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MOMENT-EQUATION APPROACH TO IMPURITY
TRANSPORT IN TOKAMAKS

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CHAPTER 1

INTRODUCTION

As tokamaks are developed to commercial status, one of the most important processes requiring understanding is the movement of non-hydrogen particles inside the plasma. These impurities with their associated electrons take up some of the plasma pressure, reducing the system efficiency. More importantly, these ions cause enhanced radiation which, if controlled, could help to control the burn. If uncontrolled, the impurities could quench the burn, possibly barring the tokamak from becoming a viable commercial power reactor. The purpose of this report is to overview neo-classical impurity transport through a fluid theory description, and experiments relevant to the topic.

In a magnetic field, ions are tied to the field lines to within a banana width. These field lines form flux surfaces of constant pressure and density in a tokamak. With no field irregularities, the only mechanism for particle transport across the flux surfaces is by Coulomb collisions between different species, like species collisions having no net effect to lowest order. In general, transport is caused by momentum exchange, collisions being the typical form. Since ion momentum exchange due to impurities is greater than ion momentum exchange due to electrons by a factor of the charge state squared and the square root of the mass of the impurity divided by the mass of the electron, the ion-electron collisions can be ignored when the impurity concentration times the charge state squared is on the order of the main ion concentration or greater.

Three theoretical forms for particle transport are discussed here. The distinctions among these are mainly historical, as all can be included within the same framework, although the resulting transport effects differ greatly. The three, in order of historical development, are classical, Pfirsch-Schlüter, which introduced geometric effects, and neoclassical, which included long mean free path effects.

The first of these, classical transport, is dependent on a particle on one flux surface colliding with a particle on a nearby flux surface. Since the particles are tied closely to the flux surfaces, the movement in each collision is small, and therefore the net particle flux is small. The first mention of collisionally driven impurity transport was given by Longmire and Rosenbluth in 1956 [1]. They found that classical transport predicted impurities being driven up main ion gradients until a central impurity peak developed. The final impurity peak was predicted to be steeper than the main ion gradient by a factor of the impurity charge state. This is due to the impurity gyro-radius being the charge state times larger than the main ion gyro-radius. A review by Braginskii [2] in 1965 summarizes classical transport.

Classical transport did not take into account the magnetic field gradients and curvature, and therefore ignored the inexact cancelation of the magnetic gradient drift. This correction leads to Pfirsch-Schlüter transport [3]. Pfirsch-Schlüter transport is dependent on collisions between particles on the same flux surface, and is driven only by poloidal flows. In general, Pfirsch-Schlüter transport leads to an enhancement of classical transport by a factor $(1+2q^2)$, where q is the safety factor, or inverse rotational transform, and varies in value from ~ 1 in the center of the plasma to $\sim 3-6$ at the outer boundary.

Pfirsch-Schlüter theory was applied to impurity transport by Rutherford [4]. With the inclusion of first order corrections to the particle and heat friction coefficients, the possibility arose that the normally negative temperature gradients could inhibit or reverse the inward impurity flow. The case studied by Rutherford was for an impurity mass much larger than the main ion mass, for which this temperature screening effect is not predicted. Boley, Gelbard, and Stacey [5] extended this treatment for an arbitrary number of species.

The most recently developed theory in this line is neoclassical transport, which takes into account the magnetic trapping of particles in a plasma in which the time between collisions is long compared with the time required for the particle to complete a poloidal transit. With long mean free paths, some of the particles become trapped in "banana" orbits. This in-

creases the diffusion step size with each collision, so that this mechanism can greatly enhance particle transport. Connor [6] and Hinton and Moore [7] applied neoclassical transport to impurity transport and predicted as before that the main ion density gradient would drive impurities inward, but at the larger, neoclassical rate. The neoclassical diffusion rate is enhanced over the Pfirsch-Schlüter rate by a factor of $(R/r)^{3/2}$, where R is the major radius and r is the minor radius.

To calculate the impurity transport in the Pfirsch-Schlüter regime, Rutherford had extended the fluid, or moment equation approach of Braginskii [2] to obtain his results for two species of disparate mass, since the geometry effects do not change the collisional friction or viscosity effects. Trapped particles, on the other hand, do substantially change the plasma viscosity. Therefore, until recently, the neoclassical, or long mean free path regime, could not be treated as a fluid. The solution for the neoclassical transport was then restricted to the solution of Föcker-Plank equations in velocity space with the choice of a tractable approximation of the collision operator. This limitation was overcome when Hirshman [8,9] derived an averaged parallel viscosity coefficient for a neoclassical plasma, which placed the long mean free path effects calculated kinetically into the moment equation framework.

A concept inherent in the calculation of the impurity fluxes is ambipolarity, or the principal of detailed balance. This was introduced by Hazeltine and Ware [10], who proved that the diffusion could not set up gross charge imbalance. The requirement indicates that a flux of one species must be offset by opposing fluxes of other species, so that the net charge flux is zero. Hirshman [8,9] later proved that collisional impurities would hold the magnitude of the differential transport of both species to the magnitude of the transport with all species collisional, regardless of the collisionality of the other species. Samain [11] had also found this for large concentrations of impurities.

To this point, the only effects considered in calculating ion-impurity transport was internal, collisional momentum exchange. Ohkawa [12] suggested that a source of external momentum, such as a neutral beam, could reverse the inward flow of impurities driven by gradients. El Derini

and Emmert [13] produced a theory based on adding terms to account for the direct momentum effects of a beam. Stacey and Sigmar [14,15] included the indirect effects of the momentum source and an external drag on the flows within the plasma, as well as the direct effects, to develop a consistent theory for transport in the presence of external momentum exchange. This subsequently led to an expression for radial impurity transport driven by a momentum source.

Ohkawa [16] also suggested the alteration of radial transport by the introduction of a poloidally asymmetric particle source. This theory was then developed by Burrell [17] and Wong [18] for a collisional two species and multiple species plasma, respectively. Chu and Rawls [19] obtained similar results through a kinetic approach in the Pfirsch-Schlüter and plateau regimes. With a fluid model in the Pfirsch-Schlüter regime, Stacey [20] generalized the theory to include both momentum and particle sources simultaneously.

Further extensions to the fluid treatment of impurity transport theory including external sources were performed. Stacey and Sigmar [14,15] extended the multiple regime formalism to include both axisymmetric and momentum sources. Bennett [21] and Stacey [22] further extended the theory to include heat sources and heat fluxes in any collisionality regime. Burrell [23] also did this for a collisional plasma, and Parks, Burrell, and Wong [24] included a momentum source for an arbitrary collisionality plasma with no external impurity drag. Also, Burrell et al [25] included inertial effects on impurity transport, which is important when injected beams drive a toroidal plasma rotation.

CHAPTER 2

IMPURITY TRANSPORT DERIVATION

A. MOMENTUM EQUATIONS

The moment equations mentioned above are the velocity moments of the Fökker-Plank equation. The first two moment equations are

$$\frac{\partial n_i}{\partial t} + \nabla \cdot n_i \bar{\mathbf{v}}_i = S_i^0, \quad 1)$$

expressing particle conservation and

$$\begin{aligned} m_i n_i \frac{\partial \bar{\mathbf{v}}_i}{\partial t} + m_i n_i (\bar{\mathbf{v}}_i \cdot \nabla) \bar{\mathbf{v}}_i + \nabla p_i + \nabla \cdot \bar{\pi}_i \\ = n_i e_i (\bar{\mathbf{E}} + \bar{\mathbf{v}}_i \times \bar{\mathbf{B}}) + \bar{\mathbf{R}}_i + (\bar{\mathbf{S}}_i^1 - m_i \bar{\mathbf{v}}_i S_i^0), \end{aligned} \quad 2)$$

expressing momentum conservation. The scalar terms are n_i , m_i , e_i , p_i , and S_i^0 ; the number density, mass, electric charge, scalar pressure, and external particle source density. The velocity, electric field, magnetic field, friction force, and momentum source are the vector quantities $\bar{\mathbf{v}}_i$, $\bar{\mathbf{E}}$, $\bar{\mathbf{B}}$, $\bar{\mathbf{R}}_i$, and $\bar{\mathbf{S}}_i$. The remaining tensor, $\bar{\pi}_i$, is the anisotropic shear force, and the subscript i refers to the species in all cases. It should be noted that each equation contains the divergence of a higher order moment, so that the set of equations is not closed.

The ordering parameter, δ , is defined as the ratio of the gyro-radius to the gradient scale length. Since this is small in tokamaks, the details of the gyromotion are not considered in these calculations. By ordering the moment equations in terms of δ , the pressure and density are found to be zeroth order, and to lie on surfaces of constant magnetic flux. The particle and heat flows are found to be first order in δ , $O(\delta^1)$, and

also lie on the flux surfaces. The flows across the flux surfaces are second order, $O(\delta^2)$. This is discussed in detail in a review of transport by Hinton and Hazeltine [26], where the first four moments are given for the case in which no external sources are present.

In order to determine the second order, cross field flows, it is first necessary to determine the first order flows on the flux surface. The vector product of the magnetic field on the first order components of the momentum balance gives the flow perpendicular to the field.

$$n_i \bar{V}_{i\perp} = \frac{1}{e_i B} \hat{n} \times (\nabla p_i + n_i e_i \nabla \phi) \quad 3)$$

$$= \frac{RB}{e_i B} (p'_i + n_i e_i \phi') \frac{n \times \nabla \psi}{|\nabla \psi|} \equiv \frac{B_p G_i(\psi)}{B_\phi B} \hat{n}_\perp ,$$

where the prime denotes a derivative normal to the flux surface, nominally radial [25]. ϕ' is the electrostatic potential set up by charge separation. This potential will eventually modify the fluxes so that a further net charge imbalance is not produced, and therefore maintains ambipolarity [28]. B_ϕ and B_p are the toroidal and poloidal magnetic fields, and $G(\psi)$ is defined by Equation 3).

The particle flow parallel to the field is determined by integration over the divergence of the velocity from the particle balance equation,

$$n_i V_{i\parallel} = -\frac{G_i}{B} \left(1 - \frac{B^2}{\langle B^2 \rangle} \right) + \langle B n_i V_{i\parallel} \rangle B / \langle B^2 \rangle \hat{n} \quad 4)$$

The brackets $\langle . . . \rangle$ denote an average over the flux surface, and $\langle B n_i V_{i\parallel} \rangle$ is a constant of integration, which remains undetermined at this point. This constant of integration is determined by parallel momentum balance, and contains the effects of the external sources in addition to the gradients. A complete derivation by this Author can be found in Ref. [21], which includes all sources and higher order moments.

The perpendicular and parallel flows may be combined to obtain

$$n_i \bar{V}_i = \left(\frac{\bar{B} \cdot n_i \bar{V}_i}{B_p^2} \right) \bar{B} - \frac{R}{e_i} (p_i' + n_i e_i \phi') R \nabla \phi \quad 5)$$

where the flow parallel to the field may be written

$$\Gamma_{ip} \equiv \left(\frac{\bar{B} \cdot n_i \bar{V}_i}{B_p^2} \right) = \frac{F}{e_i \langle B^2 \rangle} (p_i' + n_i e_i \phi') + \langle B n_i V_{i\parallel} \rangle / \langle B^2 \rangle . \quad 6)$$

It is useful to note that the second term in 5) is a solid body toroidal rotation, whereas the first term varies with the poloidal position.

B. PARALLEL FLOWS

The flux surface average of the parallel component of the momentum balance equation is,

$$\langle \bar{B} \cdot \nabla \cdot \vec{\pi}_i \rangle = \langle \bar{B} \cdot \vec{R}_i^1 \rangle + \langle \bar{B} \cdot \vec{S}_i^1 \rangle . \quad 7)$$

The gradient of the flux surface quantities disappear in the flux surface average, and the parallel component of the $\vec{V} \times \vec{B}$ term is zero by definition. Hence, the parallel component of the divergence of the pressure anisotropy, is balanced by friction, $\vec{R}_i$, and a momentum source and drag, $\vec{S}_i$.

A form for the anisotropic shear, or parallel viscosity, was presented by Braginskii [2] for a classical plasma. In this it was argued

that the shear must be proportional to the lowest order particle flows on the flux surface. The classical viscosity coefficient,

$$\eta_{\parallel} \sim n_i T_i / \nu_{ii} \quad , \quad 8)$$

was found to be negligibly small, and the parallel shear was set to zero.

If the left hand side of the anisotropic shear equation, Eqn. 7), is set to zero, the set of equations of all species i of this form may be solved simultaneously for the parallel flows, using the appropriate forms for the friction and source terms.

If any species is in a long mean free path regime, the shear for the particle viscosity is increases due to the increased particle path length, and the shear for that species is not small. The fluid formulation for the viscosity was established by Grimm and Johnson [29]. A viscosity coefficient which covers all collisionality regimes was derived by Hirshman [9, 28, 29] which embodies all the information needed for transport in the long mean free path regimes.

$$\langle \bar{\mathbf{B}} \cdot \nabla \cdot \vec{\pi}_i \rangle = 3\mu_{ii} \langle (\hat{\mathbf{n}} \cdot \nabla \mathbf{B})^2 \rangle \Gamma_{ip} \quad . \quad 9)$$

This is of course independent of the solid body rotation. Given a friction form involving the particle flows, and an external drag and momentum source, the parallel particle flows can be determined by solving the sets of equations 7), and 9), for each species simultaneously. This then allows

Note: Recent toroidal rotation experiments on PLT indicate that the perpendicular viscosity must be ~ 100 times larger than the theoretical value. This was inferred from the maximum rotation speed with a known momentum input [28].

evaluation of the constants $\langle B n_i V_{i\parallel} \rangle$ and, hence, the complete determination of the flows and transport fluxes across flux surfaces.

C. TRANSPORT FLUXES

To determine the higher order particle flows across the flux surface, the surface average of the toroidal component of the momentum balance is taken,

$$\langle R^2 \nabla \phi \cdot m_i n_i \frac{\partial \bar{V}_i}{\partial t} \rangle = \langle R^2 \nabla \phi \cdot n_i e_i \bar{V}_i \times \bar{B} \rangle + \langle R^2 \nabla \phi \cdot \bar{R}_i \rangle + \langle R^2 \nabla \phi \cdot \bar{S}_i \rangle, \quad 10)$$

to lowest order, the first term on the right hand side reduces to the particle flow normal to the magnetic surface, since the significant portions of the cross product are the radial velocity and the poloidal field. After equilibrium flows are established, the particle flux is

$$\Gamma_i = \langle \nabla \psi \cdot n_i \bar{V}_i \rangle = n_i \langle R^2 \nabla \phi \cdot \bar{V}_i \times \bar{B} \rangle = -\frac{1}{e_i} (\langle R^2 \nabla \phi \cdot \bar{R}_i^1 \rangle + \langle R^2 \nabla \phi \cdot \bar{S}_i^1 \rangle) \quad . \quad 11)$$

This shows a toroidal momentum source directly driving the cross field flux, as well as the usual collisional friction. This source effect was first suggested by Ohkaka [12], who, along with Ref. [13], included only the direct momentum effects of the fast particles. Stacey and Sigmar [14,15] included the momentum effects on the bulk particle flows, and subsequent authors extended this treatment to include higher moments and sources [20-23].

Since gross charge separation can not occur, the radial current for steady state must be zero. This is the ambipolarity constraint,

$$\sum_i e_i \Gamma_i = 0 \quad . \quad 12)$$

From this, the electrostatic potential may be determined, completing the description of the fluxes.

The second pair of moment equations may be solved in parallel with the first pair in order to derive the cross field heat flux, as well as corrections to the particle flux. Heat flows on the flux surface change both the friction and the pressure anisotropy. This has the effect of adding new transport terms due to temperature gradients instead of simply pressure gradients. Higher order sources may also be added.

This approach has ordered out classical transport. For classical transport, no magnetic field gradient is considered, causing the parallel flows to equal the parallel averaged flows. This negates the parallel viscosity, leaving only the smaller, perpendicular viscosity to consider. The classical perpendicular viscosity is small since the particles lie very close to the flux surface.

D. GENERAL APPROXIMATIONS AND RESULTS

The friction between any two species is proportional to the net velocity difference between the two species. A choice of a coefficient of friction is then needed. The obvious choice is the collision frequency momentum exchange,

$$v_{ij} \propto \frac{n_j z_j^2 z_i^2 \ln \Lambda}{\sqrt{\mu} T_2^3}, \quad (13)$$

where μ is the reduced mass and T the common temperature. Λ is the ratio of the distance at which a single collision will produce a 90° deflection to the Debye length.

Braginskii had predicted that the impurities would peak sharply at the center, roughly

$$\frac{n_z}{n_{z0}} = \left(\frac{n_i}{n_{i0}} \right)^z, \quad (14)$$

This ignores the displacement of the main ions by the impurity. Approximately correcting for this, the peak for an average concentration of iron of $n_z z_z^2 / n_j = 1$ is less than 4 instead of 26. It can be seen that the sharp impurity peaks predicted by Braginskii would only be seen for very small impurity concentrations.

Since the particles are not in a simple Maxwellian, this collision frequency is not quite accurate. The particles actually are interacting over a net velocity difference, distorting the Maxwellian. Braginskii [2] noted this for an electron-ion plasma, and introduced a correction factor. Rutherford [4], in applying the Braginskii solution form to impurity transport, extended this correction factor. Samain [11] and Hirshman [30] later extended this correction factor, C_n , to long mean free path regimes. The factor C_1 is applied to particle friction causing pressure anisotropy. C_2 is applied to both heat friction, causing pressure anisotropy, and to particle friction causing heat anisotropy. Likewise, the factor C_3 is relevant to heat friction, causing heat shear. These factors are appropriate only to a large mass ratio approximation.

When these friction coefficients are used, the impurity transport is proportional to C_1 times the density gradients and $(C_1 - C_2)$ times the temperature gradient. With the usual negative main ion gradient, the impurities are pushed to the center of the plasma. If the main ions are in the collisional regime, the temperature gradient also drives impurities inward. When the main ions are in the long mean free path regime, the temperature gradient "screens" the density gradient effect, possibly driving impurities outward [11, 30]. When the Lorentz mass ratio approximation is not valid, temperature screening may also be predicted in the collisional, Pfirsch-Schlüter, regime.

In addition to gradients, impurity transport is also caused by external sources. As mentioned above, poloidally non-symmetric particle sources can drive impurity fluxes. This effect is conceptually equivalent to the lack of cancelation of the magnetic gradient drift which causes Pfirsch-Schlüter transport, and hence this effect does not exceed collisional scaling in any case. Therefore, it is only important when impurities are collisional.

A momentum source such as a neutral beam is also found to drive impurity transport [12-15]. Only the source component parallel (or anti-parallel) to the plasma current can cause transport. Generally, from the

Stacey-Sigmar theory [14, 15] in which bulk flow modification is calculated, when the beam is injected parallel to the current (co-injection), the impurity transport effect is outward if $\beta_{zi} k_i > \beta_i k_z$. β_i is the ratio of the external drag on the species to species collision frequency with the other species, $\beta_i = v_{di}/v_{iz}$, and K_i is the momentum input to the species. This is a collisionally scaled result in the collisionally scaled regimes, and rises to neoclassical scaling when all particles are in the long mean free path regimes.

Burrell [25] et al developed an alternate beam driven flow reversal theory based on inertial effects. In this theory, when the rotation velocity approaches the impurity thermal velocity, poloidal impurity density variations develop. The poloidal friction variation then drives impurity fluxes, which is outward for co-injection if $m_z v_{z\parallel}^2 > 2T_i$. This inertial term is thought to be larger in smaller machines than the above mentioned flow reversal term driven directly by the momentum source.

Although calculations have been performed using each ionic species [18, 32-35], the huge amounts of computer time and core space indicate the need for a simplification. An approximation frequently used is known as coronal equilibrium [36]. This is a method of determining the average charge state of an impurity species as a function of electron temperature, due to an equilibrium between their ionization and recombination rates. Transport in coronal equilibrium is usually then calculated as the transport of a single ionic species with the average charge state. This process unfortunately does not take into account transport due to gradients in each charge state. Also, if the transport time is shorter than the time to re-establish coronal equilibrium, the actual transport quickly diverges from the approximation. This is a good approximation for large or dense machines, although it is weaker where gradients are steep such as in the edge region.

An approximate correction to the limitations of coronal equilibrium is by transporting a few charge states at a time in "clumps" [37]. This method does not take into account the fact that only a few charge states exist at any one location, and hence the clump generally does little better than coronal equilibrium.

Transport coefficients are summarized by Boley [38] for mixed collisionality tokamaks, and fits pertaining to the viscosity coefficients are summarized by Hirshman [39] and Hinton and Rosenbluth [40]. In addition to the works mentioned above, Hogan [41] reviewed models for impurity effects and transport. Higher order sources were also found to have only a small effect on impurity transport [20].

CHAPTER 3

EXPERIMENTAL EVIDENCE

The presence of an impurity in a tokamak plasma can be directly determined only by observation of the impurity's characteristic X-rays. For this approach to be successful, the impurity must not be fully stripped, and the concentration of the impurity must be sufficiently large so that a particular line may be observed. The observation of forbidden transition lines has been found to be particularly useful [42, 43]. The impurity content can also be indicated indirectly both through enhanced radiation and increased plasma conductivity.

Early tokamaks exhibited various impurity characteristics. In the ST tokamak, impurities were found not to form a sharp central peak at the plasma center with both hydrogen [44] and helium [45] as the working gas. Oxygen and iron were the main impurities, and the electron temperature was slightly peaked at the center. Transport theory for the Pfirsch-Schlüter regime would predict finding the impurities in a slight ($\sim$ factor of 2-3 for the impurity concentrations) peak at the center of the plasma. Since these references looked for a sharp peak, the smaller peaks were not ruled out.

The T-4 tokamak, on the other hand, had centrally peaked impurity profiles [46], usually by factors of 3-5. This machine was approximately the same size as the ST tokamak, but had a flatter electron temperature gradient. In the absence of the temperature gradient, the T-4 tokamak should have had a flatter, not more peaked impurity profile according to the theory. In the T-4 case, the impurities studied were heavy noble gases injected at the edge.

Other early machines in which impurities were studied are the ATC [47] and T-3 [48] tokamaks. Carbon had been observed to flow to the center of the T-3 tokamak, and as in the case of the T-4 tokamak, a slightly peaked profile was inferred. Observations of aluminum injected into ATC verified that the aluminum moved inward at approximately the Pfirsch-Schluter rate, but no report of a final profile was given.

Although central accumulation can be observed in some cases, the moderate accumulation is not always consistent with the steeper profiles predicted. This indicates that Pfirsch-Schlüter transport is not inconsistent with the experimental observations. The results of the early experiments are inconclusive, and a rigorous analysis of these devices including multiple charge states and other potentially important phenomena has never been performed, leaving the experimental verification of neoclassical theory for impurity transport an open question.

Similar experimental results have been obtained for more recent, larger tokamaks. ISX-A, designed for impurity study, used argon injection to trace the transport of impurities. A multiple charge state analysis was used to model the discharge [50], using more than one impurity species. Reasonable agreement with the experiment was obtained with neoclassical transport at twice the predicted value [51], while using an empirical electron heat conduction term. With large amounts of injected tungsten, the modified theory modeled an inverted electron temperature profile observed in experiments.

Experiments on other machines of this size range show similar impurity behavior. Charge state distributions of tungsten on PLT prove that this machine is in coronal equilibrium [52]. The impurities in this device tend to accumulate at the center but do not peak sharply. Results from T-10 [53] also indicate neoclassical transport rates and some of the effects, including a hollow electron temperature profile due to central radiation of accumulated tungsten. In summary, experimental observations of ohmically heated tokamaks qualitatively support neoclassical transport as a major mechanism for impurity transport. There are many unresolved aspects which complicate the picture, such as proper external source and recycling rates to use in the modelling. This greatly affects the ability to correlate theory and experiment, with the result that it is not possible to draw any firm conclusion about the validity of neoclassical theory in predicting impurity transport.

The effect of asymmetric particle sources on impurity transport was predicted prior to the experiment [15-17]. Researchers on ISX conducted the experiment by gettering inner and outer torus walls and injecting neutral gas at the bottom of the plasma. By switching the plasma current

and magnetic field directions, both impurity transport enhancement and retardation were observed at the magnitudes and directions predicted in the outer edge (~ 2 cm.) of the plasma for small quantities of injected neon [51]. However, no consistent modification of the flows of the intrinsic impurities could be observed.

Experiments to prove the validity of the beam driven flow reversal theory [15] were also performed. Data from both PLT [52] and ISX [54] indicated enhanced impurity flow to the center with counter-injection and reduced flow to the center with co-injection, as predicted. However, inward tungsten fluxes in PLT with counter-injection were much larger than this theory could account for [52]. The inertial term of Burrell was found to match better the observed flows. On the other hand, a steady state analysis based on [15] indicated that a plasma averaged impurity concentration of $\alpha = n_z Z_z^2 / n_i \sim .25$ could significantly slow and possibly reverse the inward impurity flux for both ISX and PLT [55]. This was done for the case of the initial impurity concentration being peaked at the outside of the plasma, as would be the seen with an external source of impurities. Further analysis needs to be completed to see if the theoretical explanation of flow reversal and the experimental evidence are sufficiently close to label either theory as the observed mechanism.

Overall, the experimental evidence supports the Pfirsch-Schlüter and neoclassical predictions of plasma impurity transport, at least to the extent that they are similar in magnitude and general action. Coefficients which are based on microscopic plasma quantities may account for some of this. Experimental data have significant uncertainties intrinsic to it not only for the observed flows and profiles of impurities, but also the bulk plasma profiles and other properties which dictate the impurity transport. Sawtooth and other plasma oscillations complicate the issue. With these uncertainties it is difficult to use the impurity transport models as a predictive tool, but further study may extend these models so that they can aid in the design of future power producing reactors.

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