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NEOCLASSICAL IMPURITY FLUX IN THE PRESENCE OF ANOMALOUS TRANSPORT

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Abstract

When there are fluctuations, the usual argument for automatic ambipolarity no longer applies, because the fluctuations contribute to the mean momentum balance. The ambipolar electric field E_r is now determined by ambipolarity of the neoclassical and anomalous particle fluxes. For electrostatic fluctuations the fluxes are only weakly non-ambipolar, because of the ion drift driven by their mean inertia and pressure tensor. $E_r - B_\theta U_\parallel$ is therefore close to the neoclassical prediction. However, the flow along a stochastic magnetic field is strongly non-ambipolar, and can change E_r from inwards to outwards. Impurity neoclassical fluxes are most sensitive to the electric field, because of their high Z, and the abrupt change in E_r at the onset of stochasticity can cause a rapid pump-out of impurities. The observed reduction in impurity density during periods of MHD activity, such as sawteeth or ELM's, agrees qualitatively with analytic prediction.

1. INTRODUCTION

It has frequently been stated [1-3] that neoclassical transport is automatically ambipolar, in the sense that $\sum Z_a \Gamma_a = 0$ follows from conservation of momentum during collisions, where Z_a is the charge on the ath species and Γ_a its mean radial flux. In a neoclassical plasma it is therefore unnecessary to invoke quasi-neutrality as a justification for ambipolarity. It will be shown in Sec. 2 that the usual demonstration of automatic ambipolarity no longer holds when fluctuations are present. The case of electrostatic fluctuations will be considered in detail. Here the off-diagonal component of the pressure tensor has a non-zero average value, which must be included in the mean momentum balance. This upsets the simple relation between neoclassical diffusion and collisional friction.

Whether or not neoclassical transport is automatically ambipolar has some important consequences. If it is, then the anomalous transport must also be ambipolar, to satisfy quasi-neutrality. If it is not, then quasi-neutrality requires only that the total fluxes are ambipolar. As discussed in Sec. 2, electrostatic fluctuations produce a slightly non-ambipolar flux, because the ion inertia and the off-diagonal components of the pressure tensor give rise to a difference in the mean cross-field flux of ions and electrons. This flux imbalance is small but, if it could not be balanced by an opposite neoclassical charge flux, it would result in a rapid build up of large radial electric fields. Particle flow along a stochastic magnetic field can be strongly non-ambipolar, requiring a large change in the neoclassical flux to balance it.

Sec. 3 discusses the way the radial electric field E_r adjusts itself, so that the charge imbalance in the anomalous transport is compensated by the neoclassical fluxes. Because impurities are particularly sensitive to changes in E_r, an impure plasma is considered. The result of Conner's analysis [4] for neoclassical impurity transport neglecting fluctuations will first be summarised, since much of this analysis is still valid when fluctuations are present. It is only when momentum and particle fluxes are averaged over magnetic surfaces that the neoclassical and anomalous parts must be combined.

Only a small change in $E_r - B_\theta \overline{U}_{//}$ is required for the neoclassical fluxes to compensate the small imbalance in the anomalous transport produced by electrostatic fluctuations. Magnetic fluctuations can have a much larger effect. In a stochastic magnetic field the capability of electrons to escape along field lines is much greater than that of ions, because of their larger thermal velocity. However, because it was previously thought that the anomalous transport must be ambipolar, it was concluded that the escaping electrons are held back by the ions. Consequently, a stochastic magnetic field could produce significant electron thermal flux, but very little particle flux. Since the ambipolar condition is now applied only to the total flux, the potentially large ion neoclassical flux can be balanced by a large electron flow along a stochastic field.

Neoclassical impurity fluxes are most strongly affected by E_r, because of their larger atomic charge. The comparison between prediction and experiment in Sec. 4 thus concentrates on impurities. Much of the experimental impurity behavior can be explained qualitatively by the following analysis. During periods when MHD activity is low, impurities tend to accumulate in the centre. When MHD activity is strong, such as in disruptions or sawtooth crashes, a rapid loss of impurity is observed. This is consistent with the electron flow along the stochastic field shorting out the inward radial electric field, resulting in a large outward neoclassical impurity flux. A more complete comparison with experiment, together with more detailed analysis, is given in reference 5.

2. IS NEOCLASSICAL TRANSPORT AUTOMATICALLY AMBIPOLAR?

(a) The basic equations

The usual derivation of automatic ambipolarity will now be outlined, but including the effect of the electrostatic fluctuations which are always present in real plasmas. We consider the simplest model of a tokamak, in which flux surfaces are circular and concentric. The derivation starts from the generalised fluid momentum balance equation for the a^{th} species

$$n_a m_a \frac{d\underline{u}_a}{dt} = n_a e_a\left(\underline{E} + \underline{u}_a \times \underline{B}\right) - \nabla \cdot \underline{\underline{P}}_a + \underline{R}_a \tag{1}$$

where $\underline{R}_a$ is the collisional friction with other species. Effects arising from the particle velocity distribution are hidden in the pressure tensor $\underline{\underline{P}}_a$.

The variation in density over a flux surface, r = const, can be written in the form

$$n_a(r,\theta,\varphi) = \bar{n}_a(r) + \tilde{n}_a(r,\theta) + \sum_{m,s} n_{ams}(r)\cos\left(m\theta + s\varphi - \omega t + \alpha_{ms}\right) \tag{2}$$

where $\bar{n}_a$ is the mean, $\tilde{n}_a$ is the neoclassical variation, which has the form $\tilde{n}_{ac}\cos\theta + \tilde{n}_{as}\sin\theta$, and the last term is the sum over fluctuations. The electrostatic potential variation, Φ, has the same form, except that the phase constant in the fluctuations will be written as β_{ms}.

(b) The pressure tensor

The variation in the pressure tensor can be separated into its neoclassical and fluctuation parts

$$\underline{\underline{P}} = \underline{\underline{P}}^N + \underline{\underline{P}}^F$$

When studying LF electrostatic fluctuations, the pressure tensor is expressed as a scalar part plus a tensorial part resulting from finite Larmor radius (FLR) effects

$$\underline{\underline{P}}^F = p^F \underline{\underline{I}} + \underline{\underline{\pi}}$$

For drift waves, $k_\perp \rho_a$ ~ 0.1-0.3 both theoretically and experimentally, where $\rho_a = (2T_a m_a)^{1/2} / e_a B$ is the Larmor radius, so FLR effects are expected to be significant. Such effects should be completely negligible for the neoclassical variation, since here the spatial scale is the minor radius, and $\rho_a/a << 1$. However, parallel viscosity is important and, when studying neoclassical transport, the pressure tensor is normally written as a diagonal tensor with unequal parallel and perpendicular temperatures

$$\underline{\underline{P}}^N = p_\perp \underline{\underline{I}} + (p_\parallel - p_\perp)\, \underline{b}\, \underline{b}, \quad where\ \underline{b} = \frac{\underline{B}}{|B|} .$$

(c) φ-component of momentum balance

We will now average the φ-momentum equation for the ath species over a magnetic surface. The φ-component of the pressure tensor, including the effect of the magnetic field curvature, is

$$\left(\nabla \cdot \underline{\underline{P}}\right)_\varphi = \frac{B B_\theta}{r} \frac{\partial}{\partial \theta}\left(\frac{p_\parallel - p_\perp}{B^2}\right) + \frac{1}{R}\frac{\partial p^F}{\partial \varphi} + \left(\nabla \cdot \underline{\underline{\pi}}\right)_\varphi \tag{3}$$

Integrating over φ immediately eliminates $\partial p^F/\partial \varphi$. For the simple model used for the tokamak field

$$B = \frac{B_o}{h(\theta)}, \quad B_\theta = \frac{B_{\theta o}}{h(\theta)}, \quad where\ h(\theta) = 1 + \epsilon \cos \theta, \quad \epsilon = \frac{r}{R}.$$

Multiplying the momentum equation by h^2 makes the $(p_\parallel - p_\perp)$ term an exact differential, and it vanishes in the subsequent θ-integration. We thus obtain the following flux-surface-averaged φ-momentum equation

$$m_a \left\langle n_a \frac{du_{a\varphi}}{dt} h^2 \right\rangle = e_a \Gamma_a B_{\theta o} + e_a \left\langle n_a E^A h^2 \right\rangle + \left\langle R_{a\varphi} h^2 \right\rangle$$

$$+ \frac{e_a}{2} \sum_{m,s} s n_{ams} \Phi_{ms} \sin(\beta_{ms} - \alpha_{ms}) - \left\langle \left(\nabla \cdot \underline{\underline{\pi}}_a\right)_\varphi h^2 \right\rangle \tag{4}$$

where E^A is an externally induced electric field in the φ-direction. It produces the Ware pinch but, since it does not affect the following arguments, it will henceforth be omitted for brevity. $\Gamma_a = \langle n_a u_{ar} h\rangle$ is the average particle flux across the surface, where

$$\langle A\rangle = \frac{1}{4\pi^2}\oint d\theta \oint d\varphi\, A.$$

(d) Ambipolarity

The final step is to separate off the classical and anomalous fluxes from Γ_a leaving the neoclassical flux. The anomalous and classical fluxes are readily obtained from $\underline{B} \times \text{Force}/e_a B^2$, giving

$$\Gamma_a^N = \Gamma_a - \frac{\left\langle h^2 \underline{R}_a \times \underline{B}\right\rangle_r}{e_a B^2} - \sum_{ms} \frac{\left\langle n_{ams} \underline{E}_{ms} \times \underline{B}\right\rangle_r}{B^2} + \frac{\left\langle \nabla \cdot \underline{\underline{\pi}}_a \times B\right\rangle_r}{e_a B^2} \qquad (5)$$

$$-\sum_{ms}\left\langle \frac{m_a n_{ams}}{e_a B^2} \frac{d\underline{u}_{ams}}{dt} \times \underline{B}\right\rangle_r$$

We substitute for Γ_a in Eq. (4), and write certain φ components in the form $R_{a\varphi} = (R_{a\parallel}B_\varphi + R_{a\perp}B_\theta)/B$, where ⊥ denotes the direction in the magnetic surface perpendicular to $\underline{B}$. The ⊥ components cancel, leaving

$$m_a \sum_{ms} \left\langle n_{ams} \frac{du_{ams//}}{dt}\right\rangle = e_a \Gamma_a^N B_{\theta o} + \left\langle R_{a\parallel} h^2\right\rangle$$

$$-\frac{e_a}{2B}\sum_{ms} \underline{k}\cdot\underline{B} n_{ams} \Phi_{ms} \sin(\beta_{ms} - \alpha_{ms}) - \left\langle \left(\nabla \cdot \underline{\underline{\pi}}_a\right)_{\parallel}\right\rangle \qquad (6)$$

If there are no fluctuations, and a steady equilibrium is assumed, then

$$e_a B_{\theta o} \Gamma_a^N = -\left\langle R_{a\parallel} h^2\right\rangle$$

Now sum this equation over all species (including electrons). Since $\Sigma R_{a\parallel} = 0$ from momentum conservation during collisions, one finds $\Sigma e_a \Gamma_a^N = 0$, i.e. ambipolarity occurs naturally without having to impose it as a condition [1-3].

When fluctuations are included, the third term on the right vanishes in the summation, since quasi-neutrality is satisfied for each wave-component ($\Sigma_a n_{ams} e_a = 0$). The mean pressure tensor resulting from the fluctuations does not vanish, however, giving

$$\Sigma\ e_a \Gamma_a^N = \frac{1}{B_{\theta o}} \sum_a \left[\left\langle (\nabla \cdot \pi_a)_{\parallel} \right\rangle + m_a \sum_{ms} \left\langle n_{ams} \frac{du}{dt} ams_{\parallel} \right\rangle \right] \tag{7}$$

The neoclassical transport is therefore not ambipolar when fluctuations are present.

(e) The FLR pressure tensor

Since this pressure tensor is important for the above conclusion, I will discuss it briefly. Its components are set out, for example, by SI Braginskii [6]. A typical component of $(\nabla \cdot \underline{\underline{\pi}})_\varphi$ is

$$\frac{\partial}{\partial r}\left[\frac{\tilde{n}_a T_a}{\Omega_a r} \frac{\partial \tilde{u}_\varphi}{\partial \theta} \right] \tag{8}$$

Where the tilde denotes the sum over the fluctuating waves. Eq. (8) is a nonlinear term whose flux-surface-average does not vanish. The physical origin of the pressure tensor becomes clear only after detailed analysis. For example, if the fluid velocity is separated into its guiding centre and diamagnetic parts, $\underline{u} = \underline{v} + \underline{w}$, then some of $\nabla \cdot \underline{\underline{\pi}}$ cancels the diamagnetic part of the convective inertia. Hence $(\partial / \partial t + \underline{u} \cdot \nabla)\underline{u}$ becomes $(\partial/\partial t + \underline{v} \cdot \nabla)\ \underline{v}$. The physically correct convective derivations should move with the mean velocity of a single particle, i.e. the guiding centre velocity. Since the fluid momentum equation contains the fluid convective derivative, it must include a correction to cancel the diamagnetic part. Some of the pressure tensor expresses a real physical effect, i.e. the ion averages the electric field over its fast gyration. Thus the effective electric field seen by an ion is $\left[1 + \rho_i^2 \nabla_\perp^2\right]\underline{E}$. This leads to a small difference in the ion and electron cross field fluxes.

(f). Non-ambipolarity of Electrostatic Anomalous Transport

The local cross-field flux resulting from the electrostatic fluctuations may be obtained from the vector product of the momentum equation and $\underline{B}$. Integrating over a flux surface gives

$$\Gamma_a^{an} = \left\langle n_a \frac{\underline{E} \times \underline{B}}{B^2} - \frac{\nabla p_a \times \underline{B}}{e_a B^2} - \frac{\nabla \cdot \underline{\underline{\pi}}_a \times \underline{B}}{e_a B^2} - \frac{m_a n_a}{e_a B^2} \frac{d\underline{u}_a}{dt} \times B \right\rangle_r \tag{9}$$

The second term on the right vanishes in the integration, since it is an exact differential. When the charge flux is summed over species, the first term vanishes, leaving

$$\Sigma_a \; e_a \Gamma_a^{an} = -\Sigma \left\langle \frac{\nabla \cdot \underline{\underline{\pi}}_a \times \underline{B}}{B^2} - \frac{m_a n_a}{e_a B^2} \frac{d\underline{u}_a}{dt} \times \underline{B} \right\rangle_r \tag{10}$$

Thus the same terms which contributed to the mean momentum balance in Eq. (6-7) now give rise to a non-ambipolar flux. These terms are small, being of order $k^2\rho_a^2$ and $\omega / \Omega_a \sim k\rho_a^2 / L_n$ where $\Omega_a = e_a B / m_a$ is the cyclotron frequency and L_n the density scale length. Typically the observed density fluctuations have $k\rho_a \sim 0.2$, so the net charge flux is of order 1% of the total flux. The effect of magnetic fluctuations can be much larger, but the electrostatic fluctuations have been considered first because their analysis is more straight forward.

(g) The Effect of Magnetic Fluctuations

The analysis for magnetic fluctuations follows similar lines to that for the electrostatic. The analogue of Eq. (4) now includes the term $\langle \tilde{j}_{ar} \tilde{B}_\theta - \tilde{j}_{a\theta} \tilde{B}_r \rangle$, where $j_a = n_a e_a \underline{v}_a$. This again upsets the simple relation between neoclassical flux and collisional friction and invalidates the argument for automatic ambipolarity. In Eq. (10) the stress tensor term is replaced by $\langle \tilde{j}_{\parallel} \tilde{B}_r \rangle$.

3. EVALUATION OF THE NEOCLASSICAL IMPURITY TRANSPORT

(a) Connor's analysis

In the first neoclassical analysis of impure plasma, Connor [4] solved the kinetic equation, assuming all species to be in the banana regime, and no fluctuations. His results will first be briefly summarised, before generalising then to include fluctuations. He obtained the particle flux of each species to be

$$\Gamma_a^N = 1.46\,\epsilon^{1/2}\, n_a \rho_{a\theta}^2 \bar{\upsilon}_a \left[-\frac{n_a'}{n_a} + \gamma_a \frac{T_a'}{T_a} + \frac{e_a}{T_a}\left\{ E_r - B_\theta \frac{\sum_b \bar{\upsilon}_{ab} U_b}{\bar{\upsilon}_a} \right\} \right] \qquad (11)$$

where $\bar{\upsilon}_{ab}$ is the collision frequency of species a on species b, $\bar{\upsilon}_a = \Sigma_b\ \bar{\upsilon}_{ab}$, and $\rho_{a\theta}$ is the Larmor radius in the poloidal magnetic field. The coefficient γ_a is a function of collision frequencies. Its value varies between 0.17 and 0.5, depending on the impurity content. U_b is approximately the parallel velocity of the b^th^ species. Assuming all ions to have similar parallel velocities, the last bracket will be written as $(E_r - B_\theta U_\parallel)$.

Connor eliminated E_r using $\Sigma e_a \Gamma_a = 0$. For the case of a main light ion and one dominant heavy impurity, denoted by subscripts 1 and 2 respectively, the impurity flux is

$$\Gamma_2^{NO} = 1.46\,\epsilon^{1/2} \frac{\rho_{2\theta}^2}{m_2}\left(\frac{n_1 m_1 \bar{\upsilon}_1 n_2 m_2 \bar{\upsilon}_2}{n_1 m_1 \bar{\upsilon}_1 + n_2 m_2 \bar{\upsilon}_2} \right)\left[-\frac{n_2'}{n_2} + \gamma_2 \frac{T_2'}{T_2} + \frac{T_1}{T_2}\frac{Z_2}{Z_1}\left(\frac{n_1'}{n_1} - \gamma_1 \frac{T_1'}{T_1} \right) \right] \qquad (12)$$

where $Z_a e$ is the ion charge. The first term in the second bracket is diffusive. For the normal case where the main ion density peaks on axis, the n_1'/n_1 term gives an inward convection. If $Z_2 >> Z_1$, this inward pinch is large, leading to strong impurity peaking on axis.

(b) Extension of Connor's flux to other regimes

The neoclassical impurity flux in all collisional regimes may be written in the general form

$$\Gamma_a^N = n_a D_a \left[-\frac{n_a'}{n_a} + \gamma_a \frac{T_a'}{T_a} + \frac{Z_a e}{T_a}\left(E_r - B_\theta U_\parallel \right) \right] \qquad (13)$$

The origin of the rather complex variation in the ambipolar fluxes when species are in different collisional regimes is clearer when the ambipolar condition is imposed only as a last step. The ambipolar equation is dominated by the species with the largest value of $Z_a^2 D_a$. Figure 1 illustrates a typical variation of $Z_a^2 D_a$ with collisionality for the main ions and a dominant impurity. Because of their large Z, the impurity passes into the plateau regime, and possibly the Pfirsch-Schlüter regime, while the main ions are still in the banana regime.

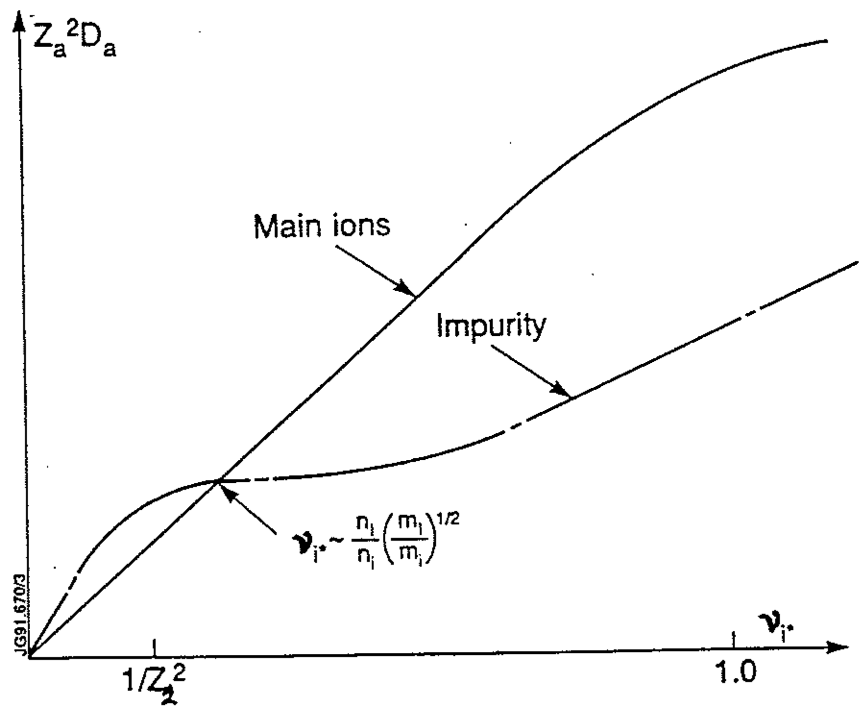

Fig. 1. Variation of the diffusion coefficient with collision frequency for main plasma and impurity ions

Ambipolarity forces E_r to adjust itself to reduce the flux of whichever species has the largest $Z_a^2 D_a$, until it is balanced by the other species, i.e.

$$\frac{Z_a e}{T_a}\left(E_r - B_\theta U_{//}\right) \approx \frac{n_a'}{n_a} - \gamma_a \frac{T_a'}{T_a} \qquad (14)$$

Substituting this E_r in Γ^N for the other species gives their actual flux.

(c) Inclusion of fluctuations in Connor's analysis

The velocity distribution can now be written, to first order, in analogous form to that used for the fluid equations

$$f(r,\theta,\underline{v}) = \bar{f}(r,\underline{v}) + \tilde{f}(r,\theta,\underline{v}) + \sum_{ms} f_{ms}(r,\theta,\underline{v})\cos\left(m\theta + n\varphi - \omega t + \alpha_{ms}\right) \qquad (15)$$

The kinetic equation can be linearised in $\in$ and $f_{ms}/\bar{f}$, and solved separately for the first order neoclassical variation and fluctuations. In the evaluation of the second order flux.

$$\Gamma_a = \frac{1}{S}\int dS \int d^3 v f(r,\theta,\underline{v}) V_r$$

where V_r is the radial guiding centre drift, only products of like terms survive.

The two effects first interact when ambipolarity is imposed

$$\sum_a e_a \Gamma_a^N(E_r) + e\Gamma_s^{an} = 0 \qquad (16)$$

where $e\Gamma_s^{an} = \sum_a e_a \Gamma_a^{an}$ is the non-ambipolar anomalous flux. This determines E_r.

(d) Effect of magnetic fluctuations

Provided the impurity content is not too large, the ambipolar electric field is determined primarily by the balance between the stochastic electron loss and the main ion neoclassical flux, denoted by subscript i.

$$1.46 \in^{1/2} \bar{\upsilon}_i \rho_{i\theta}^2 \left[-\frac{dn_i}{dr} + n_i \left\{ \gamma_i \frac{T_i'}{T_i} + \frac{e}{T_i}\left(E_r - B_\theta U_{//}\right) \right\} \right] = -\Gamma_e^{stoch.}(E_r) \qquad (17)$$

The electron loss along the stochastic magnetic field lines is also a function of E_r. If the magnetic field is highly stochastic, an outward E_r is required to hold back the electrons. This gives a large increase, by a factor of order $(m_i/m_e)^{1/2}$, in the main ion loss rate compared with the pure neoclassical case, where the ambipolar rate is determined by the slower diffusing electrons. This resolves the old question how magnetic ergodicity can significantly increase the ambipolar particle loss, since the ion parallel flow is so small. The parallel electron loss can be balanced by the potentially large ion neoclassical loss.

4. COMPARISON WITH EXPERIMENT

(a) Prediction

If the anomalous transport is ambipolar, or nearly so, the ambipolar electric field is determined by the neoclassical fluxes. As discussed in Section 3b, this is usually given by Eq. (14), evaluated for the main ions. Neglecting for simplicity the effect of temperature gradient, the impurity flux may then be written as.

$$\Gamma_a = n_a D_N \left[-\frac{n_a'}{n_a} + Z_a \frac{n_i'}{n_i} \right] - D_{an}\left[n_a' + K n_a \frac{r}{b^2} \right] \qquad (18)$$

Where subscript i denotes the main ion species (assumed to be singly charged) and subscript a the impurity. Z_a is the impurity ionic charge and b the plasma radius. The first bracket is the neoclassical flux given, for example, in Eq. (12), and the second bracket the anomalous flux, including a finch term. In the plasma core, the impurity profile evolves towards a steady state where $\Gamma_a = o$, i.e.

$$\frac{n_a(r)}{n_a(o)} = \left[\frac{n_i(r)}{n_i(o)}\right]^{\alpha} \exp\left[-\frac{K\ D_{an}}{2(D_{an} + D_N)}\ \frac{r^2}{b^2}\right] \tag{19}$$

where $\alpha = Z_a D_N/(D_N + D_{an})$. Typically D_{an} is on order of magnitude greater than D_N, making α of order unity. If the same anomalous transport law applies to the main ions, as would be expected for electrostatic fluctuations, then the steady state main ion profile in the core approximates to $n_i(o) \exp[-Kr^2/2b^2]$. Hence $n_a(r)/n_a(o) \approx [n_i(r)/n_i(o)]^{\alpha+1}$. Thus the steady state impurity profile is more peaked than the main ions, though not nearly as much as when anomalous transport is neglected ($D_{an} = o$).

When the anomalous transport is not ambipolar, the ambipolar electric field is determined by the balance between anomalous and neoclassical fluxes. If the electrons have the faster escape rate, as in a stochastic magnetic field, the inward electric field is reduced or reversed, as discussed in Section 3d. The Zn_i' / n_i in Eq. (18) is then replaced by ZeE_r/T_a, the steady state impurity profile becomes

$$\frac{n_{a(r)}}{n_{a(o)}} \approx \frac{n_{i(r)}}{n_{i(o)}} exp\left[-\alpha\frac{e}{T_e}\{\phi(r) - \phi(o)\}\right]$$

where $\phi(r) = -\int E_r dr$. If the stochasticity is strong enough to reverse E_r, the exponential factor increases with radius. The steady state profile is then either less peaked than the main ion profile, or may even increase with radius.

In practice short periods of intense MHD activity are frequently observed, separated by longer quiescent periods. The dominant loss mechanism during quiescent periods is presumably electrostatic fluctuations and, since these are nearly ambipolar, impurities are expected to accumulate steadily in the centre, the profile tending towards Eq. (19). Assuming the MHD activity to be strong enough to stochasticise the field, it should change the impurity flux to outwards, as the profile moves towards Eq.(20). The changes in E_r also affects the main ion neoclassical transport, increasing the loss rate. However, as may be seen from Eq. (13), the effect on the impurity is much stronger, because E_r is multiplied by

Z. Thus during MHD activity we expect to see a significant outflux of impurities from the centre, with a much smaller drop in the ion and electron densities.

(b) Experiment

Recent measurements of space charge potential on TEXT [8] support the predicted reversal in E_r by magnetic stochasticity. In most conditions the measured potential profile agreed reasonably with neoclassical prediction, consistent with the magnetic fluctuation being low. When operated are at relatively low values of $\bar{n}_e$ B_0 or I_p, however, the magnetic fluctuations became larger, reaching an amplitude at which magnetic stochasticity was expected in the edge plasma. In these same conditions E_r became positive over an edge region.

There are many examples where impurity accumulated on axis during MHD quiescent phases, and was expelled during MHD activity, but space permits mention of only a few. In early experiments on T-4, the central impurity density increased monotonically, but then dropped abruptly when a kink instability occurred [9]. During a disruptive instability, the central impurity density dropped by an order of magnitude, while the control electron density fell by only 5%.

A bimodal behavior, in which nominally identical discharges showed major differences in their impurity behavior, was observed in several tokamaks. For example, the more normal discharges (type S) in D-III developed sawtoothing early, and no impurity peaking occurred thereafter [10]. In type O discharges, however, there was no central MHD activity and the impurity density peaked sharply on axis. Later a large m = 1, n = 1 oscillation grew rapidly, the central impurity fell, and the discharge reverted to type S. Which mode developed depended on the early plasma-wall interaction, which influenced the initial current profile and hence the occurrence of sawteeth.

In Alcator C [11] and several other tokamaks, impurities were found to accumulate on axis between sawteeth and be ejected during each sawtooth crash. Following pellet injection in Alcator C [12] impurities accumulated on axis for about 40 msec, until they were abruptly ejected in a giant impurity disruption (GID). During a GID the central carbon density dropped by two thirds, while electron density and temperature decreased by only about 10%. A further example of impurity ejection is the ELM. In an ELM the magnetic fluctuations are localised near the plasma edge, and the impurities are ejected from the same region, without much change in the central impurity.

All the above impurity behavior is qualitatively consistent with the foregoing prediction.

5. CONCLUSIONS

1. The usual demonstration that neoclassical transport is automatically ambipolar is invalidated by the mean momentum produced by fluctuations, either electrostatic or magnetic.
2. The ion drift due to inertia and the non-diagonal pressure tensor makes the transport driven by electrostatic fluctuations weakly non-ambipolar. Electron flow along a stochastic magnetic field can give a strongly non-ambipolar flux.
3. An ambipolar electric field develops such that the non-ambipolar anomalous flux is balanced by non-ambipolar neoclassical flux. E_r is not much changed by electrostatic fluctuations, but it may be reversed by stochasticity.
4. Because of their high Z, the impurity neoclassical flux is more strongly affected by this change in E_r. Thus the onset of stochasticity is expected to produce pump-out of impurities with a much weaker reduction in the main plasma density.
5. The abrupt decrease in central impurity during sawtooth, and other strong MHD activity, and the decrease in edge impurity density during ELMs, agree qualitatively with these predictions.

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Appendix I

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