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THE EFFECTS OF IMPURITIES AND MAGNETIC DIVERTORS ON HIGH-TEMPERATURE TOKAMAKS

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ABSTRACT

A one-dimensional tokamak plasma transport code has been adapted to include impurity influx, stripping, radiation, and diffusion, as well as the usual processes of hydrogen plasma and heat transport, recycling at the boundary, and multi-generation charge-exchange. Neutral-beam heating, adiabatic compression, and divertor boundary conditions are included as optional features. Illustrative computations are given for present-day and next-generation tokamaks. The problems of impurity control are discussed, and two technical approaches are examined in greater detail: the transient cold-plasma shield, and the poloidal divertor.

1 INTRODUCTION

Nonhydrogenic ions have an important influence on the behavior of present-day tokamak discharges, and are expected to become a key element in determining feasibility of future tokamak reactor plasmas. A relatively small admixture of heavy impurity ions suffices to dominate the effective ionic charge Z_{eff} of a hydrogen plasma, thus determining the resistivity and transport processes. This effect of impurities is generally helpful in present-day devices, by increasing the ohmic-heating input power, and could be helpful even in future neutral-beam-heated tokamaks, by diminishing the anomalous transport due to trapped-particle modes. On the other hand, impurity radiation losses are already becoming appreciable in high-density present-day discharges, and may represent a significant barrier to the attainment of ignition conditions in

future reactor plasmas [1]. Here one must distinguish between impurity ions that have been stripped to helium-like states, or beyond, and those that are only partly stripped. Line-radiation and recombination radiation from the latter constitute the most important channels for significant radiation cooling. When a given ion species is partly stripped at the plasma edge, and fully stripped on the interior (as in the case of low-Z impurities such as C and O in present-day experiments), the net effect may be beneficial, since edge-cooling by radiation represents one of the least damaging forms of heat removal to the material surroundings of the tokamak discharge. On the other hand, ingestion of partly stripped ions (Fe, Mo, in present experiments, W in the future) into the plasma interior tends to depress the central plasma temperature, and thus plays a role that is in general highly undesirable. Aside from radiation cooling, impurity ions, of course, have the adverse effect of diluting the reactive ion population, thus raising the plasma β -value required to reach a given power density, and depressing the energy multiplication factor.

Impurities are injected into the tokamak discharge by two principal mechanisms: sputtering and evaporation. Plasma heat outflow through the channel of energetic ions or neutrals can sputter surface atoms from the tokamak limiter or vacuum vessel, respectively. In addition, heat outflow, particularly through the channel of superthermal electrons, can create evaporation from "hot spots"; but the plasma heat load would not be severe, in either present-day or future experiments, if spread uniformly over the entire vacuum wall. Assuming that the technical problems of local overheating can be solved, the most serious long-range problem facing high- T_i tokamaks is sputtering — especially sputtering by energetic charge-exchange neutrals from the plasma interior.

In Section 2, we describe the principal tool of the present investigations, a one-dimensional plasma transport code [2] that has been modified to include impurity stripping, radiation and diffusion. Section 3 gives illustrative results for impurity phenomena in present-day and next-generation tokamaks. Section 4 discusses special techniques for improved impurity control in future experiments.

2 THE TOKAMAK COMPUTER MODEL

The most advanced present version of our transport code incorporates the following features.

Six-Regime Plasma Transport Model. Local plasma transport coefficients are employed that describe the anomalous transport predicted for low-frequency microinstabilities [3]. In order of decreasing collisionality, the six regimes are: three regimes of collisional drift waves (a regime of Bohm diffusion, a regime in which a theory of Kadomtsev applies, and a pseudoclassical regime); two regimes of the trapped electron instability; and the trapped-ion instability. This model has been providing good fits of much of the data from present-day tokamaks, which enter the first four regimes. In computations for next-generation tokamaks, the last two regimes are dominant, although the near-axis region (where the trapping is unimportant) and the cold edge region remain in the more collisional regimes. Although the theory underlying the six-regime model provides the dependences of the transport coefficients on the various plasma parameters, there are numerical coefficients of order unity which are uncertain. For the four regimes entered by present experiments, there are empirical indications of what these coefficients should be, but for the two most collisionless regimes we can only use the crudest theoretical estimates. A quasi-linear

type of analysis does, however, indicate the relative magnitude of particle and electron and ion thermal transport, and this is incorporated into the model. In the case of the ion thermal conductivity, we also include the neoclassical contribution, which can be greater than the anomalous term. In the case where impurities are not treated explicitly, our transport equations may be summarized as follows

$$\begin{aligned} \frac{\partial n}{\partial t} &= \frac{1}{r} \frac{\partial}{\partial r} r D \frac{\partial n}{\partial r} + S_I; \quad \frac{\partial B_\theta}{\partial t} = \frac{c^2}{4\pi} \frac{\partial}{\partial r} \left[\frac{n}{r} \frac{\partial (r B_\theta)}{\partial r} \right]; \\ \frac{3}{2} \frac{\partial (n T_e)}{\partial t} &= \frac{1}{r} \frac{\partial}{\partial r} r \left(n \kappa_e \frac{\partial T_e}{\partial r} + \frac{3}{2} D T_e \frac{\partial n}{\partial r} \right) \\ &+ n j_z^2 - \frac{3m_e}{m_i} \frac{n}{\tau_e} (T_e - T_i) - W_R - W_I + W_{Be}; \\ \frac{3}{2} \frac{\partial (n T_i)}{\partial t} &= \frac{1}{r} \frac{\partial}{\partial r} r \left(n \kappa_i \frac{\partial T_i}{\partial r} + \frac{3}{2} D T_i \frac{\partial n}{\partial r} \right) \\ &+ \frac{3m_e}{m_i} \frac{n}{\tau_e} (T_e - T_i) - W_{cx} + W_{Bi}. \end{aligned} \quad (1)$$

Here S_I denotes the source due to ionization, W_R , W_I and W_{cx} denote losses due to radiation, ionization, and charge-exchange respectively, and W_{Be} and W_{Bi} denote neutral beam heating of electrons and ions. For the resistivity we take a classical value with trapping corrections:

$$\eta = \frac{v_e m_e}{n e^2 f_{TR}}; \quad v_e \tau_e = \frac{0.457 Z_{eff}}{1.077 + Z_{eff}} + 0.29 Z_{eff};$$

$$f_{TR} = 1 - [1.9(r/R)^{1/2} - 1.0r/R]/(1 + v_E^*)$$

The classical contribution to the ion thermal conductivity is

$$\kappa_i^{cl} = \frac{0.68}{1 + 0.36 v_i^*} \frac{Z_{eff}^2 \rho_i}{\tau_i} \left(\frac{r}{R} \right)^{1/2} + \frac{Z_{eff}^2 \rho_i}{\tau_i} \left(1 + 1.6 \frac{r_{Bz}^2}{R^2 B_\theta^2} \right)$$

In these formulae, we have $\rho_{e,i} = c(2m_{e,i} T_{e,i})^{1/2} / e B_\theta$ and

$$\tau_e = 3m_e^{1/2} T_e^{3/2} / 4(2\pi)^{1/2} n e^4 \ln \Lambda; \quad \tau_i = 3m_i^{1/2} T_i^{3/2} / 4\pi^{1/2} n e^4 \ln \Lambda;$$

$$v_e^* = Z_{eff} R^{3/2} B_z / \tau_e r^{1/2} B_\theta (T_e / m_e)^{1/2}; \quad v_i^* = Z_{eff} R^{3/2} B_z / \tau_i r^{1/2} B_\theta (T_i / m_i)^{1/2}$$

Turning now to the anomalous diffusion terms, the drift wave contribution is essentially the minimum of a Bohm, Kadomtsev, and pseudoclassical term:

$$D_{\text{Drift}} = \frac{D_P D_K D_B / (D_K + D_B)}{[D_P + D_K D_B / (D_K + D_B)]};$$

$$D_B = \frac{cT_e}{16eB}; \quad D_K = \frac{v_e^{1/3}}{(r_n \theta)^{2/3} (\Omega_i \Omega_e)^{1/3}} \left(\frac{cT_e}{eB} \right)^{4/3} \left(1 + \frac{T_i}{T_e} \right)^{1/3}; \quad D_P = \frac{2v_e^0{}^2}{\theta^2} \left(1 + \frac{T_i}{T_e} \right);$$

$$r_n^{-1} = \left| \frac{d \ln n}{dr} \right|, \quad r_T^{-1} = \left| \frac{d \ln T_e}{dr} \right|, \quad \theta = \frac{r r_n}{B_z} \left| \frac{d}{dr} \frac{B_\theta}{r} \right|$$

where $\rho_{e,i} = c(2m_{e,i} T_{e,i})^{1/2} / eB_z$, and $\Omega_{e,i} = eB/m_{e,i} c$. For the trapped electron and ion modes

$$D_{TE} = \frac{1}{r_n^2} \left(\frac{r}{R} \right)^{3/2} \left(\frac{cT_e}{eB} \right)^2 \frac{v_e}{v_e^2 + v_o^2}, \quad v_o = \frac{2\pi}{r_n} \frac{r}{R} \left(\frac{T_e}{m_i} \right)^{1/2} \theta^{1/2},$$

$$D_{TI} = \frac{1}{v_e r_n^2} \left(\frac{r}{R} \right)^{5/2} \left(\frac{cT_e}{eB} \right)^2 \left(1 + \frac{T_e}{T_i} \right)^{-2}.$$

The terms D_{TE} and D_{TI} are cut off if the collisionality is large enough that the modes are stable. The terms D_{DRIFT} , D_{TE} , and D_{TI} are included in the transport coefficients D , κ_e , and κ_i . In each case, the relative magnitudes of the contributions to particle diffusion and thermal conductivity are determined by quasi-linear considerations; in all cases this merely introduces numerical coefficients of order unity which, for the trapped ion and electron modes, depend on r_n/r_T . The dependence of the anomalous transport coefficients on the scale lengths r_n and r_T and on the magnetic shear θ present a degree of computational difficulty. In computing the transport coefficients, we use values for r_n , r_T , and θ obtained from simple algebraic profiles fitted at each time-step to the actual computed profiles. The errors introduced by this process are certainly less than the uncertainties in the model itself.

Neutral Hydrogen Gas. All of the computations discussed below include an influx of "cold" neutral gas at the plasma boundary, as well as a variable number (four for present tokamaks, ten for next-generation devices) of generations of "hot" neutrals arising from the cold gas by successive charge exchanges. The energy prescribed for the cold gas seems to be an unimportant variable, and we use 10 eV. The amount of cold-gas influx varies with time during the calculation, and is usually determined by the requirement that it should exactly balance the outflux of ions and charge-exchange neutrals (100% recycling). Our computations can also include a prescribed additional component of cold gas influx. The neutral gas density within the plasma is computed kinetically, generation by generation, in a two-dimensional (r, θ) treatment. The terms S_I , W_I , and W_{CX} in Eqs. (1) represent a sum over all generations. When the impurities are not treated explicitly, but modelled by Z_{eff} , the term W_I incorporates an enhanced radiated loss (usually 200 eV) per hydrogen ionization, to account for impurity radiation. When impurity radiation is explicitly treated, this loss is reduced to 20 eV per hydrogen ionization.

Neutral Injection. In some of our calculations, neutral injection is assumed to deposit into the plasma a prescribed radial distribution of mono-energetic "beam" protons. The rate of injection and the period over which it lasts, are also prescribed. The energy of each element of the beam is subsequently determined by the classical (Fokker-Planck) slowing-down formulae. The energy lost in slowing down is transferred to the background electrons and ions according to the classical formulae, and appears in Eqs. (1) in the terms W_{Be} and W_{Bi} . In this model, energy diffusion of the beam ions is neglected, as is the beam-ion "banana orbit" width and radial diffusion. The beam ion density is attenuated by charge exchanges with the background neutral gas, and the high-energy neutrals so produced are assumed to be lost entirely. Lower energy neutrals produced by charge-exchange break-up of the primary neutral beam are not included in the background gas; although this may be a serious omission near the magnetic axis, charge-exchange losses from the beam ions are not in fact very important for next-generation devices.

Impurity Ions. The most advanced version of the code includes diffusing impurity ions of a single atomic species. The local density of impurity ions n_I (summed over all the ionization levels) is determined from a fifth diffusion-type equation:

$$\begin{aligned} \frac{\partial n_I}{\partial t} = & \frac{1}{r} \frac{\partial}{\partial r} r \left\{ D \frac{\partial n_I}{\partial r} \right. \\ & - C_{iI} \left[\langle Z \rangle n_i \left(\frac{\partial n_i}{\partial r} + \alpha_1 n_i \frac{\partial \ln T_i}{\partial r} \right) - n_i \left(\frac{\partial n_I}{\partial r} + \alpha_2 n_I \frac{\partial \ln T_i}{\partial r} \right) \right] \\ & \left. + C_{II} \left[(\langle Z^2 \rangle - \langle Z \rangle^2) n_I \left(\frac{\partial n_I}{\partial r} + \alpha_3 n_I \frac{\partial \ln T_i}{\partial r} \right) - \langle Z \rangle n_I^2 \frac{\partial \langle Z \rangle}{\partial r} \right] \right\} \end{aligned} \quad (2)$$

where n_i is the proton density. The terms in C_{iI} and C_{II} are neoclassical terms; the C_{iI} terms arise from impurity-proton collisions, and the C_{II} terms arise from collisions between different ionization levels of the impurities. The coefficients C_{iI} , C_{II} , α_1 , α_2 , and α_3 are exceedingly complicated expressions taken from the most recent treatments of the banana, plateau, and Pfirsch-Schlüter regimes [4]. The coefficients α are of order unity; α_1 is negative in the banana/plateau regimes but positive in the Pfirsch-Schlüter regime [4]. The quantities $\langle Z \rangle$ and $\langle Z^2 \rangle$ denote averages over the ionization levels present. Quasi-linear-type analyses suggest that the anomalous diffusion should apply also to the impurities: hence the term D in Eq. (2). We have also made computations with this term omitted. All of the ions are assumed to be at the same temperature T_i . When Eq. (2) is included in the computations, important modifications are made to Eqs. (1). The particle diffusion equation that is integrated is that for protons; the flux term then includes a collection of neoclassical terms related to those in Eq. (2): these describe the fact that as the high-Z ions diffuse inward the protons must diffuse outward to conserve the local ion charge density. With n_i and n_I given by diffusion equations, the electron density is determined from

$$n_e = n_i + \langle Z \rangle n_I.$$

A further collection of neoclassical terms must be added to the ion energy transport; these describe the fact that more thermal energy is carried out by the protons than is carried in by the impurities. More straightforward

modifications are to employ the total ion density in the ion heat capacity term, to include impurity contributions to the electron-ion heat transfer, and to distinguish properly between n_e and n_i in all the atomic physics terms; our code incorporates all of these modifications. To determine $\langle Z \rangle$, $\langle Z^2 \rangle$, and Z_{eff} , we must know the ionization levels of the impurities. In calculating these, we make use of the fact that the time to reach approximate coronal equilibrium is typically a few times shorter than the diffusion time scale. Accordingly, we compute the fraction n_I^j of the total impurity density n_I in each of the ionization levels j from local rate equations without diffusion terms:

$$\frac{\partial n_I^j}{\partial t} = n_e n_I^{j-1} S_{j-1} - n_e n_I^j (S_j + \alpha_{j-1} + \gamma_{j-1} n_e) + n_e n_I^{j+1} (\alpha_j + \gamma_j n_e)$$

where S_j are the ionization coefficients from the level j , and α_j and γ_j the radiative and three-body recombination coefficients to the level j . The power radiated by ionization and recombination, namely

$$W_R = \sum_j [\chi_j n_e (S_j n_I^j - \gamma_j n_e n_I^{j+1}) + \frac{3}{2} T_e \alpha_j n_e n_I^{j+1}] ,$$

where χ_j are the ionization potentials, is included in the electron energy equation, as is the line radiation. For oxygen impurities, the rate coefficients and detailed expressions for the line radiation are taken from the most appropriate data [5]. We have also considered iron impurities, including only the eight most highly ionized states. For these we have used rate coefficients taken from a treatment by Hinnov [6]. For iron, the line radiation has been estimated as $2 \times 10^{-26} n_e$ watts for each impurity ion in the lithium-like state and below [6,7]. The appropriate boundary condition to be applied to Eq. (2) is of some interest. The boundary conditions that we have used are conditions on the flux of impurity ions, rather than on the boundary density n_I . For oxygen impurities, which might be expected to recycle on the limiter in much the same way as the hydrogen, the appropriate condition is that the oxygen ion flux should vanish at the boundary (neutral impurities are not included in the calculations). The iron impurities are presumed to be produced by the sputtering of the wall by energetic charge-exchange neutral bombardment. The sputtering coefficient that we have used to describe this process is

$$S = S_{\text{max}} \frac{4(E - E_C)^{3/2} E_M^{1/2}}{E_M^2 + 3E^2} , \quad E_M = 6000 , \quad E_C = 70 , \quad S_{\text{max}} = 7 \times 10^{-3}$$

where E is the energy of the bombarding neutral in eV; this formula is obtained by fitting recent data for stainless steel [8]. It is presumed that the classical inward-diffusion process is sufficiently strong in the thin edge region (penetrated by neutral impurities) that the total sputtered impurity flux enters the main part of the plasma as impurity ions. Accordingly, our code calculates the energy spectrum of the charge-exchange neutrals leaving the plasma and, by means of the sputtering coefficient, determines the impurity influx at the boundary at each instant. The impurities enter with the $\langle Z \rangle$ appropriate to the edge temperature.

Adiabatic Compression. Our code incorporates the facility for making an instantaneous major-radial compression by a factor C . The plasma density, temperature, and poloidal field profiles are compressed according

to the adiabatic laws; the energy of beam ions is raised by C^2 (for parallel beams). The discharge subsequently evolves according to the transport processes in the new configuration.

Divertor Boundary Conditions. Normal (limiter) boundary conditions are fixed low values of density and temperature. To describe the effect of a divertor on the transport in the main plasma region, we have developed alternative boundary conditions for the radial particle and energy fluxes at the separatrix, obtained by considering the parallel flow in the scrape-off region, including neutral gas effects, in a self-consistent way. These computations cannot be presented in detail; they support the conclusions presented in Section 4.

3 IMPURITY TRANSPORT COMPUTATIONS

In this section we present some results on impurity transport which have been obtained using the code described above. We present results for a small tokamak, representative of present-day devices (ST parameters: $a = 14$ cm, $R = 109$ cm), and for a several-times larger tokamak, representative of next-generation devices, now under construction (PLT parameters: $a = 45$ cm, $R = 130$ cm).

Present-Day Tokamak Model. In Table 1, we present the results of computations using several different treatments of the impurity transport. All cases have $q(a) = 5.4$, constant current, and 100% recycling hydrogen, so that the hydrogen density remains roughly constant. Plasma parameters are at 60 msec. Cases A and B are without an explicit treatment of the impurities (they use $Z_{eff} = 4$ uniform, and a radiated energy loss of 200 eV per hydrogen ionization); these cases are presented to show that our plasma transport model fits typical β_θ -values and confinement times quite well, and for comparison with the impurity cases C-F. Cases C-F

	I kA	B kG	$n_e(0)$ 10^{13} cm^{-3}	β_θ	τ'_E msec	$T_i(0)$ eV	$T_e(0)$ eV	$q(0)$	$\bar{n}_I$	$n_I(0)$ 10^{12} cm^{-3}
A	60	36	2.9	0.8	15	360	1080	1.7	$Z_{eff} = 4$	
B	100	60	5.3	0.9	25	680	1440	1.8	$Z_{eff} = 4$	
C	60	36	3.1	0.6	10	260	410	5.5	1.0	2.1
D	60	36	3.3	0.8	13	340	760	2.4	1.0	1.7
E	60	36	3.6	0.8	11	320	750	2.4	1.5	2.4
F	100	60	6.4	0.9	25	700	1300	2.1	2.0	3.6

Table 1. Results for present-day tokamak model. Case C has classical impurity transport; cases D-F also have anomalous diffusion.

have an explicit treatment of diffusing oxygen impurities. In all cases the initial oxygen ion density n_I is radially uniform, and the oxygen is assumed to recycle, so that the net flux at the boundary is set to zero. In case C, which has $\bar{n}_I = 10^{12} \text{ cm}^{-3}$ (i.e. about 5% impurities), only the neoclassical terms are included in the oxygen diffusion [the anomalous term D is omitted from Eq. (2)]. From the last two columns in

Table 1, we note that the oxygen ions become relatively strongly peaked in this case; from Fig. 1, however, we note that, as a consequence, the temperature profile has become very broad (the current profile is in fact "hollow"). In cases D-F both the neoclassical and anomalous terms are included in the oxygen diffusion. From Table 1, we note that the oxygen ions become somewhat less strongly centrally peaked in these cases. In Fig. 2 we show the radial distribution of the different oxygen ionization levels for the case E, which has 7% impurities; the radiated power in this case is only 10 kW out of an ohmic power of 120 kW.

Let us compare the the results with actual tokamak experiments. In the ST device [9], the observed T_e -peaking is typically sharper than in cases A and B. This discrepancy is clearly aggravated by both the classical (C) and classical-anomalous (D, E, F) models of impurity transport. The computed profiles have a closer resemblance to the T-3 results [10] and to "Type-B" discharges in ORMAK [11] — while the ORMAK "Type-A" profiles resemble those in ST. Further refinements in the theoretical model are clearly needed to fit this multiplicity of experimental regimes. As regards the computed degree of central impurity accumulation, this effect is weak for oxygen, especially in the classical-anomalous model, consistent with experimental observation [12].

Fig. 1. Electron temperature profiles for ST-size model. (PPPL 742139)

Fig. 2. Profiles of oxygen ionization levels for ST-size model, Case E. (PPPL 742138)

Next-Generation Tokamak Model. In Table 2, we present some results for a several-times larger tokamak device. All cases have $q(a) = 3.9$. Plasma parameters are at 200 msec. Cases A and B are without an explicit impurity treatment ($Z_{eff} = 2$ uniform, and 200 eV lost per hydrogen ionization). Case A is a pure ohmic heating discharge, whilst case B has a neutral injection beam current of 100 A of 40 keV protons, initiated at

100 msec and terminated at 200 msec, with the beam ions deposited with the profile $n_b(r) = n_b(0) [0.9 \exp(-4r^2/a^2) + 0.1]$. Cases C-F have an explicit treatment of diffusing iron impurities. In all cases the initial iron density is 10^{11}cm^{-3} and is radially uniform. The subsequent iron influx is determined by sputtering as discussed in Section 2. Cases C and D have only the neoclassical terms in the iron diffusion; cases E and F have both the neoclassical and anomalous terms. Cases C and E have only ohmic heating; cases D and F also have neutral beam heating of the same magnitude as in case B.

	I MA	B kG	$n_e(0)$ 10^{14}cm^{-3}	β_θ	τ'_E msec	$T_i(0)$ keV	$T_e(0)$ keV	$q(0)$	$\bar{n}_I$	$n_I(0)$ 10^{11}cm^{-3}
A	1.0	50	0.9	0.3	160	1.8	1.9	1.4	$Z_{\text{eff}} = 2$	
B	1.0	50	0.9	0.5	110	4.8	4.2	1.3	$Z_{\text{eff}} = 2$	
C	1.0	50	1.0	0.2	90	1.5	2.1	0.9	1.3	4.9
D	1.0	50	1.0	0.5	80	4.3	4.0	1.0	1.9	4.1
E	1.0	50	1.0	0.2	100	1.4	1.7	1.3	1.2	4.1
F	1.0	50	1.0	0.5	90	4.3	3.9	1.2	1.7	3.5

Table 2. Results for next-generation tokamak model. Cases C and E have only neoclassical impurity transport; D and F also have anomalous diffusion. Cases A, C, and E have ohmic heating; B, D, and F also have neutral beam heating.

In all cases the current and temperature profiles are quite narrow. Since the impurity density profiles differ substantially, we conclude that the temperature and current profiles are determined by a variety of other effects, namely: the trapping-effect on resistivity, the injection beam profile, and the impurity radiation losses. In Fig. 3 we show the radial distribution of the different iron ionization levels at 200 msec for the cases C and D (neoclassical); Fig. 4 shows the same data for the cases E and F (neoclassical plus anomalous). The primary peak in all the figures results from the central accumulation of the impurities assumed to be present initially. The secondary peak results from the sputtered impurities; the secondary peak is much more pronounced in cases D and F, due to the enhanced energy transport to the wall in the neutral-beam-heated cases. Both primary and secondary peaks are less pronounced in cases E and F, where the anomalous diffusion is included in the impurity transport. Nonetheless, the central accumulation is still considerable, and we conclude that the neoclassical effect might be more evident in PLT than in ST. Although, in our model, all of the sputtered influx is ingested, even in the worst cases (D and F) the total number of iron impurities produced by sputtering during 200 msec is only 0.1% of the protons. At later times, of course, the sputtered iron content is considerably more: by 500 msec in cases D and F it has typically risen to about 0.5%. Furthermore, we have included only the sputtering of the wall material by neutral bombardment, and have neglected the sputtering of the limiter by ion bombardment. At 200 msec in cases D and F the power incident on the wall in neutrals is about 0.7 MW, whereas that on the limiter in ion thermal transport is about 1 MW; this suggests that the limiter

problem might be the more severe, and that the magnetic limiter feature of a divertor might be especially valuable in next-generation tokamaks. Finally, we note that, although the iron content is only in the range 0.1-0.2%, the impurity radiation has assumed a major role in the energetics of all the cases C-F. In case E at 200 msec the ohmic power is 1.3 MW and the total radiation is 0.5 MW; in case F at 200 msec (i.e., just after termination of the 4 MW beam) the total radiation is 0.4 MW.

Fig. 3. Profiles of iron ionization levels for PLT-size model, cases C and D (neoclassical). (PPPL 742141)

Fig. 4. Profiles of iron ionization levels for PLT-size model, cases E and F (neoclassical plus anomalous). (PPPL 742140)

4 IMPURITY CONTROL

In the preceding calculations, no special techniques have been employed to interfere with the "natural" process of impurity-atom injection by charge-exchange sputtering, followed by inward diffusion of impurity ions toward the plasma center. This process could, however, be interrupted at three different points:

1. The tokamak edge heat-outflow can be switched to channels that do not cause impurity injection. This might be done by neutral-gas influx, provided that the mean energy of the resultant charge-exchange neutrals can be depressed below the sputtering threshold of the wall material, or provided that the thermal capacity of the resultant cold plasma can act as a transient absorber of heat from the central region. The latter possibility will be discussed in Section 4.1. Alternatively, the normally dominant channel of heat outflow by plasma transport can be used, but with the limiter element removed to a separately pumped chamber, to which the plasma flows by means of a magnetic divertor. This possibility is discussed in Section 4.2.

2. The impurity influx can be attenuated at the plasma edge. In the absence of the classical impurity transport phenomenon, a large fraction of the shallowly deposited impurity-atom-influx would tend to flow along the steep outward impurity-density gradient, rather than along the gentle inward gradient. This outflow could be encouraged by impairing confinement at the plasma edge, or by creating a substantial plasma edge region that communicates with a divertor chamber (Section 4.2).

3. The flow of impurity ions from the edge region to the plasma interior might be impeded or reversed by several techniques. The use of plasma waves [13], plasma flows [14], and energetic-ion beams [15] has been proposed. If the classical impurity-transport process is dominant, one could simply reverse the hydrogen plasma density gradient toward the plasma edge (e.g. by neutral-gas injection), thus creating a gradual rise of the central hydrogen plasma density, accompanied by an outflux of impurity ions.

Computer studies of some of these impurity-control techniques are discussed briefly below.

4.1 Transient Cold-Plasma Shield

One way to achieve a relatively pure tokamak plasma of reactor temperatures, is to create a hot central region, well removed from wall contact, and "isolated" — at least transiently — by a layer of colder plasma. Computer simulations show that, in sufficiently large plasmas, the cold outer layer effectively blocks the influx of neutral atoms into the central region, even when pulsed gas is used to build up the outer layer itself. It also serves to recapture energetic charge-exchange neutrals from the interior, and to delay the onset of plasma heat transport to the limiter. Finally, it serves to delay the influx of impurities, in the same manner as the outer plasma region in Figs. 3 and 4. (Impurity diffusion is not included explicitly in the computations discussed in this section.)

The cold-plasma shield has been investigated computationally for the large two-component tokamak reactor experiment for which computer studies are reported also in Ref. [16]. The limiter radius is $a_L = 90$ cm at the final plasma position of $R = 270$ cm, where we have $B = 45$ kG. The plasma is formed initially with $a_L = 66$ cm, at $R = 325$ cm, where

$B = 37.5$ kG. For two-component reactor experiments, an initial tritium plasma is neutral-beam heated and then injected with an energetic deuteron beam. Finally, it is compressed in R , away from limiter-contact. In some cases, pulsed neutral-gas injection is then used to densify the outer plasma region.

In the example of Figs. 5 and 6, the initial central plasma parameters are $n = 8 \times 10^{13} \text{ cm}^{-3}$, and $T_e = T_i = 6$ keV, and the boundary values are $n = 1.6 \times 10^{13} \text{ cm}^{-3}$, and $T_e = T_i = 10$ eV. The initial profiles are parabolic for T_e , T_i , and j_T , and cubic for n (our standard assumptions). The discharge current is $I = 1$ MA. A 240 A neutral-beam current at 120 keV is injected from 20 msec to 50 msec. The beam ions are deposited in the profile $n_b(r) = n_b(0)[0.9 \exp(-4r^2/a^2) + 0.1]$. Compression to $R = 270$ cm, and $I = 1.2$ MA takes place at 55 msec. The hot plasma radius a_p shrinks to 60 cm. The precompression boundary density $n = 1.6 \times 10^{13} \text{ cm}^{-3}$ compresses to $n = 2.3 \times 10^{13} \text{ cm}^{-3}$ at $r = a_p$. The limiter for the post-compression plasma is assumed to be at $a_L = 90$ cm. Immediately after compression, we fill up the region $a_p < r < a_L$ with additional plasma; the density falls linearly from $n = 2.3 \times 10^{13} \text{ cm}^{-3}$ at $r = a_p$ to $n = 1.6 \times 10^{13} \text{ cm}^{-3}$ at $r = a_L$, the boundary value; the temperatures fall linearly in a similar way from about 15 eV at $r = a_p$ to the boundary values 10 eV at $r = a_L$. Initially, the blanket region carries no current.

Fig. 5. Transient cold-plasma shield case. Initial 1.0-MA discharge is preheated and injected with energetic deuterons; at 55 msec it is moved to the center of a 90-cm limiter aperture. Radial profiles are at (a) 60 msec; (b) 100 msec. (PPPL 742051)

The central values and radial profiles of the important plasma parameters soon after compression (at 60 msec), and later (at 100 msec), are shown in Figs. 5(a) and (b), respectively. The central plasma parameters are $n = 9.5 \times 10^{13} \text{ cm}^{-3}$, $T_e = 6.8$ keV, and $T_i = 6.3$ keV at 60 msec, and $n = 7.8 \times 10^{13} \text{ cm}^{-3}$, $T_e = 6.6$ keV, and $T_i = 5.5$ keV at 100 msec. Figure 5(a) shows the hot front "blowing out" into the cold plasma blanket. Figure 5(b) shows that by 100 msec the blanket has essentially disappeared, although the plasma density in the edge region is quite high ($\sim 3 \times 10^{13} \text{ cm}^{-3}$). The hot plasma is in the trapped-ion regime.

Fig. 6. Energetics of the transient cold-plasma shield case of Fig. 5. On the vertical scale 1.0 equals 2.16 MJ. (PPPL 742065).

The energetics of this case are shown in Fig. 6. The unit of energy is 2.16 MJ. We imagine a fusion-energy "break-even" experiment [16] to begin just after compression. At that time, the energy in the injected energetic ions is $W_B = 0.88$ MJ, and the total nuclear energy subsequently produced is $W_N = 0.91$ MJ, giving $Q = W_N/W_B = 1.0$. Figure 6 shows a sharp increase in the rate of electron and ion energy transport, and charge-exchange transport, which occurs when the hot front reaches the limiter at about 80 msec. This can be delayed somewhat by densifying the outer plasma with neutral gas, so that the break-even experiment goes essentially to completion before substantial impurity influx can take place.

An interesting application of the cold-plasma-shield technique occurs in the case of pulsed-ignition reactor experiments.

The energy multiplication factor

is maximized by heating a small plasma core to ignition, and letting it "burn out" to the plasma edge, where the reaction is quenched by impurity influx [17].

4.2 Magnetic Divertor

A magnetic divertor channels surface plasma through a scrape-off region into a remote burial chamber. A divertor system can reduce the impurity concentration by having: (1) a high plasma capture efficiency, C_p , to reduce charged particle sputtering, (2) low neutral recycling, R_n , to reduce charge-exchange neutral sputtering, and (3) a high shielding efficiency, S , in the scrape-off region where impurity atoms are ionized and swept into the burial chamber.

A simple model with plasma diffusion across the magnetic field and an effective confinement time $\tau_{||}$ for plasma loss parallel to the field lines into the divertor provides an estimate for the width, λ , of the plasma scrape-off: $\lambda = [D_{\perp}^S \tau_{||} / (1 - n_0 \sigma_i v_e \tau_{||})]^{1/2}$ where n_0 is the neutral density, $D_{\perp}^S$ the diffusion constant in the scrape-off and $\sigma_i v_e$ is the ionization rate. Effective plasma capture requires that the physical width of the divertor entrances, W , should be $\geq 4\lambda$. The recycling due to neutral gas from the divertor is given roughly by $R_n \approx W/f_s L_w$, where f_s is the effective sticking coefficient of neutrals on a gettered surface with a length L_w in the poloidal plane. The recycling has been calculated more accurately for real geometries with a Monte Carlo code.

The shielding efficiency of the divertor scrape-off depends first on the ability of the scrape-off region to ionize incoming wall-impurity atoms. Roughly 99% of the impurities ($A = 100$, $T_z = 5$ eV) are ionized if

$n_p \lambda = 3D_i^p \tau_{ii} (\bar{n}_p/a) (1 - n_0 \sigma_i v_e \tau_{ii})^{-1} > 10^{13} \text{cm}^{-2}$, where D_i^p , $\bar{n}_p$ and a are the diffusion coefficient, average density and radius of the confined plasma, respectively. The resulting impurity ions are swept into the divertor by the plasma flow, due to either frictional or electrostatic coupling.

There are several modes of operation possible for a divertor system; an unload divertor ($C_p \approx 1$, $R_n \approx 0$ and $S \approx 0$), shielding-unload divertor ($C_p \approx 1$, $R_n \approx 0$, and $S \approx 1$), or a gas-fed-shielding divertor ($C_p \approx 1$, $R_n \approx 1$, and $S \approx 1$). In the following, we discuss the characteristics of a large poloidal divertor experiment (PDX), which has $a = 40$ cm, $R = 140$ cm, $I_p \approx 500$ kA, $B_t \approx 25$ kG, $n_e \approx 5 \times 10^{13} \text{cm}^{-3}$, and $T_i \approx T_e \approx 2$ keV, as a typical example of divertor behavior on a large tokamak. An unload divertor can be formed in this device by using a double-outer poloidal divertor (inverse D-shaped plasma). In this case, the scrape-off plasma is rapidly transported into the outer divertors at essentially the sound speed. The estimated parameters for this mode of operation are $C_p \approx 99\%$, $R_n \approx 5\%$ and $S \approx 3\%$, when the cross-field transport is $\approx 10\%$ of Bohm diffusion. A disadvantage of this system is the questionable stability of the inverse-D plasma [18].

A double-inner divertor (D-shaped plasma) may act as a shielding-unload divertor if the plasma transport into the divertors is reduced by mirror confinement of the moderately warm (0.5-1.0 keV) collisionless plasma edge. The mirror confinement causes the plasma scrape-off width and density to increase, thereby decreasing the capture efficiency ($C_p \approx 90\%$) and increasing the recycling ($R_n \approx 5-10\%$) but also increasing the shielding efficiency to near 100% for classical mirror confinement. Velocity-space instabilities may enhance the parallel loss rate for these mirror-confined plasmas, since they tend to be very long ($L/\rho_i \approx 10^3$) and narrow ($\lambda/\rho_i \approx 8$). These instabilities may be reduced by partial filling of the loss cone due to ionization of cold gas injected into the scrape-off. Mirror confinement corresponding to only 10% of classical expectation would still allow significant impurity shielding.

Another possible mode of divertor operation is a combination of the cold plasma shield approach and the magnetic divertor approach. In this case, excess plasma with sufficient density for impurity ionization is formed in the scrape-off, either by additional gas injection or by an auxiliary plasma source in one of the divertor chambers. A large fraction of this excess plasma flows directly into the opposite divertor burial chamber, thereby creating a very efficient impurity sweeping action. This shield plasma also provides a method to reduce or even reverse the classical inward transport by creating a steady-state plasma density peak in the scrape-off region. The steady-state density profile of the main confined plasma could also be altered so as to realize the combination $dn/dr > 0$ and $dT/dr < 0$, thus preventing impurities present at the edge of the confined plasma from diffusing to the center by classical transport. A further benefit of this mode is that the plasma temperature at the surface of the plasma is reduced to ~ 50 eV, thereby reducing the sputtering coefficient significantly. The power required to sustain such a linear shield plasma would be ~ 200 kW in PDX, compared with typical power losses from the confined plasma of $\sim 300-500$ kW.

Some aspects of PDX divertor operation have been modeled with the transport code. A "divertorless" PDX discharge with $R_n = 1.0$ and Z_{eff} chosen to be ≈ 8 is compared in Table 3 with an "unload divertor" discharge simulated by choosing $R_n \approx 0.2$ and $Z_{\text{eff}} = 1$. In this case we have

added 1 MW of auxiliary heating, shared equally between ions and electrons, and distributed over a parabolic profile. The plasma parameters are quoted at 200 msec; the plasma temperatures are still rising strongly at this time, especially in the divertor case.

R	Z	$n_e(0)$ 10^{13} cm^{-3}	$T_i(0)$ keV	$T_e(0)$ keV	W_J kJ	W_{aux} kJ	W_{CX} kJ	W_P kJ
1.0	8	4.9	2.7	1.8	440	300	90	180
0.2	1	2.4	3.9	3.8	60	300	10	150

Table 3.

Note that the energy input due to ohmic heating, W_J , is much larger in the divertorless case, due to the large Z_{eff} . The amount of energy lost by charge exchange, W_{CX} , and plasma flow to the limiter, W_P , is 270 kJ in the divertorless case: resultant sputtered atoms may reenter the plasma. This is to be compared with only ~10 kJ deposited in the vicinity of the discharge in the "unload divertor" case. An estimate for the reduction in impurity generation requires a knowledge of the energy distribution of the escaping particles. However, assuming that the amount of sputtering per unit of incident particle energy does not change appreciably when R_n is reduced from 1.0 to 0.2, then even a rather poor "unload divertor" would reduce the impurity generation by at least one order of magnitude.

The effects of divertor boundary conditions have also been incorporated into the transport code. In these cases, the edge values of plasma density, ion temperature, and electron temperature (depending on the assumption regarding secondary electron emission at the plate) are allowed to float, being determined by particle and energy transport processes in the scrape-off region. For a typical case, simulating the early phases of an "unload divertor" discharge with parallel ion sound flow and six-regime cross-field transport, the edge values relative to the central values were $n_e(a) \approx 5 \times 10^{11} \text{ cm}^{-3} \approx 10^{-2} n_e(0)$, $T_i(a) \approx 330 \text{ eV} \approx 0.5 T_i(0)$, with $T_e(a)$ held constant at 200 eV and $R_n \approx 0.15$. This illustrates one of the principal effects of an unload divertor, that is, the reduction of charge-exchange energy losses, which allows the ion temperature profile to rise and broaden. These rough estimates indicate that the impurity input during the quasi-static discharge can be reduced by ~10 with either an unload or unload-shielding divertor. A gas-fed-shielding divertor may allow the impurity level to be reduced even further.

In addition to impurity control during the quasi-static discharge, a poloidal divertor system can control impurity generation during plasma current start-up by: (1) forcing the plasma to form a multipole null in the poloidal field located away from the limiters; (2) ensuring that the plasma scrape-off enters the divertor throats while the plasma current rises; and (3) arranging the plasma-current cross sectional area inside the divertor separatrix to increase linearly with the plasma current. This latter function of a divertor, an expanding magnetic limiter, reduces the skin effect and can be expected to reduce the associated turbulence problems, which in large tokamaks are expected to cause large amounts of energy to be deposited on limiters, with associated impurity production during plasma start-up.

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