiter SOL.

For ECH plasmas, remarkable increase of the soft X-ray intensity and peaking of its radial profile was observed for the both polarities of the bias voltage. Since the electron temperature and the density of the core plasma were almost the same in this case, the impurity transport should be modified. For NBI plasmas, such clear peaking of the profile was not observed. In this case, however the peaking might be canceled by broadening due to MHD-activity. The detailed study of the impurity transport will be done in the next experiment.

Relating to the limiter bias, an interesting change in the slowing-down spectrum of the energetic trapped ions was observed. The limiter insertion made the depletion in the spectrum(6). This is considered to be explained by the reduction of the loss-boundary due to the limiter insertion and the resonant loss due to negative $E_r$. The negative bias, however, increased the flux at an energy range. If the limiter bias slightly affects on the intrinsic electric field, the resonant loss-cone structure of the deeply trapped particles might be changed. (The possibility of such modification of the intrinsic electric field by the limiter bias is also discussed in Ref.2) Detailed data of the electric field in the core region and a Monte Carlo simulation of the high energy particles in the limiter configuration are necessary.

5. Summary

The limiter bias experiment was performed in Heliotron E for ECH/NBI currentless plasma. The observed effects of the limiter bias are

(1) The edge plasma potential, edge electric field were controlled by the limiter bias.
(2) Limiter biasing affected on the particle transport; (a) The divertor flow was modified for shallow insertion case. (b) The reduction of the density fluctuation was observed in the limiter SOL. (c) The improvement of the particle confinement was expected for negative bias.

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References

Impurity behavior in Heliotron E

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[1] Introduction
Impurity control has been the critical issue for obtaining high temperature plasmas in various magnetic confinement devices. Biased limiters and electrodes are used to access the high confinement regime, H-mode, and to control impurity influxes [1]. In tokamaks, negatively biased limiter provides plasmas with high impurity concentration [2]. In helical systems, the ATF has been reported that biasing the limiter had no effects on the density of the intrinsic impurities [3]. In Heliotron E [4], recently a carbon limiter is used for imposing the radial electric field between the plasma edge and the vacuum chamber wall [5]. The maximum applied voltage was 300 V for both polarities. In this paper, the dependence of the impurity influxes on the radial electric field imposed by the biased limiter is described. A new method to study the screening effects for impurities is also explained.

[2] Impurities in biased limiter plasmas
A carbon limiter is biased to the vacuum chamber wall from -300 V to +300 V to control plasmas and impurities. Currentless plasmas were generated by 53 GHz gyrotrons and further heated by neutral beams with a power of 1 MW. The injection energy was 23-26 kV. The magnetic field strength is 1.9 T. The average electron density was \(3.5 \times 10^{13} \text{ cm}^{-3}\) and the central electron temperature was 500 eV. For biasing the limiter, the average electron density and the central electron temperature changed only slightly. The edge electron temperature, however, decreased in biased limiter plasmas. Figure 1 shows the VUV impurity spectra for biased(+-) and non-biased limiter plasmas. The wavelength region is from 180 Å to 400 Å. In biased limiter plasmas, the spectral intensities of oxygen, carbon and iron increase drastically compared to the non-biased limiter plasma. The dependence of the spectral intensities of
oxygen O V(193 Å) and iron Fe XVI(335 Å) on the bias voltages is shown in fig.2. Since the plasma parameters in the core region were almost same and the spectral intensities of the successive ionization stages increased, the impurity densities and influxes are large in the biased limiter plasmas. The increases of the influxes due to the enhancement of impurity production and decrease of the screening efficiency in the SOL are considered as the most probable reason. However, the charge exchange neutral fluxes on the chamber wall, measured by a neutral particle analyzer, did not increase as the bias voltage became large. The electron density measured by electrostatic probes decreased in the edge region for the biased limiter plasmas. The increase of the impurity density is considered due to the change of the screening efficiency in the scrape off layer.

[3] Si-probe to study the screening effects for impurities

It is important to measure the screening efficiency for the impurity influxes in the SOL region for understanding impurity behavior. A new trial to study the screening effects has been started in Heliotron E. Silicon, as a tracer impurity, is evaporated from a Si-probe placed near the outermost magnetic surface, which is limited by the limiter inserted from the bottom of the chamber. The source rates are estimated from the heat fluxes on the probe, which can be measured with electric probes installed in the same Si-probe. The density of the silicone in the plasma is evaluated by the VUV spectrometer. Figure 3 shows the Si XI spectral intensity for the various positions of the Si-probe in two cases of the limiter radii at 260 mm and 280 mm, respectively. The experimental results so far are limited in the non-biased limiter plasmas. The arrows indicate the positions of the outermost magnetic surfaces at the toroidal section where the Si-probe located. As the Si-probe is inserted into the plasma from the outermost magnetic surfaces, the Si XI spectral intensity increases drastically at 225 mm with the limiter radius at 260 mm. The electron density changes from $2 \times 10^{12}$ cm$^{-3}$ at 270 mm to $7 \times 10^{13}$ cm$^{-3}$ at 210 mm, and the electron temperature is about 20 eV. The variation of the electron density and temperature is small for explaining the large changes of the Si XI intensities. The ratio of Si XI to the ion saturation currents is shown in fig.4. This value approximately represents the ratio of the Si density in the plasma to the influx of Si atoms. The ratio is small outside of the outermost magnetic surface and increases as the Si-probe approaches and enters the outermost magnetic surface. This means that as Si atoms ionize near the outermost magnetic surface, it becomes easier to enter into the plasma. This experimental result confirms that the screening effects can work
in the SOL region of Heliotron E magnetic configuration. The ionization lengths of Si atoms, which corresponds the spatial resolution of this method, are estimated less than 1 cm from these plasma parameters. The advantage of this technique is to evaluate the screening effect at the desired region by surveying the Si-probe.

References
[5] T. Mizuuchi et al., in these proceedings

FIG.1. Impurity spectra in the wavelength region from 180Å to 400 Å for biased, and non-biased limiter plasmas.

FIG.2. O V(193 Å) and Fe XVI(335 Å) spectral intensities as a function of bias voltage (V).
FIG. 3. Si XI spectral intensity as a function of the Si-probe position. Two arrows indicate the locations of the outermost magnetic surfaces in the toroidal section of the Si-probe.

FIG. 4. The ratio of the Si XI spectral intensity to the ion saturation current as a function of the Si-probe position, and ion saturation current measured by electric probes installed in the Si-probe.
LOW ENERGY NEUTRAL PARTICLE ANALYSIS
AT THE STELLARATOR W7-AS

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Introduction

A Low Energy Neutral particle Analyzer (LENA) has been installed at the Stellarator W7-AS. LENA has been successfully operated at the ASDEX-Tokamak /1,2/. From LENA measurements valuable information on the $T_i$-profiles near the edge and on the plasma-wall-interaction was obtained.

The measuring principle is based on a time-of-flight (TOF) method. The charge exchange (CX) flux is mechanically chopped in bunches of 1 μs duration and the TOF distribution over 2m flight path are measured. The available energy range is from 20 to 1000 eV for $H_2$ and 40 to 2000 eV for $D_2$ discharges. Because of counting statistics is the time resolution during a discharge limited to 100 ms, when CX-energy distributions are analyzed. For the consideration of the integrated flux $\Gamma_{cx}$ or the mean energy $<E>$ the smallest possible time resolution is 2 ms.

At W7-AS LENA’s line of sight is in the midplane ~1° from the direction of the major radius. The toroidal location is between triangular and elliptical planes. Through the same port as LENA a photodiode monitoring the $H_2$ radiation is viewing the area where the CX particles detected by LENA originate. During the last campaign LENA was routinely taking data.

During W7-AS operation with 140 GHz ECRH at high density a spontaneous transition to a plasma status with improved confinement was observed which shows essential the features of the Tokamak H-mode /3/. The general conditions for and the observations at the H-mode transition are discussed in the paper of Erckmann et al. in this conference /4/. In this paper the observation with LENA during the H-mode are presented.

Results

We are observing a striking change in the shapes of the energy distributions of the low energy CX-neutrals after the H transition. Fig. 1 shows the CX spectra for a $D_2$-discharge before (+++) and after (solid line) the transition. The spectrum before the transition shows a broad maximum at ~100eV. This shape is typical for ECRH discharges also with 70 GHz and at lower densities. After the H-transition the CX flux increases considerably below 100 eV and the maximum shifts to ~50 eV (the drop at very low energy is not well
supported by the data because of the statistical errors). In the energy range of 200 to 300 eV a flattening of the spectrum is observed for D₂ discharges.

In Fig.2 the influence of the H transition on a H₂ discharge is shown. Just before the transition again the broad maximum at ~100 eV is visible (++). This is also seen earlier in the discharge (***) where the density was approximately half the value during the H-mode. Therefore the overall intensity is much lower. After the H-mode transition the spectrum is quite different: It rises cautiously down to 10 eV. (The apparent data scattering is due to rather poor counting statistics. Note that the energy scale is 2 times stretched compared to Fig.1 since E/M is plotted.)

The response of the integrated CX-flux onto the H transition is shown in Fig.3, upper trace. The H-transition occurs at 0.78 s as it shows up in the characteristic sharp drop of the Hα(3) signal, which monitors the Hα radiation from the upper limiter (lower trace). The Hα monitor viewing the plasma along the same line as LENA Hα(1) responds with an apparent drop. The signal recovers within ~50 ms, while the Hα(3) signal stays low with the exception of bursts due to ELMS. The integrated CX-flux follows closely the variation of Hα(1).

**Discussion**

The observed CX spectra are due to a line integral which involves the neutral density n°(x), the ion density n_i(x), and the ion temperature T_i(x), where x is the coordinate along the line of sight. It has been shown [1] that T_i(x) near the plasma edge can be obtained from the measured spectra, when n°(x) can be simulated and n_i(x) is known. The n°(x) simulation requires the knowledge of n_e(r) and T_e(r) and the exact geometry of the walls in the vicinity of LENA (in 3 dimensions). This is especially difficult in the present case with l>0.5 where non closed flux surfaces occur and the measured profiles cannot easily be transferred to other toroidal locations. At ASDEX we observed in all cases CX-spectra which rose continuously to the lowest observable energies. It is not known so far, what the reason for the occurrence of the maximum at ~100 eV in the W7-AS ECRH-spectra is. From our ASDEX knowledge, however, we can estimate for the present line density that the origin of the particles below 100 eV is the region from the steep gradient of n_e(r) to beyond the separatrix at r_eff=0.16 m. Hence substantial changes of the plasma at the edge must occur which cause the changes in the CX-spectra. However, it is possible that at the present l>0.5 there exists an appreciable density outside the separatrix at our toroidal location. From other experiments we know that the gas recycling at the "helical edge" is the most important neutral gas source for LENA. The helical edge passes the wall close below the LENA port. From the response of Hα(1) (Fig.3) to the H transition and the constant l it is impossible to explain the change of the CX-spectra at low
energies by a change of the neutral gas source. Hence a change of \( n_1(x) \) and/or a decrease of \( T_1(x) \) must occur in this range of \( T_{eff} \). This would shift the origin of the low energy CX neutrals further outward.

**References**

/4/ V. Erckmann et al

![Energy distribution graph](image.png)

**Fig. 1** Energy distributions of the low energy CX neutrals before (***) and after (solid line) the H-transition for a D₂ discharge
Fig. 2  Energy distribution of the low energy CX flux during the H-mode (solid lines), before the H-mode (++) and in an early phase with half the density (**) for a H₂ discharge.

Fig. 3  Time evaluation of the low energy CX flux (upper trace), the Hₐ(1) signal from the LENA location (center) and the Hₐ(3) signal from the upper limiter (lower trace). The H-transition occurs at the dashed line (H₂-discharge).
ON THE ROLE OF NEUTRAL PARTICLES ON EDGE TURBULENCE AND ELECTRIC FIELDS IN THE ATF TORSATRON

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I. Introduction

Turbulence driven by neutral particles and radiative instabilities has been considered one of the dominant processes to partially account for the observed edge turbulence characteristics [1-6]. The possible role of neutrals in determining confinement has also been discussed [7,8]. Neutrals can affect directly the ionization and the charge exchange sources. Ionization effects have been theoretically considered as a possible driving mechanism of edge fluctuations [2-5]. At the lowest temperatures (< 10 eV) charge exchange is much more probable than ionization mechanisms.

In this paper we report experimental evidence of edge turbulence and edge electric fields modified by the presence of neutrals.

II. Experimental

Edge fluctuations have been measured in the plasma edge region of the Advanced Toroidal Facility (ATF) (l = 2, M = 12 field-period torsatron with R₀ = 2.10 m and a = 0.27 m) using Langmuir probes. The probe system consists of a square array of four Langmuir probes located one field period away from the instrumented rail limiter [9]. Two tips, aligned perpendicular to the local magnetic field, are used to measure the floating potential fluctuations, from which the poloidal phase velocity of the fluctuations can be deduced. The location of the velocity shear layer is Zshear = 43 cm. Measurements were taken in the radial range Z - Zshear ≈ (5, -2) cm.

Plasma edge turbulence has been characterized by measuring the fluctuations in the ion saturation current (Iₛ = neTₑ¹/²) of a Langmuir probe, using the experimental methods previously described [9]. It has to be noted that only when temperature fluctuations are negligible fluctuations in the ion saturation current can be interpreted in terms of local density fluctuations (∆Iₛ/Iₛ = ∆n/n).
The experiments were performed in ECRH plasmas with heating power $P_{ECH} \approx 200$ kW, average electron density $n_e \approx (4-6) \times 10^{12}$ cm$^{-3}$, stored energy $S = 1$ kJ, and magnetic field $B = 1$ T.

III. Results and discussion

Fig. 1 shows the radial profiles of the electron temperature ($T_e$), the edge electron density ($n_e$), floating potential ($\phi_f$) and probe current fluctuations ($I_s/I_S$) in the plasma edge region for hydrogen plasmas with different line average densities ($\bar{n}_e$). The larger $n_e$ the lower the edge electron temperature. The edge electron density slightly increases with $n_e$ except when the electron temperature is close to 10 eV. The floating potential becomes more negative with increasing $T_e$. In the radial range $Z_{shear} - Z = (3-5)$ cm, and at the higher (35 eV) and lower (10 eV) temperatures the density profile and the radial electric field remain constant ($\approx 2$ V/cm) while $I_s/I_S$ increases when $T_e$ is close to 10 eV.

Fourier analysis of the probe current fluctuations reveals that the sharp increase of $I_s/I_S$ at the lower temperatures ($\approx 10$ eV) is due to the increment of the amplitude of the broadband fluctuations with frequencies below 100 kHz (Fig. 2).

The edge electron temperature has also been modified by means of neon puffing experiments in deuterium plasmas [10]. When due to the strong injection of neon impurities the electron temperatures stay below 10 eV, a substantial increase in $I_s/I_S$ is observed. The stored energy in the plasma is higher for the lower fluctuation levels.

The sensitivity of the turbulence levels to the local electron density is independent of the method of edge cooling: gas puffing with the working gas (hydrogen) or radiative impurity (neon) cooling.

The influence of the electron temperature on the edge electric fields (velocity shear layer [1,9]) has been studied in the temperature range (10-40 eV). The value of the edge electric field is independent of temperature when the electron temperature is higher than 20 eV, but decreases strongly as $T_e$ approaches 10 eV (fig. 3). Concurrent with this drop in the radial electric field the level of probe current fluctuations increases.

These results can be understood taking into account the temperature dependence of the ionization rate coefficient: $<\sigma v>_i$ decreases strongly when the temperature is below 15 eV, whereas it is rather constant for temperatures above 15 eV [11] (fig.4). As shown in fig.1, the edge electron density increases with increasing $n_e$ when $T_e$ is above 15 eV (i.e. in the temperature range where $<\sigma v>_i$ is basically temperature independent). For further increase of $n_e$, $T_e$ becomes close to 10 eV and $<\sigma v>_i$ decreases strongly. The neutral density at the edge increases while the electron density goes down.

The increase of $I_s/I_S$ and the decrease of $E_r$ when the electron temperature approaches 10 eV can be interpreted in terms of the influence of neutral particles on edge turbulence and electric fields.
IV. Conclusions

In the temperature range where the ionization rate coefficients are strongly temperature dependent, the electron temperature plays an important role in determining the level of edge turbulence and the value of the self-generated radial electric fields. These results provide evidence of edge turbulence and electric fields modified by neutral particles.

Acknowledgments

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References

Fig. 1. - Radial profiles of edge density, electron temperature, floating potential and probe current fluctuation levels in plasma with different line averaged densities; $n = 6.5 \times 10^{12}$ cm$^{-3}$ (A), $n = 5.5 \times 10^{12}$ cm$^{-3}$ (+), $n = 4.7 \times 10^{12}$ cm$^{-3}$ (x), $n = 4.5 \times 10^{12}$ cm$^{-3}$ (a).

Fig. 2. - Power spectra at three different temperatures.

Fig. 3. - Radial electric field and probe current fluctuations versus edge electron temperature at the shear location.

Fig. 4. - Rate coefficients for ionization and charge-exchange [11]
TEMPERATURE, DENSITY AND POTENTIAL FLUCTUATIONS BY A SWEEPED LANGMUIR PROBE IN WENDELSTEIN 7-AS

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1. INTRODUCTION

Numerous experiments using a Langmuir probe to investigate the magnitude of temperature fluctuations and their contribution to heat transport in the edge region of tokamak plasmas have been carried out (1,2). Sweeping the voltage applied to a tip fast enough to ensure that the ion saturation current, floating potential and electron temperature may be assumed to be constant during the sweep is experimentally more difficult than alternative schemes but this disadvantage is compensated by the ability to measure all three of these quantities at one spatial location (3,4). Sweep frequencies up to 600kHz have been employed to obtain the current-voltage characteristic. A radial scan in the vicinity of the velocity shear layer on W7-AS stellarator was performed. Inside and outside the shear layer the normalised magnitude of the temperature fluctuations was found to be approximately 30% larger than the magnitude of the electron density fluctuations, approaching a value of 0.12 and 0.09 respectively at a radial position 1cm inside the shear layer. An increase in the coherency of the temperature, floating potential and density fluctuations between tips with a poloidal separation of 2mm was also measured as the shear layer was crossed. Heat conduction produced by correlated temperature and poloidal electric field fluctuations is therefore possible. An increasing coherence of temperature and floating potential fluctuations leads to an increase in the coherence of temperature and plasma potential fluctuations as the shear layer was crossed.

2. EXPERIMENT

The Langmuir probe characteristic was obtained by sweeping the applied voltage between the tip and the vacuum vessel. Transformer coupling of the broadband amplifier ( Amplifier Research ; 1000W CW from 10kHz to 220MHz ) and a dummy BNC cable were used to balance the capacitive currents. The cable length from the transformer to the probe tip was 5m and so for an applied voltage at 400kHz with an amplitude of 100V a capacitive current of 125mA would flow in one of the BNC cables. With compensation this capacitive current could be reduced to 5mA. Low noise broadband buffer amplifiers were used to deliver the applied voltage monitor and current probe output to the data acquisition system ( 8-bit Nicolet digital oscilloscope ; Model Pro 60 ). For a typical sweep frequency of 400kHz, a data acquisition rate of 50 MHz was used. The 256kB available memory therefore allowed the analysis of 5ms of data consisting of 4000 individual current-voltage characteristics. A non-linear least squares fit of the current-voltage characteristic
yields the ion saturation current, floating potential and temperature. Only those points at a voltage less than kT_e above the floating potential were fitted. The mean fitting error (with T_e/T_e = 0.03 in this experiment) is required to be as small as possible as this is one of the quantities which determines the lowest level of normalised temperature fluctuations that may be analysed. The testing of recently available ferrite cores, which have low losses at a frequency of 1MHz and so can cope with the typical 100W power being transformed, is being carried out and this will allow the sweep frequency of the voltage used to obtain the current-voltage characteristic to be increased in the near future. Offset biasing of the swept voltage down to -200V avoided perturbation of the plasma encountered when sweeping the probe into the electron saturation regime (5).

3. RADIAL PROFILES

The radial profiles of the mean values and normalised fluctuation magnitudes of density, floating potential and temperature (\(\bar{n}_e/n_e, \bar{\phi}_{fl}/T_e\) and \(\bar{T}_e/T_e\) respectively) are shown in Fig. 1. The fluctuation magnitudes shown are the RMS values in the frequency range 30kHz to 350kHz. Existing measurements show a wide discrepancy with \(T_e/T_e\) greater than, comparable to and less than \(\bar{n}_e/n_e\) being reported (2,4,6). An apparent temperature fluctuation with \(T_e/T_e = 0.05\) due to \(\bar{n}_e\) and \(\bar{\phi}_{fl}\) during a sweep is estimated (3). In contrast to recently reported measurements, it was found that \(\bar{T}_e\) and \(\bar{n}_e\) were not out of phase as the shear layer was crossed (2,4). As shown in Fig. 2, both the coherency of \(\bar{T}_e\) and \(\bar{n}_e\) up to a frequency of 100kHz and the coherency of \(\bar{T}_e\) and \(\bar{\phi}_{fl}\) increased as the shear layer was crossed. The level above which the coherency is statistically significant is indicated.

4. SPATIAL CORRELATION AND TRANSPORT

Three poloidally separated tips were swept simultaneously. The spatial coherency of density, temperature and floating potential fluctuations measured at radial positions inside and outside the shear layer are shown in Fig. 3. In each case, this coherency decreases as the poloidal separation of the probe tips is increased from 2mm to 4mm and increases as the shear layer is crossed. An improvement in the spatial coherency of \(\bar{T}_e\) is expected for measurements at higher sweep frequencies. The plasma potential, \(V_p\), is given by:

\[
V_p = \phi_{fl} + \alpha T_e
\]

where \(\alpha = 2\) was used. Fluctuations of plasma potential generate poloidal electric field and radial ExB velocity fluctuations. Coherent velocity and density fluctuations induce convective heat transport, while coherent velocity and temperature fluctuations induce conductive heat transport. The ratio of heat transport by conduction and convection is given by (3):

\[
\frac{\langle \bar{T}_e \phi_{th} \rangle}{\langle \bar{n}_e \phi_{th} \rangle}
\]

Ultimately, experiments with the swept Langmuir probe are aimed to yield values of this ratio. The observed systematic increase in the power weighted coherence between \(\bar{T}_e\) and \(V_p\) as the shear layer was crossed, is based on the increasing coherence of \(\phi_{fl}\) and \(\bar{T}_e\) as the shear layer was crossed. Also a systematic decrease in the power weighted coherence between \(\bar{n}_e\) and \(V_p\) was observed as the shear
layer was crossed. The implications for the ratio of conductive and convective transport is presently under analysis.

The experiments investigating the spatial correlation of temperature fluctuations were motivated by the static Langmuir probe measurements of spatially correlated $\tilde{n}_e$ and $\tilde{\phi}_f$ and the turbulent eddy model (5,7). The extension of these results to include the cross correlation of $T_e$ with $\tilde{n}_e$ and $V_p$ will be presented.

5. H-MODE

The time development of ELM's in the recently discovered H-mode on W7-AS is another topic able to be studied by the swept Langmuir probe technique. With measurements on the microsecond time scale, the electron temperature and density pulse associated with an ELM can be resolved, enabling estimates of particle transport by coherent density and plasma potential fluctuations during an ELM to be carried out. At a position approximately 5cm away from the separatrix, an ELM is observed typically as a jump in temperature from 30eV to 50eV with a relaxation to the original value on a time scale of 1ms.

6. CONCLUSION

The spatial coherency of $\tilde{n}_e$ and $\tilde{\phi}_f$ found in static probe measurements have been reproduced with the swept probe technique. The smaller poloidal correlation length of $\tilde{T}_e$ inferred from its smaller spatial coherency could be partly explained by the contributions of an apparent $\tilde{T}_e$ induced by $\tilde{n}_e$ and $\tilde{\phi}_f$ during a sweep and the gaussian noise introduced by fitting. Combined this amounts to $\tilde{T}_e/T_e = 0.06$ which is half the measured fluctuation level at the innermost radius. An increase in the spatial correlation of $\tilde{T}_e$ and in the coherence of $\tilde{T}_e$ and $\tilde{n}_e$ is needed for consistency with the turbulent eddy model (7). Experiments at higher sweep frequencies will decide whether or not the present results remain in contradiction with this model.

(7) M.Endler et. al., 19th EPS Conference on Controlled Fusion and Plasma Physics (Innsbruck), II 787, 1992.
Fig. 2 Coherency for a single tip at a radius inside and outside the shear layer between:
(a) temperature and floating potential fluctuations
(b) temperature and density fluctuations

Fig. 3 Coherency of density, temperature and floating potential fluctuations between two tips separated poloidally by 2mm and 4mm at a radius inside and outside the shear layer.
Edge plasma profile and particle transport study on the
WENDELSTEIN 7-AS stellarator

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1. Introduction. W7-AS has toroidally varying plasma cross sections (five field periods) and a complex magnetic surface topology at the edge which is more or less governed by "natural" islands (at rationals $5/m$ of the rotational transform $\tau = n/m$) of up to several centimeters radial and poloidal extension [1]. The complexity is further increased by two movable, asymmetric local limiters introducing inhomogeneous connection lengths. Except a narrow $\tau$-range close to $1/3$, where magnetic surfaces extend to within the limiter shadow, a sophisticated plasma edge transport analysis requires three-dimensional treatment. In the case with $\tau$ just above $1/3$ (closest to tokamak limiter configurations) reasonable results were obtained by a 1D (radial) approach basing on singular flux bundles [2, 3]. What we did now is, to explore the ground by Langmuir probes, in what manner the edge plasma parameter profiles are influenced by the resonant boundary structures in the higher-\textit{\tau} case (0.44 - 0.58, including the resonances at 5/11, 5/10 and 5/9). Further we searched for some ordering parameters allowing a comparison of estimated transport coefficients in the SOL at higher $\tau$ (at least in ranges free of stronger resonant perturbations) with results found at low $\tau$. We focus on particle transport in this first attempt.

2. Boundary magnetic islands and edge plasma profiles. The chosen discharge parameters guarantee a magnetic field configuration close to the vacuum field (flat top ECRH discharges with $P_{heat} = 160$ kW and constant line density, $\langle n_e \rangle = 1x10^{19}$m$^{-3}$, net plasma current compensated to zero). The two main limiters placed at the top and bottom of an elliptical cross section (shifted by one field period) were kept fixed at the outermost position. Two fast reciprocating Langmuir probes (CFC tips, triple mode, temperatures derived by averaging after exchanging positive and floating tip) were placed at the outer tip and more towards the bottom of a triangular plasma cross section (fig. 2a). A condensed overview on the results is given in fig. 1, showing isolines of the ion saturation current density $j_\parallel$ and $j_\parallel T_e$ measured at the triangular tip (probe 1) as functions of the edge rotational transform $\tau(a)$. Multiplied by an adequate energy transmission coefficient ($=10$) the latter represents the parallel energy flow to a sink. The outstanding feature is a strong radial modulation of the profiles, with flat regions at the resonances 5/11, 5/10 and 5/9 (a shift with respect to $\tau(a)$ is explained by the fact that $\tau(a)$ refers to a fixed surface at the edge). Concerning the flat regions at the resonances, a detailed connection length analysis shows that the limiters become efficient only outside the islands in these cases. The expansion of the plasma cross section by the closed islands (relative to neighbouring $\tau$ ranges where the islands are intersected by the limiters) and long connection lengths in the vicinity of the X-point broaden the profiles along the line-of-sight of the probe. Transport simulations by field line tracing with crossfield diffusion [4] comparing this case (5/9) with a neighbouring-\textit{\tau} case, show good
qualitative agreement (fig. 2c, d). It should be remarked that, though the configurational cross section is increased at the resonances, temperature profiles in the plasma core measured by Thomson scattering generally contract in these ranges indicating short-circuit transport by the islands. Starting from the 5/9 resonance towards smaller \( t \), the 5/9 islands are shifted outward (positive shear) and become increasingly intersected by the limiters (separatrix-dominated configuration). Inside a small \( t \) range between 0.52 and 0.53, connection lengths become short (< 10 toroidal revolutions, = 120 m) close to the X-point and inside the 5/9 islands. This holds for the whole boundary because in the 5/9 case there is only one island which closes after 9 toroidal and 5 poloidal revolutions. That means in this case we have a separatrix-dominated regime with optimum plasma cross section. This is exactly the \( t \) range where we (up to now exclusively) found the H- mode. In contrast to this region the range governed by the 5/10 resonance \( t(a) = 0.47 - 0.5 \) is limiter-dominated within discrete poloidal sectors (5 separated islands), and in between the field lines end at other installations or the wall. Towards smaller \( t \) values the configuration is generally limiter-dominated.

3. Particle transport. For comparison of the particle transport at high \( t \) with the limiter-dominated "standard case" at \( t(a) \) close to 1/3 (0.345), the "optimum" separatrix-dominated configuration at \( t(a) = 0.525 \) was selected. Discharge conditions were 350 kW ECRH, maximum limiter aperture and varying densities. Radial density profiles outside the LCMS as obtained by the probes described above were exponentially fitted. Averaged particle diffusion coefficients were estimated then from the e-folding lengths by the zero-dimensional approach \( \lambda_n \beta = 2DLC/c_{S1} \) with \( D \) the diffusion coefficient, \( 2L_C \) the connection length and \( c_{S1} \) the ion sound speed. Parallel connection lengths \( L_C \) were calculated from the vacuum magnetic field and averaged along the line-of-sight of the probes and over a poloidal distance of ±2 cm from the probes. Finite \( \beta \) effects should be negligible with this respect. \( \lambda_n \) is referred to distances from the LCMS which were averaged along the parallel direction of the flux tubes intersected by the probes. The result is shown in fig. 3, where averaged particle diffusion coefficients obtained by the two probes for the two configurations are plotted versus the electron density at the LCMS. \( \lambda_n \)-values referred to the local distance from the LCMS were found to differ by up to a factor of 3 for the two probes, reflecting the strongly varying local distances of the flux tubes from the LCMS [5]. Nevertheless, after applying the averaging procedures described above, which account for the topological characteristics, a uniform scaling of \( D \) with the edge density was obtained. For the low-\( t \) range, this behaviour was found already earlier [1, 3]. We conclude that, within this approach (perhaps a factor 2 -3 of \( D \)) and referred to the edge density as ordering parameter, the particle transport does not significantly differ for the regimes and poloidal ranges investigated.

References

[5] F. Sardei et al., this conference
Fig. 1: Isolines of (a) the ion saturation current density $j_{S}^{+}$ and (b) $j_{S}^{+}T_{e}$ measured by probe 1 (see fig. 2a) versus the edge rotational trans-form $i(a)$. The dotted line indicates the LCMS defined by the limiters or the separatrix. Symbols L (fig. b) - limiter-dominated, S - separatrix-dominated configuration ("island-divertor"), IS - limiter surface outside closed islands. Rationals at the bottom indicate the type of resonances at the boundary.
Fig. 2: Cross sections of the vacuum magnetic field (a, b) at $\phi = 72^\circ$ (probe plane) showing the 5/9 boundary islands for two values of the boundary rotational transform $\imath(a)$, and respective Poincare plots taking into account diffusion (c, d). The diffusion is modelled by perpendicular displacements of the particle motion along the magnetic field after each integration step length ($D_{\text{eff}} = 1 \text{m}^2\text{s}^{-1}$, for details see ref. [4]). The plots illustrate typical constellations discussed in the text: closed islands (profile broadening) and islands intersected by the limiters (separatrix-regime, "island-divertor").

Fig. 3: Averaged edge particle diffusion coefficients $D$ (see text) versus the edge density for two different configurations (limiter-dominated at $\imath(a) = 0.345$, separatrix-dominated at $\imath(a) = 0.525$) and two different poloidal probe positions (probe 1: close to an X-point, probe 2: neighboured to an O-point in the separatrix case).
EXPERIMENTAL AND THEORETICAL STUDY ON PLASMA HEAT FLOW TO PLASMA FACING MATERIALS
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1. INTRODUCTION: The heat load to divertor plate is crucial to estimate its erosion due to radiation enhanced sublimation, chemical sputtering, evaporation, ionic impact sputtering, etc. Divertor biasing has been found to be effective for edge plasma control leading to an enhanced core confinement as well as an increased particle exhaust. However, the plasma heat flow is generally estimated by a simple sheath theory, and few systematic experimental and theoretical results on it have been obtained. Plasma heat flow is determined by not only energy transmission of sheath and presheath, but also other factors. Particularly in a cold and dense divertor plasma, ion energy reflection and surface recombination have to be brought into consideration. The former reduces divertor heat load, and the latter has the opposite role. In this contribution, the experimental results on plasma heat flow to an electrically floated as well as biased surface and the comparison of those data with the theory including ion energy reflection and surface recombination are presented. Effects of these factors are made clear.

2. THEORY: Plasma heat flow \( q(\phi) \) to plasma facing material is estimated as a function of target potential \( \phi \) with respect to the plasma potential by the following expression:

\[
q(\phi) = n_{se} C_s \left[ 2kT_i - e\phi + 0.5kT_e (C_e/C_s) \exp(e\phi/kT_e) \right],
\]

where \( n_{se} \) is the electron density at sheath edge, \( C_s \) is the ion sound velocity and \( C_e = (8kT_e/\pi m_e)^{1/2} \) is the electron thermal velocity, \( T_e \) and \( T_i \) are the electron temperature and the ion temperature. So called "energy transmission factor" \( \delta(\phi) \) is defined by the expression

\[
\delta(\phi) = q(\phi)/n_{se} C_s kT_e .
\]

In the case of electrically floated surface in a hydrogen plasma, the value of \( \delta \) is found to be 7 to 8. On the other hand, the expression of plasma heat flow including effects of ion energy reflection and surface recombination is as follows:

\[
q(\phi) = n_{se} C_s \left[ (2kT_i - e\phi)(1-R_{ie}) + 0.5kT_e (C_e/C_s) \exp(e\phi/kT_e) + e\phi_r \right],
\]

where \( R_{ie} \) is the ion energy reflection coefficient and \( e\phi_r \) is the ion recombination energy. \( R_{ie} \) depends on incident ion energy, ionic species and target material. Then, we obtain the revised energy transmission factor written by

\[
\delta(\phi) = \frac{(-e\phi/kT_e + 2T_i/T_e)(1-R_{ie}) + 2[(1+T_i/T_e)(2\pi m_e m_i)]^{1/2} \exp(e\phi/kT_e) + e\phi_r/kT_e .}{2[(1+T_i/T_e)(2\pi m_e m_i)]^{1/2} \exp(e\phi/kT_e) + e\phi_r/kT_e}.
\]
3. EXPERIMENT: The experiment has been carried out on a linear plasma device, NAGDIS-I[21,31]. The length and the inner diameter of the cylindrical vacuum vessel are 2.5m and 0.18m, respectively. The steady state plasma in NAGDIS-I is produced by PIG discharge with working gases of hydrogen, helium and argon. The present experiment was done using helium plasma. The experimental setup is shown schematically in Fig.1. Plasma heat flow was measured by the temperature increase of a tungsten target (20x20x2mm) located at the center of plasma column and set normally to the magnetic field. A calibration of this method had been done by using an electron beam. An infrared thermometer was used to measure the target temperature. The target potential relative to ground potential is variable by a bias power supply from -160V to -30V. The plasma parameters on a cross section are measured every 2mm by the fast scanning Langmuir probe with cylindrical probe tip set about 0.2m distant from the tungsten target. Figure 2 shows a typical profiles of NAGDIS-I plasma parameters. The plasma diameter was found to be about 10cm. In this experiment, electron density and temperature of center of plasma column were varied (0.7-1.2)x10^18 m^-3 and 6-10eV, respectively. The plasma parameters are assumed not to be changed largely in the area between the position of the scanning Langmuir probe and the target. These parameters were substituted into above equations to derive experimental and theoretical values of energy transmission factor δ.

4. RESULT AND DISCUSSION: Theoretical values of energy transmission factor δ was obtained based on some assumptions; (1) the probe data show the values at the so called "stagnation point" so that the value of φ in above equations is the difference between the target potential and the plasma space potential, (2) a potential
drop of presheath is $0.5kT_e$, so that the electron density at the sheath edge is $0.61n_e$, where the value at the probe location should be employed for $n_e$. (3) the electron reflection on the surface and secondary electrons were disregarded because $T_e$ is not enough high for such effects in NAGDIS-I plasma, (4) $T_i=0.2T_e$. In this experiment, $\phi$ was varied from -8V to -130V, so reduced energy $\varepsilon$ varied from $3.8x10^{-4}$ to $6.4x10^{-3}$. Ion energy reflection coefficient $R_{ie}$ was assumed to be 0.7 over the present ionic energy range. Surface recombination energy $\phi_r$ is 24.6eV for helium. Figure 3 shows the comparison between the energy transmission factor experimentally obtained and the one theoretically calculated. The effect of ion energy reflection is clearly shown at the region of $|e\phi/kT_e|>7$. Theoretical curves from eq.(4) show that energy transmission factor is reduced by the ion energy reflection indicated by the hatched part. Experimental values agree fairly well with the revised theoretical values. Theoretical curves shows that energy transmission factor comes to be minimum at near the floating potential, $e\phi/kT_e=-3.5$. However, the minimum is found to be at $e\phi/kT_e\approx -6$ in the experiment. In the region of $|e\phi/kT_e|<6$, the experimental value $\delta$ takes considerably larger value than the value given by the theory. Two reasons are considered to explain this discrepancy. One-dimensional computer simulation using P.I.C.(Particle In Cell) code suggested that it is due to a deviation of electron velocity distribution function. When the target is biased shallower, that is, the target potential approaches to the plasma potential, low energy electrons come to be able to reach the target rather than reflected. The tail of velocity distribution of electrons running away from the target is largely depopulated by the absorption of incoming electrons to the target. So at sheath edge, $T_e$ determined by dispersion of velocity distribution are estimated smaller than true ones. This phenomenon was observed by

![Fig.3: Energy transmission factor from experiment and theory. Lines are theoretical curves. The bold one is from eq.(2), and the others are from eq.(4). Circles are experimental data. $[n_e=(0.7-1.2)x10^{18}m^{-3}, T_e=6-10eV, R_{ie}$ is assumed to be...](image)
Fig. 4(a): Experimental set-up for measurement electron energy distribution function using an insulated twin probe.

Fig. 4(b): Derivatives of probe currents with respect to the probe voltages, that is considered to be proportional to electron energy distribution function. Helium plasma, $n_e = (3.1-3.4) \times 10^{16}$ m$^{-3}$, $T_{e,I} = 6.2$ eV, $T_{e,II} = 3.4$ eV, target potential $= -40.2$ V relative to ground. Probe was 0.2 m distant from biased target.

probe measurement as shown in Fig.4, where a lack of high energy electrons is observed at the probe facing the target. Another reason is the existence of intrinsic non thermal energetic electrons.

5. CONCLUSION: Energy transmission factors were measured in NAGDIS-I helium plasma and were compared with theoretical values including effects of ion energy reflection and recombination energy on the target surface. In the case of deeper biased target than the floating potential, both values agree fairly well in which the role of energy reflection was found to be important. In shallower biasing, experimental data greatly exceed theoretical values. The important reason is the deviation of electron velocity distribution function from the Maxwellian due to the lack of high energy electrons absorbed at the target as well as the presence of intrinsic non thermal energetic electrons.

REFERENCES
INTERACTION OF A HIGH POWER HOT PLASMA STREAM WITH SOLID MATERIALS

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Abstract
The GOL-3 facility was used for exploratory plasma stream target experiments under conditions rather typical for the thermal quench phase of ITER tokamak plasma disruptions. The plasma stream at the target consists of a relativistic electron beam (REB) for plasma heating and a hot plasma stream with the two components bulk plasma electrons with $T_e$ around 1 keV and suprathermal plasma electrons with characteristic energy of 10 - 20 keV. The specific energy of the hot plasma stream reaching the target can be up to 3 MJ/m² the power density up to 50 MW/cm² and the pulse duration around 6 µs.

The plasma stream target experiments allowed to study the properties of the target plasma formed from vaporized target materials in front of the target and to determine the target material erosion. In less than 1 µs after the onset of the plasma stream a cloud of evaporated material is formed. The cloud expands along magnetic force lines with velocities around 10⁶ cm/s. Line radiation is observed from the target plasma corona, continuum radiation from the bulk of the cloud. The temperature in the plasma corona is about 1 eV. The black body temperature of the bulk cloud is around 0.5 eV. The erosion for graphite increases sharply upon reaching a threshold value of 1 MJ/m².

Introduction
During the thermal quench phase of a disruption event energy densities of up to 12 MJ/m² are estimated to be dumped to the ITER divertor plates within a time interval of 0.1 to 3 ms (D.E. Post 1990). Such high heat loads result in instantaneous evaporation of the divertor plate material (Th. Klippel 1989). The amount of material evaporated per disruption could limit the lifetime of the divertor plates and thus could become a key problem in tokamak fusion technology (Kuroda 1990).

In existing tokamaks ITER typical heat loads are not achievable. Therefore the divertor material behaviour has to be studied in facilities which allow a simulation of the ITER disruption conditions. The GOL-3 facility (A.V. Arzannikov 1988) by some of it's important parameters is suitable to simulate these conditions and thus was used for exploratory hot plasma stream target experiments. Various substances (graphite, polypropylene, copper) were exposed to a plasma stream impinging perpendicularly and under different angles onto the target.

Facility and diagnostics
GOL-3 is a mirror trap of length of 7 m, with a longitudinal magnetic field of up to 6 T in the trap and up to 12 T in the mirror end plugs (see Fig. 1). A relativistic electron beam (REB) with beam energy content up to 90 kJ, beam current up to 50 kA, pulse duration 5 µs and particle energy around 0.8 MeV is generated in a quasiplanar vacuum diode and injected through one of the end mirrors into a plasma of density of $10^{15}$ cm⁻³. After the end of the beam injection the plasma cools down due to classical electron heat conductivity along magnetic field lines.

The GOL-3 facility is well equipped with diagnostics for the measurement of the operation parameters of the electron beam, of the initial plasma and of the heating
Fig. 1  GOL-3 with facility diagnostics

process of the plasma. In particular there are two systems of Thomson scattering of ruby laser light, there are diagnostics for x-ray, VUV and optical radiation, optical interferometers, bolometers, diamagnetic loops and energy analyzers (Arzannikov et al. 1990).

A special test chamber (see Fig. 2) was used for the experiments. Targets can be installed at different positions in the chamber. At the chamber specially designed diagnostics were used for studying the behaviour of the target and the target plasma. For analysis of the target plasma an optical interferometer with wavelength of 630 nm, equipment for optical spectroscopy of target plasma radiation including spectral analysis with high time and space resolution, a 9 channel VUV pinhole chamber and 4 VUV detectors were used. For analysis of the bulk target an optical and an electron microscope, a microbalance with accuracy of 0.1 mg and a profile machine for measurement of surface profiles with an accuracy of up to 1 µm are available.

Results
The plasma stream at the target consists of the REB for plasma heating in the solenoid and the hot plasma stream. Characteristics of both are listed in table 1 together with the ITER data for comparison. The mean free path of the REB electrons in graphite is a few mm. The hot plasma stream flowing out of the device consists of two components, the bulk plasma electrons and the suprathermal plasma electrons. The temperature of the bulk plasma ions is an order of magnitude smaller, their contribution to the energy density thus is negligible. As a result the specific energy of the hot plasma stream at the target can achieve 3 MJ/m² depending on the operating conditions. The power density of the hot plasma stream is up to 50 MW/cm².
Bulk targets from fine grain graphite, and foil targets from aluminum, titanium, copper and polypropylene were used. In this paper only results for graphite are given. About 1 μs after the start of the plasma stream a target plasma cloud of substantial ionization is established near the target. The diagnostic measurements have shown that the outer part of the cloud facing the hot plasma stream is a magnetized plasma corona which propagates along the magnetic force lines with velocities of up to 2 × 10^6 cm/s. The expansion velocity was determined by different methods: by time- and space dependent measurements of the maximum Hα luminescence and of the VUV brightness and by interferometry measurements.

Time and space dependent electron density distributions in the target plasma corona were determined from analysis of measured Stark broadened Hα lines. A typical example is shown in Fig. 3. The plasma temperature is around 1 eV. Observed was line radiation from the target substance (CI, CII, CIII) and from hydrogen.

A bulk plasma of higher density follows the corona. The level of ionization is small, the gas is not confined within the magnetic field but moves along and across the magnetic force lines with the same expansion velocity of about 10^6 cm/s. The radiation from the bulk plasma is mainly continuum and close to black body radiation with a temperature of 0.3 - 0.5 eV.

First erosion experiments were performed with bulk fine grain high density graphite targets. First of all it was established that irradiation with the REB beam alone causes no measurable erosion. Fig. 4 shows the results for two targets irradiated by the REB beam only and by the REB beam and the hot plasma stream. In case of REB alone no erosion was detected and no surface damage was observed. In the latter case the hot plasma stream was partially shielded with a thin graphite screen in front of the target. Under the influence of the hot plasma stream target erosion occurs. The erosion depth can achieve several hundred microns. The radial dependence of the erosion reflects the fact that the energy density of the hot plasma stream is radially inhomogeneous.

Results of a series of erosion experiments for energy densities ranging from 0.5 to 3 MJ/m² and various angles of the targets with respect to the magnetic force lines are shown in Fig. 5. There is no difference in erosion behavior for perpendicular incidence and incidence under different angles. Below 1 MJ/m² the erosion is small above 1 MJ/m² explosive erosion processes presently not understood start. Target irradiation by the REB results in volume heating of a material layer of a few mm.
Conclusions
Exploratory plasma stream target experiments have been performed to study the properties of the target plasma formed from vaporized target material in front of the target. Fine grain graphite was exposed to a hot plasma stream of specific power ranging up to 50 MW/cm² and of energy density ranging from 0.5 to 3 MJ/m². Formation and expansion of the target plasma was observed. Important parameters of the plasma were determined by using special diagnostics.

The target plasma consists of corona and bulk plasma. The corona emits line radiation, has a temperature around 1 eV, is magnetized and expands in case of perpendicular stream incidence along the magnetic force lines with velocities around $2 \times 10^6$ cm/s, typical plasma density is up to $10^{17}$ cm⁻³. The bulk plasma has higher density, emits continuum radiation and the black body temperature is around 0.5 eV. Due to its incomplete ionization the cloud expands along and across magnetic field lines with about the same velocity of around $10^6$ cm/s.

First erosion experiments yielded surprisingly high erosion values if the energy density of the hot plasma stream at the target is above 1 MJ/m². The erosion rates observed are an order of magnitude larger than the hitherto obtained rather large erosion values using laser beam facilities (van der Laan 1992) where no vapor shield effect occurs (B. Goel 1992).

Table 1  Energy and power density values

<table>
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<th>REB</th>
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<th>suprathermal electrons</th>
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<td>magnetic field (T)</td>
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<tr>
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<td>$\leq 1$</td>
</tr>
<tr>
<td>power density (MW/cm²)</td>
<td>$\leq 12$</td>
<td>600</td>
<td>$\leq 20$</td>
<td>$\leq 20$</td>
</tr>
</tbody>
</table>

References
1. D.E. Post (ed) ITER Physics, ITER Documentation Series No. 21, IAEA, Vienna (1990)
7. B. Goel et al., 17th SOFT, Rome (1992)
Energetic Electron Transport in Cold Plasma and Gas Target Divertors

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Introduction - Gas target or radiative divertors are under consideration for next generation tokamaks like ITER and TPX. A significant reduction of the parallel heat flux in a gas target has been demonstrated previously in PISCES [1] and elsewhere [2]. A reduction of the peak heat flux and the divertor electron temperature by gas puffing has been demonstrated in DIII-D [3]. Energetic electrons can carry a substantial part of the parallel electron heat flux in a divertor plasma. We have experimentally investigated a regime where up to 90% of the total electron heat flux ($Q_{eS} \leq 5 \text{MW/m}^2$) is carried by energetic electrons ($n_{eh}/n_{ec} \leq 0.1$, $U_{eh} \leq 160 \text{eV}$). The experiments were carried out in a hydrogen plasma in the PISCES-A linear plasma device at densities $n \leq 10^{19} \text{m}^{-3}$. The heat flux to the (simulated) divertor target is measured with a biased calorimeter, allowing us to separately determine the target heat fluxes due to ions as well as cold and hot electrons. It is demonstrated that the energetic electron heat flux can be attenuated by nearly two orders of magnitude with active gas injection through the target. Neutral densities $n_{H_2} \geq 10^{21} \text{m}^{-3}$ are required. Results from a 1.5/2-D fluid model based on coupled plasma/neutral equations are presented.

Experiment - A schematic of the PISCES-A linear plasma device is shown in Fig.1[4]. The plasma (source density $5 \times 10^{17} \text{m}^{-3} \leq n_e \leq 1 \times 10^{19} \text{m}^{-3}$ and bulk electron temperatures $kT_e \leq 20 \text{eV}$) is generated in a reflex arc configuration, using a hot cathode made of LaB$_6$. The plasma diameter is 0.05 m. The ion temperature ($kT_i \leq 2 \text{eV}$) is found to be close to the Franck-Condon energy. The axial magnetic field $B_z$ is generated by two sets of coils. In the plasma source region, $B_z \leq 0.1 \text{T}$, while downstream in the analysis region, $B_z = 0.19 \text{T}$. The primary hydrogen gas feed is located in the source region. The neutral hydrogen pressure in the plasma source is kept below 1 mtorr. A circular tube (length $L = 0.9 \text{m}$, diameter $2a = 0.045 \text{m}$) is located at a distance of 0.55 m downstream from the source, simulating a closed divertor channel.

Fig.1: PISCES experimental set-up for gas target experiments and diagnostics
The simulated divertor target is located at the end of the tube \((z=0)\). A second hydrogen gas feed is located close to the center of the target plate. The "neutralizer" tube is made of (insulating) anodized aluminum and is electrically floating. The neutral pressure is monitored by baratron gauges located at various axial positions. The plasma density and the electron temperature are measured by an axially moveable, water-cooled, plane Langmuir probe.

### 1 1/2-D Fluid Model

1-D fluid equations [5] are solved along the direction of the magnetic field \(B_z\) for ions, cold electrons, and atomic and molecular hydrogen neutrals. The diffusion approximation is used for the two neutral species. Some perpendicular effects are included (cross-field plasma diffusion, wall recombination, and recycling from the target and side walls). The particle balance equation can be written

\[
\frac{\partial n}{\partial t} + \frac{\partial}{\partial z} (n v) = S_n, \tag{1}
\]

where \(S_n = S_{\text{ion}} - n D_z J_0^2 a^2\) is the effective particle source \((S_{\text{ion}}\) includes plasma production by cold and hot electrons due to ionization of \(H\) and dissociation of \(H_2\)). We take the perpendicular diffusion coefficient \(D_z\) to be the sum of the Bohm diffusion coefficient and the classical diffusion coefficient due to ion-neutral collisions \(D_{\text{cl}} = kT_e (v_{\text{cx}}+v_{\text{in}}+v_{\text{ion}}) / (m_i \omega_{ci}^2)\), times some numerical factor \(\alpha\). Here, \(\omega_{ci}\) is the ion cyclotron frequency, and \(v_{\text{cx}}, v_{\text{in}}\) and \(v_{\text{ion}}\) are the momentum loss rates due to charge exchange, elastic ion-neutral, and ionizing collisions. The plasma flow velocity \(v\) is determined by the ion momentum balance:

\[
m_i \frac{\partial}{\partial t} (n v) + \frac{\partial}{\partial z} (m_i n v^2 + n kT_i + n e_{\text{cc}} kT_e + n e_{\text{eh}} U_{\text{eh}}) = -m_i n v (v_{\text{cx}}+v_{\text{in}}) - m_i n v D_z \frac{J_0^2}{a^2} \tag{2}
\]

The cold electron temperature profile is determined by the electron heat flux equation:

\[
\frac{3}{2} \frac{\partial}{\partial t} n e_{\text{cc}} kT_e + \frac{\partial}{\partial z} \left(\frac{5}{2} v n e_{\text{cc}} kT_e + q_{\|}\right) = v \frac{\partial}{\partial z} (n e_{\text{cc}} kT_e) - \frac{3}{2} S_{\text{ne}} E_{\text{ion}} - 3 \frac{n_{\text{ee}} m_i}{v_{\text{ee}} m_i} (kT_e kT_i) - Q_{\text{en}} \tag{3}
\]

Here, the terms on the left-hand side describe the convective energy flow, the first term on the right-hand side describes electron heat conduction with a harmonic flux limiter \(q_{\|} = q_s q_s / (q_f - q_s)\). The flux limit is taken to be \(q_f = 0.5 (2/m_i \pi)^{1/2} n e_{\text{cc}} kT_e^{3/2}\) and the Spitzer heat flux is \(q_s = -v_{\text{ee}} \frac{\partial kT_e}{\partial z} = -3.2 n e_{\text{cc}} kT_e / (m_e v_{\text{ee}}) \frac{\partial kT_e}{\partial z}\). The second term describes the energy losses by radiation and ionization. \(E_{\text{ion}}(n, kT_e)\) is the electron energy loss per ionization event (including excitation and dissociation) and varies from 20-150 eV depending on the local electron temperature and density [5]. At high neutral density, this term is the dominant energy loss term. The third and forth terms describes the energy transfer from electrons to ions and neutrals due to elastic collisions.

### Results and Discussion

Fig.2 shows the cold (bulk) electron density and temperature as a function of \(n_{\text{neq}}\), the neutral density at the target. Gas is admitted through the target at a rate \(\leq 2\) torr-l/s. The upstream neutral pressure is \(2 \times 10^{-4}\) torr without target gas injection and increases to \(5 \times 10^{-4}\) torr at the highest injection rate. The cold electron temperature at the source decreases somewhat with increasing \(n_{\text{neq}}\), while the upstream plasma density is not changing significantly. The target electron density first increases with the neutral density due to the increased ionization rate inside the neutralizer tube. The bulk electrons are cooled by ionization and radiation losses in the vicinity of the target. The axial plasma density profile eventually develops a maximum \((n_{\text{max}} \leq 8\) \(n_0)\) at \(0.2\) m \(< z < 0.5\) m, as reported previously [1,7]. Downstream of the density maximum, the cold electron temperature decreases to values of 2-3 eV. The ionization rate is very small at these temperatures, and particle losses dominate the particle source. Radial and axial losses are comparable in magnitude at high neutral density and the plasma density is observed to decrease towards the target. The plasma "detaches" from the target for \(n_{\text{neq}} > 8 \times 10^{20}\).
The upstream boundary conditions for the modeling calculations are given by the measured source plasma parameters; sheath boundary conditions are used at the target. Good agreement with the experimental data is found.

Fig. 3 shows the measured average energy and density of energetic electrons upstream (z=1m), and at the target (z=0), plotted versus the neutral hydrogen density at the target. The data is obtained by differentiating the characteristics of a plane Langmuir probe. The hot electron energy in the source plasma is $U_{eh} = 100$ eV and is close to the cathode sheath potential drop. $\lambda_{ee, \lambda_{ei}} \gg L_n$ and, at low neutral density, $\lambda_{en, \lambda_{ion}} \gg L_n$, where $\lambda_{ee, \lambda_{en}, \lambda_{ei, \lambda_{ion}}}$ are the hot electron collisional mean free paths for electron-electron electron-ion, elastic electron-neutral, and ionizing collisions, $L$ is the length of the plasma, and $L_n$ is the axial extent of the neutral layer at the target. The hot electrons are therefore decoupled from the cold electrons and the ions. For target neutral densities $n_0 \geq 6 \times 10^{20}$ m$^{-3}$, $\lambda_{en, \lambda_{ion}} = L_n$ so that hot electrons are thermalized and cooled by inelastic collisions in the gas layer close to the target. The hot electron density is found to decrease monotonically towards the target at high neutral densities. The hot electron heat flux is sheath limited (the hot electron flux through the sheath is given by $F_{eh} = 0.5neheh = 0.5csn$). The values of $n_{eh}$ calculated from this expression are shown in Fig. 3 and agree reasonably well with the experimental data.

Fig. 4 shows the heat fluxes to the target as a function of the neutral density at the target. The total heat flux is measured calorimetrically and normalized to $Q_0 = 2.5$ MW/m$^2$ (the value measured without target gas injection). Shown are also the components (normalized to $Q_0$) of the target heat flux due to hot electrons $Q_{eh}$ and the total plasma heat flux $Q_{pl}$ (the sum of the hot electron heat flux and the ion heat flux including the sheath potential acceleration), as evaluated from biased calorimeter current/voltage characteristics [8]. The heat flux due to cold electrons is negligible, since cold electrons are reflected by the sheath potential. $Q_{cal}/Q_0$ is always larger than $Q_{pl}/Q_0$ since energetic neutrals and radiation in the vicinity of the target contribute to the total heat flux. It is obvious that, at low target neutral densities, the hot electron heat flux is the dominant contribution, while at high neutral density, all components are attenuated by nearly two orders of magnitude.

**Conclusions** - In this paper we have demonstrated that plasma "detachment" in a gas target can be achieved in the presence of a significant fraction of energetic electrons. We have also shown that the density of energetic electrons at the target can be greatly reduced under these conditions. This reduction is probably due to the combined effects of a sheath-induced flux limit, inelastic collisions, and strong, anomalous radial transport. For the parameters of the PISCES experiment, the target heat flux is reduced by almost two orders of magnitude for neutral densities $n_0 \approx 10^{21}$ m$^{-3}$. While the magnetic field in a tokamak is at least one order of magnitude higher than in PISCES, the connection length along the magnetic field between the x-point and the divertor target is also larger by at least one order of magnitude. Therefore, the ratio of the radial diffusive speed and the poloidal flow speed of the plasma should be similar (if we assume $D_1 \sim 1/B$). Plasma detachment might be achievable in a tokamak divertor if neutral densities on the order of $10^{21}$ m$^{-3}$ can be maintained near the divertor target.

**References**

Fig. 2: Plasma density at the source ($n_s$) and at the target ($n_d$); cold electron temperature $kT_e$, and modeling results (target).

Fig. 3: Hot electron density and average energy at the source ($n_{eh s}, U_{eh s}$) and at the target ($n_{eh d}, U_{eh d}$). The target density calculated from the sheath flux limit is also shown (+).

Fig. 4: Measured target heat fluxes, normalized to the target heat flux $Q_o$ without target gas injection. Shown are the hot electron heat flux $Q_{eh}/Q_o$, the plasma heat flux $(Q_{pl}/Q_o)$ and the total heat flux $Q_{cal}/Q_o$. 
SURFACE DEFORMATION EFFECTS ON STAINLESS STEEL, Ni, Cu AND Mo PRODUCED BY MEDIUM ENERGY He IONS IRRADIATION

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To investigate dose and energy dependence of surface deformation effects (blistering and flaking), different kinds of candidate CTR first wall materials as 12Kh18N10T, W-4541, W-4016 and SS-304 stainless steels, Ni, Cu, Mo were irradiated at room temperature with 3.0, 4.7 and 6.8 MeV He$^+$ ions at IAP Cyclotron. The effects were investigated by means of a TEMSCAN 200 CX electron microscope and two metallographic Orthoplan Pol Leitz and Olympus microscopes. The results are summarized in the Table.

We observed two dose dependent main phenomena: blistering and flaking (craters). So, blisters occurrence on the irradiated surface is almost instantaneous when a critical dose (number of He ions accumulated in the region at the end of alpha particles range) is reached. Examples of blisters are presented in figs. 1, 2, 3. Increasing irradiation dose, we reached flaking stage. So, isolated submicronic fissures along grain boundaries were observed on the blister skin (fig. 4), chronologically followed by large (5–20 μm) deep cracks of hundreds of microns in length, blisters opening and, finally, flaking appearance.

To investigate their inner morphology, the blisters were mechanically opened by a stainless steel pin. We observed a lot of interesting irradiation phenomena both on the bottom and on the inner side (fig. 5) of the reversed skin: sponge- and wave-like structures (figs. 6, 7), microcraters (fig. 6), rupture secondary small blisters (fig. 5, top right), multilayer flaking (figs. 5, 8). To explain these effects, possible reasons are the important He$^+$ energy spread - e.g. for multilayer flaking, local microstructural damages (especially for stainless steel), local high temperatures values on the skin and, also, possible high gas pressure values inside the blisters.
Concerning potential explanation of the obtained results using the stress-induced and gas-driven (bubble coalescence) models, it seems that our average blister diameter data (see Table) apparently support the stress-induced model for medium energy He ions (d ∼ t^{3/2}, where t is the blister skin thickness). However, other observed effects—as microstructure changes—suggest the validity of gas-driven model. For the energy dependence of critical dose for blistering (see Table) we can assume a dependence D ∼ t^{3/2}. Therefore, not only the resistance of the blister skin (which is proportional to t) is important (as gas-driven model predicts), but also other internal resistances (e.g. inter-submicronic bubbles fractural mechanism, which proceeds the gas accumulation at the end of He ions range) play also an important role. So, it is necessary to enlarge the area of samples (new metals and alloys with various metallurgical preparations) and ions (types and energies) to positively conclude between these two models.

<table>
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<th>Material</th>
<th>Energy (MeV)</th>
<th>Dose rate (x10^{13}He/cm^2s)</th>
<th>Critical dose for blistering (x10^{18}He/cm^2s)</th>
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</table>

* single blister
Fig. 1 Blisters on W-4541, 3.0 MeV, 0.5x10^{13} He/cm^2, X 60

Fig. 2 Blisters on Ni, 3.0 MeV, 0.9x10^{13} He/cm^2, X 120

Fig. 3 Blisters on 12KH18N10T, 6.8 MeV, 5.2x10^{13} He/cm^2, X 60

Fig. 4 Submicronic fissures W-4541, 4.7 MeV, X 3000 5.4x10^{13} He/cm^2
Fig. 5 Reversed blister skin, 12KH18N10T, 4.7 MeV, 3.6x10^18 He/cm^2, X 1500

Fig. 6 Wave-like structure and microcraters, 12KH18N10T, 4.7 MeV, 3.6x10^18 He/cm^2, X 5000

Fig. 7 Sponge-like structure, Mo, 3.0 MeV, 1.5x10^18 He/cm^2, X 4000

Fig. 8 Multilayer flaking and sponge-like structure, W-4016, 4.7 MeV, 2.4x10^18 He/cm^2, X 500
DENSITY PROFILE VARIATION IN A HIGH RECYCLING DIVERTOR IN A NEXT STEP DEVICE: COMPARISON OF RESULTS FROM ANALYTIC AND MONTE CARLO NEUTRAL MODELS AND INFLUENCE ON CONVERGENCE

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Introduction

The demonstration of viable regimes for high power operation of next step devices requires 2-D plasma calculations coupled to an appropriate neutral particle treatment. Demonstration of convergence of the coupled code for the relevant parameter range is also required. Previous work [1], carried out with the B2 code [2] coupled to an analytic model, demonstrated exponentially decreasing residuals as well as plasma particle and energy balance to $10^{-6}$ of the particle flux to the plate and the input power, respectively. To permit comparison with previous work, the geometry chosen in the present paper is that of the ITER CDA divertor (22 MA, R=6 m, [3]), double null and having a poloidal X-point to strike-point distance of 1.4 m. All results refer to one outer divertor channel from midplane to divertor plate; the input power $P$ to one outer divertor is 0.4 of the total power to the SOL and 0.05 of the fusion power, and the power per unit area, $f_P$, is given without safety and peaking factors. The mesh has a strongly nonlinear distribution of cells. Results from analytic and Monte Carlo recycling models (see below) are compared. In both cases, only DT ions, atoms and molecules are treated but collision frequencies are corrected for impurities. Radial transport coefficients are uniform in space, with $\chi_e=2 \text{ m}^2/\text{s}$ and $D=\chi_i=\chi_e/3$. The results apply to the high recycling regime of next step devices.

Density profiles with analytic neutral model

Previously [4,5], we had used the B2 code coupled to an analytic neutral particle model [6] which includes atoms and molecules neglecting sideways motion, but adjusting the mean free path along the field such that the correct projection perpendicular to the inclined divertor plate is obtained. Radiation from hydrogenic species was included. While charge exchange processes were included to determine the population of thermal atoms, energy and momentum sources and sinks due to these processes could not be considered. Using this model, we had performed a simultaneous variation of the upstream density (at the intersection of separatrix and midplane) and of the input power, reported in [4]. These scans show a striking variation in the shape of the density profile along the inclined (15° to the magnetic surfaces) outer divertor plate when the input power was varied. An example is shown in Fig. 1. At low power, the density profile is single-humped. As the power increases, the profile becomes double-humped, i.e. peaks develop on the private and SOL sides of the separatrix. As the power is increased further, the density profile is again single-humped; a strong peak survives on the SOL side, displaced from the point of maximum power flow. In this example, and a number of other scans we have examined, the density dip always occurs when the peak electron temperature at the divertor plate is in the range 25 - 50 eV, the temperature range for which the ionization rate is maximum. The density peak on the SOL side coincides with a local $T_e$ of several eV (2-5 eV; charge exchange dominates over ionization).

Initial results with Monte Carlo neutral model

The analytic model has now been replaced by a full Monte Carlo neutral model,
EIRENE, described in [7]. In the initial runs with the coupled code, momentum sinks due to the plasma-neutral interaction were not included. Convergence was found to be more difficult than with the analytic model, in part because of the statistical variation introduced by succeeding Monte Carlo runs. (From one B2 timestep to the next, the sources are rescaled to account approximately for the changes in plasma parameters and a full Monte Carlo call is performed once the variation of the sources exceeds a specified value.) For a power of 47 MW into 1/2 the outer SOL (corresponding to a fusion power of about 1 GW), for an upstream density of $3.10^{19}$ m$^{-3}$ and for a divertor plate perpendicular to the flux surfaces, the peak $T_e$ at the plate was 25-50 eV but the results were initially found to be non-stationary, varying by $\pm$ 35% in the course of the iterations. Closer examination revealed that, although the integral of the sources varied little, strong local variations in plasma parameters resulted. In particular, the density profile evolved from a double-humped to a single-humped state in the course of the B2 iterations with source rescaling between two full EIRENE calls (Fig. 2). Full EIRENE calls are separated by at most 30 ms. The next full EIRENE call triggered the return to a double-humped profile. It appears that, in this temperature range and in the absence of momentum sinks (see below), particle source distributions having the same integral within 10% can lead to either single-humped or double-humped density profiles.

**Improved Convergence of the Monte Carlo model**

Following these initial runs, convergence of the coupled B2-EIRENE code was improved by introducing source mixing: after the full Monte Carlo call, the new sources are taken to be a linear combination of the newly calculated rescaled sources and the sources obtained just before the last Monte Carlo call. This attenuates the effect of the statistical variations of succeeding Monte Carlo calls. Furthermore, to ensure response of the coupled system to changes in input parameters, we require that a minimum number of B2 timesteps (typically 50-100 at 5 $\mu$s per timestep) be performed before a new full Monte Carlo call. This avoids a lengthy initial period during which many full EIRENE calls but few B2 timesteps (little variation of plasma parameters) would otherwise be performed. Full Monte Carlo calls are now separated by at most 1 ms.

With these modifications, convergence is improved and stationary solutions are obtained both at standard power (47 MW into 1/2 the outer SOL) and for a lower power case at one-tenth this value. Solutions for these cases are obtained both without and with momentum sinks. Exponential convergence of the particle and power balance (Fig. 3b) is obtained down to 1% of the particle flux to the plate and of the power input to the SOL respectively. The level at which the evolution levels off appears to be determined by the remaining perturbations imposed by successive full Monte Carlo calls, which typically occur every 200 B2 timesteps. Variations of local quantities such as the peak electron temperature at the divertor plate are typically $\pm$ 5% (Fig. 3a), with each half cycle of the oscillation triggered by a new full Monte Carlo call. For the lower-power case, the B2 evolution can also be followed with frozen sources (no rescaling, no Monte Carlo calls) and then exponential convergence to a level of $10^{-6}$ is obtained without a significant change in parameters. This is an additional demonstration that the solution is well-converged.

**Results with Monte Carlo model**

When momentum sinks due to plasma-neutral interactions are included, a power of 47 MW to 1/2 the outer SOL and an upstream density of $0.35.10^{20}$ cm$^{-3}$ yield a peak electron temperature at the divertor plate of 33 eV and a peak power per unit area on the perpendicular divertor plate of 23 MW/m$^2$. (These cases are given for perpendicular divertor
plates: inclination to 15° would probably reduce this value to the range 5-10 MW/m² without safety and peaking factors, but this remains to be calculated). With momentum sinks included, the double-humped type of density profile described above is no longer observed for the range of parameters we have studied, whereas a test calculation with momentum sinks removed again exhibits this form (Fig. 4).

The electron temperature at the point of peak power on the perpendicular divertor plate is given as a function of upstream density $n_s$ in Fig. 5. A relatively high upstream density, $0.5 \times 10^{20}$ m$^{-3}$ is required to attain 15 eV. This density is considerably higher than previously calculated with the simple analytic model (3-4 times higher than the perpendicular plate calculation and 2x higher than the inclined plate calculation [5]). A part of this difference can be attributed to the (more realistic) inclusion of momentum sinks since the test calculation without momentum sinks requires a density lower by a factor 1.3 to give the same temperature (Fig. 5). The main effect, however, must tentatively be attributed to a reduced recycling due to the more realistic inclusion of 2-D effects in the Monte Carlo model both in the neutral source distribution and in the subsequent neutral particle motion, both of which tend to spread out the particle sources radially. The peak power per unit area is given in Fig. 6 versus the electron temperature at the divertor plate. As in the previous studies, it varies roughly as $T_e^{1/3}$ for the lower values of temperature at the plate (higher values of upstream density). It is about 2/3 the value given by the analytic model for the perpendicular plate case, i.e. the power scrape-off width at the plate is now found to be somewhat broader.

Discussion

Stationary solutions with good convergence characteristics have been obtained with the coupled B2-EIRENE code for conditions of a next step device (1 GW ITER CDA) in the high recycling regime at electron temperatures at the plate in the range 10-30 eV. A local feature of the density profiles (switch from double-humped to single-humped shape), found when momentum sinks are not included, may render convergence more difficult. This feature is no longer present for the cases studied when momentum sinks are introduced. High upstream densities ($\sim 0.5 \times 10^{20}$ m$^{-3}$) are predicted to be necessary to attain 10-15 eV at a perpendicular divertor plate. Coupled B2-EIRENE calculations to investigate lower-temperature regimes at ITER EDA parameters are starting. Development of a 2-D multispecies analytic neutral model to supplement the Monte Carlo calculations for parameter variations is also continuing.

Acknowledgments

D. Reiter and collaborators of KFA Jüllich developed the EIRENE code and the coupling to the B2 code. We would like to acknowledge their help in implementing the coupled code at NET as well as discussions with R. Schneider of IPP Garching.

References

Fig 1. Density profiles along inclined divertor plate, analytic neutral model, for power to ½ outer SOL from 37 to 93 MW

Fig. 2 Density profile along perpendicular plate during initial (nonstationary) coupled B2 - Monte Carlo runs (see text).

Fig. 3 a) $T_e$ at point of peak power flux on plate and b) power balance normalized to input power vs. time during iteration

Fig. 4 Density profile along perpendicular plate for stationary B2-EIRENE runs with and without momentum sinks

Fig. 5 $T_e$ at peak power flux on plate versus density at separatrix and midplane, stationary B2-EIRENE, ⊥ plate

Fig. 6 Peak power flux on ⊥ plate versus $T_e$ at point of peak power flux on plate, stationary B2-EIRENE, ⊥ plate
A MODEL FOR DETACHED SCRAPE-OFF LAYER PLASMAS IN A TOKAMAK DIVERTOR

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1. INTRODUCTION

In JET, X-point discharges with divertor temperatures of \( \leq 5 \) eV, often show a substantial decrease in both plasma pressure and the particle flux density \( I^+_{sat} \), while the upstream plasma conditions in the SOL remain relatively unchanged (detachment). Detachment is observed in both low power ohmic discharges and high power, high density discharges ("gas target" discharges [1]). Analysis has been focussed on the "gas target" discharges. They show strong divertor radiation and total radiative fractions of about 90% without marling and are of potential interest for the problem of power exhaust in next-generation devices.

2. CONDITIONS FOR THE EXISTENCE OF A NEUTRAL CUSHION

In a differentially-pumped, linear magnetic plasma simulator, Hsu [2] demonstrated an externally sustained gas target, where plasma contact with the solid target was prevented by increasing the neutral hydrogen density to sufficiently high levels that ion-neutral \((i-n)\) collisions formed a "neutral cushion". In Ref. [3], the criteria were considered for a self-sustained gas target to form in a tokamak divertor, based on the \(i-n\) collisions associated with the natural recycling flux at the target (see Fig. 1). (The recycling flux is assumed to be large compared with any external fueling rate.)

One criterion is that \(T_e\) be sufficiently low \((T_e \leq 5 \) eV\) that there can be many \(i-n\) collisions before the recycling neutral is ionized. A further criterion is that the \(i-n\) collisions be effective at removing the plasma momentum, which requires \(\lambda_{n-i} \geq \Delta\) \((\lambda_{n-i} \) the neutral mean free path, \(\Delta\) the SOL width\) to hold, so that in general each \(i-n\) collision is followed by a neutral-wall collision. In this situation the \(i-n\) collisions can be expected to have three principal effects on the edge plasma:

(a) In the collisional region (CR) \(i-n\) collisions transfer momentum to the neutrals which then transfer it kinetically to the solid surface. As a result the plasma pressure will drop along \(B\).

(b) In the CR, because of the favourable mass ratio, also ion energy is effectively transferred to neutrals and thus to the solid surfaces, and the ion temperature decays over a distance of about \(\lambda_{i-n}\).

(c) Assuming that the plasma flow velocity at the target sheath is sonic, the effect of \(i-n\) collisions is to reduce the plasma Mach number \(M_C\) at the entrance to the CR and \(M_C \ll 1\) holds for sufficient \(i-n\) collisionality [3]. Since, by definition, there is no ionization source in the CR, also the plasma outflux rate to the target (thus the ion saturation current \(I^+_{sat}\) to target-mounted Langmuir Probes) will decrease for given upstream plasma conditions.

At the entrance of the CR, because of the low Mach number, convective energy
transport is small. Heat conduction is dominated by \( i - n \) induced ion heat conductivity and approximately given by \( (z) \) is the coordinate along \( B \)

\[
q_{\parallel,c} = -\kappa_{i-n} n_C \frac{\partial T_C}{\partial z} \approx v_t \lambda_{i-n} n_C \frac{T_C}{\lambda_{i-n}} \approx n_C T_C c_s(T_C).
\]

where \( C \) denotes values at the entrance of the CR and \( c_s \) is the sound speed at this point. (Note that \( T_e \approx T_i \approx T \) may be a good approximation outside the CR.) Equation (1) is reminiscent of the usual sheath condition with the transmission factor \( \gamma = 8 \) replaced by \( \gamma = 1 \).

3. MODELLING OF DETACHED SOL PLASMAS

In the upstream region, i.e., between stagnation point and entrance of the CR, the physics basically does not differ from that in an attached SOL. Modelling of this part is achieved by applying a standard 1-D two point SOL model with the usual sheath boundary condition of the attached case replaced by proper boundary conditions at the entrance of the CR.

A version of the basic equations of a two point model are given in Refs [4, 5] for the case of Bohm-like perpendicular heat transport \( (S, X \) denote stagnation point and downstream quantities, respectively)

\[
n_X = \frac{n_S T_S}{T_X} \tag{2}
\]

\[
\Delta = \frac{5}{32} \frac{c \ n_S T_S^2}{q_\perp B_t} \tag{3}
\]

\[
T_S = \left( \frac{49 q_\perp L^2}{4k \Delta} \right)^{2/7} \left( 1 - \left( \frac{T_X}{T_S} \right)^{7/2} \right)^{-2/7} \tag{4}
\]

\[
\frac{7 L q_\perp}{2 \Delta} = \int_0^L dz Q_{rad} + \gamma T_X c_s(T_X) n_X \tag{5}
\]

Here \( \Delta \) is the temperature SOL thicknesses and \( q_\perp \) the mean power flux across the separatrix. \( f \) reflects a possible pressure drop along the field lines. Otherwise the notation is conventional. For a detailed discussion of Eqs (2) to (5) see Ref. [4, 5].

Equation (5), which is basically the global energy balance in the SOL, can be in various ways brought into a more tractable form. In the present study we are aiming
at the modelling of discharges where the total radiation from the SOL is known from measurements. In that case one conveniently introduces \( q_{\perp,rad} = \frac{2}{T_x} \int_0^L dz Q_{rad} \), so that

\[
\frac{2 \Delta}{T_x} (q_{\perp} - q_{\perp,rad}) = \gamma n_x T_x c_s(T_x)
\]

Note that \( q_{\perp,rad} = \frac{P_{rad}^{div}}{\pi^2 R_a \sqrt{a}} \), where \( P_{rad}^{div} \) is total SOL radiation power and \( s \) the plasma elongation. Manipulating Eq. (6) by analogy with Refs [4,5], one gets

\[
n_S = C \left( \frac{f}{\gamma T_x^{1/2}} \right)^{11/16} \left( \frac{q_{\perp} - q_{\perp,rad}}{B_t^{11/16}} \frac{1}{L_{1/16}} \right)^{1/8} \left( \frac{T_x}{T_s} \right)^{7/2} \]

Equations (2) to (4) and (7) now alternatively describe attached or detached SOLs if the following specifications/differences are taken into account:

- In the attached case \((X \equiv D)\) one has \( f = 2 \) (because of \( M = 1 \) at the sheath entrance [6]), \( \gamma \approx 8 \) and the divertor temperature is a variable while \( L = \text{const} \).
- In the detached case \((X \equiv C)\) one has \( f = 1 \) (because of \( M << 1 \) at the entrance of the CR), \( \gamma = 1 \) (as outlined above) and the "effective" connection length \( L_C \) (see Fig. 1) is now a variable while \( T_C = \text{const} \) (\( \approx 5 \text{ eV} \)).

Hence in both cases we have four equations for seven quantities \( n_S, T_S, n_D, \Delta, q_{\perp}, \) \( q_{\perp,rad} \) and \( T_D/L_C \).

From Eq (7) one concludes that in the detached case \( n_S \) depends on power essentially through \( q_{\perp} - q_{\perp,rad} \approx P_{heat} - P_{rad}^{bulk} - P_{rad}^{div} \), i.e., the power flux into the CR. \((1 - (T_x/T_s)^{7/2})^{3/8} \approx 1 \). Hence, for comparison with empirical data specific knowledge about the splitting of radiation into its bulk and divertor fractions is not required.

It is illuminating to compare the detached version of Eq. (7) \((T_X \equiv T_C \approx 5 \text{ eV}, \gamma \approx 1, f = 1)\) with the version of an attached SOL with a divertor temperature close to the threshold where ion-neutral collisions become effective (marginally attached case, \( T_D \approx 5 \text{ eV}, \gamma \approx 8, f = 2 \)). Comparing Eq. (7) for these two cases one concludes that marginally attached cases and detached cases differ considerably in the external parameters \((n_S, P_{heat}, P_{rad}^{bulk}, P_{rad}^{div})\) that control the discharge. In particular, transition into the detached state requires finite changes of these external control parameters. In JET discharges, in general, these parameters evolve slowly (as compared to the characteristic time scales of the SOL) [1]. Hence transition into detachment is a gradual process in these discharges.

Model validation requires measurements of at least four of the basic variables of Eqs (2) to (4) and (7). The JET database provides \( P_{heat} - P_{rad}^{bulk} (q_{\perp}), P_{rad}^{div} (q_{\perp,rad}) \) and \( n_S \) from bolometer and Lidar measurements. The dependence on \( L_C \) is so weak in Eq. (7) that one can take \( L_C \approx L \). With this approximation Eq. (7) provides \( n_S \) from \( q_{\perp} \) and \( q_{\perp,rad} \) which can be compared with the measured values. (Note that, though Eq. (7) is basically freestanding, its derivation involves Eqs (2) to (4).)

This is done in Fig. 2 for shot 26389, which is detached after formation of the X-point \((t \approx 12 \text{ s})\), then shortly attaches when the heating power is ramped up and detaches again when in the later stage \( P_{rad}^{bulk} \) and \( P_{rad}^{div} \) increase with increasing \( n_S \). Figure 2 illustrates the difference \( n_S \) shows in attached and detached cases for the same power flux into the CR. It also gives an example for the gradual transition.
In Fig. 3 the special power dependence according to Eq. (7) is checked by comparing measured and calculated $n_S$ values for a variety of detached discharges with different input and radiation powers.

One has to be aware of the rather poor accuracy of the Lidar and bolometer measurements, since otherwise the good agreement as shown in Figs 1 and 2 might be somewhat misleading. Also most gas target discharges show transition into a poor H-mode. This seems to have little effect on the SOL properties, but that remains to be confirmed by an analysis of transition free discharges.

A Numerical Study of CX and Radiation Losses in a Divertor Channel

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JET Joint Undertaking, Abingdon, Oxon OX14 3EA, U.K.

INTRODUCTION
A new 2-dimensional code, EDGE2D/N, has been set up to investigate the physics of a simple divertor configuration. The aim is not model validation by means of detailed simulations of particular shots but rather the study of basic physics phenomena playing a role in the divertor. Special interest is given to the relative importance of transport, power losses due to hydrogenic atomic processes (charge-exchange (CX) and ionization) and hydrogen radiation in divertor relevant plasma regimes. In particular we analyse quantitatively the relevance of hydrogen atomic losses, following the ideas of Rebut and Watkins /1/. Calculations for JET as well as for ITER relevant cases have been performed.

THE MODEL
The FORTRAN source of EDGE2D/N is produced by a pre-processor written in REDUCE. The code solves the 2-dimensional plasma transport equations in a channel simulating the SOL in a deep slot divertor leg (Fig.1). For JET the length and width of the channel are 0.4 m and 0.1 m respectively, for ITER the corresponding values are 4 m and 0.1 m. Transport is classical in the parallel direction and anomalous in the perpendicular (radial) direction. Hot and cold neutrals are treated with a two-group diffusion approximation. Source terms due to CX, ionization and recombination have been included. Different possibilities for neutral recycling are examined. Following /1/ the plasma is neutralized and pumped at the target and refed from the separatrix-side (private region) or the wall (Fig.1). Partial recycling at the target has also been considered. Standard plasma boundary conditions are used at the target. An input power of 10 MW (with 5 MW in ion and electron channels) for JET is given at the channel entrance with an exponential fall-off away from the separatrix. The powers are increased by a factor of 10 in the case of ITER. Reference values for the thermal and particle diffusivities $\chi_0$, $\chi_*$ and $D_1$ are: $\chi_0 = \chi_* = 1.0 \text{ m}^2\text{s}^{-1}$, $D_1 = 0.3 \text{ m}^2\text{s}^{-1}$. The hot neutrals have the local ion temperature and the cold neutrals are assigned a fixed temperature ($\leq 3 \text{ eV}$). The average energy loss $\xi_i$ in eV per electron ionization event (including hydrogen radiation) is given by /2/: $\xi_i = 17.5 + (5 + 37.5/T_i) \log_{10}(10^{21}/n)$ with $n_i$ in $\text{m}^{-3}$ and $T_i$ in eV. Following /1/ we assume that, on average, ionization and CX energy transfer due to the hot neutrals balance.

RESULTS AND CONCLUSIONS
The following tables give a summary for JET and ITER cases. $n_i$ is taken as the maximum value at the channel entrance. The “net CX loss” takes into account the contribution to ions of cold neutral ionization. All powers are given in $\text{MW}$.

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ITER private region-recycling:

<table>
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<th>cold neutral influx $s^{-1}$ $[10^{23}]$</th>
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<th>ionizat. + rad. losses ele.</th>
<th>total rec. losses</th>
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<th>ele. power flux to target</th>
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ITER wall imposed flux:

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JET private region recycling:

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<td>6.0</td>
<td>7.0</td>
<td>0.32</td>
<td>1.07</td>
<td>0.0007</td>
<td>1.39</td>
<td>4.57</td>
<td>4.03</td>
<td>0.002</td>
<td>0.003</td>
</tr>
<tr>
<td>20.5</td>
<td>12</td>
<td>0.12</td>
<td>2.07</td>
<td>0.003</td>
<td>2.19</td>
<td>4.41</td>
<td>3.39</td>
<td>0.003</td>
<td>0.005</td>
</tr>
<tr>
<td>60</td>
<td>18</td>
<td>-0.24</td>
<td>3.27</td>
<td>0.01</td>
<td>3.04</td>
<td>4.24</td>
<td>2.71</td>
<td>0.007</td>
<td>0.006</td>
</tr>
</tbody>
</table>

In the tested range of densities recombination remains negligible. Recycling of neutrals from either the separatrix or wall side presents two rather different scenarios. In the case of separatrix-side recycling the outflow of plasma at the target equals the inflow of cold neutrals. The outflow of neutrals to the wall is negligible in comparison. On the other hand if a cold
neutral influx is imposed from the wall the outflow of neutrals to the wall is in general a large fraction of the injected flux and larger than the plasma flux to the target. The considerably higher flux of neutrals from the wall compared to that from the private region (even though plasma densities and total losses don’t differ much), accounts for the fact that only a fraction of the cold neutrals gets ionized, the rest returning as neutrals. For high densities (ca. $1.5 - 2 \cdot 10^{20} \text{ m}^{-3}$) the temperature along the separatrix is about a factor $2 - 4$ lower for separatrix-side recycling than for wall recycling and decreases in the case of ITER from $280 \text{ eV}$ to $60 \text{ eV}$ (along the separatrix). For lower densities (ca. $3 - 5 \cdot 10^{19} \text{ m}^{-3}$) the variations are smaller, decreasing from $360 \text{ eV}$ to $240 \text{ eV}$. CX is more effective where the plasma temperature is high and therefore to exploit this to a maximum, neutrals have to be injected from the separatrix-side. In contrast, the atomic losses in the case of wall recycling are mainly due to ionization.

Trends of results are similar for JET and ITER cases (Figs 2-4 refer to one particular ITER case). If ITER recycling from the private region is considered, for the range of densities tested, the atomic losses vary around 22-29% of the input power. At very high density (ca. $4.8 \cdot 10^{20} \text{ m}^{-3}$, not in the table) the atomic losses increase again up to 35% of which about two thirds are due to CX. It is more favourable to have some recycling at the target in addition to separatrix-side recycling. In this case the density of the plasma is considerably lower for the same atomic losses. For example in the JET case having 30% separatrix and 70% target recycling lowers the density along the separatrix by a factor of two (from $1.4 \cdot 10^{20} \text{ m}^{-3}$ to $7 \cdot 10^{19} \text{ m}^{-3}$) while increasing the atomic losses from 18.5% to 22%. At the same time the neutral flux decreases from $1.7 \cdot 10^{23} \text{ s}^{-1}$ to $1.2 \cdot 10^{23} \text{ s}^{-1}$.

Concentrating the recycling in the upstream half of the channel is not beneficial: at $n_i = 1.4 \cdot 10^{20} \text{ m}^{-3}$ the losses go from 18.5% to 16.8%. Similarly doubling the channel length gives only a slight increase of the total atomic losses (at $1.4 \cdot 10^{20} \text{ m}^{-3}$ from 18.5% to 21.7%). In order to maximize atomic losses and improve the efficiency of the scheme, the neutral influx (from the separatrix-side) would have to be tailored (e.g. increasing the plasma density as the temperature decreases along the channel). A better scheme probably would imply injection of impurities upstream. In this case recirculation of the neutrals is intended mainly to trap the impurities in the divertor region.

For the range of densities considered here, which extends to rather high values, about 20-30% of the incoming power can be dissipated via atomic losses (including hydrogen radiation). When the density is of the order of $3 \cdot 10^{19} \text{ m}^{-3}$ or less the temperature remains high even if the atomic losses are comparable to cases at high density. To lower the plasma temperature throughout the channel a very high density ($> 2 \cdot 10^{20} \text{ m}^{-3}$) and separatrix-recycling are necessary. These calculations indicate that hydrogen atomic losses are probably not efficient enough to lead to plasma extinction for realistic values of the plasma density and channel length. The energy losses scale well from JET to ITER, which allows a test of the concept at JET.
REFERENCES


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Figure 1. The computational plasma channel: a slab geometry with a non-uniform mesh and a constant magnetic pitch of $\beta = 0.1$ is used.

Figure 2. Plasma density: ITER separatrix-side recycling for $n_i = 1.5 \cdot 10^{20} \text{m}^{-3}$ - $\Delta n_i = 1.15 \cdot 10^{19} \text{m}^{-3}$

Figure 3. Ion temperature: ITER separatrix-side recycling for $n_i = 1.5 \cdot 10^{20} \text{m}^{-3}$ - $\Delta T_i = 12.4 \text{eV}$

Figure 4. Electron temperature: ITER separatrix-side recycling for $n_i = 1.5 \cdot 10^{20} \text{m}^{-3}$ - $\Delta T_e = 6.7 \text{eV}$
2-D MODELLING OF THE JET DIVERTOR

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1. Introduction

The design of the first version (Mk I) of the JET pumped divertor is fixed, but various "advanced" geometries are being studied for the second phase (Mk II).

The two main characteristics that determine a good divertor are the divertor's ability to handle power exhaust at acceptable erosion rates and its control of the impurity production and retention in the scrape-off layer (SOL). A high recycling regime is favourable for both aspects.

For modelling studies, the JET code EDGE2D [1] has been updated (EDGE2D/U). The new code treats impurities self-consistently and allows for targets with an arbitrary inclination with respect to the magnetic field lines. We have carried out studies of divertor performance using EDGE2D/U for different geometrical configurations, showing that the geometry of the divertor chamber plays an essential role, because i) it controls the accessibility to the high recycling regime by influencing the return of neutrals to the main plasma, and ii) it can increase the wetted area on targets, thus reducing peak heat load.

2. A Summary of the Model

In EDGE2D/U, fluid equations for the conservation of particles, momentum and energy are solved for hydrogenic and impurity ions. The electron density is evaluated from quasi-neutrality. The model allows for arbitrarily high impurity concentrations, with the full non-coronal distribution of impurity charge states and the corresponding energy losses being determined. A single impurity temperature, set equal to the hydrogen ion temperature, is assumed. Particle and energy sources due to neutrals recycled and sputtered at the target and chamber walls are computed by a full 2D Monte Carlo module (NIMBUS).

The metric coefficients needed for the transport equations are computed from a two-dimensional mesh derived from the magnetic flux surfaces obtained from MHD computation of JET equilibria, taking the material wall into consideration.

In the calculations reported here we have assumed full plasma recycling at the divertor targets, but other schemes have also been considered. Constant transport coefficients across magnetic surfaces have been used. Matching of experimental decay lengths at the midplane for JET [2] has suggested the choice \( \chi_r = \chi_\perp = 2.0 \text{m}^2\text{s}^{-1}, D_\perp = 0.1 \text{m}^2\text{s}^{-1} \), increased by a factor of five in the divertor region. No inward pinch has been considered so far. We have assumed the input power \( P = P_r = 10MW \). The wall material is Carbon.

3. Comparison of results for two divertor configurations

We have continued studies of several divertor geometries for which preliminary results were reported in [3]. Here we report only the results for two of them, since they represent the extreme cases. The first is the Mk I configuration currently being installed at JET. The
second one could be achieved in a Second Phase of the JET divertor operations; namely, a baffled vertical plate design where recycling neutrals are directed more toward the private region (Mk II). The two configurations are sketched in Fig. 1 and 2, where “S”, “1” and “2” refer to separatrix and flux surfaces 1 and 2 cm beyond separatrix at the midplane.

The possibility of reaching high density and low temperature regimes has been studied by means of a scan of the density at the separatrix from 1 to \(2 \times 10^{19} m^{-3}\), which is expected to be reached in JET at the power level considered. A summary of relevant results is given the following table. Here \(n_s\) is the density at the separatrix on the outer midplane (\(\times 10^{19} m^{-3}\)), \(n_T\) is the density at the outer strike point, \(T_{es}\) is the electron temperature at the separatrix on the outer midplane (eV), \(T_{ir}\) is the electron temperature at the outer strike point, \(P_h\) is the power loss due to the atomic processes of charge exchange and ionization (with its associated radiation loss) of the hydrogen (MW), \(P_r\) is the radiation from impurities, \(R_{S}^{SOL/LIT}\) is the ratio of the hydrogen source in the SOL to that in the SOL plus divertor, \(R_{Z}^{SOL/LIT}\) is the corresponding ratio for Carbon, \(S_Z\) is the total source of impurities (\(\times 10^{23} m^{-3}\)), \(n_{Z,S}\) is the Carbon density at the separatrix on the outer midplane (\(\times 10^{23} m^{-3}\)).

<table>
<thead>
<tr>
<th>Mk</th>
<th>(n_s)</th>
<th>(n_T)</th>
<th>(T_{es})</th>
<th>(T_{ir})</th>
<th>(P_h)</th>
<th>(P_r)</th>
<th>(R_{S}^{SOL/LIT})</th>
<th>(R_{Z}^{SOL/LIT})</th>
<th>(S_Z)</th>
<th>(n_{Z,S})</th>
</tr>
</thead>
<tbody>
<tr>
<td>I</td>
<td>1</td>
<td>1.3</td>
<td>123</td>
<td>116</td>
<td>3.7</td>
<td>0.2</td>
<td>23%</td>
<td>18%</td>
<td>3.1</td>
<td>0.6</td>
</tr>
<tr>
<td></td>
<td>1.6</td>
<td>7.6</td>
<td>77</td>
<td>29</td>
<td>6.4</td>
<td>1.9</td>
<td>13%</td>
<td>6.1%</td>
<td>5.1</td>
<td>3.8</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>16</td>
<td>75</td>
<td>15</td>
<td>8.3</td>
<td>2.7</td>
<td>12%</td>
<td>4.8%</td>
<td>3.2</td>
<td>2.3</td>
</tr>
<tr>
<td>II</td>
<td>1</td>
<td>20</td>
<td>81</td>
<td>15</td>
<td>9.1</td>
<td>1.1</td>
<td>0.5%</td>
<td>0.5%</td>
<td>4.3</td>
<td>2.4</td>
</tr>
<tr>
<td></td>
<td>1.6</td>
<td>35</td>
<td>79</td>
<td>1.3</td>
<td>11</td>
<td>1.8</td>
<td>0.4%</td>
<td>0.1%</td>
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<td>2.7</td>
</tr>
<tr>
<td></td>
<td>2</td>
<td>47</td>
<td>77</td>
<td>0.6</td>
<td>14</td>
<td>1.6</td>
<td>0.2%</td>
<td>0.3%</td>
<td>1.9</td>
<td>1.4</td>
</tr>
</tbody>
</table>

The principal effect of changing the geometrical configuration is to affect the distribution of ionisation sources due to plasma recycling. This in turn affects the profiles of temperature, density and flow. Mk II is more closed to neutrals than Mk I, as shown by the much smaller relative sources in the SOL \(R_{S}^{SOL/LIT}\) and \(R_{Z}^{SOL/LIT}\). Higher target densities and lower temperature are obtained in Mk II than in Mk I for all values of \(n_s\). Thus the total impurity production is also lower except at very low density. Impurities are better retained in Mk II.

The power load due to conduction and convection to the targets is lower in Mk II than in Mk I, not only because the wetted area is larger (Figs. 1, 2), but also because the total radiation loss is larger (see Table). This is illustrated in Fig 3 for the low-density case \(n_s=10^{19} m^{-3}\) which is the most severe case. The near-orthogonal case Mk I has a relatively peaked distribution which requires sweeping to reduce loads to an acceptable level, while the vertical geometry Mk II can be operated without sweeping. Mk I shows an inversion of the temperature profile at the target, the temperature being lower at the separatrix (Fig 4). This effect is due to the vertical targets that push neutrals toward the separatrix, thus increasing the ionisation source there. The effect is overall beneficial because the high temperature on the chamber side (\(\approx20 eV\)) is compensated by the very low density (\(\approx10^{17} m^{-3}\)), thus limiting
sputtering to acceptable levels. In Mk I, instead, the temperature is higher where the density is higher. As the density \( n_s \) increases to \( 2 \times 10^{19} \text{m}^{-3} \), in the case of Mk II the plasma enters a regime of very low temperature and high density at the target, and the production of impurities is reduced (Fig.5 and Table). In this regime the power conducted and convected to the plates is only 20% of the input power. Of course, the power to the target should also include the recombination energy and some fraction of the divertor radiated power. The performance of Mk I also improves with \( n_s \) as expected, but at a lower rate than Mk II (Fig.5 and Table for \( n_s \approx 2 \times 10^{19} \text{m}^{-3} \)). It is worthwhile to note that, while the decay lengths of \( T_e \) and \( T_i \) in the SOL at the midplane are relatively insensitive to \( n_s \) and to the geometrical configuration, the decay length of the density depends on both (for a given power input). This is illustrated in Fig.6, which shows the radial density profiles at the outer midplane for various cases. The density profiles actually depend on the ratio of the perpendicular and parallel transport of particles. However, the parallel transport depends on the flow velocity \( v_{||} \) which in turn depends strongly on the distribution of sources and on the temperature at the targets, under the assumption that \( v_{\parallel} = v_{\text{reac}} \) at the targets. This result should be tested experimentally and may modify results of predictions based on simpler models. As far as impurity radiation \( P_R \) is concerned, we have found that in all of the configurations considered \( P_R \) is less that 15% of the total input power if sputtered carbon is considered. \( P_R \) is marginally larger in Mk I than MK II because more impurities are sputtered in MK I. The situation is opposite at low density. Studies of increasing impurity radiation by injecting impurities at some distance from the target are being carried out at present.

4. Conclusions

Geometrical effects in the divertor can affect significantly the plasma density and temperature profiles, the radiation pattern and the distribution of heat loading on the targets. These effects on the density profile may be important even at the midplane, if \( n_s \) is large enough. In general, a closed divertor such as Mk II is more effective than the open Mk I in alleviating the heat load on the targets, thus reducing or eliminating the need of sweeping, and in retaining impurities. Total atomic losses (mainly from hydrogen) can amount to a large fraction of the input power for Mk II, and increase rapidly with density. The wall material seems to be relatively unimportant since most of the radiation comes from the hydrogen. It is expected that geometry effects would play less of a role at higher densities than those considered here, but such conditions might not be attainable in JET.

5. References

Analysis of Cold Divertor Concepts by 2-D Simulations

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(**) Katholieke Universiteit Leuven, 3001 Heverlee, Belgium
(+) Courant Institute New York, USA

Introduction
The most challenging issue in the design of future fusion experiments is a divertor concept which simultaneously provides tolerable power loading at the target plates and sufficiently large pumping rates. With respect to the target power load a reduction by at least one order of magnitude seems to be required. Actually, two candidates for the required upstream losses are discussed: neutral recycling losses and impurity radiation. This results in the concepts like the radiative slot divertor [1] using impurities to radiate most of the power, or the cold gas divertor variants. Along this latter line, a concept has been proposed for ITER [10] in which the energy is mainly transferred to the divertor side walls by CX-neutrals, before it reaches the target plate.

Obviously, the problem as a whole is very complex involving a lot of atomic physics processes over an extreme range of temperatures and densities. A reasonably accurate description requires at least a 2d treatment, and must include a large amount of atomic and wall interaction data. In the following we present and discuss exploratory 2d simulations of divertors in present and future tokamaks including possible design variants. The main focus will be on neutral gas effects with only a few remarks on impurity radiation.

2-D divertor simulation results
For the simulations we use the two-dimensional multi-species fluid code B2 [2,8] coupled to the three-dimensional linear Monte-Carlo code EIRENE [7]. Exploratory modelling runs were done for ASDEX, ASDEX-Upgrade, JET, and ITER, including some possible variants such as an ASDEX-Upgrade slot divertor. The power into the divertor was chosen according to typical heating powers, assuming in some cases already some bulk energy losses. These simulations cover a large range of parameters. Fig.1 shows 100 neutral particle trajectories launched on both target plates according to the particle and energy fluxes hitting the plates for four in geometry as well as in plasma conditions different cases in order to give a qualitative idea of the neutral behaviour. ASDEX with its symmetrized double null geometry has a nearly transparent divertor, whereas ASDEX-Upgrade and especially the JET case has an opaque divertor and extremely confined neutrals. In contrast, the ITER case [6] is again nearly transparent for neutrals as will be discussed below. Table 1 shows the typical quantities for cold divertor (maximum target temperatures about 10 eV) solutions for ASDEX, ASDEX Upgrade and JET. Obviously, because the mean free path of the neutrals get shorter from ASDEX to JET compared with the divertor
dimension, the relative CX-losses get smaller, whereas total volume losses due to the interaction of electrons with neutrals, e.g., with hydrogen atoms (radiation and ionization) and molecules (dissoziation), increase. Also listed is the recombination energy flux, because every ion that hits the wall deposits at least 13.6 eV recombination energy.

<table>
<thead>
<tr>
<th></th>
<th>ASDEX</th>
<th>Upgrade</th>
<th>JET</th>
</tr>
</thead>
<tbody>
<tr>
<td>$P_{Div}/MW$</td>
<td>0.27</td>
<td>2.15</td>
<td>7.34</td>
</tr>
<tr>
<td>$P_{H,CX}/MW$</td>
<td>0.09 (33%)</td>
<td>0.25 (12%)</td>
<td>0.90 (12%)</td>
</tr>
<tr>
<td>$P_{H,Rad}/MW$</td>
<td>0.08 (30%)</td>
<td>1.08 (50%)</td>
<td>2.44 (33%)</td>
</tr>
<tr>
<td>$P_{Target,e+i}/MW$</td>
<td>0.10 (37%)</td>
<td>0.82 (38%)</td>
<td>4.00 (55%)</td>
</tr>
<tr>
<td>$P_{recomb}/MW$</td>
<td>0.03</td>
<td>0.44</td>
<td>1.53</td>
</tr>
<tr>
<td>$n_{Sep,Mid}/(10^{19}m^{-3})$</td>
<td>2</td>
<td>4</td>
<td>7</td>
</tr>
<tr>
<td>$T_{Target}/eV$</td>
<td>11</td>
<td>10</td>
<td>10</td>
</tr>
</tbody>
</table>

Table 1
Power into the divertor $P_{Div}$, power transferred by CX neutrals $P_{H,CX}$, total volume losses $P_{H,Rad}$ due to the interaction of electrons with neutrals, e.g., with hydrogen atoms (radiation and ionization) and molecules (dissoziation), total power fluxes of electrons and ions to the targets $P_{Target,e+i}$, power flux due to the recombination of ions at the wall $P_{recomb}$, separatrix density at the midplane $n_{Sep,Mid}$ and maximum target temperature $T_{Target}$ for three ASDEX, Upgrade and JET cases. In brackets the normalized power fluxes ($P/P_{Div}$) are given.

<table>
<thead>
<tr>
<th></th>
<th>ITER-1</th>
<th>ITER-2</th>
</tr>
</thead>
<tbody>
<tr>
<td>$P_{Div}/MW$</td>
<td>141.3</td>
<td>35.1</td>
</tr>
<tr>
<td>$P_{H,CX}/MW$</td>
<td>21.0 (15%)</td>
<td>6.1 (17%)</td>
</tr>
<tr>
<td>$P_{H,Rad}/MW$</td>
<td>15.7 (11%)</td>
<td>17.6 (50%)</td>
</tr>
<tr>
<td>$P_{Target,e+i}/MW$</td>
<td>104.6 (74%)</td>
<td>11.4 (33%)</td>
</tr>
<tr>
<td>$P_{recomb}/MW$</td>
<td>8.1</td>
<td>6.3</td>
</tr>
<tr>
<td>$n_{Sep,Mid}/(10^{19}m^{-3})$</td>
<td>3</td>
<td>3</td>
</tr>
<tr>
<td>$T_{Target}/eV$</td>
<td>45</td>
<td>7</td>
</tr>
</tbody>
</table>

Table 2
Same as Table 1 for two ITER cases.

Table 2 summarizes the results for two ITER cases. ITER-1 is a case with high input power for which we get high CX losses, but only because we have a hot therefore for neutrals transparent divertor (see also Fig.1). The (in absolute magnitude) high CX losses are obtained by the very effective transformation of hot CX neutrals into cold molecules at the open baffle structures in the transparent divertor. As a consequence, the suggested cold gas divertor concept for ITER does not seem to be able to work without additional radiation losses, because the neutral recycling losses for realistic separatrix densities (up to $10^{20}m^{-3}$) are limited to at most 30 to 40%. Simulating these additional radiation losses in the simplest way by just reducing
the input power (case ITER-2 in Table 2, e.g. by radiating away a fraction of the input power from closed flux surfaces in a photosphere) we end up with the desirable cold divertor solution and reduced target power fluxes. In this case the divertor is again opaque for the neutrals and the absolute magnitude of CX-losses is reduced.

Discussion and conclusions
The neutral recycling losses are effective if the plasma fan in the divertor is transparent for the neutrals. However, if the mean-free path of the neutrals gets shorter compared to the divertor dimensions the amount of energy transferred by neutrals is reduced drastically.

It seems to be impossible to build a divertor for ITER based only on neutral recycling losses, which are limited for realistic ITER parameters to about 30% for transparent divertor fans of the SOL input power, which are clearly then hot solutions. For typical ITER geometry and plasma data for a cold divertor the mean free path for charge-exchange is very short compared with the divertor dimensions [3], what is confirmed by the results of the ITER-2 case. In this case, the energy loss through neutrals have rather surface-like nature than volumetric ones. The total amount of energy transferred by CX neutrals to the side walls is reduced, because they are not able to penetrate into hot plasma regions. However, in combination with additional radiation losses a cold divertor can be realized. Simple estimates show the possibility and limits of these radiation losses. One can hope that accounting for non-coronal effects, like finite lifetime of the impurities in the divertor fan and CX recombination due to hydrogen background neutrals [1], the radiation losses by low-Z materials are high enough for an operating ITER divertor [4,5]. In fact, modelling runs confirms the possibility that for additional radiation losses it is possible to reduce the power load on the targets to tolerable values or even achieve gas blanket solutions [9]. However, the question of stability appears, because the radiation zone tends to form close to the X-point where the available volume is the largest one. The possibility of stabilizing a kind of divertor marfe is highly questionable from this point of view, because there seems to exist a bistable situation with the stable states of a x-point marfe and an attached state at the plate and with only transient highly dynamically states of divertor marfes.

References
Fig. 1
100 neutral particle trajectories launched on both target plates according to the particle and energy fluxes hitting the plates.
Reversal of plasma flow in tokamak divertors

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* Chemical & Nuclear Engineering Dept., University of New Mexico, Albuquerque, USA

1. Introduction

In a magnetic divertor, retention of impurity ions is expected to be dependent on an expulsive thermal force directed up the gradient of ion temperature being opposed by frictional entrainment in a plasma flow towards the target. Preferred conditions of high recycling, however, can induce a reversal of usual plasma flow, with consequent reinforcement of thermal forces potentially leading to damaging contamination of the core. Backflow in diverted plasmas was first anticipated theoretically by Nedospasov & Tokar', subsequently observed experimentally in DITE and JET, and has been seen in a number of numerical studies.

We report briefly on a systematic investigation of steady-state divertor flow reversal for ITER-relevant conditions, by detailed numerical modelling. The BRAAMS “B2” edge plasma transport code is used, with both analytic approximations and EIRENE Monte Carlo simulation of neutral particle recycling. The flexibility of numerical models regarding physics admitted is exploited to expose the key rôles of redistribution of recycling sources across magnetic surfaces in flow reversal. Concomitant amplification of cross-field ion diffusion in the SOL is also examined.

2. Divertor computational model

Calculations are performed for the lower half of an assumed up-down symmetric double-null poloidal divertor, pertaining to the ITER CDA physics phase. The LINDA discretiser was used to generate an accurate numerical representation, depicted in Fig.1, with high refinements close to each target (cell lengths ≈1 mm along the field) and about the separatrix (equatorial widths ≈1/2 mm). Note some departures from orthogonality occur close to each X-point. For ease of viewing, large variations in metric functions (incorporated in calculations) are suppressed in plots of results below, which are referred to a computational map of the outside lower divertor branch.

Since tests of principle only are intended, strict quantitative rendition of latest design is not attempted. Similarly, consistent impurity transport is not yet considered, but rather bulk flow patterns of mixed $^2$H & $^3$H ions. The “B2” 2-D edge plasma transport model assumes classical parallel coefficients, and we adopt anomalous cross-field diffusivities of $D = 3 \chi_i = \chi_e = 2 \text{ m}^2\text{s}^{-1}$. Interior boundary conditions, set on a closed surface marginally inside the separatrix, fix a density of $n = 3.9 \times 10^{19} \text{ m}^{-3}$, and power efflux (for the whole configuration) of 116 MW, distributed in ratios 3:1 between electrons & ions and 4:1 outside to inside branches. Conventional target sheath conditions, and exterior pedestal values, are imposed. Also physically disjoint but numerically adjacent regions are
separated by purely insulating grid cuts \( \text{BB}'; \text{EE}' \). This might affect flow fields, but a much greater effect probably accompanies exclusion of real inclined target sections, as reconsidered later.

Treatments of neutral particle recycling include (i) full description in automatic linkage\(^8\) to the EIRENE Monte Carlo simulator, (ii) analytic approximation by the so-called Hotston model\(^3,7\), and (iii) an artificial 1-D form, designated "O-minimal", which constrains atoms too to move wholly on magnetic surfaces until reionized. Pumping is by a neutral particle albedo over a plausible duct entrance (G F') with EIRENE, and by chosen target ion reflectivities with analytic models. Recall steady states are sought, and these are ensured through-out by securing very low unsteady residuals in plasma transport balances\(^1,2\).

3. **Mechanism of flow reversal**

A simple visualization of the process of plasma flow reversal may be deduced from the approach of Nedospasov & Tokar\(^3\), and a later similar discussion by Cooke & Prinja\(^8\). For recycling in a 2-D Cartesian space \((x, y)\), denote ion flux density by \( \dot{I}(x, y) \) m\(^{-2}\) s\(^{-1}\) and reionization source strength by \( S_n(x, y) \) m\(^{-3}\) s\(^{-1}\). Let \( y \) be directed normal to magnetic surfaces, and \( x \) be an along-field co-ordinate from a target at \( x = 0 \); then for a Bohm sheath \( \dot{I}_x(x = 0, y) < 0 \). Steady-state particle continuity now yields :

\[
\frac{\partial \dot{I}_x(x, y)}{\partial x} + \frac{\partial \dot{I}_y(x, y)}{\partial y} = S_n(x, y) \, ,
\]

\[
\Rightarrow \dot{I}_x(x, y) = \int_0^x S_n(x', y) \, dx' - |\dot{I}_x(x = 0, y)| - \frac{\partial}{\partial y} \int_0^x \dot{I}_y(x', y) \, dx' \, . \tag{1}
\]

The last term in (1) is generally diffusive, and so of lesser importance. Thus \( \dot{I}_x(x, y) \) changes sign when the net integrated source along a magnetic tube of force exceeds that flux being recycled at its target end. For self-consistent recycling, this fundamentally can occur through motions of neutral particles across magnetic surfaces\(^3,9\), and their reionization in more intense plasma regions displaced from their magnetic flux tube of origin.

A typical flow pattern at high recycling is illustrated in Fig.2(a). Arrows represent only a local direction of ion flux (s\(^{-1}\)) through each cell in the left-hand frame, and are

![Fig.1 Discrete ITER geometry & computational map](image-url)
scaled in proportion to its magnitude on the right. Note this latter involves some spillage of symbols over cells and the target border. Regimes may be characterized by a flux amplification \( \mathcal{F} \), defined as the ratio of total ion flux (s\(^{-1}\)) on the outside lower target (F F') to that over its related innermost boundary (D E). In Fig.2(a), Hotston recycling at \( \mathcal{F} = 100 \) clearly leads to an extensive reversed flow close to the separatrix. Superimposed are contours of reionization “excess”, namely the accumulated source minus transverse ion losses along each magnetic flux tube from the target, normalized by that ion flux to its respective target element (cf (1)). Onset and layer width of backflow in the SOL are strongly correlated with the “excess” front. Also apparent is induced plasma diffusion back towards the core over a portion of the separatrix from the X-point, which could promote intrusion of impurities. A corollary is restriction of core efflux to a narrower separatrix arc, an effect designated7 “choking”, possibly also impeding core exhaust.

Fig.2 Ion flux (s\(^{-1}\)) through each cell of outside lower branch

direction only — left  direction & magnitude — right

(a) Hotston (\( \mathcal{F} = 100 \)) : “excess” contours 1 — 1.0 2 — 1.05 3 — 1.1

(b) 0–minimal (\( \mathcal{F} = 49.5 \))
A more quantitative summary is shown in Fig. 3, for a Hotston case at $F = 50.0$ and EIRENE description giving $F = 288$. In each magnetic flux tube for which peak backflow is at least 5% of its target-element ion flux, this maximum is normalized by total target ion flux and plotted against maximum reionization “excess” in that tube. For both results, a similar step-like growth of flow reversal occurs as net source deposition in the tube exceeds unity.

The O-minimal treatment of recycling allows any reionization “excess” artificially to be excluded. A corresponding flow field involving $F = 49.5$ is depicted in Fig.2(b). As revealed in Fig.3, this degree of recycling in the presence of neutral particle redistribution across magnetic surfaces would produce a large reversed flow; purely in its absence, Fig.2(b) shows it to have disappeared.

Additionally, backflow can produce increased cross-field density gradients$^{10}$ in the SOL, so magnifying diffusive ion flux, fed by recycling sources. Integrated flux ($s^{-1}$) through each magnetic surface is presented in Fig.4 for results involving EIRENE ($F = 288.$), Hotston ($F = 100.$), O-minimal ($F = 49.5$), and weaker Hotston ($F = 20.0$) recycling. Core efflux, of course, gradually rises for effectively larger pumping. The first two cases each exhibit amplification of cross-field ion flux, after an initial fall at the separatrix due to diffusion into the divertor private region. No amplification occurs for O-minimal sources, abolishing flow reversal, although it does when cross-field neutral particle motions are allowed at the same level of recycling (result referred to in Fig.3). In the last instance of low recycling, sources are no longer strong enough under these plasma conditions to produce backflow, and cross-field ion flux again falls monotonically across the SOL.

4. Conclusions

Divertor plasma flow reversal could degrade impurity retention and impede core exhaust. It seems intrinsic to strong source conditions. Detailed ITER-related calculations support its connection for self-consistent recycling with motions of neutral particles across magnetic surfaces and their heterogeneous reionization. Flow results presented, however, must be qualified by omission of charge-exchange momentum losses, and particularly exclusion of actual oblique target sections. These would tend to redirect recycling across the SOL, possibly shifting, diminishing or even enhancing reversed flows.

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References

1 B Braams
NET rpt EUR-FU/XII-80/87/68

2 H Pacher... Nucl Fusion 12(1982)373

3 A Nedospasov & M Tokar
Proc XIIIth EPS Conf, Aachen, '83

4 P Harbour... J Nucl Mat 128/9(84)359

5 L De Kock...
Proc XIIth IAEA Conf, Nice, '88

6 M Petravic... J Nucl Mat 128/9(84)91

7 in NET rpt EUR-FU/86/91-100

8 D Reiter... J Nucl Mat 128/9(82)60

9 P Cooke & A Prinja
Nucl Fus 27(87)1165

10 S Krasheninnikov
Nucl Fus 32(92)1927
INTRODUCTION

A two-dimensional fluid code called UEDGE is used to simulate the edge plasma in tokamak divertors and to evaluate methods for reducing the heat load on divertor plates by radiating some of the power before it reaches the plates. UEDGE is a fully-implicit code being developed jointly by us, R. B. Campbell and D. A. Knoll. For these studies, UEDGE uses a banded matrix solver [1-2] and a fixed-fraction impurity model. Work is presently underway with Campbell and Knoll to include a memory-efficient iterative solver [3] and a model of impurity transport [4]. Simulations of the proposed TPX device [5] show that a few percent nitrogen concentration in the scrape-off layer can radiate up to 80% of the divertor power, thus reducing the peak heat flux and electron temperature at the divertor plate to acceptable values. A comparison of the neutral gas distribution from UEDGE with results from the DEGAS Monte Carlo neutrals code [6] confirms the validity of our fluid neutrals model.

MODEL DESCRIPTION

The UEDGE model includes classical collisional transport [7] along the magnetic field, classical cross-field plasma flow and anomalous cross-field transport. Neutral hydrogen that evolves from the divertor plate as a consequence of recycling or gas puffing is modeled as a two-species fluid; one species represents Franck-Condon neutrals generated by molecular breakup, and another species represents energetic neutrals due to charge-exchange processes in the plasma. Rate parameters for ionization, recombination and charge-exchange processes are obtained from an atomic physics package [8] that includes density-dependent multi-step ionization processes for hydrogen [9].

The fluid model equations are implemented numerically as described in References [2,10,11]. Orthogonal mesh surfaces in our model are based on MHD equilibria from the TEQ code [12]. Figure 1 is a schematic representation of the double-null geometry for TPX that illustrates the configuration for results presented here. We assume up/down symmetry, but allow particles and energy to flow between the inboard and outboard halves of the core and private flux regions.

Fixed density and temperature boundary conditions are used at the innermost core plasma flux surface in our model. The boundary conditions at the outermost open flux surface are zero radial particle flux for ions and fixed temperature for both ions and electrons. For neutrals, the wall albedo (reflection coefficient) is fixed, typically at a value of 0.95 to simulate pumping of neutrals. Boundary conditions at the innermost flux surface in the private flux region under the x-point are the same as above except that fixed temperatures are replaced by zero energy flux conditions.

We use sheath boundary conditions at the divertor plates with energy transmission factors \( \delta_i = 2.5 \) and \( \delta_e = 2.0 + \kappa \) where \( \kappa \) is the sheath potential drop in units of the electron temperature. We allow the ion parallel flow velocity at the plate to be supersonic if it approaches at a supersonic rate. Separate recycling coefficients for the Franck-Condon and charge-exchange neutral components specify the neutral sources at the divertor plates as a fraction of the incident ion flux.
Non-equilibrium radiative losses due to impurities are included in the electron energy balance equation; the energy loss rate is of the form \( n_e \cdot n_Z \cdot L_Z(T_e, n_o, n_e, \tau_{res}) \), where \( n_Z \) is the impurity density and the radiative loss rate \( L_Z \) depends on electron temperature \( T_e \), impurity charge exchange on neutral hydrogen of density \( n_o \), and the residence time \( \tau_{res} \) of the impurity atoms in the plasma. We use a table look-up for the radiative loss rate in UEDGE. The table data are obtained from a series of runs with the MIST code [8] which solves for the impurity charge state distribution, including the effect of charge exchange on neutral hydrogen and finite residence time in the plasma. The radiation rates for nitrogen are shown in Figure 2. One sees that the loss rate is enhanced over the coronal equilibrium rate (lowest curve) for electron temperatures greater than 10 eV. For the simulations presented here we assume a fixed spatially uniform impurity concentration, \( n_Z/n_e \), and we set \( \tau_{res} = 1 \) sec, so non-equilibrium radiation effects are mainly due to charge exchange.

RESULTS

We present TPX simulation results on the effect of various levels of impurity radiation. We neglect drifts and net electrical current flow between divertor plates because these effects destroy the assumed up/down symmetry of the TPX plasma, requiring simulation of the full double-null geometry. The core plasma boundary values were fixed at \( n_{core} = 1.69 \times 10^{19} \text{ m}^{-3} \) and \( T_{core} = T_{core}^i = 300 \text{ eV} \). Anomalous perpendicular transport coefficients are the standard ITER-like values, \( \chi_n^x = 3 \chi_n^1 = 3D_n = 2.0 \text{ m}^2\text{sec}^{-1} \). These values yield 6.4 MW of power crossing the lower half of the separatrix from the core into the scrape-off layer; 5.3 MW via electrons and 1.1 MW via ions. We assume 100% recycling at the divertor plates and 95% recycling of neutrals at the radial boundaries. Results are given for various values of the impurity (nitrogen) fraction.

Figure 3 shows results from a scan of the impurity fraction. The total heat flux entering the sheath at the divertor plates decreases by about 80% as the impurity fraction is increased to 6% of the plasma density. Correspondingly, the radiated power due to the nitrogen impurity increases to 4 MW. The Hydrogen radiation decreased from 1.2 MW to 0.8 MW, and the radial ion heat loss was 0.2 MW. The peak heat flux and temperature at the outboard plate are both strongly reduced by the impurity radiation as shown by the radial profiles in Figures 4 and 5. In these figures, radial distance is measured from the separatrix, positive outward and negative in the private flux region.

Plasma profiles along a flux surface just outside the separatrix in the outboard leg of the divertor are shown in Figures 6 and 7. The left edge of each figure (at \( \approx 2.7 \text{ m} \)) is the x-point position and the right edge (at \( \approx 3.3 \text{ m} \)) is the divertor plate. The solid lines show the results with no impurities and the dashed lines results with 6% nitrogen. As shown here, at high impurity fraction, the peak of the radiation zone moves upstream away from the divertor plate to where the electron temperature is above 10 eV. This is due to the strong dependence of the radiation rate on electron temperature shown in Figure 2. We also find that the peak of the hydrogen ionization source moves upstream, and the ion flow becomes supersonic near the plate for this case.

We test the validity of our 2-fluid neutrals model by comparing the neutral hydrogen density distribution with results from the DEGAS Monte Carlo code. The DEGAS code uses the geometry and plasma distribution from UEDGE, with the same TPX input parameters as described above, but on a radially extended mesh with no impurities present. In Figure 8 we plot the atomic neutral density along a flux surface just outside the separatrix in the outboard divertor leg. We see that the 2-fluid neutrals model, with flux limits imposed, tracks the DEGAS neutral density reasonably well along the entire
length of the divertor leg, especially in the recycling region near the divertor plate. The statistical error bars on the DEGAS results are about ±20%, with some additional uncertainty due to grid resolution.

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REFERENCES
5. THOMASSEN, K. I., this conference (1993)

Fig. 1: TPX mesh configuration

Fig. 2: Nitrogen radiative rates
Heat Flux ($Wm^{-2}$)

Impurity Radiation ($Wm^{-2}$)

Fig. 3: Effect of nitrogen impurity

Fig. 4: Outboard divertor plate $T_e$

Fig. 5: Outboard divertor plate heat flux

Fig. 6: Outboard divertor separatrix profiles

Fig. 7: Outboard divertor separatrix radiation

Fig 8: Neutrals model comparison
ISLAND DIVERTOR CONCEPT FOR THE STELLARATOR WENDELSTEIN 7-X

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Introduction

The proposed Advanced Stellarator Wendelstein 7-X, is a Helias configuration which has been optimized with respect to several criteria concerning plasma performance and coil geometry. The main data are: \( R_0 = 5.5 \) m, \( a = 0.52 \) m, \( B_0 = 3 \) T. The aim of steady-state operation and a large heating power require a divertor configuration for control of plasma-wall interaction and minimization of impurity inflow. The goal is to optimize the magnetic field in the boundary region without losing any of the other favourable properties, to locate suitable target plates in the outflowing plasma avoiding too excessive power loads and to hold back the re-emitted neutral gas for high recycling probability and pumping efficiency. The concept described in this paper utilizes the existence of large magnetic islands or their remnants in the plasma boundary.

Island divertor

In the standard configuration of Wendelstein 7-X (W7-X), five islands at the boundary exist corresponding to the rational value of the rotational transform \( t_0 = 5/5 \). Increasing \( t \) to 5/4 at the edge yields four islands which are larger than those at \( t = 1 \). Usually they exist only in the vicinity of the O-points surrounded by an ergodic layer. The ergodisation can be enhanced by finite \( \beta \) effects or by increased shear at the boundary /1/. The separatrix of the islands provides the diversion of the field lines and takes over the role of the X-line in torsatrons or in tokamaks. The outflowing plasma crosses the separatrix by diffusion and streams along the field lines towards the rear of the island where the target plates are located. The toroidal and poloidal position of the plates is chosen where the radial dimension of the islands is maximal and where the diverted field lines attain the largest distance from the main plasma. The viability of the island divertor concept depends on the stability of the island structure against external and internal magnetic field perturbations, and the capability to guide the outflowing plasma to the target plates, i.e. a certain range of parallel to perpendicular transport is required. Several studies have been made to check this concept and to optimize the geometry of the target plates. These include:

- Studies of the magnetic field structure, localization of the O-points and X-points outside the last magnetic surface, effect of perturbation fields on the islands, and geometry of flux bundles in the vicinity of islands and ergodic regions;
- investigation of particle diffusion and heat conductivity in and around islands by mapping and Monte Carlo techniques;
- optimization of divertor target plates and calculations of heat loads on them;
- design of sweep coils /2/ for control and modification of the edge field structure and for compensation of symmetry breaking error fields by adjusted DC currents, and for reduction of the heat load by AC currents.
Target plates

For a given power flux along the field lines, the intersection angle with the plates determines the power density on the plates. If the power density is held below the technical limit of 10MW/m², an intersection angle of about 2°–3° results with an estimated width of the scrape off layer (SOL) of about 2 cm and a heating power of 10MW. The total length of the plates is then determined by the intersection angle, the requirement that all field lines in the SOL hit the plates, and the requirement that the leading edge of each successive plate is protected by the "shadow" (see Fig. 1) of the previous one. In each of the 5 field periods two target plates are installed with a toroidal length of about 5 m and an average width of 0.5 m, resulting in a total surface of 25 m². These target plates follow roughly the O-point of the ς = 1 islands in a toroidal range of 54 degree (Figs. 2 and 3).

The problem arising from the different configurations created by variation of the rotational transform, ς, is solved by adjusting parts of the plates to each specific case. In the standard case (ς₀ = 1) the field lines hit more the center part of the plates and in the high and low iota case (ς₀ = 5/4 and 5/6, resp.) the target plates are loaded more towards one or the other end.

A Monte-Carlo technique is used to obtain the intersection patterns on the target plates. The SOL is simulated by calculating the particle motion parallel to the field together with perpendicular displacements after each step length (e.g. displacement 0.1cm, step length 30cm, for an anomalous thermal conductivity $\chi_\perp = 3$ m²/s in a plasma which has a parallel conductivity of $10^7$ m²/s). The starting points of the field lines are statistically spread on a magnetic surface close to inside the bounding separatrix; the integration is stopped at the target plates. The intersection points of the separatrix and the adjacent field lines form stripes on the target plates in a width reflecting the width of the SOL (see Fig. 4) separated in the directions parallel and antiparallel to the magnetic field.

The intersection pattern differs for the various cases examined:

Fig. 1. Schematic arrangement of segmented target plates. The length of the plates is chosen to protect the leading edge of each successive plate with the "shadow" of the previous one.

Fig. 2. Magnetic surfaces (one field period, just inside the SOL), target plates, and sweep coils of W7-X viewed from the radial outside.

Fig. 3. Cross-section of the "bean-shaped" toroidal plane of W7-X. Shown are the system of magnetic surfaces, the SOL represented by Monte-Carlo calculation, and target plates. The first wall is shown as a dashed line and the radial dimension of the coil winding pack as the hatched area.
- Standard case (HS5V10N)— variation of \( t_o \) (0.83,0.86,0.89), variation of island size using the sweep coils, variation of diffusions coefficient (\( \chi_x = 3 \text{ m}^2/\text{s} \) and 0.6 \( \text{ m}^2/\text{s} \)).
- Low iota case — variation of island size and radial position
- High iota case — variation of island position and ergodisation

No deterioration of the divertor action occurs in the high \( \varphi \) case as long as perturbations do not spoil the general field structure within the short length of about 3-5 toroidal transits, which takes the field line from the front to the rear side of the island. The stripes are narrower when the diffusion coefficient is reduced, and because the SOL width depends on the connection length and this length on the island size, the width of the stripes decreases with increasing size of the islands. Consequently, the smallest stripes and largest power densities are found in the high \( \varphi \) case.

**Neutral gas behaviour**

The plasma parameters in front of the target plates depend on the plasma surface interaction, most significantly by the interaction with the neutral gas re-emitted by the targets. Low plasma temperatures and high densities in the SOL reduce the impurity source (sputtering) and reduce the impurity inflow if they are ionized close to the target plates. These favourable edge plasma parameters reduce also the power density on the target plates by radiative losses (recombination) and diffusive broadening of the SOL. In this context the aim of W7-X is to establish a high recycling region in front of the target plates. This is attainable by a proper design of target plates, baffles, gas-puff equipment, an efficient pumping system, and the possible operation at a high edge density without a disruptive density limit.

The behaviour of the neutrals is studied with the neutral gas transport code EIRENE /3/. The geometry part of the code is adapted to the complex geometry of the magnetic surfaces and the island structure; see Fig. 5. The intersection pattern on the target plates from the Monte Carlo calculations serve as particle source regions for the neutrals. At present, the code is used as a stand-alone code with fixed plasma parameters. The profile functions for the electron and ion temperatures and the density in the plasma column are: \( T_e(r) = 4.6/[1+(r/20)] \) [keV], \( T_i(r) = 2.4/[1+(r/36)] \) [keV], \( n_e(r) = 1 \cdot 10^{14}/[1+(r/38)] \) [cm\(^{-3}\)]; \( r \) is the average radius of the magnetic surfaces.
The parameters are adapted to code predictions on the basis of transport analysis. In the SOL (outside the separatrix) the field lines contact the target plates; a radial decay length of about 1.5 cm is assumed. The plasma pressure is assumed to be constant along the field lines, however the density towards the target plates is increased by a factor of 4 to a value of $8 \cdot 10^{13}$ cm$^{-3}$ in conformity with a temperature drop to values of 15 to 40 eV. These parameters lead to a re-ionisation of 80% to 90% of the neutral particles just in front of the target plates and nearly all are re-ionized in the SOL; see Fig. 6.

The first interest of the calculations was to explore the pumping efficiency. In W7-X, cryo-pumps are planned behind or near the target plates in a chamber separated from the main plasma by baffle plates, which prevent the neutrals from moving away, especially in those regions where the plasma has a small radial extension. Two alternatives are considered in the calculations. The first uses transverse slits in the target plates of about 1.5 cm width and a target/slit ratio of 3.4 to get the neutrals into the pumping chamber. In this case about 1% of the source flux on the target plates enters the chamber, composed mainly of charge-exchange particles (Fig. 5). A factor of 3 larger flux into the chamber is obtained in the second case. Here the target plates have no slits, however one side of the chamber is left open (Fig. 6). The neutral particles coming from the target plates and from the edge plasma arrive at the pumping chamber after a few reflections.

Conclusions

The island divertor concept presented in this paper adapts the axisymmetric open divertor to the 3-dimensional stellarator configuration of W7-X. The proposed target plates are fit to various values of $z$, permitting a variation of the island size and position; their operation is rather insensitive to magnetic field perturbations. They are located in regions where the islands have a maximum radial extent which implies a large distance from the main plasma. The toroidal length of the target plates is sufficient to keep the power density within technical constraints; a nearly constant intersection angle of the field lines at the highly loaded regions provides a smooth power load and avoids hot spots. Leading edges are avoided by making use of the "shadow effect". There is a good chance to enter the high-recycling regime in a wide parameter range and to collect a substantial fraction of the neutralized particle with the cryo-pumps.

References

SOL MODELING FOR THE W7-X ERGODIC DIVERTOR CONCEPT

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Helias configurations [1, 2] as optimized [3] for the Wendelstein 7-X (W7-X) stellarator [4] exhibit interesting properties also for physics issues which were not original objectives of the optimization procedure, for example the divertor operation, which plays an important role with respect to power and particle exhaust, impurity control, and minimization of erosion and impurity production.

Investigations of the diversion properties of the magnetic field of W7-X (major radius \( R_0 = 5.5 \) m, number of periods : \( N = 5 \)) allowed a divertor concept to be developed [5]: location of five helical troughs at a distance of \( 1/5 \) of the plasma radius outside the last closed magnetic surface (lcms) along the so-called helical edges [6]. These helical troughs fulfill two important conditions: i) The troughs lead to a complete separation of the magnetic field lines starting at the lcms and ending at the divertor troughs from the field lines starting at the first wall, so that the interaction of the charged particles leaving the plasma is completely removed from the first wall and managed on the troughs. This separation is independent of the detailed island positions. ii) There are no leading edges - i.e. the vertical impingement of scrape-off-layer field lines on material objects, which results in excessive heat and particle deposition - is avoided because the deposition patterns do not extend to the boundaries of the troughs.

In order to study the flow of energy and particles outside the confinement region onto the plasma facing surfaces of these troughs, field line tracing is considered to be the basic method. Starting magnetic field lines at the plasma facing surfaces of the helical troughs and tracing them until they end at the troughs yields a map of the troughs onto themselves [7, 8]. Characterization of these field lines according to their lengths shows that i) the foot points on the troughs form ordered areas up to \( \mathcal{O}(10^2 \) m) length (see Fig. 1); ii) there is a clear correlation between field line lengths and minimum distance to the lcms so that an ordered-layer structure prevails up to lengths of \( \mathcal{O}(10^2 \) m) (see Fig. 2), in particular, there are no short field lines which connect the troughs and approach the plasma; iii) field lines of several hundred meter lengths form an inner ergodic area enclosing the lcms and are mapped to line-like regions on the troughs which cover an area of \( \mathcal{O}(1 \) m²).

Because of the layer structure of magnetic field lines in the edge region it appears possible to use these field lines as coordinate lines for plasma edge modeling: in a first attempt the density and temperature distribution in the edge region are calculated by solving 1D particle and energy transport equations along the 3D field lines. In order to solve the 1D transport equations it is necessary - analogous to tokamaks [9] - to define the
region which contains the scrape-off-layer (SOL) plasma, i.e. the region in which volume sources of heat and matter - representing the flow of heat and matter from the main plasma - load field lines (in contrast to the divertor channel plasma (DCP)). Considering the layer structure suggests how to define this scrape-off-layer region, which should show the same geometrical behaviour as the layers, while the thickness of the SOL is a parameter which has to be determined. Such a SOL surface is calculated with the help of the free-boundary MHD-equilibrium code NEMEC [10] which constructs surfaces nearly tangential to the field lines in the region beyond the lcms [7].

FIG. 1. Characterization of the field line map between the troughs. Of the five helical troughs the one centered at \( v = 0.5 \) is shown. Only one half of it is plotted, the other half is given by the stellarator symmetry at \( v = 0.5 \). The field line foot points on the troughs are ordered according to the field line lengths: \( (0 \leq 9 \text{ m}) \), \( (9 \leq 36 \text{ m}) \), \( (36 \leq 81 \text{ m}) \), \( (\geq 81 \text{ m}) \). The white areas are formed by the foot points of the longest field lines, while the darkest ones belong to the shortest field lines. Discretization of the trough area: \( \approx 0.02 \text{ m} \) (poloidal), \( \approx 0.04 \) (toroidal).

FIG. 2. Characterization of the field line map between the troughs. Poincaré plot in part of the triangular cross-section. The small and thick points characterize the different layers which are formed by field lines with different lengths: from the outermost to the innermost layer \( (0 \leq 9 \text{ m}) \), \( (9 \leq 36 \text{ m}) \), \( (36 \leq 81 \text{ m}) \), \( (\geq 81 \text{ m}) \) correspond to \( (*) \) \( (*) \) \( (*) \). The hatched area marks the helical trough, the solid line the lcms, and the dashed line represents the first wall.

Together with the information on the field line positions relative to the SOL plasma region, one-dimensional fluid solutions along the field lines have been obtained [8]. The simplest possible choice for computation corresponding to a high-density situation in front of the targets is energy transport by pure electron heat conduction along the field
lines with boundary conditions corresponding to small temperatures at the plates. This approach yields temperature, density and velocity profiles along the field lines. As an example Fig. 3 shows the temperature along various field lines.

![Temperature Profile](image.jpg)

**FIG. 3.** Temperature along various field lines.

The results for temperature, density and velocity along the field lines together with their three-dimensional structure yield the distributions of these quantities in the SOL.

Within this framework of approximation energy and particle transport across the field lines have been neglected. Combining the 1D calculations with complementary Monte-Carlo methods which make use of the standard assumptions of the SOL modeling [11], and which are described in detail in [5, 6], transport across the field lines is simulated by a ‘diffusion’ of field lines. This diffusion is achieved by random displacements during field line tracing after characteristic mean free paths.

![Field Line Characterization](image2.jpg)

**FIG. 4.** Characterization of the field line map between the troughs taking diffusion into account. Of the five helical troughs the one centered at \( v = 0.5 \) is shown. Only one half of it is plotted, the other half is given by the stellarator symmetry at \( v = 0.5 \). The field line foot points on the troughs are ordered according to the field line lengths: \( (0 \leq 9 \text{ m}), (9 \leq 36 \text{ m}), (36 \leq 81 \text{ m}), (\geq 81 \text{ m}) \). The white areas are formed by the foot points of the longest field lines, while the darkest ones belong to the shortest field lines. Discretization of the trough area: \( \approx 0.02 \text{ m} \) (poloidal), \( \approx 0.04 \text{ m} \) (toroidal).

Taking diffusion into account the areas on the troughs formed by the foot points of the long field lines (see Fig. 4) increase, which corresponds to a more uniformly distributed
particle and energy load on the divertor plates. With zero-dimensional SOL parameter estimates [6, 8] as initial guess the temperature distribution shown in Fig. 5 has been obtained, which can be considered as a first step of an iterative procedure for determining the local SOL parameters.

FIG. 5. Characterization of the temperature distribution in the SOL. Poincaré plot in part of the triangular cross-section. The points of 4 different thicknesses mark the temperature along the field lines: (...) refer to $0 \leq 35$ eV, $35 \leq 70$ eV, $70 \leq 105$ eV, $\geq 105$ eV. The solid line shows the lcms. The SOL width is $\approx 0.02$ m.

REFERENCES

1. Introduction
Predictions for the plasma edge and divertor behaviour in future experiments require realistic, 2-dimensional numerical models, which have to be validated against present experiments. For ASDEX-Upgrade we use the B2-EIRENE code package to describe the scrape-off layer plasma and the neutral gas dynamics. Many exploratory runs have already been made to study exhaust optimization, mafies or deep divertors.[1–3] Here we present the first attempts for a quantitative comparison with experimental results from Ohmic ASDEX-Upgrade discharges.

2. Codes
For modelling of the ASDEX-Upgrade scrape off-layer plasma, the Garching version of the two-dimensional multi fluid code B2[4] and the three-dimensional linear Monte-Carlo code EIRENE[5] have been coupled in a fully self-consistent way[3]. The first steps of model validation discussed here, however, were done in a single-fluid calculation. This version includes the elastic collisions[6], and in addition to the targets, the vessel wall is also treated as a particle source.

3. Experiment
As a starting point we have taken an ohmic discharge in deuterium with $I_p = 800$ kA, $B_T = -2$ T (i.e. ion $\nabla B$ drift downwards), $q_{95} = 3.8$, $n_e = 3 \times 10^{19}$ m$^{-3}$ (# 2421–2424). The vacuum vessel was boronized. With the measured loop voltage of about 1 V these discharges have an ohmic input power of about 800 kW, with about 200 kW radiated by the main plasma. The magnetic configuration was a standard divertor with a single null divertor at the bottom of the plasma. In this low-$\beta_p$ configuration the separatrix on the lower outer target plate sits relatively far out on the plate (about 3/4 of the plate length). This limits the radial extension of our computational grid which has to be “fixed” to the target plate.

During this experiment a large number of diagnostics were used to study the divertor, including visible spectroscopy, thermography, a moveable Langmuir probe, neutral gas measurements and $H_\alpha$ diagnostic. Unfortunately, the edge diagnostics in the main chamber have not yet been fully available.

Later a second experiment was performed using the same magnetic configuration, but involving a density scan (#3033–3053).

4. Model validation
4.1 Divertor plasma profiles
First, we concentrate on the main, power conducting part of the outer divertor fan. Divertor plasma profiles were mainly taken from the moveable in-vessel probe[7], which gives horizontal
profiles of $n_e$ and $T_e$ roughly 10 cm above the separatrix intersection with the lower outer target plate. Since at the respective density no strong gradients are expected, and since the probe itself is quite massive, we project the probe positions onto the target plate and treat the data as if it were target plate data.

The best agreement of the B2-EIRENE results with the probe data is achieved using a quite small particle diffusion coefficient $D_L = 0.1 \text{ m}^2/\text{s}$, and $\chi_e = 1.5 \text{ m}^2/\text{s}$, $\chi_i = 0.5 \text{ m}^2/\text{s}$. No inward drift velocity has been used, and these transport parameters are assumed to be poloidally constant. The comparison in figure 1 shows good agreement with regard to $T_e$, while the density profile shows some discrepancies. The decay length of both profiles looks similar, but the maximum of the measured density profile is shifted with respect to the temperature profile and is inside the separatrix (which is known within ± 1 cm), in the private flux region. Such a shift of the density profile has been observed already in ASDEX[8]. In ASDEX and in ASDEX-Upgrade this shift changes sign with reversal of the toroidal field direction[7]. This might be due to classical drift effects, which are not yet included in the present version of the code. These drift terms are presently being implemented into the B2-code[9], and their influence will be checked in further calculations.

On the outer side of the divertor plate the measured $n_e$- and $T_e$-profiles show a wide shoulder which extends out to the divertor baffle, but was not included in the present numerical grid as mentioned before. The abrupt drop in the calculated profiles is a consequence of the boundary conditions at the grid edge. Similar shoulders in $T_e$ and $n_e$ have already been observed in ASDEX and were attributed to enhanced turbulent transport at the outer edge region[10].

The power load onto the target plates measured with the thermography system[11] shows an in-out asymmetry (90 kW on the inner target, 230 kW on the outer target), as it is usually the case with the ion $\nabla B$-drift towards the divertor. Without the drift terms mentioned above, the code calculations show a much smaller asymmetry with about 175 kW on the inner target and 200 kW on the outer target. The power load on the outer target agrees reasonably with the measurement, but on the inner target the code calculates a power load which is too high by a factor of 2. This might change with the implementation of the drift terms[11].

4.2 Neutral gas measurements

The flux density of the neutral gas in the divertor chamber is measured with two different diagnostic systems, namely with a residual gas analyzer (RGA) close to one of the turbo pumps, and with in-vessel ionization gauges (ASDEX gauges)[12]. Both systems are absolutely calibrated with a capacitance gauge measuring the pressure inside the vessel in runs with a toroidal field pulse, but without plasma. From that pressure the particle flux density is derived assuming room temperature. The ionization gauges are installed at three different positions in the vessel: the “divertor gauge” behind the lower passive stabilizing loop (PSL), the “X-point gauge” in the private flux region, between and below the target plates and the “main chamber gauge” in the top of the vessel.

In this discharge the main chamber gauge measures a flux density of $2.3 \times 10^{15} \text{ cm}^{-2}\text{s}^{-1}$, in very good agreement with the code, while the divertor gauge measures a flux density of about $1.5 \times 10^{17} \text{ cm}^{-2}\text{s}^{-1}$, which is nearly 2 orders of magnitude higher, and in reasonable agreement with the code calculations which yielding a flux density of $1.1 \times 10^{17} \text{ cm}^{-2}\text{s}^{-1}$ behind the PSL. Since the plasma extends further out than simulated in the code, as is indicated by the probe data discussed before, this discrepancy is understandable.

For the modelling discharges the RGA measurement was not available, but for the second experiment involving discharges at different densities, all the neutral gas measurements were available, and their results are shown in figure 2. All the flux densities show a strong increase with
line-averaged density, while the absolute values are different. At densities of $n_e \leq 2.5 \times 10^{19} \text{ m}^{-3}$ the flux onto the divertor gauge is a factor of about 2.8 higher than the flux onto the RGA, then this ratio decreases to about 2 at $6 \times 10^{19} \text{ m}^{-3}$. B2-EIRENE calculates a flux ratio of only 1.1 at $n_e = 3 \times 10^{19} \text{ m}^{-3}$. However, a contribution of 20% hydro-carbons in the divertor region, as measured with the RGA[13], would double the ionization gauge signal.

The measured flux density at the X-point gauge is a factor 2 higher than the divertor flux density for all densities $n_e \leq 4 \times 10^{19} \text{ m}^{-3}$, and only at the highest density where the plasma is detached[14] they are the same. The code package calculates the flux density at the X-point to be only half the one in the divertor. This again can probably be related to the power load onto the inner target plate being too high in the code.

At $n_e = 3 \times 10^{19} \text{ m}^{-3}$ the discharge should be the same as the discharge discussed in the previous section. However, the divertor gauge shows a flux density which is about 40% higher, probably due to different wall conditions. Nevertheless the calculated flux density behind the PSL is also shown in figure 2, and using the same transport parameters we have also simulated discharge with a lower and a higher edge density corresponding to $n_e = 2.6 \times 10^{19} \text{ m}^{-3}$ and $n_e = 3.3 \times 10^{19} \text{ m}^{-3}$ respectively, also shown in figure 2. As discussed before the calculated results are lower than the measured ones, but the density dependence is the same.

5. Summary

First attempts to compare simulations of the B2-EIRENE code package with experimental data from ASDEX-Upgrade have been presented. They yield a very small diffusion coefficient, but quite typical values for $\chi$. The central, power conducting part of the scrape-off layer is already described well. More details have to be included, such as the shift of the density profile and shoulders on the $n_e$- and $T_e$-profiles which have also been observed on ASDEX. As next step also impurities have to be included (multi-fluid simulation), and spectroscopic measurements (impurities and $H\alpha$ light) can be checked against the simulations.

6. References

[2] KASTELEWICZ, H. et al., these proceedings.
[13] POSCHENRIEDER, W. et al., these proceedings.
[14] LAUX, M. et al., these proceedings.
Figure 1. Comparison of the in-vessel probe data (projected onto the divertor target plate) with the results from B2-EIRENE, using \( D = 0.1 \text{ m}^2/\text{s}, \chi_e = 1.5 \text{ m}^2/\text{s}, \chi_i = 1.5 \text{ m}^2/\text{s} \).

Figure 2. Neutral gas flux densities measured at different positions in the vacuum vessel as a function of the electron density (#3033-3053). The dark circles show the results calculated with B2-EIRENE for the similar discharge #2424, with the transport parameters as in figure 1.
Carbon Transport in the Plasma Edge of TEXTOR

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Introduction:

There are a number of experimental observations at TEXTOR concerning the impurity distribution in the plasma periphery, which require a 2D-modelling. Examples are the spectroscopic measurements of the spread of \( Li^0 \) and \( Li^+ \) into the plasma edge after \( Li^0 \)-beam injection from a laser irradiated target with a well-known space-time behaviour of the source [1], and measurements of the poloidal [2] and radial [3] distribution of intrinsic carbon arising from chemical or physical sputtering of the limiter and wall structures. The solution of the carbon impurity transport problem, which we will present here in a preliminary form, extends to a rather large integration domain and requires numerical methods (see fig.1 for the triangulation of the applied finite element method). We solve the problem in the trace-impurity limit for a given hydrogen background plasma in contact with the ALT-II limiter assumed to be the dominant source of carbon atoms (due to physical sputtering by hydrogen ion bombardment). We account for drift motion, anomalous cross-field diffusion and classical transport of the \( C^{z+} \)-ions along \( \vec{B} \) with a strong collisional coupling to the electrons and hydrogen ions. The hydrogen plasma parameters are taken from SOLXY-code calculations [4] with appropriate extrapolations to the plasma interior.

Analysis:

The analysis of the density profiles of \( C^{z+} \)-ions starts with the particle balance equation in toroidal coordinates \( (r, \theta, \phi) \) assuming axisymmetry \( \partial / \partial \phi = 0 \) and the standard magnetic field structure \( \vec{B} = (0, -\Theta, 1)B_0 / h \) with \( \Theta^2 \ll 1 \) and \( BR \simeq \text{const.} \), where the major radius \( R = R_0 + r \cos \theta = \)
$R_{oh}$ governs the strength of the curvature effects:

$$
\frac{1}{r} \frac{\partial}{\partial r} \left[ r (\Gamma_r^A + V_r^E n) \right] + \frac{1}{r} \frac{\partial}{\partial \theta} \left[ -\Theta \Gamma_\parallel + V_\theta^E n \right] + \frac{\Gamma_R}{R} = S_n - \nu_{1n} \tag{1}
$$

$S_n - \nu_{1n}$ describes the effective particle source due to ionization and recombination processes. $\Gamma_r^A = -D_\perp \partial n/\partial r$ and $V_r^E$ are respectively the radial particle flux due to anomalous diffusion and the components of the electric drift velocity. $\Gamma_R = \Gamma_r^A \cos \theta + \Theta \Gamma_\parallel \sin \theta + \Gamma_\parallel^D$ is the particle flux in the direction of the major radius. Its drift contribution $\Gamma_\parallel^D$ is connected with the vertical electrostatic and pressure forces (along the coordinate $\zeta$ at constant $R$) induced by the centrifugal and grad B drifts as well as by the vertical asymmetry in the source function (poloidally localized limiter source). Neglecting viscosity effects and presupposing the same temperature for longitudinal and cross-field thermal motion for all ions, we have:

$$
\Gamma_\parallel^D = \frac{E_\zeta}{B} n - \frac{\partial}{\partial \zeta} \left[ \frac{\rho V_\parallel^2 + 2p}{zeB} \right] \tag{2}
$$

where $\rho = M n$. Note that in the limit $R \to \infty$ the diamagnetic terms disappear from the particle balance equation. $\Gamma_\parallel = n V_\parallel$ is the particle flux along $\vec{B}$ induced by the neutralizer action of the limiter and by the short-circuit of $E_\zeta$ (Ohmic discharges). Since it is especially pronounced at the plasma edge, where the inequalities $(D_\perp/\Delta_\perp, V_\parallel^E/\Delta_\perp) \ll \alpha \nu_C$ hold ($\Delta_\perp$ being a characteristic gradient length in radial direction), we may calculate $V_\parallel$ from the 1D approximation of the parallel momentum equation:

$$
\frac{1}{r} \frac{\partial}{\partial r} \left[ \rho (V_\theta^E - \Theta V_\parallel) V_\parallel - \Theta p \right] - zen E_\parallel = R_\parallel + S_{p\parallel} - \nu_{1p} V_\parallel \tag{3}
$$

where $S_{p\parallel}$ denotes the momentum source function along $\vec{B}$. $R_\parallel$ accounts for momentum transfer from elastic collisions (friction and thermal forces):

$$
\frac{R_\parallel}{\rho} = \alpha \nu_C (V_\parallel^+ - V_\parallel) - \Theta \frac{z^2}{M} \left[ 2.21 \frac{\partial kT_+}{\partial \theta} + 0.71 \frac{\partial kT_-}{\partial \theta} \right] \tag{4}
$$

Neglecting the pressure gradient along $\vec{B}$ (see also [6]), eq.(3) can be brought into the form $\rho (V_\theta^E - \Theta V_\parallel^0) \partial V_\parallel^0/\partial \theta \simeq zen E_\parallel + R_\parallel^0$ corresponding to the auxiliary condition for the velocity parameter in the shifted Maxwellian used in the moment method of ref.[5b]. Adopting this method, we consider the neglected terms in the moment equation for the first order expansion coefficient of the distribution function.
In a first computation campaign we neglected $\Gamma R/R$ in equ.(1) and treated the extreme collisional case $v_T/\Delta_|| \ll a v_C$ thus avoiding the complications inherent to the nonlinear convection and the boundary condition for $V_\parallel$ [7]. The above equations were supplemented by boundary conditions similar to those discussed in [5a] (see fig.1) and an analytical source function due to hydrogen ion sputtering of $C^0$-atoms from ALT-II with yield data taken from ref.[8]. Being bound to the limiter geometry as used for the SOLXY-code calculations we assumed a cosine angular (but 4eV monoenergetic) distribution of the sputtered $C^0$-atoms around a preferential direction into the plasma (under 45° with respect to the surface normal of the limiter) thus simulating to some extent the real limiter shape.

Results:

Figs.2-3 show the radial density profiles $n_z$ of $C^{2+}$ for $z = 3,4$ at inboard and outboard poloidal positions, as calculated for $D_\perp = 4 \times 10^{18}/n_p$ in MKSA-units using the SOLXY-plasma model [4]. As expected from $\nu_\parallel \tau_\parallel \geq 1$ with $\gamma_\parallel \equiv q R_0/v_T$ and in qualitative agreement with experimental data from ref.[2] the profiles of the first three charge states of carbon (up to the Li-like state) exhibit a marked poloidal dependence with an ion concentration in the vicinity of the ALT-II limiter. $C^{4+}$, appearing deeper in the plasma, is rather uniformly distributed in poloidal direction due to $\nu_\parallel \tau_\parallel \leq 1$. At the outer midplane the calculated $C^{4+}$-profile has a maximum of $\approx 1.6 \times 10^{11}/cm^3$ at $r \approx 42cm$. In view of the assumptions made in the theory the agreement with experimental data from ref.[3] must not be overestimated. It should, however, be noted that in contrast to the 1D modelling of ref.[3] the source function has been calculated from the plasma edge parameters. Further studies are necessary to improve the calculation of $V_\parallel$ and to explore systematically the processes involved in establishing the density profiles.

References:
[1] G.Kocsis et al., this conference
Fig. 1 Poloidal Plasma Cross Section: Triangulation and Boundary Conditions

Fig. 2

Fig. 3

density of $C^{3+}$
density of $C^{14+}$
The magnetized sheath in front of a flush-mounted Langmuir probe is studied with two-dimensional particle-in-cell simulations (2d real space, 3d velocity space). The dependence of the sheath structure and the probe current on the angle between magnetic field and probe surface, the probe size and voltage is determined. Non-saturation of the ion-current in case of small angles is explained by an increased effective area owing to the bending of ion trajectories in the sheath.

1. Introduction

Recent measurements with Langmuir probes which are flush-mounted in divertor target plates, and thus tilted against the magnetic field lines, deviate from the standard probe theory, e.g. no saturation of the ion current is observed in case of grazing magnetic field incidence[1].

In front of an absorbing wall which is at lower potential than plasma potential, a Debye sheath with positive space charge is formed. Particle simulations showed that, if a tilted magnetic field is present, plasma and Debye sheath are seperated by a quasi-neutral magnetic sheath with length scale equal to the ion gyro radius[2, 3]. In the sheath the streaming velocity of the ions, which is parallel to the magnetic field at the entrance and equal to ion sound speed $c_s$, is bent towards the normal to the wall, so that the normal velocity component can finally exceed sound speed. The ion gyro radius is strongly increased. Consequently, for a detailed study of the sheath a kinetic model is necessary. Therefore particle simulations of a suitable slab model of a flush-mounted Langmuir probe (Fig. 1) were performed.

2. 2D PIC Code and Numerical Model

Since the sheath of a finite size probe in an oblique magnetic field is essentially two-dimensional, an existing 1d-3v (variables $x, v_x, v_y, v_z$) particle-in-cell (PIC) Code[2] was extended to two spatial dimensions (2d-3v). The full particle motions in the self-consistent electrostatic potential and a homogeneous magnetic field are calculated using Boris' algorithm[4] with splitting of the electric force. In some runs electrons are treated in guiding center approximation to speed up the calculation. The Poisson equation is solved on a rectangular grid, and area-weighting is used for charge deposition and electric field interpolation. Since the main part of the PIC Code is intrinsically parallel, the code was ported to and run efficiently on the NCube parallel computer at IPP.

The numerical model consists of the collisionless sheath in front of an absorbing wall containing a long narrow probe oriented parallel to the ignorable coordinate $z$ (Fig. 1). The magnetic field and the electric field are both in the $x$-$y$ plane. Particles are created with probability corresponding to a Maxwellian flux at the boundary opposite to the
Fig. 1: Geometry of the sheath model.

Fig. 2: Contours of the electric potential (a), and ion flux vectors (b) from a simulation with $\alpha = 30^\circ$, $mi/me = 900$, $\Omega_e = 2$, $V = -20$. 
3. Results and Conclusions

Many simulations with the following parameters have been made, \( \frac{m_i}{m_e} = 900, \ \Omega_e = 1 \) or \( 2, \ T_e = T_i, \ d = 60, \ 5^\circ \leq \alpha \leq 90^\circ, \) and \( -20 \leq V \leq 0. \) The dimensionless variables

\[ \Omega_e = \omega_{ce}/\omega_{pe}, \ d = \text{probe width}/\lambda_D, \ \Delta = \text{sheath thickness}/\lambda_D, \ V = e(\Phi - \Phi_{fl}) \]

are used, where \( \Phi_{fl} \) denotes the floating potential. In Fig. 2 results of a simulation with \( \alpha = 30^\circ \) and \( V = -20 \) are presented. Fig. 2a shows the contour plot of the electric potential. The two-dimensional Debye sheath with narrowly spaced equipotential lines is clearly to be seen; the direction of \( B \) is indicated by the arrow. There are two main effects of this sheath:

First, the electrons are completely reflected (if \( -V \gg 1 \)) so that a shadow region at one side of the probe (top in Fig. 2) exists which receives no electrons. And second, additional ions are focused onto the probe surface by the strong electric field as can be seen in Fig. 2b which shows the macroscopic ion flow towards the wall. This focusing, which is strongest towards the edges of the probe, leads to a hollow current distribution on the probe (Fig. 3). The Debye sheath growth is well described by the Child-Langmuir type relation \( \Delta = 6 + 0.8 (-V)^{3/4}. \)

The current-voltage characteristic for the parameters of Fig. 2 is shown in Fig. 4. The single points are from simulations, the dashed curve is \( 1 - e^V \); clearly the current does not saturate at large values of \( -V \). The solid curve represents the relation \( I/I_0 = 1 + 0.023 (-V)^{3/4} - e^V \), which is easily explained by assuming that the effective probe area for the ions is increased by the factor \( 1 + 0.023 (-V)^{3/4} \). The same result is obtained from the Child-Langmuir relation, if the effective probe width (perpendicular to the field...
lines) is given by $d_{\text{eff}} = d \sin \alpha + (\Delta - 6) \cos \alpha$. Experimental I-V characteristics which exhibit non-saturation are of the same form $1 + c (-V)^{a} - e^{V}$ with $a \approx 0.75$.

However, in experiments the sheath is thin compared with the probe dimension so that it can cause a considerable effect only if the magnetic field is almost parallel to the wall ($\alpha \approx 1^{\circ}$). Since two-dimensional simulations with small $\alpha$ ($< \approx 5^{\circ}$) need a very big grid, and consequently a very large number of particles, 1d simulations with $\alpha$ down to $1^{\circ}$ were made to determine how the sheath extension scales at small angle. The result is that the Child-Langmuir type relation remains valid, but the coefficient of $(-V)^{3/4}$ strongly increases with decreasing angle owing to the change of the length scale from the Debye length to the ion gyro radius. Thus the results of the 2d simulations with medium-size angles can at least qualitatively be applied to the small-angle case, too.

In conclusion, the sheath inhibits saturating of the ion current to the probe by increasing the effective probe area.

References


Numerical Marfe Simulations at ASDEX and ASDEX-Upgrade

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Introduction

Marfes (or radiation instabilities) are observed as precursors of the density limit in tokamak discharges. They are ploidally localized, axisymmetric structures in the tokamak edge region with strongly enhanced plasma density, reduced temperature and high radiative energy losses due to impurities. Their appearance is accompanied by a strong decrease of the plasma power flux to the divertor plates or even by complete plasma detachment. Marfes are observed as quasi steady state configurations or moving phenomena. The basic mechanism driving the marfe formation has been discussed in various papers /1, 2, 3/.

Marfes and their relation to the global density limit have been extensively studied in ASDEX /4, 5/. In order to confirm the validity of the basic marfe model, numerical simulations of the marfe dynamics in realistic ASDEX double null geometry have been done with the B2 code stand-alone /5/, as well as with the coupled B2–EIRENE package /5, 6/. Because of up-down symmetry and toroidal curvature, a marfe is formed always at the high field side midplane above a critical edge density. With an experimentally determined carbon density of $0.5 \times 10^{18} \text{m}^{-3}$, the critical density for marfe onset agrees roughly with the experimental marfe limit. The anomalous transport coefficients in these runs were adjusted in such a way that experimental edge and divertor quantities below the marfe limit were reasonably fitted /7/.

In the following, a similar simulation study is presented in some detail for ASDEX Upgrade single null discharges. Though the physics should be rather similar, we expect strong modifications in the marfe pattern, since the geometry and the source and sink distributions are completely different.

Numerical Model

The edge and divertor plasma is described in the hydrodynamic approximation by the 2d multifluid code B2 (for the ions of D, He, C), which is self-consistently coupled to the 3d Monte-Carlo-code EIRENE for the neutrals. EIRENE takes into account all essential details of the discharge vessel. The numerical grid is produced from an ASDEX Upgrade equilibrium using the Sonnet grid generator.

In addition, for the extensive parameter scans (particularly for obtaining the limit of marfe onset), we also used the B2 code in connection with a simpler neutral recycling model and an adhoc ansatz for impurity radiation losses which, however, have been calibrated against the full version. The - a priori unknown - anomalous cross field transport is modelled by an 1/n transport law using coefficients which have been obtained from a first fit at ASDEX:

- Diffusion coefficients: $D = D_0 \cdot (3 \cdot 10^{19} \text{m}^{-3}/n_i)$, $D_0 = 0.5 \text{m}^2\text{s}^{-1}$.
- Thermal conductivity: $\kappa_{\text{e},i} = n_{\text{e},i} \chi_{\text{e},i}$, $\chi_{\text{e}} = 3 \cdot D$, $\chi_i = 2 \cdot D$.

Other transport laws may be used in the future once more detailed data are available.

Marfe Formation

Marfe formation is initiated by successively enhancing the plasma density for sufficiently high impurity concentration. At a critical density a marfe is formed (usually near the (lower) x-point) which is connected with a drastic change of the whole edge plasma state. As a consequence, strong changes
are experimentally observed, e.g., an enhancement of the impurity radiation and bremsstrahlung in the vicinity of the marfe, and a decrease of the total power flux and change of the power deposition profile at the target plates. Figs. 1a,b compare calculated contour plots of the electron density just before and just after marfe formation. The location of the marfe is clearly to be seen by the strongly enhanced density inside the separatrix near the x-point. The shape of the marfe, its location and the changes of plasma parameters caused by the marfe may, however, strongly vary with the discharge parameters.

Figs. 1 Contour plot of the electron density just before (a) and just after (b) marfe formation (n_{e,g} = 9.5 \times 10^{18} m^{-3} (a) and n_{e,g} = 1.0 \times 10^{19} m^{-3} (b); P = 1 MW, q = 4.45)

Marfe Limits

The limits for marfe onset and marfe disappearance have been determined in dependence on the input power P, the safety factor q and the plasma density n_e (simplified recycling and radiation model, impurity concentration c_{imp} = 5\%). Since the bulk plasma density (or line average density) is not available within the present model we used the electron density at the inner boundary of our computational region (outer midplane), n_{e,g}, as the characteristic quantity presuming that this plasma region is less influenced by the particular shape and location of the marfe itself.

Figs. 2 Limit for marfe onset (full lines) and the disappearance of marfes (broken lines) as function of the input power P for q = 4.5 (a) and as function of the safety factor q (b) for P = 1 MW (lower line) and P = 3 MW
With the above assumptions we find (Fig. 2a) that the density limit for marfe onset increases linearly with the input power $P (q = 4.5)$. The density limit where the marfe vanishes for decreasing density (Fig.2b) behaves more complicated and shows irregularities, which may be partly due to the non-linear nature of the marfe itself which changes the plasma parameters strongly when approaching the very limit. For higher powers ($P \geq 2.5$ MW) a remarkable hysteresis exists, i.e., once the marfe is formed, it may exist down to densities lower than the limit for marfe onset.

Contrary to experimental findings at ASDEX / 8/, practically no dependence of the marfe onset on the safety factor $q$ is found here (Fig. 2b). However, it may be seen from Figs.3 that marfes formed at larger $q$ have a much higher relative density peak than those at smaller $q$ and, hence, are experimentally easier to detect. Again, a hysteresis is obtained for all $q$ values in the case $P = 3$MW.

![Fig. 3 Marfe for the same conditions as in Fig.1b, but for smaller safety factor $q = 2.8$; $n_{e,g} = 1.0 \times 10^{19}$.](image)

Figs. 4 a) Calculated $C^{1+}$ radiation field of a marfe with double structure; b) Numerically calculated optical picture of this radiation field as it would be seen by the CCD referred to in Fig.4c; c) Photograph of a marfe obtained by a CCD camera (see text)

**Radiation and Plasma Profiles**

The pattern of impurity radiation emitted by the marfe depends sensitively on the shape and intensity of the marfe. As to the considered carbon dominated plasma, mainly the ions $C^{3+}$ and to a lower extend $C^{2+}$, $C^{1+}$ and $He^+$ contribute to the radiation losses of the marfe. $C^{4+}$ and $C^{5+}$ radiate
almost homogeneously around the core plasma. For strong marfes (produced at high plasma density and impurity concentration) the temperature decrease in the center of the marfe becomes so large that line radiation ceases despite the high electron density. In this case radiation comes mainly from the steep shell surrounding radially the marfe. Less strong marfes have their radiation maximum at its density center.

First measurements of marfe radiation have already been done in ASDEX Upgrade. Fig. 4c shows the picture of a marfe with double structure obtained with a CCD camera in the total light of the plasma. The camera is located at the height of the lower x-point and looks somewhat upwards with respect to the horizontal. This may qualitatively be compared with Fig. 4b which is numerically obtained from the calculated $C^{+1}$ radiation field of Fig. 4a assuming the same position of observation as the camera.

Marfe Dynamics

The time dependent multifluid (D-He—C) calculations show a highly dynamic behaviour of the marfe plasma state. It is found that marfes may move in a complicated way around the lower x-point region upwards along the edge plasma (preferable at the inner, but also at the outer side of the torus), they may deform, even split into separate parts and reconnect again.

Summary

It has been shown by numerical simulations that for certain ranges of the discharge parameters (particularly for high plasma density and impurity concentration) marfes will occur in the edge plasma of ASDEX Upgrade.

The limit for marfe onset has been determined in dependence on the plasma density ($n_e$), the power influx into the edge plasma ($P$) and the safety factor ($q$). Enhancing the plasma density, for sufficiently high impurity concentration, there exists a sharply defined critical value above which the plasma state abruptly reorganizes and a marfe is formed. For fixed impurity concentration, this density increases for increasing input power but is nearly independent of the safety factor although the relative density enhancement of the marfe increases with $q$ (i.e., the marfe becomes more pronounced for high $q$ and flattens for decreasing $q$). Starting, in return, from the marfe plasma state and lowering the density, the plasma returns to the marfe-free state at another critical density which is usually lower than the first one. This hysteresis increases with the energy input and nearly vanishes below about 1 MW. It must be mentioned, however, that marfe formation as a highly nonlinear process also depends crucially on the very details of the model, so at this stage the results have only a qualitative meaning. Further investigations, in particular improvements of the radial transport and inclusion of drift terms are, therefore, in progress.

The time dependent multi-fluid calculations show that marfes are usually not in steady state, but deform and move around the lower x—point region with an intensity which depends on how far the density is beyond the limit for marfe onset. This behaviour is coupled to the motion of the lower impurity charge states which provide effective radiative energy losses and thus cool the plasma.

References

/ 4 / A. Stäbler et al., Nucl. Fusion 32 (1992) 1557
/ 5 / A. Stäbler et al., IAEA Würzburg 1992, paper CN-56/C-2—6
/ 7 / McCormick et al., this conference
SIMILARITY SOLUTION FOR "PLASMA SHIELD" IN HARD DISRUPTIONS

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ABSTRACT

The behaviour of the material ablated from the wall during thermal quench is investigated in terms of a similarity solution approach. An application to an ITER-type device is carried out.

I. INTRODUCTION

Reported damages due to hard disruptions are never apparently so large as the total amount of energy involved would allow them to be (see, e.g., Refs [1] and [2]). In the literature, Ref. [3] has been the first paper to attempt a theoretical explanation of such feature, offering the suggestion that wall damage is ultimately determined by the effectiveness of radiative energy transport mechanisms acting within the ablation material. Such a suggestion was actually pursued further only a decade later [4]. Amidst recent work [5–7], Ref. [7] is in particular notable, in that it specifically includes magnetic field effects in describing the dynamics of the ionized ablation cloud (i.e., of the so-called "plasma shield"). Reference [8] has a similar scope, except that it further includes finite-β effects. Moreover, the latter reference shows that atomic physics and radiation physics features can be modelled in a relatively streamlined fashion without much losing in terms of representativeness. The resulting one-dimensional model system of equations is of such nature as to allow for a similarity solution approach. Then, still within the same paper, the similarity solution equations are further simplified, to the point of eventually allowing the derivation of the final result in closed analytical form. In the present work, we take such procedure actually one step back, and directly address the numerical solution of the system of the similarity equations – thus avoiding the error that has been possibly introduced in Ref. [8] while pursuing further approximations.

II. MODEL EQUATIONS

In Ref. [8] a one-dimensional mathematical model has been derived and justified, which describes resistive magnetohydrodynamical features and radiative energy transport features within the plasma shield. Such model system of equations writes as follows:

\[
\left( \frac{\partial}{\partial t} + \mathbf{V} \cdot \nabla \right) N = 0 , \tag{1}
\]

\[
\left( \frac{\partial}{\partial t} + \mathbf{V} \cdot \nabla \right) \left[ \frac{3}{2} (1 - Z) NkT + Nk W \right] + (1 - Z) NkT \frac{\partial V}{\partial x} + \frac{\partial F}{\partial x} - \eta j^2 = 0 . \tag{2}
\]

\[
F = -1.1 \times 10^{56} \frac{T^{15/2}}{Z^4 Z N^2} \frac{\partial T}{\partial x} \tag{3}
\]
The dependent variables in the above equations are the density $N$ of ions plus neutrals, the temperature $T$, the radiative energy flux $F$, the "frozen-in" macroscopic velocity $V_F$ and the "diffusive" macroscopic velocity $V_D$ (with the cumulative macroscopic velocity $V$ being defined as $V_F + V_D$). The diamagnetic current density $J$ must be calculated from the derivative of the magnetic field $B$, which is in turn defined in terms of $N$ and $T$ by the pressure balance relation:

$$(1 + \bar{Z}) NkT + \frac{B^2}{2\mu_0} = \frac{B^2}{2\mu_0},$$

where $B_o$ is the vacuum magnetic field. The quantities $\bar{Z}$, $\bar{W}$ and $\eta$ must be expressed as functions of $N$ and $T$ by solving the system of Saha's ionization equilibrium equations (see Ref. [8] for details). In the equations MKS units are used - with however temperatures and single particle energies being measured in eV's.

III. SIMILARITY SOLUTION

In Ref. [8] it was shown that by introducing the similarity variables:

$$\zeta = x^2 t^{-1}, \quad \bar{F} = x^{-1} t F,$$

$$\bar{V}_F = x^{-1} t V_F, \quad \bar{V}_D = x^{-1} t V_D,$$

the problem simplifies to a system of ordinary differential equations:

$$\left[ \frac{d}{d\zeta} (2\bar{V}-1) + \frac{\bar{V}}{\zeta} \right] N = 0,$$

$$\left[ \frac{d}{d\zeta} (2\bar{V}_F-1) + \frac{\bar{V}_F}{\zeta} \right] Nk \left[ \frac{3}{2} (1 + \bar{Z}) T + \bar{W} \right] + \left( \frac{d}{d\zeta} + \frac{1}{\zeta} \right) [(1 + \bar{Z}) NkT \bar{V} + \bar{F}] +$$

$$- 2 \bar{V}_F \frac{d}{d\zeta} [(1 + \bar{Z}) NkT] = 0,$$

$$\left[ \frac{d}{d\zeta} (2\bar{V}_F-1) + \frac{\bar{V}_F}{\zeta} \right] B = 0,$$

$$\bar{V}_D = - 2 \frac{\eta}{B^2} \frac{d}{d\zeta} [(1 + \bar{Z}) NkT].$$
\[
\tilde{F} = -2C \frac{T^{15/2}}{Z^4 \tilde{Z} N^2} \frac{dT}{d\tilde{\zeta}},
\]  

(12)

where \( V = V_D + V_F \) and \( C = 1.1 \times 10^{56} \). Thus, the scope targeted with the present paper was directly addressing the numerical solution of the latter system of equations. To this end, through a further change of variables the system of equations has been cast into canonical form (namely, a form such that anyone equation contains only a single differential term) and then a standard, Runge-Kutta type integration procedure has been applied. Boundary conditions have been prescribed in the form of the requirement of a "physically reasonable" behaviour at either end of the ablation matter cloud (where the simultaneous involvement of both ends has of course called for a suitable iterative approach). In any case, among boundary requirements the controlling role is surely played by power balance at the interface between disrupting plasma and ablation cloud.

IV. RESULTS
The obtained profiles of the temperature \( T \), radiative energy flux \( F \), ablation plasma "beta" and ablation plasma velocity \( V \) are shown respectively in Figs. 1 to 4 - for input data that are representative of the ITER device, namely: thermal quench

![Fig. 1 - Temperature versus spatial coordinate](image1)

![Fig. 2 - Radiative flux versus spatial coordinate](image2)

![Fig. 3 - Plasma shield "beta" versus spatial coordinate](image3)

![Fig. 4 - Macroscopic velocity versus spatial coordinate](image4)
duration 100 μs, average heat load during thermal quench 120 GWm⁻², vacuum magnetic field 5 Tesla, wall material graphite. The values plotted are referred to pulse end. The behaviour at earlier times can be inferred from the fact that profiles remain self-similar at all times, while the spatial coordinate changes as t¹/₂ and the magnitudes of F and V change as t⁻¹/₂.

Note the marked change of steepness of the profiles near the wall (except for the "beta", that is). This is in part due to the error that is affecting the model at low temperature. It can be shown, however, that such error does not sensitively affects the conclusions as to the overall plasma shield thickness, and cumulative depth of ablation. Note also that the radiative energy flux has a (shallow) maximum away from the hotter end of the ablation cloud. This can only be explained as the effect of a local energy source - which can indeed be traced, in the model, to the Joule heating produced by diamagnetic currents.

In Table I a number of other results are summarized, and compared with the corresponding values obtained instead from the approximate analytical solution derived in Ref. [8]. The data in the Table are referred to pulse end, and the suffix "0" denotes boundary values on the side facing the disrupting plasma - with x₀ being thus a measure of the plasma shield thickness. Moreover, βₜₚ is the value of the plasma shield "beta" near the wall, and Δ is the produced erosion thickness. Finally, of the total energy that is poured in during the pulse, fᵥ is the fraction that is reflected back into the vessel cavity by blackbody radiation, fₜₚ is the fraction that is transmitted to the solid wall, and fₚₚ is (by difference) the fraction that is deposited inside the plasma shield. The comparison between the two sets of data indeed shows that the analytical approximation - although not greatly accurate - is nevertheless fairly representative of the exact result.

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REFERENCES

Study of the Nonlinear Evolution of the Gyro-Kinetic Plasma in the Guiding Center Approximation

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Recent experimental and theoretical studies of the plasma edge of tokamaks have revealed the importance of turbulence in the physics of the plasma edge. In this context, one of the most important basic set of equations relevant to the physics of the plasma edge of a tokamak is the two-dimensional finite Larmor radius guiding-center equations. A study of the asymptotic state of this system of equations has been presented in Ref. [1]. This system has three "rugged" quadratic invariants. A camonical-ensemble probability distribution characterized by three temperature states was derived [1], and it was shown that this system can have possible negative temperature states leading to an inverse energy cascade, with the equilibrium spectral energy density condensing in the low \( k \) modes. Once this stage has been reached, the energy remains in the lowest modes. For the higher modes less energy is available, so that the level of turbulence is significantly reduced. This inverse cascade manifests itself through the appearance of vortices with scale sizes at the maximum that can be accommodated within the system, and which are important for the transport of energy and particles across the magnetic field. The relevance of these results to the physics of the turbulence at the edge of the plasmas has been recently pointed [2, 3].

In the present work, we extend this model by including the polarization drift, which has a different sign for ions and electrons and thus gives rise to a charge separation in a time varying electric field. The pertinent equations have been recently reported [2, 3]. We consider a straight and constant magnetic field.

\[
\frac{\partial n_e}{\partial t} + (V_D + V_{Pe}) \cdot \nabla n_e + n_e \nabla \cdot V_{Pe} = 0
\]
\[
V_D = -\nabla \times \vec{e}_z / B
\]
\[
\frac{\partial n_i}{\partial t} + (\vec{V}_D + \vec{V}_{Pi}) \cdot \nabla n_i + n_i \nabla \cdot \vec{V}_{Pi} = 0,
\]
\[
\vec{V}_D = g \otimes V_D = -\nabla \phi \times \vec{e}_z / B ; \quad \vec{V}_{Pi} = g \otimes V_{Pi}
\]
\[
\nabla^2 \phi = -\frac{q}{\varepsilon_0} (\vec{n}_i - n_e)
\]
\[
\phi = g \otimes \phi, \quad \vec{n}_i = g \otimes n_i
\]
\[
V_{Pi} = \frac{1}{\Omega_i^2} \frac{q}{m_i} \frac{dE}{dt} \quad V_{Pe} = -\frac{1}{\Omega_e^2} \frac{q}{m_e} \frac{dE}{dt}
\]

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As described in [1], $g$ is a convolution integral operator, so that we have, for instance

$$\bar{n}_i = g \otimes n_i = \int g(\vec{r} - \vec{r}') n_i(\vec{r}') \, dxdy$$

(9)

In the Fourier $k$-space, it becomes a filter operation, which is easily incorporated in the different Fourier modes. Each coefficient $a_k$ of the mode $e^{i k \cdot \vec{r}}$ is multiplied by a factor $g = \exp(-\frac{i}{2} k^2 r_i^2)$, where $k$ is total wave vector and $r_i$, is the ion Larmor radius. We emphasize that the densities are not particle but guiding center densities and this explains the absence of any pressure gradient terms in the present model.

The numerical code used to solve equations (1-9) uses a method of fractional steps and has been presented in Ref. [3]. These equations are solved in a slab geometry periodic in $x$, and finite in the $y$ direction. In dimensionless units, we have the following steps:

1) Solve for $\Delta t/2$:

$$\frac{d\bar{n}_e}{dt} = \frac{\partial \bar{n}_e}{\partial t} + V_D \cdot \nabla \bar{E}$$

(8)

$$\frac{\partial \bar{n}_e}{\partial t} + \left( E_y - \beta_e \frac{\partial E_x}{\partial t} - \beta_e E_y \frac{\partial E_x}{\partial x} + \beta_e E_x \frac{\partial E_y}{\partial y} \right) \frac{\partial \bar{n}_e}{\partial x} = 0$$

(10)

(with a similar equation for the ions, $\beta_e = m_e/m_i$).

The solution of Eq.(10) is calculated by the shift

$$n_e^{n+\frac{1}{2}} = n_e^n \left( y, x - \left( E_y^n - \beta_e \frac{\partial E_x^n}{\partial t} - \beta_e E_y^n \frac{\partial E_x^n}{\partial x} + \beta_e E_x^n \frac{\partial E_y^n}{\partial y} \right) \frac{\Delta t}{2} \right)$$

(11)

(The shift in Eq. (11) is effected using a cubic spline interpolation.)

Then solve for $\Delta t/2$

$$\frac{\partial \bar{n}_e}{\partial t} + \left( -E_x - \beta_e \frac{\partial E_y}{\partial t} - \beta_e E_y \frac{\partial E_x}{\partial x} + \beta_e E_x \frac{\partial E_y}{\partial y} \right) \frac{\partial \bar{n}_e}{\partial y} = 0$$

(12)

(with a similar equation for the ions)

The solution of Eq. (12) is calculated by the shift

$$n_e^{n+\frac{1}{2}} = n_e^n + \frac{1}{2} \left( x, y + \left( E_x^n + \beta_e \frac{\partial E_y^n}{\partial t} + \beta_e E_y^n \frac{\partial E_x^n}{\partial x} - \beta_e E_x^n \frac{\partial E_y^n}{\partial y} \right) \frac{\Delta t}{2} \right)$$

(13)

2) Solve Poisson equation to update the electric field. $\bar{E} = -\nabla \phi$.

Solve for $\Delta t$ the equation:

$$\frac{1}{n_e} \frac{\partial n_e}{\partial t} = \beta_e \left( \frac{\partial E_x}{\partial x} - \frac{\partial E_y}{\partial y} \right) + \beta_e \left( E_y \frac{\partial E_x}{\partial x^2} - E_x \frac{\partial E_y}{\partial y^2} \right)$$

(14)

(with a similar equation for the ions). The solution of Eq.(14) is given by:

$$n_e^+ = n_e^- e^{+\beta_e} \left[ \frac{\partial E_x^+}{\partial x} - \frac{\partial E_x^-}{\partial x} + \frac{\partial E_y^+}{\partial y} - \frac{\partial E_y^-}{\partial y} \right] + \beta e \Delta t \left( E_y \frac{\partial E_x^+}{\partial x^2} - E_x \frac{\partial E_y^+}{\partial y^2} \right)$$

(15)
The subscripts + and - denote the values at the present and previous time steps respectively.

3) Repeat Step 1

The $y$ length of the box is 17 (in units of $v_{thi}/\omega_{ci}$, where $v_{thi}$ is the ion thermal velocity and $\omega_{ci}$ is the ion cyclotron velocity). The number of points in the $x$ direction is 64 and in the $y$ direction is 128. The initial profile is of the form

$$n_e = n_i = \frac{1}{2} \left( 1 + \tanh \left( \frac{y}{2} \right) \right)$$

Electron density profile as function of $y$ at $t = 0$ and $t = 500 \, \omega_{ci}^{-1}$ is shown in Figs. 1 and 2. Figs. 3 and 4 show contour plots of the equipotential at $t = 700 \, \omega_{ci}^{-1}$ and $t = 800 \, \omega_{ci}^{-1}$. They show a tendency for a sheared flow to form on which is superimposed a vorticity structure, corresponding at $t = 800 \, \omega_{ci}^{-1}$ to the lowest mode in the periodic $x$ direction. The drift velocity field is shown in Fig. 5 at $t = 800 \, \omega_{ci}^{-1}$ and Fig. 6 at $t = 1300 \, \omega_{ci}^{-1}$. The shear flow on which vorticities are superimposed is apparent in Fig. 5, the vorticity appearing in the region where the flow is changing sign. The dominance of the sheared flow is apparent in Fig. 6 at $t = 1300 \, \omega_{ci}^{-1}$ with a reduced vorticity structure appearing in a reduced layer in the transition where the velocity field is changing direction. A study of the ratio of energy to enstrophy shows this ratio is asymptotically close to 10. A recent study [4] for a 2-dimensional set of equations close to what we present in Eqs. (1 - 3), but without the ion gyro-radius effect, predicts a ratio of $\pi^2$.

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References

$t = 700\omega_{ci}^{-1}$  

![Fig. 3](image1.png)  

$t = 800\omega_{ci}^{-1}$  

![Fig. 4](image2.png)  

$t = 800\omega_{ci}^{-1}$  

![Fig. 5](image3.png)  

$t = 1300\omega_{ci}^{-1}$  

![Fig. 6](image4.png)
Bifurcation of Electron Temperature in the High Recycling Regime

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Introduction. Thermal instability in the boundary region of tokamaks has been recently studied numerically by Capes et al.[1], who considered the one-dimensional thermal conduction between the mid-plane of a tokamak and the divertor target plates. Multiple solutions and bifurcation of the electron temperature are caused by impurity radiation and its non-linear dependence on the temperature. In this paper it will be shown that the non-linearity of heat flux in the high-recycling regime in front of the divertor target plates can also trigger a bifurcation even in absence of impurity radiation. In this paper we consider a magnetic flux tube which is bounded by target plates on two sides. Such a case exists in tokamaks and stellarators in the scrape-off layer outside the last magnetic surface. Here, the high recycling regime is of particular interest since it allows to keep the temperature on the target plates at a low level and to minimize sputtering and other damaging effects on the target plates. To achieve this state, the density in front of the target plates must be sufficiently high so that ionisation of recycling neutrals occurs in a short distance from the target plates. Thus, the heat flow into the high recycling layer must cover the ionisation and radiation losses and the thermal flux onto the wall. Lackner et al.[2] have modelled this effect and proposed the following ansatz for the heat flow into the high recycling layer, or — since this layer is considered as very small — into the target plate.

$$q_\parallel = (\delta_t \cdot kT_t + R \varepsilon_{\text{eff}})n_t \sqrt{\frac{kT_t}{m_i}(\gamma_e + \gamma_i)}$$  \hspace{1cm} (1)

$\delta_t$, $\gamma_e$ and $\gamma_i$ are constants, $n_t$, $T_t$ density and temperature at the target plate and $\varepsilon_{\text{eff}}$ the ionisation energy enhanced by the radiated energy. $R$ is the recycling coefficient.

On the plasma side of the high recycling layer these losses must be covered by the heat conduction into this layer. One of the standard approximations of the plasma boundary is constant pressure along the flux tubes $p_0 = 2n_t kT_t = n_e kT_e$; $n_e$, $T_e$ are the plasma parameters at the stagnation point where the parallel flow velocity is zero. Introducing the pressure as a parameter on the flux tube yields the boundary condition

$$\frac{(\delta_t \cdot kT + R \varepsilon_{\text{eff}})}{2kT} \frac{p_0}{m_i}(\gamma_e + \gamma_i) = -n_\parallel \frac{B}{B^2} \cdot B \cdot \nabla T \quad ; \quad T = T_b$$  \hspace{1cm} (2)
term in eq.(1) makes the heat flux $g(T)$ non-monotonic in $T$. With $R \to 0$ the threshold in heating power also shrinks to zero.

**Numerical results.** After introducing the dimensionless variable $X = T/T_{eff}$ and $kT_{eff} = R\varepsilon_{eff}/\delta_t$ eqs.(5) have been solved numerically. The source term is modelled by $h(y) = Q_0y/(0.01 + (y - y_0)^2)$. $Q_0$ is proportional to the total power input into the flux tube and $y_0$ describes the location of the maximum power input. Further dimensionless constants are: $C_a = [7C_1/2\kappa_aB(a)T_{eff}^3]^{1/7}$, $C_b = [7C_1/2\kappa_aB(b)T_{eff}^3]^{1/7}$, $D_a = F'\kappa_aB(a)/C_1\sqrt{T_{eff}}$, $D_b = F'\kappa_aB(b)/C_1\sqrt{T_{eff}}$. The constant $C_1$ is proportional to the pressure in the flux tube $C_1 = \delta p_0/2(\gamma_e + \gamma_t/m_i)$. $F'(a)$ and $F'(b)$ depend on the source function $h(y)$, both are proportional to the total power input $Q_{tot}$ or $Q_0$. In these dimensionless variables the boundary conditions are

$$\left(\frac{X_a}{C_a}\right)^{7/2} = \left(\frac{X_a}{C_a}\right)^{7/2} + \sqrt{X_a} + \frac{1}{\sqrt{X_a}} - D_a$$

$$\left(\frac{X_b}{C_b}\right)^{7/2} = \left(\frac{X_b}{C_b}\right)^{7/2} + \sqrt{X_b} + \frac{1}{\sqrt{X_b}} + D_b$$  \hspace{1cm} (7)

The eqs.(7) are two coupled polynomials of $8^{th}$ order, in general they have more than one solution. In case of symmetry $B(a) = B(b), C_a = C_b, D_a = D_b$ two solutions are on the symmetry line $X_a = X_b$, however, there are also two solutions off the symmetry line. Fig.1 shows these 4 solutions depending on the control parameter $Q_0$.

![Figure 1: Bifurcation of the temperature on the target plates. Solution of eqs.(7) with $C_a = C_b = 4.0$ and $kT_{eff} = 10$ eV. The x-axis is $Q_0 \propto$ heating power. $X = T/T_{eff}$ is the norm. temperature. The curves represent the temperature on one of the target plates.](image-url)

Below a threshold in $Q_0$ there are no solutions of eqs.(7); above the threshold a small region exists with two solutions followed by a region of four solutions. In this region the
where the parallel thermal conductivity is \( n \chi_\parallel = \kappa_0 T^{8/2} \). The essential feature of this boundary condition is its non-linearity and the non-monotonic dependence on the temperature. This will give rise to bifurcation and multiple solutions as ref.[1] by volume radiation.

The heat conduction equation. In our model we consider the flux tube between two target plates outside the high recycling layer. We neglect convective energy transport and perpendicular thermal conduction. Therefore, the temperature is the solution of an one-dimensional conduction equation along the flux tube. Let be \( h(x) \) a source term which describes the power input into the flux tube; \( x \) is the length coordinate along the flux tube. The temperature is the solution of the following equation \(-\nabla \cdot n \chi_\parallel \vec{B}/B^2 \vec{B} \cdot \nabla T = h(x)\), which by introducing the magnetic potential \( y(\text{dy} = B \text{dx}) \) as a new independent variable and \( f = 2/7T^{7/2} \) instead of the temperature can be transformed to

\[
- \frac{d^2 f}{dy^2} = \frac{h}{\kappa_0 B^2}. \tag{3}
\]

This transformation eliminates the geometry of the flux tube and leads to a standard Poisson equation which can be solved analytically. The solution which satisfies the boundary condition \( F(a) = F(b) = 0 \) is

\[
F(y) = \int_a^b G(y, x) \frac{h}{\kappa_0 B^2} dx; \quad G(y, x) = \begin{cases} 
(y - a)(b - x) & : y \leq x \leq b \\
(b - a) & : a \leq x \leq y 
\end{cases} \tag{4}
\]

The general solution of eq.(3) is \( f(y) = F(y) + A_1 + A_2(y - a)/(b - a) \), where \( A_1 \) and \( A_2 \) are correlated to the boundary values of \( f \) by \( A_1 = f_a; A_2 = f_b - f_a \). Inserting this ansatz into the boundary conditions (2) yields

\[
f_b = f_a + \frac{g(T(f_a))}{\kappa_0 B(a)} - F'(a); \quad f_a = f_b + \frac{g(T(f_b))}{\kappa_0 B(b)} + F'(b) \tag{5}
\]

where the function \( g(T) \) is the left hand side of eq.(2) and \( T(f) = (7/2f)^{2/7} \). These non-linear equations determine the boundary values \( f_a \) and \( f_b \). Having found \( f_a, f_b \) the temperature profile is calculated using \( F(y) \). With fixed values of \( \varepsilon_{\text{eff}}, \delta_t \) and \( \gamma_0, \gamma_1 \) there are three control parameters which determine the solution of eq.(5). These are the total input power \( Q_0 \) which yields a factor to \( F' \), the pressure \( p_0 \) and the factor \( \kappa_0 \) in the thermal conductivity. The symmetry of the boundary conditions is disturbed by the magnetic field, which in general is different on both ends of the flux tube. Furthermore, the constants in the function \( g \) may also differ on both sides. Adding the two equations in (5) yields

\[
\frac{g(T_a)}{B(a)} + \frac{g(T_b)}{B(b)} = \int_a^b \frac{h(y)}{B^2} dy. \tag{6}
\]

The left hand side of this equation is positive and has a positive minimum since the function \( g(T) \) has a positive minimum at \( T = T_{\text{eff}} \). As a consequence the integrated source term on the right hand side must be larger than a threshold value, otherwise a solution does not exist. The reason is the ionisation in the recycling layer. The second
temperatures at the boundaries $y = a$ and $y = b$ are: case 1) $X_a = T_1$ and $X_b = T_1$; case 2) $X_a = T_2$ and $X_b = T_2$; case 3) $X_a = T_3$ and $X_b = T_4$; case 4) $X_a = T_4$ and $X_b = T_3$. The four temperature profiles are shown in Fig. 2a. A remarkable result is the existence of two asymmetric solutions $X_a \neq X_b$ although the boundary conditions and the source term are symmetric to the midplane of the flux tube. As expected, asymmetric boundary conditions lead to asymmetric solutions (Fig. 2b). These asymmetries arise from the magnetic field which is asymmetric along the flux tube or from the asymmetry of the source function $h(y)$.

![Figure 2: Temperature profiles along the flux tube. Fig. 2a (left): Symmetric case $C_a = C_b = 4.05, Q_o = 2500, y_o = 0.5$. Fig. 2b (right): Asymmetric power input $C_a = C_b = 4.05, Q_o = 2500, y_o = 0.6, a = 0, b = 1$.](image)

Summary and conclusions. The model of a high recycling layer in front of two target plates at the ends of a flux tube leads to non-linear boundary conditions of the heat conduction equation. The heat conduction equation can be transformed to a linear Poisson equation and be solved analytically. In general four solutions exist leading to four temperature profiles. Asymmetric solutions with different temperatures on the target plates ($T_a \neq T_b$) can exist although all boundary conditions are symmetric to the midplane of the flux tube. The thermal stability of the solutions has not yet been considered. The bifurcation point of the multiple solutions depends on the power input and the pressure $p_o$ on the flux tube — as in ref. [1]. However, the bifurcation also depends on the magnetic field on the target plates ($B(a) \neq B(b)$) and the asymmetry of the power input.

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References.
EFFECT OF NEUTRAL PARTICLES ON DENSITY LIMITS
IN TOKAMAKS

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1 Introduction

The global stability and confinement of a tokamak plasma are significantly influenced by the boundary plasma parameters. The onset of density disruptions, which limit the maximum plasma density, is triggered by impurity radiation in the edge plasma and can be connected with the radiative thermal instability [1]. At the density $n_c$, the total radiative power $P_{rad}$ is equal to the total input power $P_{in}$ into the plasma ($S := P_{rad}/P_{in} = 1$). Above $n_c$ ($S > 1$) no steady state of the plasma column exists. Contrary to predictions made in [1], where neutral particle kinetics is not taken into consideration, experimental results show that disruptions can occur for $S < 1$ [2], [3]. It was shown in [4] that the carbon impurity radiation cooling of the plasma is strongly affected by charge exchange between carbon ions and hydrogen atoms which penetrate into the plasma as a result of recycling or gas puffing.

We present analytical expressions for the cooling rate $Q_R$ as a function of the plasma temperature $T$, $\xi_N := N/n$ and $\xi_i := n_i/n$, where $N, n, n_i, n_P$ are the densities of hydrogen atoms, impurity ions and the plasma, respectively. We investigate the influence of the neutral particles on the critical densities and the stability of the system, taking into account ionization, charge exchange and impurity cooling.

2 Model Equations

Averaged over the magnetic field lines in the SOL, the dynamics of the plasma and neutral particles perpendicular to the magnetic field is described by the following set of coupled equations:

$$\frac{\partial n}{\partial t} + \frac{\partial \Gamma_n}{\partial x} = k_i(T)nN - \theta n/\tau_1(T),$$  (1)

$$\frac{\partial N}{\partial t} + \frac{\partial \Gamma_N}{\partial x} = -k_i(T)nN - \theta R n/\tau_1(T),$$  (2)

$$\frac{3}{2}\frac{n\partial T}{\partial t} + \frac{\partial \Gamma_T}{\partial x} = \sigma(T)E^2 - Q_R(T, N, n) - \theta n T/\tau_2(T) - \theta \kappa_{||}(T) T/L^2,$$  (3)

$$\Gamma_n = -D_\perp(n, T) \partial n/\partial x,$$  (4)

$$\Gamma_N = -(1/2k_{ex})\partial (v^2 N)/\partial x,$$  (5)

$$\Gamma_T = -\kappa_{\perp}(n, T) \partial T/\partial x.$$  (6)
\( D_\perp \) and \( \kappa_\parallel \) are the coefficients of the perpendicular anomalous plasma diffusion and heat conduction, and \( \kappa_\parallel \sim T^{3/2} \) and \( \sigma \sim T^{3/2} \) are the classical heat and electric conduction coefficients, resp.; \( \tau_m := \beta_m L/v_s \) is the lifetime of the plasma \((m = 1)\) in the SOL and its energy \((m = 2)\) due to streaming along the field lines to the limiter \((\beta_m > 1, L = \pi q R \text{ - connection length})\), and \( v_s = \sqrt{2T/m} \) is the ion sound velocity; \( k_i \) and \( k_{\text{ex}} \) are the rate coefficients for ionization and charge exchange, resp., and \( R \) is the recycling coefficient at the limiter. \( x := r - r_s, r_s \) is the separatrix radius, \( \theta = \theta(x) \) the Heaviside function, and \( E \) the electric field of the tokamak.

The analytical expression for the cooling rate \( Q_R = n^2 \xi G(T, \xi_N) \) \([\text{W/cm}^2]\) (see Fig. 1) in the temperature range \(5eV \leq T \leq 50eV\) with

\[
G = A(T - T_1)e^{-B(T - T_1)},
\]

\[
A = 10^{-28}(3.1e^{-0.024(\ln \xi_N + 11)^2} + 0.7), \quad B = 0.6e^{-7.810^{-4}(\ln \xi_N + 14)^3} + 0.055
\]

\((T \text{ in } eV, n \text{ in } \text{cm}^{-3}, T_1 = 4.4eV)\) may be simplified in two steps:

\[
A = 0.78 \times 10^{-26} \xi_N^{-0.18}, \quad B = 0.041 \xi_N^{-0.31}, \quad \xi_N > 10^{-4},
\]

(see [5]) and according to [1]

\[
G = G_0 \theta(T_L - T),
\]

\[
L_0 = \frac{A \left(1 + B(T_W - T_1)\right) e^{-B(T_W - T_1)}}{T_L}, \quad T_L = T_1 + 2/B,
\]

where \( T_W \) is the wall temperature.

### 3 Equilibrium and Critical Densities

The system of eqs. (1)-(6) was solved under the steady-state condition \( \partial / \partial t = 0 \) for \( x \in [x_c, x_w] \) \((x_c \text{ - position in the central plasma, } x_w \text{ - distance from the separatrix to the wall})\). For the fluxes \( \Gamma_A \) with \( A = n, N, T \) it is of the following form:

\[
\frac{d\Gamma_A}{dx} = H_A - L_A
\]

with the corresponding source \((H_A)\) and loss \((L_A)\) terms. An equilibrium state exists if the condition

\[
I_A(x) := \Gamma_A(x) + 2 \int_{x_c}^x dx' \Gamma_A(H_A - L_A) \geq 0 \quad \forall \; x \in [x_c, x_w]
\]

is fulfilled. Otherwise critical parameters may exist. For \( x = x_w \) the general 6-parameter solution of our problem can be reduced down to a 3-parameter one.

For the case of complete recycling at the limiter \((R = 1, \Gamma + \Gamma_N = 0)\) we obtain with \( D_\perp = D_0 n^p \)

\[
N = \frac{1}{v_s^2} \left(C - \frac{2k_{\text{ex}} D_0 n^{2+p}}{2 + p}\right), \quad C = \text{const.}
\]

For \( p = 0, x_c \to -\infty, x_w \to \infty, n(x_c) = n_0, N(x_c) = 0, n(x_W) = 0, N(x_w) = N_w \) and averaged temperature it follows (cf. [6]) that

\[
N_w \leq N_w^* = 1/k_i r_1, \quad n_0^2 = N_w v_s^2/k_{\text{ex}} D_0.
\]
Next we investigate the energy balance equation with averaged densities, assuming the radiation term as the dominant loss term in the approximation (10), (11), and consider the two cases of (i) a transparent \( (l_N > l_L) \) and (ii) a non-transparent \( (l_N < l_L) \) SOL for the neutrals. 

\[ l_N = \frac{v_N}{\sqrt{2} \sqrt{k_B T}} \] \( k_B \) is the mean free path of the neutrals and \( l_L \) the width of the radiation function. The condition \( I_T(x_w) = 0 \) reads 

\[ \Gamma_T^N(x_c) = 2 \kappa_1 n_\xi_G (T_L - T_w, \text{ } l_N > l_L) \]

\[ \Gamma_T^N(x_c) = T_N - T_w, \text{ } l_N < l_L \]

with \( T_L \) as given in eq.(11) and \( T_N = T_w + \Gamma_T(x_c)/2 \kappa_1 \), which is equivalent to the above-mentioned condition \( S = 1 \) \( (S \equiv Q_R \delta / \Gamma_T(x_c), \delta = N, L) \). Using eq.(14), we obtain the following scaling relations \( (D_0 \sim q^p, B(T_w - T_1) \ll 1, A, B \text{ given in } (9)): \)

\[ n_c \sim \left( \frac{P_{in}}{q^N \xi_i} \right)^{1/(2.56 + p)} \text{ } (i), \quad n_c \sim \left( \frac{P_{in}}{\xi_i q^{0.13}} \right)^{1/(1+0.13(1+p))} \text{ } (ii). \]

An example of a self-consistent solution of the problem is displayed in Fig. 2. The radiation density localized in the SOL strongly depends on the profile functions.

4 Radiative Thermal Instability

The stability of the equilibrium state with respect to poloidally and toroidally homogeneous perturbations is investigated by assuming all quantities \( \vec{A} \) to be of the form \( \vec{A}(x, t) = \vec{A}_0(x) + \delta \vec{A}(x) \exp(\gamma t) \). \( \vec{A}_0 \) is the equilibrium solution (section 3) and \( \delta \vec{A} \) is determined by a 6-dimensional linear system of equations with coefficients depending on \( \vec{A}_0 \). A numerical solution to the problem is in preparation. Here we give a preliminary analytical estimate for case (i) considering only the neutral gas density variation in eq. (3) which leads to

\[ \kappa_1 \frac{\partial^2 \delta T}{\partial x^2} - \frac{3}{2} n_\gamma \delta T = \frac{\partial Q_R}{\partial T} \delta T + \frac{\partial Q_R}{\partial N} \delta N. \]

Using approximation (10), this results in the dispersion relation

\[ k_\perp + k_\parallel c t g \kappa_\perp = S(1 + 0.31 \eta_N), \quad k_\perp^2 = \frac{3 n \gamma T}{2 \kappa_\perp}, \quad k_\parallel^2 = \frac{0.26 S \eta_N}{2 - S} - k_\perp^2, \quad \eta_N = \langle T \frac{\partial T}{\partial N} \rangle l_L. \]

The threshold for this thermal instability \( S_c \) \( (\gamma = 0) \) is shown in Fig. 3. The essential result is that neutral particles can cause the stable equilibrium state to be unstable to thermal perturbations.

Fig. 1: Impurity radiation for different values $\xi_n$: 1 - $\xi_n = 10^4$, 2 - $\xi_n = 10^3$, 3 - $\xi_n = 10^2$, 4 - $\xi_n = 10^1$

Fig. 2: Profiles of 1 - particle flux [$10^{18}$ s$^{-1}$cm$^{-2}$], 2 - energy flux [Wcm$^{-2}$], 3 - temperature [450 eV], 4 - plasma density [$8.5 \cdot 10^{19}$ cm$^{-3}$], 5 - neutral density [$10^{10}$ cm$^{-3}$], 6 - radiation density [$10^4$ Wcm$^{-2}$], ($x_w = 15$ cm, $R = 1$, $\xi = 10^2$, $D_w = 10^8$ cm$^{-2}$s$^{-1}$, $L = 2 \cdot 10^9$ cm)

Fig. 3: Radiation density threshold as a function of $n_w$.
LINEAR THEORY OF ION VISCOSITY EFFECTS ON EDGE TEARING MODES.

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I. Introduction
The inviscid drift tearing modes dispersion relation is a singular limit of the viscous theory. In analogy with the theory of the Orr-Sommerfeld equation for viscous fluids, it is shown that linear magnetic field reconnection perturbations may be destabilized by ion viscosity, at large drift frequency.

In the regime of operation of the largest tokamaks with ion and electron temperature of several keV it is expected that finite ion temperature and pressure anisotropy modify the dynamics of the tearing modes of helicity numbers \((m,n)\) introducing both a real part of the mode frequency, of the order of the electron diamagnetic frequency, \(\omega^* = mT_e^*/e\nu_b^2\) and a kinematic viscosity \(\nu^* \) [1] which tends to damp the fluid vorticity field. In the outer regions of a tokamak plasma, viscosity may compete with collisional resistivity, which is responsible for magnetic field lines reconnection. A classical result of the analytic theory of drift tearing modes is a reduction of the linear growth rate at high drift frequency [1] while viscous effects have been considered negligible in the parameter range prevailing at the time. A discrepancy between the analytical results and numerical calculations, which did not show a decrease of the growth rate at large \(\omega^*\), was found to be due to the effects of formation of standing drift waves in a finite size plasma.

We extend the analytic linear theory of drift tearing modes to investigate whether similar effects can be due to finite ion viscosity, in analogy with the expectations of Reynolds theorem in viscous hydrodynamics.

II. Drift tearing equations
We start from the customary reduced MHD equations, valid for large aspect ratio \(e = a/R < 1\) and for the tokamak ordering \(B_\theta / B_\phi = O(e), \beta = O(e^2), \text{div} v = O(e^3)\). A cylindrical helical coordinate system \((r,\chi,z)\) is used, related to the cylindrical system \((r,\theta,z)\) by \(\chi = (\theta-kz)\), with \(k = m/R\), \(g = (1+(\nu kr)^2)^{-1/2}\) and

\[
\begin{align*}
\eta &= -e_r, & \xi_r &= g(e_{\chi_r} - kr e_z), & \xi_z &= g(e_{\chi_z} + kr e_\phi) \\
\end{align*}
\]

To leading order in \(kr\) the relevant RMHD equations in S.I. units are:

\[
\frac{\partial \psi}{\partial t} + (\nu + \nu^*) \cdot \nabla \psi = \frac{n_i}{\rho_s} \left( \nabla^2 \psi - 2k \xi_z \right) \tag{1}
\]

\[
\frac{\partial \xi_z}{\partial t} + (\nu + \nu^*) \cdot \nabla \xi_z = \frac{1}{nm_i} \left[ B \cdot \nabla \chi_z \right] + \frac{\mu^*}{nm_i} \nabla^2 \xi_z \tag{2}
\]

where \(\xi_z = \xi_z (\nabla \cdot \psi)\) and the \(B\) and \(\psi\) vector fields are expressed in terms of the scalar helical flux function \(\psi(r,\chi,z)\) and velocity stream function \(\phi(r,\chi,z)\) as \(B = \nabla \psi \times \chi\xi + B\chi\xi\) and \(\psi = \nabla \phi \times \chi\xi\). In eqs (1,2) the electron and ion responses generate the drift velocities \(\nu^* = \nu^* \chi \nabla \phi \times \chi \xi\) and the scalar pressure fulfills \(dp_0/\partial t = 0\). \(\psi^* = \psi^* e \chi \xi^*\) and the viscous term in eq (2) \(\nu^*/n_i m_i\) is of the order of \(0.3 p^2 / \nu^2\). These are the leading order \(p_i\) contributions assumed to be sufficient to describe the finite ion temperature effects addressed in this work.
The system (1-2) is linearized around an equilibrium \( q=q(r), \quad \psi=\psi_0(r), \quad p=p_0(r), \quad \phi_0=0, \quad n_0=\text{const.} \), considering single helicity perturbations of \( \psi_1, \phi_1, p_1 \) of the type \( f_1=\text{Re}\{f_1(r)e^{-im\gamma t+im\chi t}\} \). It is assumed that the external boundary value problem in the ideal MHD region is solved and provides the value \( \Lambda' \) of the jump of \( \ln \psi_1/\ln r \) across \( r=r_\text{s} \) to be matched to the analogous quantity calculated in the inner layer.

The linearized equations, in the thin layer approximation \( lr_\text{s}/a<<1 \) are:

\[
\begin{align*}
i(\omega-\omega^*)\left(\psi_1 + (F_0/\omega)x\phi_1\right) = & -i(\eta/\mu_0)x\psi_1'' \quad (3)\\
\mu\phi_1'''+i(\omega-\omega^*)\phi_1'' = & i\left[(F_0/\rho_0\mu_0)x\psi_1'' + (J_{\phi}/\rho_0)\psi_1\right] \quad (4)
\end{align*}
\]

where the following definitions are used:

\[
F_0 = k'_{\phi}B_{00}(r_\text{s}) = nq'(r_\text{s})B_{00}(r_\text{s})/\tau_\text{s} ; \quad q(r_\text{s}) = m/n ; \quad x = r - r_\text{s} ; \quad \mu = \mu_0/\rho_0 ; \quad \tau_\text{s} = \mu_0a^2/\eta
\]

and \( v_\text{A} = B_00/(m\rho_0)^{1/2}, \tau_\text{A} = a/v_\text{A}, \tau_\text{V} = a^2/\mu, \rho_0 = n_0m_1 \). It is clear from eqs. (3,4) that a vanishing viscosity limit is singular because the differential equation for the velocity stream function changes by two orders. Viscous effects can compete with resistive ones especially in the outer regions of a plasma with \( T_e > T_i \) since the ratio of resistive and viscous diffusion times \( \tau_\text{R}/\tau_\text{V} \) increases as \( T_e^{3/2}/T_i^{1/2} \).

### III. Viscous drift tearing equation.

We set \( \Omega = \omega + \omega^* \), \( \delta^* = (\omega-\omega^*)/\omega \), \( \gamma = \text{Im} \omega \), and introduce the relevant scale lengths: the inertial layer width, \( x_\text{A} = \omega/(k'/v_\text{A}) \), and the resistive and viscous layer widths \( x_\text{R} = \eta/\mu_0|\omega|, x_\text{V} = \mu/|\omega| \). Equations (3,4) with the constant \( \psi_1-\psi_\text{s} \) approximation give rise to the Orr-Sommerfeld type equation [3] for the stream function \( \psi_1 \):

\[
(\omega\delta^*/F_0)x\psi_1 = -\delta^* x^2 \phi_1 + i(\Omega/\omega)x^2 x_\text{R} \phi_1'' + x^2 x_\text{R}^2 x_\text{R} \phi_1'''' (5)
\]

The formal treatment of the problem reveals a viscous effect consisting in the coupling of different branches, thereby affecting the stability regime.

The inner \( \Lambda' \) to be matched with a given exterior \( \Lambda' \) is defined as:

\[
\Lambda' = \lim_{c \to 0} \int_C \psi_1'(x)dx = -i\lambda \delta^* x^2 \int_C [1 - zY(z)]dz \quad (6)
\]

where \( C \) is a suitable contour in the complex \( z \) plane, and \( 1, z, Y(z) \) are defined as:

\[
\lambda' = -ix^2 x_\text{R}^2 (\Omega/\omega\delta^*) = -i\omega(\omega + \omega^*)(\omega - \omega^*)^{-1}(\eta\rho_0/F_0^3) \\
z = x/\lambda ; \quad Y(z) = (\omega\lambda/k'_{\phi}B_{00})\phi_1(z) \quad and \quad v = \mu(\delta^* e^{i3\pi/2}/\Omega^3)^{1/2}(F_0^2/\eta\rho_0)^{1/2}
\]

In dimensionless form eq.(5) is \( vY''''-Y''+z^2Y = z \). Since \( v \) multiplies the highest derivative a boundary layer technique can be used, introducing a stretched "inner" variable, \( \zeta(z,\nu) = v^{-1/2}z \). An "outer" solution in \( z \) is constructed as an expansion in \( v^{1/2} \):

\( y_\text{out}(z) = y_0(z) + \Sigma_k v^{k/2}y_k(z) \) where \( y_0(z) \) is given by the usual second order equation obtained by setting \( v=0 \) and the leading term of the expansion obeys a parabolic cylinder equation.

An "inner" solution in \( \zeta \) is also constructed as an asymptotic expansion \( Y_\text{in}(\zeta) = \Sigma_k v^{k/2}Y_k(\zeta) \) where the \( Y_k(\zeta) \) obey a fourth order equation. The four free constants of the first non zero term \( Y_1(\zeta) = A\exp(-\zeta) + B\zeta + C + D\exp(\zeta) \) are fixed by requiring non divergent or secular behaviour at large \( \zeta \), and vanishing at \( \zeta=0 \). Finally the constant \( A \) is determined by the Prandtl
asymptotic matching with the outer solution up to order $v^{1/2}$. A uniform solution with the correct asymptotic behaviour for large and small $|z|$ is then constructed:

$$y(z) = \frac{z^2}{\pi} \left( \sin \theta \right)^{1/2} d\theta + v^{1/2} \frac{z^{1/4}}{\sqrt{\pi}} \left[ z^{1/2} K_{1/4}(z^{1/2}) - \frac{1}{\beta r} \Gamma(1/4) e^{-\beta r} \right]$$

(7)

where $K_{1/4}$ is a modified Bessel function.

From the definition of $A'$ given by eq.(6), the dispersion relation is obtained:

$$\omega(\omega + \omega_c) \left( \omega - \omega_c \right)^3 \left[ 1 + v^{1/2} \frac{\ln a}{\pi^{1/2} \sigma} \left( \frac{T_0}{\eta \rho_0} \right)^{1/4} (\omega - \omega_c)^{-1/4} (\omega + \omega_c)^{-3/4} \right] = i\gamma^T$$

(8)

where on the right hand side there is the classical tearing mode growth rate.

The results of the numerical solution of this dispersion relation are shown in Figs.1-3 for the viscosity coefficient $\mu=0$ and for $\mu>0$. In the first case (Fig.1) we reobtain the results of Ref. [1]. With parabolic profiles of $q(r)$ and temperature profiles of the type $T_\alpha(r) = T_0 \alpha[(1-(r/a)^2)^{1/2}]$, with $\alpha = 0.5$, $\beta = 0.25$, $T_0 = 1$ keV, $T_{10} = 1$ keV, $n_s = 3.10^{19}$ m$^{-3}$, $B = 2.2$ T, $q(a) = 4.17$, and with $\mu > 0$ the previous root of the dispersion experiences a reduction of the growth rate $\gamma$ at low drift frequencies while a new root appears with a low growth rate for $\omega^*/\gamma_T < 1$ and $\gamma^*/\gamma_T$ for $\omega^*/\gamma_T > 1$, where the old root tends to be stable (Fig.2). By reducing $T_{10}/T_0$, thereby increasing the viscous effects relative to the resistive ones, the new and old branches couple into a single branch (Fig.3) which has no decreasing trend for $\omega^*/\gamma_T > 1$ and has a minimum at the crossing point of two branches, just as obtained in Ref [1] from numerical integration.

In conclusion for high $T_e$ different from $T_i$ the viscous dissipative coupling of dispersion branches may favor their growth into the nonlinear stage where they are expected to become stable. For $m=4$, $n=1$, $q=4.17$ (Fig.3) a large viscous destabilization with $\gamma^* - 0.7 \gamma_T$ can occur; also the reduction of the rotation frequency observed is not due to viscous damping but to the exchange of two branches of the dispersion relation which in the inviscid case are uncoupled.

The linear precursors of ELMS have been shown [4] to have similar tearing mode characteristics, in real frequency and growth rate. In conclusion the new analytic dispersion shows that the mechanism of viscous dissipation of vorticity affects the magnetic resistive instability by allowing a coupling of the branch of the dispersion relation dominant in the low $\omega^*/\gamma_T$ range with another branch, which is subdominant for low $\omega^*/\gamma_T$ but dominant for large $\omega^*/\gamma_T$. The ensuing phenomenology is similar to what can be expected from Reynolds theorem.

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References.
Fig. 1 $\gamma' / \gamma T$ and $\omega / \omega^* \times \omega^* / \gamma T$ for inviscid case.

Fig. 2 Viscous case. Full line: $\gamma' / \gamma T$ vs $\omega^* / \gamma T$. Dotted line: $\omega / \omega^* \times \omega^* / \gamma T$. for $m=4, n=1$, $Te(0)=1.5 \text{ keV}, Ti(0)=1 \text{ keV} B=2.2 \text{ T}$

Fig. 3 Viscous case. Full line: $\gamma' / \gamma T$ vs $\omega^* / \gamma T$. Dotted line: $\omega / \omega^* \times \omega^* / \gamma T$. for $m=4, n=1$, $Te(0)=1.5 \text{ keV}, Ti(0)=1 \text{ keV} B=2.2 \text{ T}$
A HYDRO-DYNAMIC DESCRIPTION OF SCRAPE-OFF PLASMAS CONTAINING SEVERAL ION SPECIES.

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Introduction. An adequate description of plasmas with several ion species is of importance for transport modelling of hydrogen isotopes and impurities in the scrape-off layer (SOL) of thermonuclear devices and a hydro-dynamic approximation is often used for this purpose[1]. A common ion sound velocity along the magnetic field is usually taken as boundary conditions for the motion equations of all species at the material surfaces - limiters or divertor plates. In principal such a choice is arbitrary and may come in contradiction with the structure of equations as it occurs, for example if boundary conditions are posed in a region where the plasma flow is supersonic one [2]. The consideration performed in the present paper shows that only one condition being an analog of the Bohm criterium can be posed at the material surface; other conditions are not as a matter of fact the boundary ones and must be posed at a certain intermediate point in the flow, which position is not known a priory.

Equations of motion. We consider a one-dimensional non-viscous flow along the magnetic field of a plasma with two ions species of different masses (a generalization on the case of a larger number of ions of different masses and charges is straightforward and will be done elsewhere). Zero electric current is assumed and perpendicular transport of momentum is neglected. Under such assumptions the momentum equations for ions and electrons are as follows:

\[ \frac{d}{dt}(m_{1,2}n_{1,2}V_{1,2}^2 + n_{1,2}T) = e n_{1,2}E|_l - R_{1,2}, \]  

\[ \frac{d n_e T}{dl} = -en_eE|_l, \]  

where \( T \) is the plasma temperature taken the same for ions and electrons and invariable with \( l \), \( n_{1,2}, V_{1,2} \) are the densities and velocities of ions, \( n_e \) is the electron density, \( R_{1,2} = n_1n_2m_{1,2}a(V_{1,2} - V_{2,1}) \) are the friction forces between different ions owing to coulomb collisions, \( a = 4(2\pi/m_{1,2})^{1/2}\Lambda_e/(3T^{3/2}) \) and \( m_{12} = m_1m_2/(m_1 + m_2) \)[3].

Bohm criterium. In order to derive conditions being an analogous of the Bohm criterium we apply the Poisson equation

\[ \frac{d^2\varphi}{dl^2} = -4\pi e(n_1 + n_2 - n_e), \]

to the distribution of the electric potential \( \varphi \) in the Debye sheath. In order to provide zero electric current to the material surface \( \varphi \) changes here by several \( T \)
on the Debye length and $E_\parallel$ is much larger than outside the sheath. Therefore the condition

$$4\pi neE_\parallel (n_1+n_2-n_e)=\frac{d\varphi}{dl} \frac{d^2\varphi}{dl^2} = \frac{1}{2} \frac{d}{dl} E_\parallel^2 \geq 0,$$

must hold at the boundary between the sheath and pre-sheath region.

In order to calculate the l.h.s. we summarize Eqs.(1),(2) assuming that in the sheath the forces $R_{1,2}$ and particle sources are negligible so that fluxes $\Gamma_{1,2}=n_{1,2}V_{1,2}$ are invariable with $l$ (and are assumed equal). Taking into account that the plasma is quasineutral at the border between sheath and pre-sheath ($n_1+n_2=n_e$) we obtain the conditions:

$$G_{1,2}(M_1,M_2)=F(M_1,M_2)/(2M_{1,2}^2 -1) \geq 0,$$

where

$$F=2(M_{1,2}^2-1)(2M_{1,2}^2-1)\alpha_1 + 2(M_{1,2}^2-1)(2M_{1,2}^2-1)\alpha_2,$$

$$\alpha_{1,2}=(1+V_{1,2}/V_{2,1})^{-1}, M_{1,2}=V_{1,2}/V_s^{1,2}$$

are the Mach numbers with $V_s^{1,2}=(2T/m_{1,2})^{1/2}$ being the ion sound velocities.

These conditions generalize the Bohm criterium $M_{1,2} \geq 1$ for one ion species. For the case $m_2/m_1=2$ Fig.1 shows the region in the plane $(M_1,M_2)$, where (3) hold, and the branches of the curve $F(M_1,M_2)=0$ are also presented there.

Equations for Mach numbers. Now we consider the SOL part beyond the Debye sheath, where the plasma is quasi-neutral. The SOL transparency for neutrals is assumed so that the transverse diffusion from the discharge core is the main source of particles - Fig.2 - and the particle flux densities increase linearly with the distance from the symmetry plane: $\Gamma_{1,2}=\Gamma_0 l/L$. In such a case Eq.(2) gives the following expression for $E_\parallel$:

$$E_\parallel=eT(\alpha_1 \frac{d\ln M_1}{dl} + \alpha_2 \frac{d\ln M_2}{dl} - \frac{1}{l}).$$

Rewriting Eqs(1),(2) for Mach numbers one obtains:

$$\frac{dM_{1,2}}{ds} = \frac{M_{1,2} F_{1,2}(s,M_{1,2})}{s F(M_{1,2})},$$

where $s=l/L$,

$$F_{1,2}=(1-2M_{2,1}^2)(2+\alpha_{2,1} + 2M_{1,2}^2) - (1-2M_{1,2}^2)\alpha_{2,1} +$$

$$+ 2\varrho s^2(1-M_{2,1}^2)(\alpha_{1,2}^{-1} -2),$$

and $\varrho=\Gamma_0 m_{1,2} dL/T$. 


The approach to the numerical solution and results. The solutions $M_{1,2}(s)$ of Eqs.(4) correspond to a certain integral trajectory $M_2(M_1)$ in the plane $(M_1, M_2)$. This trajectory starts at the point $(0,0)$ and according to the generalized Bohm criterium finishes at the second (right) branch of the curve $F=0$ or passes through it into the region $G_{1,2}>0$. Here $F_{1,2}(M_1, M_2)$ are negative: it follows from the behaviour of these functions by large $M_{1,2}$ and from Fig.3, where the curves $F(M_1,2)=0$, $F_{1,2}(s=1,M_1,2)=0$ are shown for $m_2/m_1=2$, $\rho=1$. Hence in the region in question the r.h.s. of Eqs.(4) are negative, because $F$ is positive, and integral trajectory can not pass into the region $G_{1,2}>0$ and finishes at the right branch of the curve $F(M_1,2)=0$.

On the way from the point $(0,0)$ to the right branch of the curve $F(M_1,2)=0$ the integral trajectory intersects by a certain $s_*<1$ the left branch of this curve at a certain point $(M_1^*, M_2^*)$. For a smooth transition through this point $F_{1,2}(s_*,M_1,2)=0$ must also be zero. This means that the curves $F_{1,2}(s_*,M_1,2)=0$ and the left branch of the curve $F_{1,2}(M_1,2)=0$ intersect simultaneously in one point as it takes place in the case of $s=1$ presented in Fig.3. Calculation shows that it is an intrinsic feature of the functions $F, F_{1,2}$: for all $s, \rho$ and mass ratios the curves $F=0$ (the left branch), $F_1=0$ and $F_2=0$ intersect simultaneously (of course the point of intersection is not the same for different $s, \rho$ and $m_2/m_1$).

Such conclusions make it possible to propose the following "recipe" for the numerical solution of Eqs.(4). Initially we take a certain $s_*$ from the interval $(0,1)$ and find $M_{1,2}^*$ solving the equations $F(M_{1,2}^*)=0$ (left branch), $F_1(s_*,M_{1,2}^*)=0$ or $F_2(s_*,M_{1,2}^*)=0$. Then Eqs.(4) are integrated at the interval $s_* \leq s \leq 1$ with the initial conditions $M_{1,2}(s_*)=M_{1,2}^*$, where the latter are a little bit larger than $M_{1,2}^*$. $s_*$ is determined from the requirement: at $s=1$ the integral trajectory must come to the right branch of the curve $F=0$. After this Eqs.(4) are solved at the interval $0 \leq s \leq s_*$ with the initial conditions $M_{1,2}(s_*)=M_{1,2}^*$, where $M_{1,2}^*$ are a little bit smaller than $M_{1,2}^*$. In order to avoid the singularity at $s=0$ it is convenient to proceed to new variables $N_{1,2}=M_{1,2}/s$. Fig.4 shows the solutions $M_{1,2}(s)$ found in such a way for $\rho=1$, $m_2/m_1=2$. The solid parts of curves correspond to the interval $s_*=0.94 s \leq 1$ where the integral trajectory passes between two branches of the curve $F=0$.

Conclusion. The present study shows that a consistent description of plasmas with several ions must be in principal different from those used in running codes with boundary conditions posed at material surfaces. This situation reminds that taking place in the "gas target" mode of a divertor operation where strong plasma cooling causes a super-sonic motion near the plates[2].

References.
Figure Captions.

1. The curves $F(M_1, M_2) = 0$ and the region where the generalized Bohm criterium $G_{1,2} \geq 0$ is satisfied ($m_2/m_1 = 2$).
2. The SOL geometry and the picture of particle fluxes.
3. The curves $F(M_1, M_2) = 0$, $F_{1,2}(s = 1, M_{1,2}) = 0$ for $\rho = 1$.
4. Computed distributions of Mach numbers for $m_2/m_1 = 2$, $\rho = 1$.
The experimental observations in the boundary layer of the TEXTOR tokamak with toroidal belt limiter ALT-II at \( \theta_{\text{lim}} = -45^\circ \) show marked poloidal asymmetries of the plasma profiles. Most of them could be explained by modelling with the 2d-two-fluid-code SOLXY, which manifested the crucial role of drift motions and nonambipolarity to be included selfconsistently in the calculations /1,2,3/. There remained, however, two major discrepancies with experimental evidence in ohmically heated discharges. Firstly, the radial density profiles at the top position (\( \theta = +90^\circ \)) are always clearly increased (\( \Delta \) shifted outwards) relative to those at the outboard midplane (\( \theta = 0^\circ \)). Secondly, the radial profiles of the poloidal velocity \( v_p \) measured at \( \theta = 0^\circ \) exhibit a shear layer with flow reversal near the separatrix and high velocities in the electron diamagnetic direction (+\( \theta \)) just inside the separatrix connected with large inward directed electric fields /2,4/. It is the purpose of this paper to develop a physical picture for the explanation of these phenomena.

We consider the TEXTOR boundary layer in the region \( 44 < r < 50 \) cm with the separatrix radius \( a = 46 \) cm. The major radius is \( R = R_0 + r \cos \theta \) with \( R_0 = 175 \) cm. The magnetic field is \( B = B_\phi = B_0 R_0/R \) with \( B_0 = 2.25 \) T, and the plasma current is \( I_p = 400 \) kA. The transformed radial coordinate \( y \) is defined by \( dy = Rdr/R_0 \) with \( 0 < y < 6 \) cm and the separatrix radius \( y_s = 2 \) cm. The electron side (e-side) of the idealized limiter is at \( \theta_{\text{lim}} + \frac{1}{2} \Delta \theta \), the ion side (i-side) at \( \theta_{\text{lim}} - \frac{1}{2} \Delta \theta \), with \( \Delta \theta \approx 3^\circ \), and the bisectrix is at \( \theta_{\text{lim}} + 180^\circ \).

In order to simulate quite realistic plasma profiles, the code SOLXY is now routinely run with spatially dependent anomalous cross-field transport coefficients. We use \( D_\perp = 0.6 f(y) g(\theta) \) \( R/R_0 \) \( m^2 \text{s}^{-1} \), see Pl. 1. The heat conductivities and viscosity are \( \chi_i = nD_\perp, \chi_e = 3nD_\perp \), \( \eta_\perp = \frac{1}{2} mnD_\perp \). Conforming to experiment we have \( T_e = 2T_i \), see Pl. 5, 6. The magnetic field reversal \( (I_p, B_\theta, B_\phi) \rightarrow - (I_p, B_\theta, B_\phi) \) can be simulated in the model calculations by mirroring the limiter position at the midplane (\( \theta_{\text{lim}} = -45^\circ \rightarrow +45^\circ \)) and mirroring back the final results. We have marked all plots referring to the mirrored position (reversed operation) by a plus sign behind the plot number, e.g. Pl. 4+.

For the discussion we use Fig. 1 and the poloidal component of the equation of motion (slightly simplified by omitting some correction terms related to field line curvature):

\[
\rho \left( \frac{\partial v_\theta}{\partial t} + v_\theta \frac{\partial \rho}{\partial \theta} + \frac{\partial F_{\text{vsec}}}{\partial \theta} + j_i B_\theta + I_0 \right) = I_1 - \rho v_r \frac{\partial v_\theta}{\partial r} \quad \text{with} \quad I_1 = -\rho v_r \frac{\partial v_\theta}{\partial r}
\]

and

\[
F_{\text{vsec}} = -h_0 (\nabla \cdot \mathbf{v}) \mathbf{v} = h_0 \left[ \frac{\rho}{\rho_0} \left( \eta_\perp v_\perp + \frac{\partial (\eta_\parallel v_\parallel)}{\partial r} \right) + \frac{\partial \rho}{\partial r} \left( \eta_\parallel v_\parallel \right) \right], \quad v_\perp = \frac{1}{B} \left( \frac{\partial P_i}{\partial \theta} + \frac{\partial \Phi}{\partial r} \right)
\]

\[
\rho = mn, \quad h_0 v_\perp = v_\theta - v_\perp, \quad h_0 = B_\phi/B, \quad I_0 = I_{\text{eo}} = S_{\text{eo}} - mv_\theta S_\perp
\]

In the scrape-off layer (sol) the electric potential \( \Phi \) is defined by the Langmuir sheath potential and the parallel Ohm's law, which yields also the rotational drift velocity \( v_\perp \). In the adjacent transition layer, however, we must prescribe \( v_\perp \) somewhere, say at the bisectrix, and then proceed backwards in order to infer \( \Phi \). Hitherto we simply assumed an exponential decay, see \( f_{\text{exp}} \) of Pl. 7. But as a consequence of the neutralizing action of the limiter onto the sol plasma drifting with \( v_\perp \) towards the e-side we get very large poloidal pressure gradients (and electric
Schematic Plasma Edge Structure for $\theta_{\text{lim}} = -45^\circ$

fields) on the i-side, whereas the momentum transfer $I_0$ from the reemitted neutrals to the ions is more important on the e-side, because the particle influx $\Gamma_r = n v_r$ is concentrated near $\theta \approx 0^\circ$. The resulting poloidal force $F_{\lim}$ in Fig. 1 is not confined to the sol, but extends also into the transition layer just inside the separatrix (partly due to the angular distribution of the emitted neutrals), where it must lead to charge separation and hence to the build-up of a negative radial electric field $E_r = -\partial \Phi / \partial r$ everywhere around the periphery. The corresponding electric drift must create a global circulation layer (GC-layer) around the minor circumference. We take this into account by prescribing $v_\perp$ to be proportional to the cubic parabola $f_{\text{GCL}}$ of Pl. 7, with the additional constraint that $\int v_\perp \, dy = \int v_\theta \, dy$ (integrated from $y = 0$ to $y = y_{\text{lim}}$). This minimizes the parallel flow velocity $v_\parallel$ and thus the viscous damping by the force $F_{\text{visc}}$, which allows the plasma within the GC-layer to spin up to poloidal velocities of $v_\theta \lesssim 4 \text{ km/sec}$, see Pl. 2. The spinning motion must be limited by a global force balance (dynamic equilibrium) provided by counteracting forces from $I_1$, which describes the momentum transfer by particles diffusing with $v_r = -D_{\perp} \partial n / \partial r$ from the plasma core across the GC-layer to the sol. The outboard centered particle influx $\Gamma_r$ is swept towards the top while diffusing through the GC-layer, forming there a characteristic density mountain and finally entering the sol upstream with respect to the $v_\theta$-velocity prevailing in the sol, Pl. 3, 11. In this way the significant increase and shift of the density profile at the top position can now be largely reproduced and easily explained. The situation is somewhat different for $\theta_{\text{lim}} = +45^\circ$ because in this case the outboard centered $\Gamma_r$ is transported via the GC-layer directly from the i-side of the limiter to the e-side, which means that it finally enters the sol downstream with respect to the sol-$v_\theta$. As a consequence the counteracting force $I_1$ becomes in this region near the e-side a propulsive force so that the viscous damping $F_{\text{visc}}$ must now play a non-negligible role in order to ensure dynamic equilibrium. For $\int v_\perp \, dy = 0.9 \int v_\theta \, dy$ we find a considerable parallel flow in the transition layer (Pl. 13+), an even more pronounced density mountain (Pl. 4+ , 12+), a slightly enhanced global circulation (Pl. 9+, 10+) and corresponding electric fields of $E_r \approx -105 \text{ V/cm}$ (Pl. 14+) compared to $\approx +20 \text{ V/cm}$ in the sol, both numbers nicely agreeing with experiment /4/. The density mountain in the GC-layer and the concomitant density valley are connected by poloidal pressure gradients to be compensated by Lorentz forces $j_r B_\theta$ from radial currents, so that the current density lines are sinuslike deformed (Pl. 8+).

Whereas the rotational drift within the sol is in the ion diamagnetic direction, the circulation in the GC-layer is in the electron one, which establishes quite naturally a highly sheared flow near the separatrix. We conclude that the GC-layer should be of special relevance also for impurity transport and the L- to H-mode transition in tokamaks with additional heating and faster characteristic velocities.

Diffusion Coefficient $D^t [m^2/sec]$

Poloidal Velocity $v_B [km/sec]$

Density $n [m^3]$ $5.6 \times 10^{18}$

Density $n [m^3]$ $5.6 \times 10^{18}$

Electron Temperature $T_e [eV]$ 52

Ion Temperature $T_i [eV]$ 86
Rotational Drift Velocity at Bisection

\[ v(y) \in \text{fun~expy-y.} \]

Global Current Flows off:

- Density
- Flux Density
- Electric Potential

\[ n \in [10^{18} \text{m}^{-3}] \]

\[ m \in [10^{19} \text{m}^{-3}] \]

\[ \phi [V] \]

\[ V \in [kT/\text{sec}] \]
Kinetic Treatment of Space Plasma Edge Structure in $\mathbf{E} \times \mathbf{B}$ Fields

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1. Introduction. The paper presents the results of investigation into the problem of describing with kinetic rigour the bounded plasma being in the external magnetic field. Two cases have been considered: (i) transition layer formed through the contact of the magnetized plasma with the material wall and (ii) formation of the plasma edge in plasma-vacuum transition with the external magnetic field applied. The physical model is based on the following basic model postulates: (1) the stationary and one-dimensional problem of the bounded plasma in the external magnetic field is considered; (2) the plasma in the transition area is assumed to be collisionless; (3) beyond the transition layer the plasma is assumed to be Maxwellian through collisions. The statements of the boundary kinetic problems are in close agreement with the postulates, with the exact analytical solutions found to them being presented in work /1–4/. These solutions underlie the investigation into the transition layer structure and the boundary of the plasma with vacuum. The condition has been found which must be satisfied to correctly simulate the transition layer and also the spatial distribution of electric potential and the density of magnetizing current for a nonisothermal plasma.

2. Statement of problem

The plasma area in question is bounded by the planes: $x = 0$ (arbitrary boundary in the plasma) and $x = d$ (coordinate of the material wall). An inhomogeneous external magnetic field is present: $\mathbf{B} = \{0, B_y, 0\}$ where $B_y = \text{const}$. Vector potential $\mathbf{B} = \{\mathbf{v} \times A\}$ has only one component $A_z = \{0, 0, A_z\}$. 

In our case: $A_z(x) = -B_y x$. It can be shown that addition of the magnetic field component along the axis $x$ does not affect the solution to the kinetic problem and so our statement of the problem covers the statement of the problem with an inclined magnetic field /5/.

As indicated in /1/ the set of kinetic equations for the plasma components and the plane geometry of the problem is transformed into the form

$$V_x \frac{\partial f_\alpha}{\partial x} - \frac{dU_\alpha}{dx} \frac{\partial f_\alpha}{\partial x} = 0$$

where the generalized potential for plane geometry has the form:

$$U_\alpha(x, S) = \frac{1}{2} \left[ S - \left( \frac{e}{m_\alpha} \cdot A_z(x) \right) \right]^2 + \left( \frac{e}{m_\alpha} \right) \Phi(x);$$

$A_z(x)$ and $\Phi(x)$ are magnetic and electric potentials; $S$ is an independent invariable; $\alpha$ is a plasma component (ions or electrons); $e$ is a particle charge with sign "+" for ions, "−" for electrons.

The nonequilibrium distribution functions (DF) satisfying the Maxwellian boundary conditions and being a solution to the kinetic equation (1) as represented in /1/ have the form

$$f_\alpha = H(\lambda_\alpha) a_\alpha b_\alpha \cdot \exp \left[ -\frac{m_\alpha}{kT} \left( \frac{V_x^2}{2} + \frac{V_y^2}{2} + \frac{1}{2} \left( S - \left( \frac{e}{m_\alpha} \cdot A_z(x) \right)^2 \right) \right) \right];$$

$$a_\alpha = N_\alpha \left( \frac{m_\alpha}{2 \pi kT} \right)^{3/2}$$

is a normalizing multiplier;

$$b_\alpha = \exp \left[ -\frac{e}{kT} \left( \Phi(x) - \Phi(0) \right) \right]$$

is a Boltzmann multiplier;

$$\lambda_\alpha = V_x - \text{sign}(x - \bar{x}_\alpha) [\bar{U}_\alpha - U_\alpha]^{1/2}$$

is an argument of Heavyside function $H$;

$$\bar{U}_\alpha = \max_x \{ u_\alpha \}$$

and $U_\alpha(\bar{x}) = \bar{U}_\alpha$ is maximum of total potential and its coordinate;
In /3/ it is shown that type function (2) is not a unique solution to the stationary boundary kinetic problem. These function proved to be the solutions with accuracy to the arbitrary distribution function of trapped particles. Work /4/ presents a mathematical apparatus devised and designed for filling the potential wells with trapped particles arbitrarily. Write the distribution function of the plasma components as

\[ f_a = f_a^P + f_a^T, \tag{3} \]

where \( f_a^P \) is the distribution function of the passing particles (i.e., the particles passing over the boundary \( x = 0 \)); \( f_a^T \) is the distribution function of the trapped particles (i.e., the particles not passing over the boundaries of the plasma area in question).

An explicit form of the distribution function of the trapped particles cannot result from solving the stationary boundary kinetic problem. In the current work the distribution function of the trapped particles is specified so that in phase space there should be no break between the distribution functions of the passing and the trapped particles. Proceeding form this write the distribution function for the passing and the trapped particles in line with physical statement of the problem accepted in the work for the transition layer in the external magnetic field.

\[ f_a^P = \alpha_a \cdot \alpha_a \cdot \exp \left( - \left( \frac{m}{kT} \right) \left( \frac{V_x^2}{2} + \frac{V_y^2}{2} + \frac{1}{2} \left( \frac{e}{m} \right) \cdot A_a(x) \right)^2 \right) \times \left( H(\alpha_a) - \Psi(H) \right), \tag{4} \]

\[ f_a^T = \alpha_a \cdot \alpha_a \cdot \exp \left( - \left( \frac{m}{kT} \right) \left( \frac{V_x^2}{2} + \frac{V_y^2}{2} + \frac{1}{2} \left( \frac{e}{m} \right) \cdot A_a(x) \right)^2 \right) \times \psi(H), \tag{5} \]

where \( \Psi(H) = H(x - \bar{x}_{la}) - H(\bar{x}_{ra} - x) \cdot H(\bar{U}_a - U_a - \frac{V_y^2}{2}) \).

\( \bar{x}_{la} \) and \( \bar{x}_{ra} \) are the coordinates of the left and right boundaries of the potential well; \( \bar{U}_a \) is local maximum value of the total potential at the point \( \bar{x}_{la} \).

Write the Poisson equation conforming to our problem

\[ \varepsilon_0 \Delta \Phi(x) = - \sum_a \varepsilon_a \int f_a \, d\mathbf{V}. \tag{6} \]

This equation was solved numerically for DF of the form (3) and (4) with zero boundary conditions for electric potential. DF (4) together with expressions (5) and (6) permit describing the trapped particles all over the interval \( 0 < x < d \), thereby we obtain a solution for the plasma-wall transition layer. In simulating the plasma-vacuum transition area we made use of DF of the form (3) since it permits bounding the trapped particle availability by the coordinate \( x = d/2 \) so that in the area \( d/2 < x < d \) we have vacuum spacing.

The magnetizing current was calculated from the expression

\[ j_z = \sum \varepsilon_a \int V_a f_a \, d\mathbf{V}. \]

For the plasma with density smaller by a factor of \( 10^{-20} \, m^{-3} \) the intrinsic magnetic field produced by this current is considerably smaller than the external one and its effect was neglected.

3. Results of numerical calculations

Calculations were carried out for the isothermal and nonisothermal hydrogenic plasma with particle density \( n = 10^{17} \, m^{-3} \) over the range of electron temperatures \( 5 \cdot 10^4 \leq T_e \leq 2 \cdot 10^8 \, K \). After solving the equation (7) the density distributions were calculated of ions and electrons of the densities of electric charge and the electric field for the plasma-wall (Fig.1 and 3) and plasma-vacuum (Fig.2 and 4) transition area. The transition area has two characteristic scales equal to the average Larmor radii of the ions \( (R_i) \) and electrons \( (R_e) \). The analysis of the results presented in Fig.1 and 2 shows that in order to find the right solutions the condition \( d > \max (R_i, R_e) \) should be satisfied.
otherwise there occurs at the plasma-layer boundary an abrupt change in electric potential, the electric field and the charge density being discontinuous. No physical processes occurring at the plasma-layer boundary can account for this discontinuity. The only interpretation of the results obtained can be the fact by that only in case of $d > R_l$ we correctly simulate the transition layer with the availability of an external magnetic field. From the calculations it proceeds that the whole transition area represents a double electric layer and the plasma boundary has a sharply outlined edge whose width is determined by the density of nonisothermal plasma and by the magnetic field quantity. The intrinsic magnetic field together with the external one gives rise to electric drift in the crossed fields along the boundary and because of the inhomogeneity of the charged particle concentration of the plasma in the transition area there arises a magnetizing current (cyclotron current). Fig.2 and 4 present the space structure change in the transition area and the blurring of the plasma edge as the nonisothermal plasma increases. Distribution of all characteristics is qualitatively different for the cases $T_e/T_i < m_i/m_e$ (curves 1, 2, 3) and $T_e/T_i > m_i/m_e$ (curves 5 and 6) and in the case $T_e/T_i = m_i/m_e$ in the transition region the double electric layer vanishes, the electric field and potential being equal to zero (curve 4).

4. Conclusion

The results obtained permit drawing the conclusion that the statement of the stationary one-dimensional boundary kinetic problem enables one to thoroughly investigate the transition layer structure and the plasma edges in the availability of the external magnetic field. Thereby pre-requisites have been provided for investigating the plasma system structure in the availability of plasma flow components along the material wall. Solution to this problem enables one to investigate the pinching of the plasma cylinder by the intrinsic magnetic field with the electric current flowing along the plasma cylinder.

/2/ Zhykharsky A.V. International Conference on Plasma Physics, 1992, p.1553-1556;
/5/ Chodura R., Contributed Papers, 1992 ICPP, Part II, V. 16c, 871-873)
Fig. 3. Space structure of plasma-vacuum edge.
Fig. 4. Space structure of plasma-wall transition layer

Nonizothermality: \( \frac{T_e}{T_i} \) = 1 - 1; 2 - 200; 3 - 1000; 4 - 1800; 5 - 2000; 6 - 4000.

a) ion density; b) electron density; c) electric potential; d) electric field; e) electric charge density;
f) magnetizing electric current density.
THE IMPACT OF THE BIASING RADIAL ELECTRIC FIELD ON THE SOL IN A DIVERTOR TOKAMAK.
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Strong radial electric field can be induced within the SOL in a divertor tokamak by applying a voltage to divertor plates with respect to the first wall. This biasing scheme results in the strong radial electric field which is much larger than the natural electric field, usually of the order $T_e/e$. Experiments employing this biasing scheme were carried out on the tokamak TdeV [1]. Many interesting effects such as - modifications of the density profile and radial transport of impurities as a function of the polarity and the magnitude of the biasing voltage, the generation of the flux surface average toroidal rotation proportional to the applied voltage, redistribution of the plasma outflow onto divertor plates and so on - were demonstrated to result from the biasing. Furthermore, in contrast to studies carried out employing a different biasing scheme which primarily results in a poloidal electric field, the strong radial electric field impacts more significantly within SOL than the poloidal electric field [1]. Here, we aim to show that the main effects observed experimentally follow from the analysis, provided continuity and momentum balances are employed invoking anomalous viscosity and inertia [2,3].

Forsimplicity, we consider the double - null configuration with circular magnetic surfaces. Radial current is obtained from the toroidal and parallel components of the momentum balance

$$\frac{m_1}{r} \frac{\partial (r \mu_r u_{\phi})}{\partial r} + \frac{m_1}{r} \frac{\partial (r \mu_\theta u_{\phi})}{\partial \theta} = - (\nabla \cdot \vec{v}) \phi + \frac{1}{c} j_{B \theta}$$

(1)

$$\frac{m_2}{r} \frac{\partial (r \mu_{r \|} u_{\|})}{\partial r} + \frac{m_1}{r} \frac{\partial (r \mu_{\theta \|} u_{\|})}{\partial \theta} = - (\nabla \cdot \vec{v}) \| - \nabla \| p$$

(2)

Parallel current may be obtained from

$$\frac{1}{r} \frac{\partial j_{\theta}}{\partial \theta} + \frac{\partial (n u_{\theta})}{\partial r} = 0$$

(3)

Parallel and toroidal velocities are related through

$$u_{\|} = u_{\phi} + \Theta u_{\theta}$$

Besides, the following relation is maintained within SOL

$$u_{\phi} = V_0 + u_{p1} + \Theta u_{\theta}$$

(4)

which follows from the radial component of the momentum balance. Note that Eq.(4) is valid locally. Assuming that the potential brought about by biasing significantly exceeds $T_e/e$, the radial distribution of plasma potential is governed by the potential profile on plates, thereby determining the velocity of the poloidal rotation $V_0$. For a given value of $V_0$, the radial electric field is stronger than the poloidal electric field. Weak effects owing to toroidicity are neglected.

Combining (1) and (2) and neglecting a weak parallel viscosity compared with the pressure gradient, one obtains

$$\frac{1}{r} j_{B \theta} = - \frac{\Theta}{r} \frac{\partial p}{\partial \theta}$$

(5)

Integrating (5) over a magnetic surface, the total current is obtained as

$$I = \frac{2\pi R_e}{B} \left( p_- - p_+ \right)$$

(6)

where $p_-$ and $p_+$ are plasma pressures at the line - tying either to plates or to the first wall. Eq.(6) describes the case, when a strong radial electric field and current are induced only at the outer half of the torus. If the same voltage is applied at the inner half, the current given by
Eq. (6) has to be amplified by factor 2. Qualitatively, Eq. (5) yields opposite currents at the upper and lower halves of a magnetic surface, resulting in the total current to be governed by the pressure asymmetry in the vicinity of a material surface.

Flux-surface averaging Eq. (2) and neglecting the classical viscosity, one obtains

\[ \pi m_i \frac{\partial (\mu r(A) u_{||})}{\partial r} + \pi \left( \frac{m_i}{r} \right) \left( n^+ u^+ + n^- u^- \right) = \frac{\Theta}{r} (p_+ - p_-) - \pi \left( \nabla \times \mathbf{A} \right) \parallel (\nabla \times \mathbf{A}) > \] (7)

Here, \( f = \int_0^{\pi/2} \frac{\rm d\theta}{\pi} \) stands for the averaging over the outer half of a magnetic surface, \( u_{||} \) is the average toroidal velocity and \( \nabla \times \mathbf{A} \) is the anomalous viscosity. At the interface with the material boundary plasma flows with the sound speed along a field line according to the Bohm criterion. Thus

\[ u_{\theta} = \pm \Theta C_S \pm \quad C_S = \sqrt{(T_e + T_i)/m_i} \]

Accounting for Eq. (4), one obtains

\[ V_0 + u_{\theta} + \Theta u_{||} = \Theta C_S^+ \]

\[ V_0 + u_{\theta} + \Theta u_{||} = -\Theta C_S^- \]

It is worthwhile to emphasize again that \( V_0 \) is a constant along a field line due to a strong radial electric field applied externally. Substituting boundary conditions into Eq. (7), one obtains

\[ p^+ - p^- = \frac{m_i}{2\Theta} (n^+ C^+_S + n^- C^-_S) - \frac{\pi m_i}{2\Theta} \frac{d(m_r(A) u_{\phi})}{dr} + \frac{\pi}{2\Theta} \frac{d(m_r(A) u_{\phi})}{dr} \] (9a)

Substituting the current given by Eq. (6) into Eq. (9a), one obtains

\[ I = -\frac{\pi R c m_i (V_0 + u_{\phi})}{B\Theta} (n^+ C^+_S + n^- C^-_S) + \frac{\pi R c m_i}{B\Theta} \frac{d(m_r(A) u_{\phi})}{dr} + \frac{\pi}{B\Theta} \frac{d(m_r(A) u_{\phi})}{dr} \] (9b)

In order to close the system we employ the continuity equation reading

\[ \frac{1}{r} \frac{\partial}{\partial r} (n u_{\theta}) = Q = S - \frac{1}{r} \frac{\partial (\mu r(A))}{\partial r} \]

(10)

Here, \( S \) is a source of particles and we have neglected the neoclassical radial transport.

Integrating over the field line, one obtains

\[ n(V_0 + u_{\phi} + \Theta u_{\phi}) = \int_0^{\pi/2} Q(\theta) \sin \theta + C \] (11)

A constant \( C \) is determined by boundary conditions at the ends of a field line intersecting a material surface

\[ n^+ (V_0 + u_{\phi} + \Theta u_{\phi}^+) = \int_0^{\pi/2} Q d\theta + C = n^+ \Theta C_S^+ \] (12)

\[ n^- (V_0 + u_{\phi} + \Theta u_{\phi}^-) = -\int_0^{\pi/2} Q d\theta + C = -n^- \Theta C_S^- \]

The sum of Eqs. (12) yields the constant \( C \) provided the symmetry over the midplane is maintained

\[ C = \frac{n^+ C_S^+ - n^- C_S^-}{2} \Theta \] (13)

Assuming temperatures to be equal at ends of a magnetic field line \( T_e^\pm = T_i^\pm = T_c \), one obtains \( C \) as a function of the radial current \( I \) from Eq. (6)
Now Eqs. (11 & 14) yield the relation between the toroidal velocity at the equatorial plane \( u_\phi(\theta=0) \), which is measured experimentally [1], the radial electric field and the total current \( I \).

Substituting Eq. (14) into Eq. (11) and taking \( \theta = 0 \), one obtains

\[
V_0 + u_{\text{pi}} + \Theta u_\phi(0) = -\frac{IB\Theta}{4\sqrt{2\pi R_n\Theta T_c} \frac{1}{2} m_i^{1/2}} \tag{15}
\]

In low recycling divertors, the toroidal velocity at the midplane \( u_\phi(0) \) is of the same order as the toroidal rotation velocity \( u_\phi \) in Eq. (9). Hence, the system of Eqs. (9 & 15) yield the magnitude and the profile of the toroidal rotation velocity and the current for a given radial profile of the electric field. The exact solution may be obtained numerically by solving the system of Eqs. (1 & 11) for a given distribution of \( Q \).

If the radial scale of the toroidal rotation profile \( \delta_\phi \) is larger than the SOL width \( \delta \) (this is the case in experiments on TdeV [1]), then one obtains from Eqs. (9 & 15) assuming \( u_\phi = u_\phi(0) \)

\[
V_0 + u_{\text{pi}} + \Theta u_\phi(0) = 0 \tag{16}
\]

However, all the terms in Eqs. (9 & 15) are of the same order if \( \delta_\phi = \delta \) and the system has to be solved numerically. Eq. (16) yields the dependence of toroidal velocity on the applied voltage, which is in good agreement with the measured on TdeV [1]. Note that in absence of the external bias, Eq. (16) is satisfied exactly. This implies that the toroidal rotation velocity emerge within SOL coparallel to the Ohmicheating current even when the radial current through separatrix is equal to zero. The same effect has been shown to emerge for a poloidal limiter. The main reason stems from the positive natural electric field existing within the SOL.

The value of current is obtained from (9 & 15) and reads for \( \ln^+ - n^+ \ll n^+ = n_c \)

\[
I = \left[ 1 - 0.5 \exp(-0.5) \right] \frac{-1}{\Theta B} \left[ \rho_{\text{tr}}(A) \right] \left( \frac{du_{\phi}}{d\rho} \right) \left( \frac{d\rho}{dr} \right) \tag{17}
\]

Here, the relation reading \( n_c/n(0) = \exp(-0.5) \), which follows from the parallel momentum balance, has been employed.

Note that Eq. (17) is qualitatively the same as Eq. (7) in Ref. [4], which yielded the profile of the toroidal rotation velocity caused by a biased electrode inside the separatrix. However, quantitatively, the crossfield radial conductivity emerges to be much larger outside the separatrix within the SOL.

Indeed, employing in (9b) \( u_\phi = - (V_0 + u_{\text{pi}}) / \Theta \), one obtains

\[
I = \frac{k \pi^2 R n m_i^2 / T_i^{1/2}}{\Theta B^2} \frac{dr}{d\rho} \left( E_r \right) \tag{18}
\]

where \( k \) is the coefficient of the order 1. Comparison with voltage - current characteristics measured on TdeV yields \( k = 2 \).

The linear dependence of current on the electric field is maintained unless the pressure asymmetry on ends of the magnetic field line becomes significant \( n^+ - n^- \approx \max(n^+,n^-) \).

Eq. (6) yields the maximum value of the total radial current

\[
I_{\text{max}} = 4\pi R n T_i / B \tag{19}
\]

which is obtained for the electric field.
Moreover, when these large values of the electric field are reached, the substantial redistribution of the plasma outflow between divertor plates occurs. Since $E_r > 0$ the preferential outflow takes place into the low divertor chamber for the direction of the magnetic field on TdeV. In the high recycling regime, the situation remains essentially the same.

Maximum velocity averaged over the magnetic surface is obtained from $u_r = j_r/R$ as

$$u_r^{(\text{max})} = \frac{2cT_i}{erB}$$

which yields the ratio of $u_r^{(\text{max})}$ to the anomalous transport $u_r^{(A)}$ of the order 1.

$$\frac{u_r^{(\text{max})}}{u_r^{(A)}} = \frac{\pi \rho_c i}{\Theta \delta} \lesssim 1$$

Therefore, the following important effect results from this estimate. The option to control the density profile within the SOL emerges by imposing the strong electric field of the order $|V_0| = \Theta C_S$. The positive polarity of the bias results in a more uniform profile and the negative in a steeper one. This effect has also been observed on TdeV.

Finally, addressing the impurity transport the toroidal and the parallel components of the momentum balance are employed

$$m_r \frac{\partial (m_r u_r^{(A)} u_{\phi}^{(t)})}{\partial r} + m_\perp \frac{\partial (m_\perp u_\perp^{(A)} u_{\phi}^{(t)})}{\partial \theta} = - (\nabla \pi)_r + \frac{1}{c} j_r B_\theta - m_r n_i \nabla \cdot (u_{\phi}^{(t)} - u_{\phi})$$

$$m_r \frac{\partial (m_r u_r^{(A)} u_{\perp}^{(t)})}{\partial r} + m_\perp \frac{\partial (m_\perp u_\perp^{(A)} u_{\perp}^{(t)})}{\partial \theta} = - (\nabla \pi)_\perp - \nabla |p_\perp| - Z n_T \nabla \cdot (u_{\perp}^{(t)} - u_{\perp})$$

Neglecting viscosity terms compared to pressure term and the small difference between $u_{\phi}$ and $u_\perp$, one obtains for the current due to impurities

$$j_r = - \frac{c}{B_r} \left( \frac{\partial p_i}{\partial \theta} + Z n_T \frac{\partial \ln n}{\partial \theta} \right)$$

Assuming that $T_e = T_i = T = \text{constant} (\theta)$ and the Boltzmann distribution of impurities within a magnetic surface reading $n_I/n = \text{const} (\theta)$, one obtains after averaging

$$I_r = \frac{1 - Z}{2} \frac{n_I}{n}$$

where $I$ is the current carried by ions. Thus, the average outflow velocity of impurities is equal with accuracy of the coefficient $(1 + Z)/2$ to the radial transport of main ions driven by the electric field.

Finally, we conclude that the biasing of divertor plates with respect to the first wall provides the powerful tool to control particle, impurity transport and the outflow into a divertor chamber.

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REFERENCES.

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