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Low Z Impurity Transport in Tokamaks

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ABSTRACT

Low Z impurity transport in tokamaks was simulated with a one-dimensional impurity transport model including both neoclassical and anomalous transport. The neoclassical fluxes are due to collisions between the background plasma and impurity ions as well as collisions between the various ionization states. The evaluation of the neoclassical fluxes takes into account the different collisionality regimes of the background plasma and the impurity ions. A limiter scrapeoff model is used to define the boundary condition for the impurity ions in the plasma periphery. In order to account for the spectroscopic measurements of power radiated by the lower ionization states, fluxes due to anomalous transport are included. The sensitivity of the results to uncertainties in rate coefficients and plasma parameters in the periphery are investigated. The implications of the transport model for spectroscopic evaluation of impurity concentrations, impurity fluxes, and radiated power from line emission measurements are discussed.

1. INTRODUCTION

Low Z impurities in a tokamak discharge affect not only the plasma composition and hence the resistivity but also the local power balance and the maximum plasma density attainable. The importance of low Z impurities is well recognized and their effects are qualitatively well understood. Spectroscopic measurements of line emission are commonly used to evaluate the concentration, influx rate and power radiated by low Z impurities. In the initial phase of the discharge, these measurements can determine the impurity density and in the quasi-steady state, the influx rate. In general, a transport model must be used in the quasi-steady state to infer the impurity concentration from the measurements of the influx rate. Spectroscopic measurements of low Z impurities on ATC [1] and PLT [2] have been analyzed using a one-dimensional transport model. The general implications of this model in interpreting spectroscopic measurements of line emission, and in evaluating impurity concentrations, radial profiles of the fully stripped state, $Z_{\text{eff}}(r)$ and power radiated will be discussed.

The impurity density, n_j , in the j th ionization state can be calculated by solving the impurity transport equation:

$$\frac{\partial n_j}{\partial t} = -\frac{1}{r} \frac{\partial}{\partial r} (r \Gamma_j) + n_e (n_{j-1} S_{j-1} - n_j S_j + n_{j+1} \delta_{j+1} - n_j \delta_j) \quad (1)$$

$$j = 1^+ \dots Z^+$$

where Γ_j is the radial impurity particle flux, S_j the ionization rate coefficient, and δ_j is the recombination rate. In Eq. 1, n_j is expressed in cylindrical coordinates and has been averaged over a flux surface, which is equivalent to assuming poloidal and toroidal symmetry. Though recent experiments [3,4] indicate poloidal asymmetries can exist, this effect will not be explicitly considered in the present one-dimensional formulation of the problem. The transport for low Z impurities has been studied theoretically and in conjunction with several experiments. Since different models and approximations have been developed, the recent tokamak literature will be reviewed.

1.1 EXPERIMENTAL MEASUREMENTS OF IMPURITY INFLUX AND TRANSPORT

Gervids and Krupin [5] on TM-3 compared the measurements of carbon line intensity with calculations which neglected diffusion, whereupon they concluded that transport was important after "burnout" of the various carbon lines. (In the initial phases of the discharges, as the temperature increases, the impurities are ionized to successively higher ionization states and thus "burned out".) The flux of CV ions was determined by integrating Eq. 1.

$$\Gamma_j(r,t) = \frac{1}{r} \int_0^r \left[\frac{\partial n_j}{\partial t} - n_e (n_{j-1} S_{j-1} - n_j S_j + n_{j+1} \delta_{j+1} - n_j \delta_j) \right] r' dr', \quad (2)$$

for $j = 4^+$ using the experimentally determined densities and temperatures. The flux of CV was found to be inward. While determining Γ_j for the lower states of ionization according to Eq. 2 is useful as a measure of the influx of impurities, it does not determine the accumulation and hence total concentration of impurities unless Γ_j and n_j are evaluated for all states (including the fully stripped).

Using a similar procedure, the TFR group calculated the flux of incoming oxygen impurities [6,7] in the quasi-steady state. In this experiment, the observed flux was so large as to be inconsistent with observed impurity accumulation. This led to the conclusion of significant recycling in the plasma periphery [6,8]. Though the TFR group [6] compared the observed flux with an estimate of Pfirsch-Schlüter diffusion and found reasonable agreement, a coherent transport model which accounts for the large outward flux required to achieve a quasi-steady state was not presented.

In PLT, the observed ionization rate (and hence influx) of oxygen as determined from both the brightness of the OVI resonance line [2], as well as the behavior of the higher states of ionization, is also not consistent with impurity accumulation. Typically in the quasi-steady state of the discharge, the intensity of not only the OVI but also highly excited OVII lines are observed to remain constant. While the intensity of the OVI lines is determined primarily by the influx of oxygen, the intensity of the OVII lines is determined also by recombination. Since, during the discharge, the intensity of OVII typically attains a quasi-steady state, the total concentration of oxygen does not increase (accumulate) with time. Unlike the reported TFR experiments [6,7,8], in PLT the Z_{eff} as deduced from Thomson scattering measurements was low (~ 2) for the high density helium discharges studied [2]. Thus, modification of the background ion density profile by low Z impurity ions would have been relatively less important. In Alcator, the

observed fluxes of oxygen were sufficiently low such that impurity accumulation cannot be readily excluded. In the analysis of these measurements, the inward flux was modeled as due to Pfirsch-Schlüter diffusion, though the total oxygen density was estimated by assuming a finite particle confinement time [9].

Two other experiments specifically designed to study impurity transport were conducted at Princeton on the FM-1 spherator [10] and on the ATC [11,12,13] tokamak by injection of impurities for a time short compared with the transport time. On FM-1, the confinement of barium impurity ions was found to be in agreement with a model which included both classical as well as anomalous transport [10]. Subsequently, Marmar has pointed out that recombination may have been significant in the FM-1 experiment [13]. Recombination may explain why at the highest densities the impurity particle confinement time decreased with increasing density, whereas the transport model predicted it should have increased. Nonetheless, recombination does not explain the principal result of an increase in confinement time with density at lower densities. On ATC, space and time resolved measurements of various aluminum lines in the volume outside the $q = 1$ surface were in reasonable agreement with neoclassical transport theory [11]. The authors [12] have also pointed out that the diffusion velocity is also consistent with pseudoclassical anomalous transport, for in these experiments neoclassical transport was of the same order as anomalous transport. Nonetheless, from measurements of the flux of Al X and Al XI, the authors conclude that an inward flux not proportional to V_n exists [13].

In addition to spectroscopic measurements, impurity transport has been studied by X-ray techniques and most recently with a neutral beam. Vershkov and Mirnov [14] on T3 concluded that impurities accumulate in the center and furthermore that MHD activity can disperse the impurities. Von Goeler, et al., on ST also using X-ray detectors concluded that Z_{eff} was relatively constant within the inner half of the discharge and that impurity accumulation was not significant [15]. Goldston measured the q profile on ATC using a neutral beam and, in contrast, concluded that Z_{eff} was peaked on axis [16]. In a recent series of experiments conducted on TM-3 [17] and PLT [18], the existence of an outward impurity flux was determined on the basis of X-ray measurements of argon [17,18] and oxygen [17] impurity radiation. So far, detailed comparison with theory has not been reported for either experiment.

In summary, the applicability of any impurity transport model has not yet been definitively established experimentally. Hence, it is necessary to consider neoclassical as well as anomalous transport. Significant recycling of the lower ionization states has been observed on a number of experiments. It appears that further measurements of the Z_{eff} profile are required to put in perspective the present experimental results.

1.2 COMPUTATIONAL MODELS

In recent years, one-dimensional transport models have been developed to calculate the radial distribution of the different ionization states. Tazima, et al. [19], calculated the radial

distribution of oxygen for ST. The transport model was similar to Pfirsch-Schlüter diffusion but neglected the inward transport of oxygen due to proton-oxygen collisions. Thus, steady state solutions with an influx of impurities were obtained. The results indicated the presence of significant concentrations of the fully stripped state in the plasma periphery in addition to the lower ionization states. The relationship of the fully stripped state in the plasma periphery to the influx of oxygen atoms was not discussed. Amano and Okamoto [20] assumed Pfirsch-Schlüter transport and no forced neutral influx and found that only the highly ionized states exist in the plasma periphery in any significant concentrations. Dnestrovsky, et al. [21] performed similar calculations assuming neoclassical fluxes appropriate to the banana regime even though, as the authors point out, the impurities are not in the banana regime. This choice was made to enhance the ion and electron fluxes which the authors point out experimentally are not consistent with neoclassical theory. This approach does not clearly identify the relative importance either of neoclassical or anomalous transport. A slab model for the lower ionization states in PLT has been developed by Rutherford, et al. [22]. This model considered both the effects of anomalous and Pfirsch-Schlüter transport and indicated the significance of anomalous transport.

Another approach, used by Popoular [23] and in other transport simulations [24,25], is to sum over all ionization states and assume coronal equilibrium or small deviations from

coronal equilibrium [23]. Sardei and Duchs [26] used a compromise approach in which six transport equations were solved representing the background plasma, and different oxygen and iron ionization states. In this way, the effects of high Z impurities on low Z transport were simulated. They also investigated the effects of an enhanced electron collision frequency to simulate anomalous transport. This study examined how relatively large concentrations of impurities affect the ion density profiles and hence Z_{eff} . Both of these approaches [23-26] made detailed comparison with experiments difficult since the calculated impurity fluxes are not the measured flux.

Another approach which does not consider the self-consistent problem of impurity profile evolution but which clarifies the role of impurity transport on power radiated has been used by Breton, et al., [27]. In this model a constant inward or outward flow velocity was assumed. The form of this model was influenced by experiments conducted on TFR. While this approach is physically intuitive, it does not explore the full implication of their transport models or of the imposed flux.

In this paper, a more complete treatment of the neoclassical fluxes will be used taking into account the different collisionality regimes of both the background and impurity ions. In addition, the effects of anomalous transport on impurity fluxes will be presented. The purpose of the calculations is to provide a framework in which to interpret and analyze spectroscopic measurements on tokamaks in the context of various transport models. In Section 2, the transport model and the relevant atomic physics cross-sections

will be presented. The results of the transport model will be discussed in Section 3 and the implications of this study for the interpretation of spectroscopic measurements will be addressed in Section 4.

2. MODEL

2.1 TRANSPORT MODEL

As discussed in the introduction, present experimental results indicate that both neoclassical as well as anomalous impurity transport should be considered. Thus, $\Gamma_j = \Gamma_j^{\text{col}} + \Gamma_j^{\text{A}}$ where Γ_j^{col} is the total collisional neoclassical flux and Γ_j^{A} is the anomalous flux.

In most current tokamak experiments, the impurity ions are in the Pfirsch-Schlüter regime over most of the plasma volume, while the background ions are in the banana-plateau regime. Depending upon the parameter $\omega_{\text{ta}}^2 / \nu_{\text{S}}^{\text{a}} \nu_{\text{E}}^{\text{a}}$ (where ω_{ta} is the transit frequency and $\nu_{\text{S}}^{\text{a}}$ and $\nu_{\text{E}}^{\text{a}}$ are the momentum and energy exchange collision frequencies of species "a") the impurities at different minor radii can be in either the extreme Pfirsch-Schlüter regime corresponding to the results of Rutherford [28] and Hirshman [29] or in an "intermediate Pfirsch-Schlüter regime" [30,31]. The principal difference between these two regimes is the flux due to gradients in the ion temperature profile. In the extreme Pfirsch-Schlüter regime, except for cases corresponding to very high impurity concentrations [32], the flux due to ion temperature gradients is typically small. On the other hand in the intermediate Pfirsch-Schlüter regime, the flux due to the ion temperature gradient can be significant. The previous results of Ref. [29] have been extended to the intermediate Pfirsch-Schlüter regime. In our present calculations, the neoclassical flux is composed of three terms corresponding to

a) flux due to impurities in the Pfirsch-Schlüter regime (including the transition from the intermediate to the extreme regime), b) anisotropy-driven neoclassical fluxes, and c) classical fluxes. Impurity collisions between different ionization states are taken into account. For typical parameters of interest, since strong temperature equilibration [29,32] is more common, the calculations assume that all of the ions are strongly equilibrated. In this study, fluxes due to low Z impurity collisions with high Z impurities are neglected. Since the expression for the neoclassical fluxes are rather complicated, they are presented in Appendix A.

Whereas the calculations for the neoclassical fluxes are well established the same cannot be said of the anomalous fluxes. The model used in this paper is influenced by the modeling of experiments on FM-1 where

$$\Gamma_j^A = -D_A(r,t) \frac{\partial n_j}{\partial r} \quad (3)$$

was assumed in addition to the fluxes due to neoclassical theory [10]. Though $D_A(r,t)$ is related to the plasma parameters, and therefore likely to be spatially varying, in this paper, D_A will be assumed to be equal to a constant. Since in Section 3.3 the sensitivity of these calculations to changes in plasma parameters in the periphery will be examined, it is desirable not to have the anomalous diffusion coefficient changing simultaneously, especially since the functional dependence on plasma parameters of D_A is not known. The greatest uncertainty in the anomalous flux is probably not the magnitude of D_A (which can be estimated

on the basis of comparison with experiments) but whether a simple relationship such as that employed in Eq. 3 can adequately describe the flux. For example, in Eq. 3 the effects of impurity temperature gradient terms are ignored. While we are not aware of any data which contradicts our transport model, the experimental and theoretical uncertainties are still quite substantial.

In solving the transport Eq. 1, boundary conditions must be imposed. On axis the boundary condition is

$$\frac{\partial n_i}{\partial r} = 0 \quad . \quad (4)$$

In the plasma periphery, the choice of boundary condition is more difficult. Merely forcing the density to zero or a very small density can result in an artificial inward flux when only neoclassical fluxes are considered [21]. In order to avoid the artificial inward flux, a different approach has been used employing a limiter scrapeoff model. In the region of the discharge, $r > a_{lim} = 40$ cm for PLT, the impurity ions are assumed to be lost to the limiter from the plasma at a velocity $v_S = \sqrt{(T_e + T_i)/m_b}$, where m_b is the background ion mass and T_e and T_i are the electron and ion temperatures. Thus in the scrape-off region, this model introduces a loss term n_i/τ_S in Eq. 1 where $\tau_S = \pi R/v_S$ and R is the major radius. This model for the impurities in the plasma in the scrapeoff is clearly simplified and further experimental measurements are required to determine the actual impurity flow velocities [33]. Fortunately, except for the

lowest ionization states, the results of the calculations are insensitive to a factor of 3 uncertainty in τ_S . This model very naturally introduces boundary conditions on the plasma. Almost all of the flux out of the main body of the discharge (typically > 99%) goes to the limiter as opposed to the wall which is located at 50 cm. Thus uncertainties in the boundary conditions at the wall are not important.

The flux of neutral atoms is calculated from

$$\Gamma_O = \left[\Gamma_{\text{ext}} + R \sum_j \Gamma_j \right], \quad (5)$$

where Γ_{ext} is the imposed flow of atoms into the discharge (typically from the walls), R is the recycling coefficient taken to be unity for the impurity ions and Γ_j is the flux of ions to the limiter and to the vacuum vessel wall. In the calculations presented here $\Gamma_{\text{ext}} = 0$. The neutral atom profile is

$$n_O(r) = n_O(a) e^{-\frac{1}{v_O} \int_r^a dr' n_e(r') S_O(r')}, \quad (6)$$

where v_O is the neutral velocity taken to be 1 eV and $n_O(a)$ is determined by Γ_O . Singly ionized oxygen has a large cross-section for charge exchange with hydrogen atoms [34]. Charge exchange collisions have not been included but have been approximately simulated by varying the neutral atom velocity. Varying the oxygen neutral energy from 0.2 to 10 eV, results in a negligible change in the higher ionization states 3^+ through 8^+ but does affect the lower ionization states. The neutral

oxygen radiation decreased by a factor of 4.4, O^{1+} radiation by 3.3, O^{2+} by 2.1, O^{3+} by 1.5 and O^{4+} by 1.1. As the neutral energy of oxygen decreases it becomes ionized in the plasma scrapeoff and undergoes greater recycling. As will be discussed in Section 3, the density and radiation of these very low ionization states are also very sensitive to plasma parameters in both the periphery and the scrapeoff region.

In solving the transport equations, the electron and ion temperature and electron density profiles are assumed known and the hydrogen density profile is calculated from the assumption of quasineutrality. All of the calculations were done for parameters typical of the PLT tokamak using the electron density and temperature profiles in Fig. 1 [35]. In the plasma periphery, the density and temperature profiles are poorly known especially in the region 40-45 cm [35]. In order to examine the sensitivity of the calculations to small profile changes, the profiles shown in Fig. 2 and 3 have also been studied for the same peak parameter shown in Fig. 1. The calculations are for the profiles corresponding to an edge density of $1. \times 10^{13} \text{ cm}^{-3}$ and temperature of 30 eV unless otherwise stated. The principal advantage of using the measured density and temperature profiles is to separate as much as possible the impurity transport effects from the uncertainties in the transport of the background plasma. One disadvantage is that this limits the validity of the calculation to conditions where the impurity density is small compared with the main plasma density. This is not especially restrictive since many tokamaks presently operate in this range of parameters.

2.2 RATE COEFFICIENTS

Except for the cases discussed in Section 3.4, ionization rate coefficients of Lotz [36] were used. The radiative recombination rate coefficients were evaluated using the results of Breton, et al. [27] which are based on the calculations of Seaton [37] and Von Goeler, et al. [15]. The dielectronic recombination rate coefficients were calculated on the basis of Burgess' formula [38] using the tabulated results of Donaldson and Peacock [39]. The excitation rate coefficients of Vainshtein, et al. [40] were used in calculating the power radiated. The above rate coefficients (presented in Appendix B), especially the ionization and excitation rate coefficients, were chosen on the basis of comparison with experimental data. Nevertheless, the uncertainties in the rate coefficients are a factor of ~ 2 .

In order to analyze the influence of the uncertainties in rate coefficients on the radial impurity density distribution and power radiated, calculations were performed using ionization rate coefficients of Drawin [41] and excitation rate coefficients of Van Regemorter [42] and Seaton [43] (see Appendix C). Depending upon the ionization state, these cross-sections can differ from the preceding ones by factors of 3 to 4.

For oxygen and carbon impurities and $T_e < 500$ eV, line radiation is the dominant radiation loss. For higher temperatures, bremsstrahlung and radiative recombination tend to dominate but for typical tokamak parameters are not very significant. Thus, only line radiation loss was evaluated from

$$P_j^{na} = 1.602 \times 10^{-19} n_e n_j S_j^{nm} E_j^{nm} \text{ (Watts)} , \quad (7)$$

where $E_j^{nm} = |E_j^n - E_j^m|$ is the transition energy, S_j^{nm} is the excitation rate coefficient where j indicates the ionization state and n and m the upper and lower levels of the transition, respectively. The different oxygen and carbon transitions evaluated in this study are tabulated in Appendix D. In addition the power loss by ionization was evaluated:

$$P_j^{ion} = 1.602 \times 10^{-19} n_e n_j S_j E_j \text{ (Watts)}, \quad (8)$$

where E_j is the ionization energy. Typically, the power loss by ionization is small compared with the power radiated.

3. DISCUSSION

3.1 NEOCLASSICAL TRANSPORT

The density profiles of oxygen and carbon ionization states and the corresponding resonance line radiation assuming only neoclassical transport ($D_A = 0$.) is shown in Figs. 4 and 5. These calculations were evaluated after 300 msec for the electron and background ion temperature and density profiles shown in Fig. 1 and correspond to a quasi-steady state equilibrium for the lower ionization states, though the fully stripped state is still evolving with a time constant > 1 sec. Because $\Gamma_{\text{ext}} = 0$ (no forced flux of neutrals), the total oxygen and carbon content remains constant corresponding to 2% of the electron content. The density of the lower ionization states is related to the outward flux to the limiter and walls of the higher ionization states. This is because in steady state, the flux of neutral atoms is directly proportional to the outgoing flux of ions (Eq. 5). The density of the lower ionization states (oxygen states less than 5^+ , and carbon states less than 4^+) is thus not determined by recombination of the higher ionization states which exist in the plasma periphery due to diffusion from the center of the discharge.

These neoclassical calculations differ significantly from similar calculations assuming that both the background plasma and the impurities are in the extreme Pfirsch-Schlüter regime. Using the transport coefficients of [29], the densities of the lower ionization states were computed to be considerably less by a factor of ~ 100 than shown in Fig. 4. In the present neoclassical calculations, the enhanced recycling of the lower

ionization states is due in part to the improved treatment of the contribution to the impurity fluxes due to the ion temperature gradient.

To see how this enhancement comes about consider the expression for the neoclassical fluxes:

$$\Gamma_a^{\text{col}} = \sum_{b \neq H} C_{ab} \frac{\partial n_b}{\partial r} + C_{aH} \frac{\partial n_H}{\partial r} + C_{aT} \frac{\partial T_i}{\partial r} . \quad (9)$$

The radial fluxes of O^{5+} and O^{8+} due to the density and ion temperature gradients are shown in Fig. 6. The flux due to the ion temperature gradient is important over a significant fraction of the profile except near the periphery which corresponds more nearly to the extreme Pfirsch-Schlüter regime. The effect of the ion temperature gradient term is to modify the profiles of the highly ionized states so that the outward flux of these ionized states increases substantially, in comparison with calculations assuming that the background ions as well as the impurity ions are in the Pfirsch-Schlüter regime [29]. The ∇T_i terms also significantly modifies the resulting Z_{eff} profile as shown in Fig. 7. The Z_{eff} profile is not peaked on axis as might be suggested by a simple estimate of neoclassical effects but is relatively constant across the profile. Because of the importance of the ∇T_i terms, as well as the other fluxes, unless a relatively detailed calculation is performed, simple estimates of neoclassical effects in predicting both $Z_{\text{eff}}(r)$ as well as the particle recycling time can be significantly in error.

Experimentally, the power radiated by the CIV and OVI resonance lines is typically a factor of 3-10 greater than that predicted by

the present neoclassical theory, and a factor ~ 500 greater than that assuming transport in the extreme Pfirsch-Schlüter regime [29]. It appears that neoclassical transport alone can not explain the power radiated by the lower ionization states. In the succeeding calculations, the effects of anomalous transport together with neoclassical transport will be considered.

3.2 ANOMALOUS AND NEOCLASSICAL TRANSPORT

The inclusion of anomalous transport substantially decreases the particle recycling time of the lower ionization states as shown in Fig. 8 for oxygen. The particle recycling time is defined in three ways. The first is a generalization of the particle recycling time for a hydrogen ion which is the total oxygen density (n_{TOT}) divided by the neutral oxygen flux required to maintain steady state (Γ_{TOT}). Since this definition of recycling time is very sensitive to the plasma parameters in the scrapeoff region, an alternative is to consider the density of ions in the upper three ionization states divided by the flux of these ions to the limiter and wall. This is not an arbitrary choice since spectroscopic measurements of the Li-like ions (O^{5+} and C^{3+} lines) are commonly used and, as will be shown below, can be related to this flux. The third choice was the density of the fully stripped state divided by the flux of the fully stripped state to the limiter and walls. Since recombination can be significant, this definition can result in a recycling time which is not necessarily characteristic of a transport time. Nonetheless, it is apparent from Fig. 8 that the lower ionization states recycle much faster than the higher ionization states.

This observation supports the TFR conclusion that recycling of the fully stripped state in the interior must be considerably less than the lower ionization states in the periphery. In comparing these calculations with measurements on PLT of the power radiated by OVI and CIV resonance lines, anomalous diffusion coefficients in the range 2000 to 8000 cm^2/sec appear to be reasonable. Theoretical predictions of anomalous diffusion coefficients due to microinstabilities are typically also in this range [22,24]. In order to illustrate the sensitivity of these calculations to changes in temperature and rate coefficients, D_A was chosen equal to 4000 cm^2/sec .

As shown in Fig. 9 and 10 for $D_A = 4000 \text{ cm}^2/\text{sec}$, the location of the different ionization states is very nearly the same as in Figs. 4 and 5 respectively. There are small differences in the profile shape, though the largest difference is in the amplitude of the ionization state. The relative intensity of the different ionization states also changes and, as will be described below, is also sensitive to changes in the temperature and density profiles as well as changes in the rate coefficient.

3.3 EFFECT OF EDGE PROFILE MODIFICATION

Figures 11, 12 and 13 illustrate the effects of modifying the plasma density and temperature profile in the periphery using the profiles shown in Figs. 2 and 3. As the density and temperature increases, the neutral mean free path decreases and the recycling of the lower ionization states within the limiter scrapeoff increases as shown in Fig. 13. The radiation by the very low ionization states O^{1+} , O^{2+} , and O^{3+} is very sensitive to changes in

the plasma scrapeoff condition and thus is not necessarily representative of conditions corresponding to the bulk of the discharge. If the temperature in the plasma periphery including the scrapeoff increases, the power radiated by the very low ionization states increases because of the increased recycling of these states. On the other hand, the power radiated by OVI decreases as shown in Fig. 12. In the ATC experiments, the neutral mean free path was longer than in PLT because the electron density and temperature in the scrapeoff was less. Hence recycling within the limiter-scrapeoff was not very important and a decreasing edge temperature was, in part, responsible for the increase in power radiated observed with the application of rf power [1,44,45].

Figure 14 illustrates how the power loss (radiation by both resonant and non-resonant lines and ionization) varies with anomalous diffusion coefficient, D_A , as well as with the density and temperature profiles in the plasma periphery. As expected, increased density in the plasma periphery results in an increased power loss, though for PLT parameters the effect of modest temperature profile variations are not very significant. The changes in the edge profile parameters would affect not only the impurity density profiles and the power radiated as discussed here but also D_A . Because of the dependence of D_A on edge profile parameters (which in general are poorly known) the impurity density and power radiated by the lower ionization states would be further modified.

Since the density of the lower ionization states is not determined by recombination, estimating the power radiated by assuming coronal equilibrium in transport simulations of the energy balance is not only a poor approximation but obscures the underlying physical mechanism responsible for the radiation (rapid impurity recycling).

3.4 SENSITIVITY TO RATE COEFFICIENTS

As was discussed in Section 2.2, the ionization, recombination and excitation rate coefficients are often known to only within a factor of ~ two. In order to examine the sensitivity of these calculations to uncertainties in the rate coefficients, oxygen profiles and power radiated for the case corresponding to Fig. 9 ($D_a = 4000 \text{ cm}^2/\text{sec}$) have been reevaluated using different rate coefficients. The total number of ions in each species for two different ionization rate coefficients (Lotz [36], as in Fig. 9 and Drawin [41]) are compared in Table I. Though the densities in the upper ionization states agree quite well, there is a considerable difference in the lowest three states O^{1+} , O^{2+} , and O^{3+} as well as the oxygen neutral influx. For the lower ionization states, Drawin's cross-sections [41] are larger by a factor of 2-3 than Lotz's, resulting in the enhancement of the lower states. Though the density and hence emission of the various ionization states changes substantially, the location and profile shape does not, due to the strong temperature dependence in the cross-sections.

In Table II, the calculation of oxygen density for the various ionization states shown in Fig. 9 is compared with two cases in which the recombination rate was varied by a factor of 5 for all

ionization states. In this case the densities of the lower states do not change very appreciably since recombination is less important than ionization and transport, but the densities as well as location and profile shape of the higher ionization states do change substantially. In smaller machines, such as ATC, variations in the recombination rate result in a smaller change in the calculated oxygen profiles because the particle transport time is shorter.

Resonance and non-resonance line radiation was evaluated using the oxygen density profiles shown in Fig. 9 and the corresponding temperature and density profiles using three different excitation rate coefficients, Vainshtein, et al. [40], Van Regemorter [42] and Seaton [43] (Table III). The resulting variation in the power radiated by any given ionization state is quite large even though the variation in the total power radiated is much smaller. Clearly the uncertainties in excitation rate coefficients also influence the calculations, resulting in an overall uncertainty for any given ionization state of a factor of 2-4.

4. IMPLICATIONS FOR SPECTROSCOPIC MEASUREMENTS

These results indicate both the limitations as well as the usefulness of detailed code calculation in obtaining a quantitative interpretation of spectroscopic measurements. The limitations arise from both the experimental uncertainty in plasma parameters in the periphery as well as uncertainties in ionization, recombination and excitation rate coefficients and in the basic formulation of the transport equations. The usefulness of such calculations is found in interpreting changes in emission of various impurity lines (for instance during plasma heating experiments) and in relating the emission both to the influx of impurities and changes in total concentration.

On the basis of the calculations presented here as well as others related to the ATC rf heating experiment [1,45,46], it is apparent that measurement of several lines from different ionization states as well as measurements of temperature and density profiles are required to understand the behavior of low Z impurities. Furthermore, by choosing the appropriate transition, it is possible to obtain a more reliable interpretation. One example will be treated in detail. In the quasi-steady state, the outgoing flux of oxygen $6^+, 7^+$ and 8^+ equals $\int n_e (n_{5^+} S_{5^+} - n_{6^+} \delta_{6^+}) dv$. Typically, recombination of oxygen 6^+ is not very important as shown in Table II. Thus, the outgoing flux is balanced by ionization of oxygen 5^+ . Power radiated by oxygen 5^+ is proportional to $\int n_e n_{5^+} S_{5^+}^{mn} E_{5^+}^{mn} dv$. Thus, if the temperature dependence

of S_{5+}^{mn} is similar to S_{5+} , measurements of the power radiated will be proportional to the influx. In Fig. 15, the power radiated by both the OVI resonance line as well as the OVI 150 Å (3^2P-2^2S) line are plotted as a function of the outgoing flux (to the limiter and walls) of oxygen 6^+ , 7^+ and 8^+ . Due to the different temperature dependence of the ionization rate coefficient of 5^+ and the excitation rate coefficient for resonance line radiation, the power radiated by OVI resonance line cannot be simply related to the flux. On the other hand, the 150 Å line and also the OVI 3811 Å (3^2P-3^2S) line appear to be more reliable for evaluating changes in oxygen influx. In analyzing the spectroscopic measurements conducted during the ATC rf heating experiment, it was possible, using measurements of the line OVI 3811 Å, to determine that the change in emission was due in part to a change in the recycling of oxygen impurities.

Since a large fraction ($\geq 1/2$) of the oxygen line radiation is due to resonance line radiation, measurements of the OVI resonance line appear to be a good indicator of the variation in the total radiated power. Table IV is a summary of the cases considered here. Except for the case of only neoclassical transport ($D_A = 0$) in which the flux of the lower states appears to be too low and not to agree with the experiment, the ratio of the total power loss (radiation and ionization) to that by OVI resonance line does not vary appreciably. Except for the case of Seaton's excitation rate coefficients [43] (which are not considered especially reliable), the uncertainty in rate

coefficient does not appreciably change this factor. It appears that by monitoring OVI resonance radiation, it is possible to estimate the total power loss to within ~ 50%. Similarly, CIV resonance line radiation can be used to monitor the power loss due to carbon impurities as shown in Table V.

It does not appear possible to reliably evaluate the total density of oxygen from measurements only of the lower ionization states in the quasi-steady state since the diffusion coefficients in the plasma periphery are not well known. The density of the lower ionization state is intimately related not only to the absolute impurity content in the discharge but also to the recycling rate. However, in the early phase of the discharge during "burnout", measurements of the lower ionization states can be related to the impurity concentration at that time. The difficulty in quantitatively evaluating the oxygen content in the quasi-steady state from the early time measurements is that during the time in which the discharge evolves to a quasi-steady state, an unknown forced influx of neutrals can exist in addition to the recycling flux.

5. SUMMARY

The simulation of low Z impurity transport in tokamaks indicates the need for the inclusion of anomalous transport in addition to neoclassical transport. However, the present neoclassical model is in closer agreement with experiment than previous neoclassical transport calculations which assumed that both the impurity and background plasma ions are in the Pfirsch-Schlüter regime. One consequence of the calculation is the relatively rapid recycling of lower ionization states compared with the higher ionization states in qualitative agreement with the TFR experiment [6,8]. The density and hence the power radiated by the very low ionization states are relatively sensitive to uncertainties in rate coefficients, and plasma conditions in the periphery. Thus, it is necessary to study the behavior of various impurity lines both in the plasma periphery and nearer the center of the discharge to obtain a comprehensive understanding of the variation in impurity flux and content during the quasi-steady state.

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REFERENCES

- [1] ELLIS, R.A., JR. DAUGHNEY, C., HAWRYLUK, R.J., POST, D.E.,
SUCKEWER, S., TAKAHASHI, H., Bull. of Amer. Phys. Soc. 21
(1976) 1089. HAWRYLUK, R.J., SUCKEWER, S., POST, D.E.,
Conf. on Atomic Processes in High Temperature Plasmas, G-1,
Knoxville, Tenn. (1977). SUCKEWER, S. and HAWRYLUK, R.J.,
Proceedings of the Third Topical Conf. on Radio Frequency
Plasma Heating, C9-1, Pasadena, Ca. (1978).
- [2] HINNOV, E., SUCKEWER, S., BOL, K., HAWRYLUK, R.J., HOSEA, J.,
MESERVEY, E., to be published in Plasma Phys.
- [3] TERRY, J.L., MARMAR, E.S., CHEN, K.I., MOOS, H.W., Phys. Rev.
Lett. 39 (1977) 1615.
- [4] SUCKEWER, S., HINNOV, E., SCHIVELL, J., PPPL 1430 (1978).
- [5] GERVIDS, V.I., KRUPIN, V.A., Sov. Jour. Fiz. Plazmy 1
(1975) 357.
- [6] TFR Group, Phys. Rev. Lett. 36 (1976) 1306.
- [7] Equipe TFR, Nucl. Fusion 15 (1975) 1053.
- [8] TFR Group, EUR-CEA-FC-868 (1976).
- [9] TERRY, J.L., CHEN, K.I., MOOS, H.W., MARMAR, E.S., Nucl.
Fusion 18 (1978) 4.
- [10] SCHMIDT, J.A., SAUTHOFF, N.R., HAWRYLUK, R.J., Phys. Rev.
Lett. 18 (1974) 1272.
- [11] COHEN, S.A., CECCHI, J.L., MARMAR, E.S., Phys. Rev. Lett. 35
(1975) 1507.
- [12] MARMAR, E.S., COHEN, S.A., CECCHI, J.L., PPPL MATT-1246 (1976).
- [13] MARMAR, E.S., Ph.D. dissertation, Princeton University (1976).

- [40] VAINSHTEIN, L.A., SOBELMAN, I.I., YUKOV, E.A., Cross Sections of Excitation of Atoms and Ions by Electrons, Nako, Moscow (1973).
- [41] DRAWIN, H.W., Z. Physik 225 (1969) 483.
- [42] VAN REGEMORTER, H., Astrophys. J. 138 (1962) 906.
- [43] SEATON, M.J., Atomic and Molecular Processes (ed. D. R. Bates), Academic Press, New York and London (1962).
- [44] SUCKEWER, S., HAWRYLUK, R.J., Phys. Rev. Lett. 40 (1978) 1649.
- [45] HSUAN, H., HAWRYLUK, R.J., SUCKEWER, S., TAKAHASHI, H., Proceedings of the Third Topical Conference on Radio Frequency Plasma Heating, Pasadena, CA, C8, (1978).
- [46] HIRSHMAN, S.P., to be published.
- [47] HIRSHMAN, S.P., Phys. Fluids 21 (1978) 224.
- [48] KUNZE, H.J., Spac. Science Rev. 13 (1972) 565.
- [49] DATLA, R.U., BLAHA, M., KUNZE, H.J., Phys. Rev. A12 (1975) 1076.
- [50] DATLA, R.U., NUGENT, L., JR., GRIEM, H.R., Phys. Rev. A14 (1976) 979.
- [51] BEJGMAN, I.L., VAINSHTEIN, L.A., Sov. J.—Zh. Exp. Theor. Fiz. 52 (1967) 185.
- [52] KUNZE, H.J., JOHNSTON III, W.D., Phys. Rev. A3 (1971) 1384.
- [53] JOHNSTON III, W.D., KUNZE, H.J., Phys. Rev. A4 (1971) 962.
- [54] KUNZE, H.J., GABRIEL, A.H., GRIEM, H.R., Phys. Rev. Lett. 39 (1977) 1256.
- [55] ELTON, R.C., KÖPPENDORFER, W.W., Phys. Rev. 160 (1967) 194.

- [56] TAYLOR, P.O., GREGORY, D., DUNN, G.H., PHANEUF, R.A.,
GRANDALL, D.H., Phys. Rev. Lett. 39 (1977) 1256.
- [57] HINNOV, E., Private communications.
- [58] BELL, O., Proc. Phys. Soc. 88 (1966) 587; Ann. Astrophys. 29
(1966) 683.
- [59] SPITZER, L., JR., Astrophys. J. 107 (1948) 6.
- [60] WISE, W.L., SMITH, M.W., GLENNON, Atomic Transition
Probabilities, National Bureau of Standards, Washington (1966).
- [61] MOORE, C.E., Atomic Energy Levels, VI, National Bureau of
Standards 467, Washington (1949).
- [62] STRIGANOV, A.R., SVENTITSKII, N.S., Tables of Spectral
Lines of Neutral and Ionized Atoms (translated from Russian),
New York-Washington (1968).
- [63] KELLY, R.L., PALUMBO, L.J., Atomic and Ionic Emission Lines
Below 2000 Angstroms, NRL Report 7599, Washington (1973).

FIGURE CAPTIONS

FIG. 1. Density and temperature profiles for the PLT Tokamak corresponding to a limiter temperature of 30 eV and an edge density of $1 \times 10^{13} \text{ cm}^{-3}$ at the limiter location ($r = 40 \text{ cm}$).

FIG. 2. Electron and ion temperature profiles corresponding to an electron and ion temperature of 15, 30 and 60 eV at the limiter ($r = 40 \text{ cm}$).

FIG. 3. Electron density profiles corresponding to an edge density of 5×10^{12} , 1×10^{13} , $2 \times 10^{13} \text{ cm}^{-3}$ at the limiter ($r = 40 \text{ cm}$).

FIG. 4a,b. Profiles of oxygen ionization states and resonance line radiation assuming neoclassical transport ($D_A = 0$) for the conditions corresponding to Fig. 1.

FIG. 5a,b. Profiles of carbon ionization states and resonance line radiation assuming neoclassical transport ($D_A = 0$) for the condition corresponding to Fig. 1.

FIG. 6a,b. Neoclassical fluxes of O^{5+} and O^{8+} , respectively (dashed curves) due to the gradients of the impurity (j^+), background plasma (D) and ion temperature (T_i) for the profiles corresponding to Fig. 4.

FIG. 7 Z_{eff} profiles corresponding to impurity profiles in Figs. 4 and 5 ($D_A = 0$).

FIG. 8. Particle recycling time as a function of anomalous diffusion coefficient.

FIG. 9a,b. Profiles of oxygen ionization states and resonance line radiation assuming neoclassical and anomalous transport ($D_A = 4000 \text{ cm}^2/\text{sec}$) corresponding to Fig. 1.

FIG. 10a,b. Profiles of carbon ionization states and line radiation assuming neoclassical and anomalous transport ($D_A = 4000 \text{ cm}^2/\text{sec}$) corresponding to Fig. 1.

FIG. 11a,b. Resonance line radiation for neoclassical and anomalous transport ($D_A = 4000 \text{ cm}^2/\text{sec}$) using the density profiles shown in Fig. 3 and the temperature profile in Fig. 1.

FIG. 12a,b. Resonance line radiation for neoclassical and anomalous transport ($D_A = 4000 \text{ cm}^2/\text{sec}$) using the temperature profiles shown in Fig. 2 and the density profiles in Fig. 1.

FIG. 13. The particle recycling time ($\Gamma_{\text{TOT}}/n_{\text{TOT}}$) as a function of anomalous diffusion coefficient for the temperature and density profiles shown in Figs. 1, 2 and 3.

FIG. 14. Power loss (resonant and non-resonant line radiation and ionization loss) as a function of diffusion coefficient for the density and temperature profiles shown in Figs. 1, 2 and 3.

FIG. 15a,b. Power radiated by the OVI 150 $\text{\AA}$ line and OVI resonance line as a function of outgoing flux in the 6^+ , 7^+ and 8^+ states.

TABLE I

Calculation of Ion Content Using Lotz and Drawin
Ionization Rate Coefficients ($D_A = 4000 \text{ cm}^2/\text{sec}$)

Ionization State	Lotz [36] (ions 10^{16})	Drawin [41] (ions 10^{16})
1+	3.6	12.0
2+	2.4	5.4
3+	3.2	5.5
4+	5.6	4.4
5+	7.4	6.6
6+	75.0	65.0
7+	50.0	50.0
8+	201.0	201.0

TABLE II

Calculation of Oxygen Impurity Ions For
Three Different Recombination Rate
Coefficients ($D_A = 4000 \text{ cm}^2/\text{sec}$)

Ionization State	$\delta \times .2$ (ions 10^{16})	$\delta \times 1$ (ions 10^{16})	$\delta \times 5$ (ions 10^{16})
1+	3.6	3.6	3.4
2+	2.4	2.4	2.4
3+	2.6	3.2	4.0
4+	5.4	5.6	8.8
5+	7.4	7.4	11.3
6+	6.2	75.0	152.0
7+	33.0	50.0	50.0
8+	231.0	201.0	116.0

TABLE III

Calculation of Line Radiation Assuming Different
Excitation Rate Coefficients Using the Impurity
Profiles in Fig. 9

Ionization State	Vainshtein, et al Ref. [40]		Van Regemorter Ref. [42]		Seaton Ref. [43]	
	Res (kW)	Total (kW)	Res (kW)	Total (kW)	Res (kW)	Total (kW)
0	1.0	1.3	0.26	0.88	0.55	2.10
1+	11.0	27.0	16.0	35.0	22.0	32.0
2+	20.0	45.0	5.0	22.0	6.0	27.0
3+	68.0	102.0	19.0	58.0	8.4	26.0
4+	75.0	82.0	64.0	101.0	73.0	121.0
5+	58.0	76.0	61.0	121.0	24.0	53.0
6+	30.0	54.0	64.0	67.0	85.0	97.0
7+	33.0	40.0	91.0	110.0	44.0	53.0
Total	296.0	427.0	320.0	515.0	263.0	411.0

TABLE IV

Comparison of Power Loss and Power Radiated
of the OVI Resonance Line

D_A (cm^2/sec)	$n_e(a)$ (10^{13} cm^{-3})	$T_e(a)$ (eV)	Power Loss (kW)	OVI Resonance Radiation (kW)	Power Loss OVI Resonance Line Rad.
0	1.0	30	120	7.1	17.0
$2. \times 10^3$	1.0	30	304	37.0	8.2
$4. \times 10^3$	1.0	30	464	57.0	8.1
$6. \times 10^3$	1.0	30	588	71.0	8.3
$8. \times 10^3$	1.0	30	687	81.0	8.5
$1. \times 10^4$	1.0	30	768	89.0	8.6
$2. \times 10^3$	0.5	30	185	19.0	9.7
$4. \times 10^3$	0.5	30	294	36.0	8.2
$8. \times 10^3$	0.5	30	478	58.0	8.2
$2. \times 10^3$	2.0	30	723	89.0	8.1
$4. \times 10^3$	2.0	30	933	105.0	8.9
$8. \times 10^3$	2.0	30	1230	132.0	9.3
$2. \times 10^3$	1.0	15	312	39.0	8.0
$4. \times 10^3$	1.0	15	462	59.0	7.8
$8. \times 10^3$	1.0	15	675	84.0	8.0
$2. \times 10^3$	1.0	60	333	26.0	12.8
$4. \times 10^3$	1.0	60	506	41.0	12.3
$8. \times 10^3$	1.0	60	743	62.0	12.0

TABLE V

Comparison of Power Loss by Carbon Impurities and
Power Radiated by the CIV Resonance Line

D_A (cm^2/sec)	$n_e(a)$ (10^{13} cm^{-3})	$T_e(a)$ (eV)	Power Loss (kW)	CIV Resonance Radiation (kW)	<u>Power Loss</u> CIV Resonance Line Rad.
0	1.0	30	71.0	6.6	10.8
2×10^3	1.0	30	330.0	42.0	7.8
4×10^3	1.0	30	467.0	63.0	7.4
8×10^3	1.0	30	611.0	89.0	6.9

Table VI. Carbon Transition ($Z = 6$)

Ion, I_1 (eV) Ground Configuration	Transitions ($m \rightarrow n$)	E^m (eV)	E^{mn} (eV)	λ (Å)
CI 11.27 $1s^2 2s^2 2p^2 3p_0$	$2p^2 3p - 3s 3p^0$ (1 $\rightarrow$ 2)	0.0	7.48	{ 1658.12 - 1656.27 (6 lines)
	$2p^2 1s - 3s 1p$ (1* $\rightarrow$ 2*)	2.68	5.00	2478.55
	$2p^2 3p - 2p^3 3p^0$ (1 $\rightarrow$ 3)	0.0	9.33	{ 1329.59 1329.10 1328.10
	$2p^2 3p - 3d 3d^0$ (1 $\rightarrow$ 4)	0.0	9.71	{ 1277.73 1277.55 1277.28
CII 24.38 $1s^2 2s^2 2p^2 3p_{1/2}^0$	$2p^2 3p^0 - 2p^2 2d$ (1 $\rightarrow$ 2)	0.0	9.29	{ 1335.68 1334.52
	$2p^1 2p^0 - 2p^2 2s$ (1 $\rightarrow$ 3)	0.0	11.90	{ 1037.02 1036.33
	$2p^2 3p^0 - 2p^2 3p$ (1 $\rightarrow$ 4)	0.0	13.71	{ 904.47 - - 903.61 (4 lines)
	$2p^2 3p^0 - 3s 2s$ (1 $\rightarrow$ 5)	0.0	14.45	{ 858.56 858.09

Table VI. Carbon Transition ($Z = 6$) Continued

Ion, I_1 (eV) Ground Configuration	Transitions ($m \rightarrow n$)	E^m (eV)	E^{mn} (eV)	$\lambda(\text{\AA})$
	$2p^2 P^0 - 3d^2 D$ (1 + 6)	0.0	18.04	$\left\{ \begin{array}{l} 687.36 \\ 687.06 \end{array} \right.$
C III 47.88 $1s^2 2s^2 1S_0$	$2s^2 1S - 2p^1 P^0$ (1 + 2)	0.0	12.69	977.03
	$2s^2 1S - 3p^1 P^0$ (1 + 3)	0.0	32.10	386.20
	$2s^2 1S - 3s^1 P$ (1 + 4)	0.0	38.43	322.58
	$2s^2 1S - 4p^1 P$ (1 + 5)	0.0	39.97	310.17
	$2p^1 P^0 - 2p^2 1D$ (2 + 6)	12.69	5.40	2296.87
	$2p^3 P^0 - 2p^2 3P$ (1* + 2*)	6.49	10.55	$\left\{ \begin{array}{l} 1176.37 - \\ - 1174.93 \end{array} \right.$ (6 lines)
	$2p^3 P^0 - 3s^3 S$ (1* + 3*)	6.49	23.04	$\left\{ \begin{array}{l} 538.31 \\ 538.15 \\ 538.08 \end{array} \right.$
	$2p^3 P^0 - 3d^3 D$ (1* + 4*)	6.49	26.98	$\left\{ \begin{array}{l} 459.63 \\ 459.52 \\ 459.46 \end{array} \right.$

Table VI. Carbon Transition ($Z = 6$) Continued

Ion, I_1 (eV) Ground Configuration	Transitions ($m + n$)	E^m (eV)	E^{mn} (eV)	λ (Å)
CIV 64.49 $1s^2 2s^2 2s_{1/2}$	$2s^2 S - 2p^2 P^0$ (1 + 2)	0.0	8.00	{ 1550.77 1548.19
	$2s^2 S - 3p^2 P^0$ (1 + 3)	0.0	39.68	{ 312.45 312.42
	$2s^2 S - 4p^2 P^0$ (1 + 4)	0.0	50.62	244.91
	$2p^2 P^0 - 3s^2 S$ (2 + 5)	8.0	29.55	{ 419.71 419.53
	$2p^2 P^0 - 3d^2 D$ (2 + 6)	8.0	32.28	{ 384.18 384.03
CV 392.067 $1s^2 1S_0$	$1s^2 1S - 2p^1 P^0$ (1 + 2)	0.0	307.87	40.27
	$1s^2 1S - 3p^1 P^0$ (1 + 3)	0.0	354.49	34.97
	$1s^2 1S - 4p^1 P^0$ (1 + 4)	0.0	370.90	33.43
	$1s^2 1S - 5p^1 P^0$ (1 + 5)	0.0	378.51	32.75
	$2s^3 S - 2p^3 P^0$ (2* + 3*)	298.94	5.44	{ 2270.91 2277.25 2277.92

Table VI. Carbon Transition ($Z = 6$) Continued

Ion, I_1 (eV) Ground Configuration	Transitions ($m \rightarrow n$)	E^m (eV)	E^{mn} (eV)	λ (Å)
C VI 489.946 $1s^2 S_{1/2}$	$1s^2 S - 2p^2 P^0$ (1 $\rightarrow$ 2)	0.0	367.44	33.73
	$1s^2 S - 3p^2 P^0$ (1 $\rightarrow$ 3)	0.0	435.51	28.46
	$1s^2 S - 4p^2 P^0$ (1 $\rightarrow$ 4)	0.0	459.33	26.99
	$1s^2 S - 5p^2 P^0$ (1 $\rightarrow$ 5)	0.0	470.35	26.35

Table VII. Oxygen Transition ($Z = 8$)

Ion, I_1 (eV) Ground Configuration	Transitions ($m \rightarrow n$)	E^m (eV)	E^{mn} (eV)	λ (Å)
OI 13.61 $1s^2 2s^2 2p^4 3p_2$	$2p^4 3p - 3s^3 S^0$ (1 + 2)	0.0	9.52	{ 1306.03 1304.87 1302.00
	$2p^4 3p - 3d^3 D^0$ (1 + 3)	0.0	12.08	{ 1028.16 1027.43 1025.77
	$2p^4 3p - 3s^3 D^0$ (1 + 4)	0.0	12.54	{ 990.81 - 988.58 (6 lines)
	$2p^4 3p - 3s^3 P^0$ (1 + 5)	0.0	14.12	{ 879.55 - 877.80 (6 lines)
	$2p^4 3p - 3d^3 P^0$ (1 + 6)	0.0	15.28	{ 812.16 - 810.67 (6 lines)
OII 35.15 $1s^2 2s^2 2p^3 4s_{3/2}^0$	$2p^3 4s^0 - 2p^4 4p$ (1 + 2)	0.0	14.97	{ 834.46 833.33 832.75
	$2p^3 4s^0 - 3s^4 p$ (1 + 3)	0.0	22.98	{ 539.85 539.55 539.09

Table VII. Oxygen Transition ($Z = 8$) 'Continued

Ion, I_1 (eV) Ground Configuration	Transitions ($m \rightarrow n$)	E^m (eV)	E^{mn} (eV)	$\lambda(\text{\AA})$
	$2p^3 4s^0 - 3d^4 p$ (1 $\rightarrow$ 4)	0.0	28.82	{ 430.18 - - 429.56 (6 lines)
OIII 54.93 $1s^2 2s^2 2p^2 3p_0$	$2p^2 3p - 2p^3 3d^0$ (1 $\rightarrow$ 2)	0.0	14.88	{ 835.29 - - 832.93 (4 lines)
	$2p^2 3p - 2p^3 3p^0$ (1 $\rightarrow$ 3)	0.0	17.65	{ 703.85 - - 702.33
	$2p^2 3p - 2p^3 3s^0$ (1 $\rightarrow$ 4)	0.0	24.43	{ 508.18 507.68 507.39
	$2p^2 3p - 3s^3 p^0$ (1 $\rightarrow$ 5)	0.0	33.16	{ 374.44 - - 373.81 (6 lines)
	$2p^2 3p - 3d^3 d^0$ (1 $\rightarrow$ 6)	0.0	40.57	{ 305.88 - - 305.60 (6 lines)
	$2p^2 3p - 3d^3 p^0$ (1 $\rightarrow$ 7)	0.0	40.84	{ 303.80 - - 303.41 (6 lines)

Table VII. Oxygen Transition ($Z = 8$) Continued

Ion, I_1 (eV) Ground Configuration	Transitions ($m \rightarrow n$)	E^m (eV)	E^{mn} (eV)	$\lambda(\text{\AA})$
OIV 77.39 $1s^2 2s^2 2p^2 2p_{1/2}^0$	$2p^2 P^0 - 2p^2 D$ (1 $\rightarrow$ 2)	0.0	15.74	$\left\{ \begin{array}{l} 790.20 \\ 790.10 \\ 787.71 \end{array} \right.$
	$2p^2 P^0 - 2p^2 S$ (1 $\rightarrow$ 3)	0.0	20.38	$\left\{ \begin{array}{l} 609.83 \\ 608.40 \end{array} \right.$
	$2p^2 P^0 - 2p^2 P$ (1 $\rightarrow$ 4)	0.0	22.39	$\left\{ \begin{array}{l} 555.26 - \\ - 553.33 \end{array} \right.$ (4 lines)
	$2p^2 P^0 - 3s^2 S$ (1 $\rightarrow$ 5)	0.0	44.34	$\left\{ \begin{array}{l} 279.94 \\ 279.63 \end{array} \right.$
	$2p^2 P^0 - 3d^2 D$ (1 $\rightarrow$ 6)	0.0	52.00	$\left\{ \begin{array}{l} 238.57 \\ 238.36 \end{array} \right.$
OV 113.87 $1s^2 2s^2 1s_0$	$2s^2 S - 2p^1 P^0$ (1 $\rightarrow$ 2)	0.0	19.69	629.73
	$2s^2 S - 3p^1 P^0$ (1 $\rightarrow$ 3)	0.0	72.01	172.17
	$2p^1 P^0 - 3p^1 S$ (2 $\rightarrow$ 4)	19.69	68.04	182.21
	$2p^3 P^0 - 3p^3 S$ (1* $\rightarrow$ 2*)	10.20	74.62	$\left\{ \begin{array}{l} 166.23 \\ 166.15 \\ 166.11 \end{array} \right.$

Table VII. Oxygen Transition ($Z = 8$) Continued

Ion, I_1 (eV) Ground Configuration	Transitions ($m \rightarrow n$)	E^m (eV)	$-E^{mn}$ (eV)	$\lambda(\text{\AA})$
OVI 138.08 $1s^2 2s^2 2s_{1/2}$	$2s^2 S - 2p^2 P^0$ (1 + 2)	0.0	12.0	{ 1037.61 1031.91
	$2s^2 S - 3p^2 P^0$ (1 + 3)	0.0	82.58	{ 150.12 150.09
	$2p^2 P^0 - 3s^2 S$ (2 + 4)	12.08	79.35	{ 184.11 183.94
	$2p^2 P^0 - 3d^2 D$ (2 + 5)	12.0	83.64	{ 173.08 172.94
	$1s^2 1S - 2p^1 P^0$ (1 + 2)	0.0	573.91	21.60
	$1s^2 1S - 3p^1 P^0$ (1 + 3)	0.0	665.58	18.63
OVIII 871.39 $1s^2 2s_{1/2}$	$1s^2 S - 2p^2 P^0$ (1 + 2)	0.0	653.5	18.97
	$1s^2 S - 3p^2 P^0$ (1 + 3)	0.0	774.5	16.01
	$1s^2 S - 4p^2 P^0$ (1 + 4)	0.0	816.9	15.18
	$1s^2 S - 5p^2 P^0$ (1 + 5)	0.0	836.4	14.82

Fig. 1. 783821

Fig. 2. 783949

Fig. 3. 783818

Fig. 4a,b. 786044

Fig. 5a,b. 783932

Fig. 6a,b. 783930

Fig. 7. 783819

Fig. 8. 783824

Fig. 9a,b. 783928

Fig. 10a,b. 783842

Fig. 11a,b. 786045

Fig. 12a, b. 783933

Fig. 13. 783802

Fig. 14. 783820

Fig. 15. 783948

APPENDIX A: NEOCLASSICAL FLUXES

The collisional radial particle flux for an impurity ion "a" in a large aspect ratio tokamak plasma can be expressed as:

$$\Gamma_a^{\text{col}} = \Gamma_a^{\text{cl}} + \Gamma_a^{\text{BP}} + \Gamma_a^{\text{PS}} \quad (\text{A1})$$

Here

$$\Gamma_a^{\text{cl}} = \left\langle \frac{\vec{R}_a \cdot \vec{B} \times \vec{i}_r}{e_a B^2} \right\rangle \quad (\text{A2a})$$

is the classical flux driven by the perpendicular component of the friction force R_a ;

$$\Gamma_a^{\text{BP}} = - \left\langle \frac{\vec{B} \cdot \nabla \cdot \vec{\pi}_a}{e_a B B_p} \right\rangle \quad (\text{A2b})$$

is the long mean free path (banana-plateau regime) contribution to the neoclassical flux, driven by the parallel component of the viscous stress $\nabla \cdot \vec{\pi}_a$;

$$\Gamma_a^{\text{PS}} = - \left\langle \frac{\vec{R}_a \cdot \vec{B}}{e_a B B_p} \left(1 - \frac{B^2}{\langle B^2 \rangle} \right) \right\rangle \quad (\text{A2c})$$

is the Pfirsch-Schlüter flux, driven by the variation of the friction force on a magnetic surface. The surface average is defined as $\langle A \rangle \equiv \int_0^{2\pi} (1 + \epsilon \cos \theta) A d\theta / 2\pi$, where $\epsilon = r/R$ is the inverse aspect ratio.

These fluxes were calculated by generalizing the results of [29-31, 46,47]. In these calculations, collisions between different ionization states as well as with the (hydrogenic)

background plasma (denoted by H) were taken into account.

The fluxes are expressed as linear combinations of the thermodynamic forces A_j^a , where in cylindrical coordinates:

$$A_1^a = \frac{\partial}{\partial r} \ln p_a \quad (A3a)$$

$$A_2^a = \frac{\partial}{\partial r} \ln T_a \quad (A3b)$$

The impurity Ware fluxes, which are smaller by $\beta_p^{-1} (m_e/m_i)^{1/2}$ than the gradient-driven fluxes, have been neglected.

CLASSICAL FLUXES

The classical flux is:

$$\Gamma_a^{cl} = \sum_{j=1,2;b} \frac{1}{2} \rho_a^2 \frac{T_b e_a}{T_a e_b} \frac{\tilde{\ell}_{1j}^{ab}}{m_a} A_j^b, \quad (A4)$$

where $\rho_a = c \sqrt{2m_a T_a} / e_a B$ is the Larmor radius and $\tilde{\ell}_{ij}^{ab}$ are the classical friction coefficients. Using the results of Ref. [46], the classical friction coefficients can be expressed as:

$$\begin{aligned} \tilde{\ell}_{11}^{ab} = & - \delta_{ab} \sum_k \frac{m_a n_a}{\tau_{ak}} \left(1 + \frac{m_a}{m_k} \right) (1 + x_{ak}^2)^{-3/2} \\ & + \frac{m_a n_a}{\tau_{ab}} \left(1 + \frac{m_a}{m_b} \right) (1 + x_{ab}^2)^{-3/2} \end{aligned} \quad (A5a)$$

$$\begin{aligned} \tilde{\ell}_{12}^{ab} = & \frac{3}{2} \delta_{ab} \sum_k \frac{m_a n_a}{\tau_{ak}} \left(1 + \frac{m_a}{m_k} \right) (1 + x_{ak}^2)^{-5/2} \\ & - \frac{3}{2} \frac{m_a n_a}{\tau_{ab}} \left(1 + \frac{m_a}{m_b} \right) x_{ab}^2 (1 + x_{ab}^2)^{-5/2} \end{aligned} \quad (A5b)$$

In these expressions,

$$\tau_{ab} = \frac{3\sqrt{m_a} T_a^{3/2}}{n_a e_a^2 e_b^2 4\sqrt{2\pi} \ln \Lambda} \quad , \quad (A6a)$$

$$x_{ab} = \left(\frac{T_b m_a}{T_a m_b} \right)^{1/2} \quad (A6b)$$

Typically, the classical fluxes are small compared with the Pfirsch-Schlüter fluxes except near the plasma axis.

ANISOTROPY-DRIVEN NEOCLASSICAL FLUXES

The long mean free path contribution to the neoclassical particle flux has been obtained by evaluating the expressions derived in [47].

$$\Gamma_a^{BP} = \sum_{j=1,2;b} \left[.73 q^2 \left(\frac{R}{r} \right)^{3/2} \frac{e_b}{e_a} \rho_b^2 \frac{\hat{\mu}_{bj}}{m_b} \left(\frac{\hat{\mu}_{aj}}{\sum_k \hat{\mu}_{kl}} - \delta_{ab} \right) \right] A_j^b \quad , \quad (A7)$$

where $q = rB_T/RB_p$ is the safety factor and $\hat{\mu}_{aj}$ is the normalized viscosity coefficient:

$$\hat{\mu}_{aj} \equiv \frac{n_a m_a}{\tau_{aa}} \left\{ \left(x^2 - \frac{5}{2} \right)^{j-1} v_T^a(x) \tau_{aa} \right\} \quad . \quad (A8)$$

Curly brackets denote integration over a Maxwellian [47]. The collision frequency integrals can be expressed as:

$$\left\{ v_T^a \tau_{aa} \right\} \equiv \beta_{a1} \quad (A9a)$$

$$\left\{ v_T^a \tau_{aa} \left(x^2 - \frac{5}{2} \right) \right\} \equiv \beta_{a2} - \frac{5}{2} \beta_{a1} \quad , \quad (A9b)$$

where β_{aj} are positive constants. Approximate formulae for β_{aj} are obtained by matching to the asymptotic behavior of β_{aj} in various collision frequency regimes, yielding:

$$\beta_{aj} = \frac{\beta_{aj}^B}{1 + \beta_{aj}^B/\beta_{aj}^P} \frac{1}{1 + \beta_{aj}^P/\beta_{aj}^{PS}}, \quad (A10)$$

where:

$$\beta_{a1}^B = \sum_b \frac{n_b \epsilon_b^2}{n_a \epsilon_a^2} \left[(1 + x_{ab}^2)^{1/2} + x_{ab}^2 \ln \left\{ \frac{x_{ab}}{1 + (1 + x_{ab}^2)^{1/2}} \right\} \right] \quad (A11a)$$

$$\beta_{a2}^B = \sum_b \frac{n_b \epsilon_b^2}{n_a \epsilon_a^2} (1 + x_{ab}^2)^{-1/2} \quad (A11b)$$

$$\beta_{a1}^P = \left(\frac{\sqrt{\pi}}{2} \right) \left(\frac{1}{1.46} \right) v_{*a}^{-1}, \quad (A11c)$$

$$\beta_{a2}^P = 3 \beta_{a1}^P, \quad (A11d)$$

$$\beta_{a1}^{PS} = \frac{3}{4} \left(\frac{1}{1.46} \right) \frac{l_{a1}}{v_{*a}^2 \epsilon^{3/2}}, \quad (A11e)$$

$$\beta_{a2}^{PS} = \frac{3}{4} \left(\frac{1}{1.46} \right) \left(\frac{l_{a2} + (5/2) l_{a1}}{v_{*a}^2 \epsilon^{3/2}} \right). \quad (A11f)$$

Here, $\epsilon = r/R$, $v_{*a} = (\tau_{aa} \omega_{ta})^{-1} \epsilon^{-3/2}$ is the collisionality parameter, and $\omega_{ta} = (2T_a/m_a)^{1/2} (Rq)^{-1}$ is the transit frequency. The quantities l_{aj} are the collisional viscous coefficients given by:

$$l_{a1} = \frac{5}{2} I_{11}^a \left[I_{11}^a I_{00}^a - \left(I_{01}^a \right)^2 \right]^{-1}, \quad (A12a)$$

$$l_{a2} = \frac{5}{2} \left(I_{11}^a + \frac{7}{2} I_{01}^a \right) \left[I_{11}^a I_{00}^a - \left(I_{01}^a \right)^2 \right]^{-1}. \quad (A12b)$$

Here:

$$r_{jk}^a \equiv \sum_b \left(\frac{n_b e_b^2}{n_a e_a^2} \right) \left[\left(v_{jk}^{ab} - w_{jk}^{ab} \delta_{ab} \right) \right] \quad (A13)$$

and

$$v_{00}^{ab} = \left(3 + 5 x_{ab}^2 \right) \left(1 + x_{ab}^2 \right)^{-3/2} \quad (A14a)$$

$$w_{00}^{aa} = 1/\sqrt{2} \quad (A14b)$$

$$v_{01}^{ab} = \frac{3}{2} \left(3 + 7 x_{ab}^2 \right) \left(1 + x_{ab}^2 \right)^{-5/2} \quad (A14c)$$

$$w_{01}^{aa} = \frac{3}{2} \frac{1}{\sqrt{2}} \quad (A14d)$$

$$v_{11}^{ab} = \left(35 x_{ab}^6 + 38.5 x_{ab}^4 + \frac{185}{4} x_{ab}^2 + \frac{51}{4} \right) \left(1 + x_{ab}^2 \right)^{-7/2} \quad (A14e)$$

$$w_{11}^{aa} = \frac{15}{4} \frac{1}{\sqrt{2}} \quad (A14f)$$

Except near the center of the discharge, the anisotropy driven fluxes are also typically much less than the Pfirsch-Schlüter fluxes.

Pfirsch-Schlüter Fluxes

The Pfirsch-Schlüter fluxes were evaluated by extending the results of Ref. [29] to the case where the background plasma ions are not in the Pfirsch-Schlüter regime. The results of this calculation are valid in both the "intermediate" [30,31] as well as "extreme" Pfirsch-Schlüter regime. In this section, the notation of Ref. [29] will be employed.

In order to calculate r_a^{PS} it is necessary to determine the parallel friction coefficients ℓ_{ij}^{ab} , which relate the friction forces R_a , H_a to the parallel particle and heat flow:

$$R_a = \sum_b \left[\ell_{11}^{ab} u_{\parallel b} + \frac{2}{5} \ell_{12}^{ab} \left(\frac{q_{\parallel b}}{P_b} \right) \right] \quad (A15a)$$

$$H_a = \sum_b \left[\ell_{21}^{ab} u_{\parallel b} + \frac{2}{5} \ell_{22}^{ab} \left(\frac{q_{\parallel b}}{P_b} \right) \right] \quad (A15b)$$

In general, however, the following relations pertain for all collision frequencies:

$$R_a = \sum_b \left[\tilde{\ell}_{11}^{ab} u_{\parallel b} + \frac{2}{5} \tilde{\ell}_{12}^{ab} \frac{q_{\parallel b}}{P_b} + \tilde{\ell}_{13}^{ab} u_{2b} \right] \quad (A16a)$$

$$H_a = \sum_b \left[\tilde{\ell}_{21}^{ab} u_{\parallel b} + \frac{2}{5} \tilde{\ell}_{22}^{ab} \frac{q_{\parallel b}}{P_b} + \tilde{\ell}_{23}^{ab} u_{2b} \right] \quad (A16b)$$

Here, $\tilde{\ell}_{ij}^{ab}$ are the classical (perpendicular) friction coefficients for $i \leq 2$, $j \leq 2$, $\tilde{\ell}_{j3}^{ab} \propto \{L_j^{3/2} C_{ab} (L_2^{3/2} f_o)\}$, and $u_{2b} = \int v_{\parallel} L_2^{3/2} f_b dv$, where $L_j^{3/2}$ are the Sonine polynomials [29]. Thus, in order to obtain ℓ_{ij}^{ab} , it is necessary to express u_{2b} in terms of the lower order parallel flows ($u_{\parallel}$, $q_{\parallel}$). The relation $u_{2b} = f(u_{\parallel}, q_{\parallel}; v_*)$ depends on collision frequency regime. Indeed, it is shown [31] that u_{2b} satisfies the coupled equations:

$$u_{a2} \left[1 + \frac{\omega_{ta}^2}{v_a v_E} \right] - \sum_b \Delta_{ab} u_{b2} = - \sum_b \left[\alpha_{ab} u_{\parallel b} + \beta_{ab} \frac{2}{5} \frac{q_{\parallel b}}{P_b} \right] \quad (A17)$$

where α_{ab} , β_{ab} , and Δ_{ab} are defined in [29], and v_S^a , v_E^a are momentum and energy relaxation rates:

$$v_E^a \tau_{aa} = 2 \left(\frac{8}{35} \right)^2 (m_a^{22} + n_{aa}^{22}) \quad (A18a)$$

$$v_S^a \tau_{aa} = (M_a^{22} + N_{aa}^{22}) \quad (A18b)$$

where

$$m_a^{22} + n_{aa}^{22} = 3 \sum_b x_{ab}^2 \frac{n_b^2 e_b^2}{n_a^2 e_a^2} \left[\left(\frac{5}{2} x_{ab}^4 + 2 x_{ab}^2 + \frac{13}{4} \right) \left(1 + x_{ab}^2 \right)^{-7/2} - \delta_{ab} \frac{15}{32} \frac{1}{\sqrt{2}} \right] \quad (A19a)$$

$$M_a^{22} + N_{aa}^{22} = \sum_b \frac{n_b^2 e_b^2}{n_a^2 e_a^2} \left[\left(21.9 x_{ab}^8 + 28 x_{ab}^6 + 57.4 x_{ab}^4 + 17 x_{ab}^2 + 6.8 \right) \left(1 + x_{ab}^2 \right)^{-9/2} - 1.81 \delta_{ab} \right] \quad (A19b)$$

Thus, solving for u_{a2} in the limit of small Δ yields:

$$u_{a2} = - \sum_b \left(\hat{\alpha}_{ab} u_{\parallel b} + \frac{2}{5} \hat{\beta}_{ab} \frac{q_{\parallel b}}{R_b} \right) \quad (A20)$$

where

$$\hat{\alpha}_{ab} = w_a \left(\alpha_{ab} + \sum_c w_c \Delta_{ac} \alpha_{cb} \right) \quad (A21a)$$

$$\hat{\beta}_{ab} = w_a \left(\beta_{ab} + \sum_c w_c \Delta_{ac} \beta_{cb} \right) \quad (A21b)$$

and

$$w_a \equiv \left(1 + \frac{w_a^2}{v_S^a v_E^a} \right)^{-1} = \left[1 + 9.57 \epsilon^{-3} v_{*a}^{-2} (m_a^{22} + n_{aa}^{22})^{-1} (M_a^{22} + N_{aa}^{22})^{-1} \right]^{-1} \quad (A21c)$$

With these values for $\hat{\alpha}$, $\hat{\beta}$, Eqs. (20a-c) and all subsequent formulas in [29] pertain. In the limit $\Delta \rightarrow 0$, note that

$$\begin{pmatrix} \hat{\alpha}_{ab} \\ \hat{\beta}_{ab} \end{pmatrix} = W_a \begin{pmatrix} \alpha_{ab}^{\infty} \\ \beta_{ab}^{\infty} \end{pmatrix} \quad (A22)$$

where $\alpha^{\infty}, \beta^{\infty}$ are the collisional values in the extreme Pfirsch-Schlüter regime. As the collisionless regimes are approached, $W_a \rightarrow 0$ and $u_{2a} \rightarrow 0$, so the parallel friction coefficients tend to the perpendicular friction coefficients.

For heavy-on-light ion collisions, the prescription is simply to replace the $C_j^*(Z_a^*)$ functions in Eqs. (50a-c) of Ref. [29] with the following functions:

$$C_1^*(x; W_a) = 1 - \left(\frac{0.52x}{0.59 + x} \right) W_a \quad (A23a)$$

$$C_2^*(x; W_a) = \frac{3}{2} - \left(\frac{.29 + 1.20x}{0.59 + x} \right) W_a \quad (A23b)$$

$$C_3^*(x; W_a) = \sqrt{2} + \frac{13}{4} x - \left[\frac{(0.41 + 1.66x)^2}{0.59 + x} \right] W_a \quad (A23c)$$

where $x = Z_a^*$.

For like particle interactions, the D-functions in Eqs. (53a-d) of Ref. [29] are replaced by:

$$d_0 \equiv \sum_{k \neq H} D_0(\bar{Z}_k, W_k) \quad (A24a)$$

$$D_0(x, W_a) = \left(12 + \frac{1677}{96} x \right)^{-1} W_a \quad (A24b)$$

$$D_1(x, W_a) = 1 - \frac{15}{32} x D_0(x, W_a) \quad (A24c)$$

$$D_2(x, W_a) = \frac{3}{4} - (1.5 + 3.26x) D_0(x, W_a) \quad (A24d)$$

$$D_3(x, W_a) = 2 + 3.69x - (2.19 + 4.76x)^2 D_0(x, W_a) \quad (A24e)$$

Here, $x = \bar{z}_a$. These formulas can also be used to compute the ΔD functions in [29].

Using the new C^* , D and ΔD functions, it is possible to evaluate ℓ_{ij}^{ab} . The strong temperature equilibration approximation corresponding to Eqs. (33-40) in [29] has been used here to calculate r_a^{PS} .

Typically, very near the plasma periphery where the impurities are closer to the extreme Pfirsch-Schlüter regime, the ion temperature gradient diffusion coefficients are small compared with the background plasma gradient. Nonetheless, within a few centimeters of the limiter the ion temperature gradient terms become significant and outward [30] (see Fig. 7). As discussed above, the inclusion of the transition from the intermediate to the extreme Pfirsch-Schlüter regimes does make an important difference in the recycling of the lower ionization states.

APPENDIX B: IONIZATION, RECOMBINATION, AND EXCITATION RATE COEFFICIENTS

The semi-empirical ionization rate coefficients, S_j , of Lotz [36] are given by:

$$S_j = 6.7 \times 10^9 \sum_{i=1}^2 \left(\frac{A_j^i \xi_j^i}{(kT_e)^{3/2}} \frac{1}{u_j^i} E_1(u_j^i) - \frac{B_j^i \exp(C_j^i)}{u_j^i + C_j^i} E_1(u_j^i + C_j^i) \right) (\text{cm}^3 \text{sec}^{-1}) \quad , \quad (\text{B1})$$

where

$$u_j^i = I_j^i / kT_e \quad (\text{B2})$$

and

$$E_1(U) = \int_U^{\infty} \frac{\exp(-x)}{x} dx \quad . \quad (\text{B3})$$

I_j^i and ξ_j^i are the ionization energy and the number of equivalent electrons of the i^{th} subshell, respectively. The coefficients A_j^i , B_j^i , and C_j^i are tabulated by Lotz, where for $j \geq 3$, $A_j^{1,2} = 4.5 \times 10^{-14} \text{ cm}^2 \text{ev}^2$, $B_j^i = C_j^i = 0$. According to experimental results of Kunze, Griem, et al., [48,49,50], the rate coefficients of Lotz for lithium and beryllium-like light ions are a factor of 1.5-2.0 too high.

Vainshtein, Sobelman, and Yukov [40] evaluated the excitation rate coefficients by the Coulomb-Born Approximation for $\Delta S = 0$ transitions. For intercombination transitions ($\Delta S = 1$), the coefficients were evaluated using the modified Born-Openheimer

method [51]. The resulting expression for the excitation rate coefficients is:

$$S_j^{mn} = 10^{-8} \left(\frac{I_H I_j^n}{E_j^{nm} I_j^m} \right) \exp(-u_j^{nm}) \left(\frac{g}{2\ell_0 + 1} \right) g(u_j^{nm}) \quad (\text{cm}^3 \text{sec}^{-1}) \quad , \quad (\text{B4})$$

where

$$u_j^{nm} = E_j^{nm} / kT_e$$

and

$$E_j^{nm} = E_j^{mn} = |E_j^n - E_j^m|$$

is the excitation energy (in eV) with E_j^m and E_j^n being the energy of the n and m levels with respect to the ground state of ionization j . I_j^n is the ionization energy from level n of an ion in the j stage of ionization ($j = 0$ corresponds to a neutral) and I_H is the energy of ionization of hydrogen (both in eV). The function $g(u)$ is given by

$$\left. \begin{aligned} g(u) &= \frac{[(u+1)u]^{1/2}}{u+\chi} \quad , \quad j = 0 \\ g(u) &= \frac{(u+1)u^{1/2}}{u+\chi} \quad , \quad j \geq 1 \end{aligned} \right\} \quad \Delta s \rightarrow 0 \quad (\text{B5})$$

and

$$\left. \begin{aligned} g(u) &= \frac{u^{3/2}}{(u+1)^{1/2}(u+\chi)} , \quad j = 0 \\ g(u) &= \frac{u^{3/2}}{u+\chi} , \quad j \geq 1 \end{aligned} \right\} \Delta s = 1 \quad . \quad (B6)$$

The coefficient q depends on the azimuthal quantum numbers l_0 and l of the electron before and after the transition. In [40] the method of calculation of q is given. For example, for a transition without a change in spin ($\Delta s = 0$) and $l_0^N + l_0^{N-1}$, q ; $q = N$. In [40] graphs and tables are given for A and χ for the most important transition starting from the principal quantum numbers $n = 1$ and $n = 2$ for isoelectronic sequences hydrogen to neon.

Vainshtein, et al. [40] excitation rate coefficients were compared with experimental measurements by Kunze, et al. [48, 52-54], Elton and Koppendorfer [55], Taylor, et al. [56], and Hinnov [57] for oxygen and carbon. In general Vainshtein, et al. excitation rate coefficients are in better agreement with the experiment than those by Van Regemorter [42], Seaton [43], Belly [58], or Drawin [41], especially for non-resonance transition.

Collisional de-excitation S_j^{nm} and three-body recombination S_j^n rate coefficients were calculated using the principle of detailed balance

$$S_j^{nm} = S_j^{mn} \frac{g_j^n}{g_j^n} \exp(u_j^{nm}) \quad (\text{cm}^3 \text{sec}^{-1}) \quad (\text{B7})$$

$$S_{j+1}^n = S_j^n \frac{g_j^n}{2g_{j+1}^n} \left(\frac{h^2}{2\pi m_e k T_e} \right)^{3/2} \exp(u_j^n) \quad (\text{cm}^6 \text{sec}^{-1}) \quad (\text{B8})$$

where g_j^n is the statistical weight of the level n .

Radiative recombination rate coefficient was calculated by using a modification [15, 27] of Seaton's formula [37],

$$\alpha_{j+1}^n = 2.6 \times 10^{-14} (1_{j+1}^n + 2_{j+1}^n) \quad (\text{cm}^3 \text{sec}^{-1}) \quad (\text{B9})$$

where

$$1_{j+1}^n = (j+1)^2 (u_H)^{1/2} \frac{u_j^n}{n^3} \exp(u_j^n) E_1(u_j^n) \quad (\text{B10})$$

$$2_{j+1}^n = 2j^2 (u_j^n)^{1/2} \phi_{n+1}(j^2 u_H) \quad (\text{B11})$$

The functions ϕ_{n+1} have been tabulated by Spitzer [59].

Dielectronic recombination rate coefficients were calculated by using Donaldson and Peacock evaluation [39] of Burgess's formulae [38]:

$$\alpha_{d,j+1} = 7.9 \times 10^{-12} (kT_e)^{-3/2} (j+1)^4 \theta(u_j^1, n_t) \psi(u_{j-1}^1, n_t) \quad (\text{B12})$$

where $n_t \approx 4$ for carbon and oxygen for $10^{11} \text{cm}^{-3} n_e \leq 10^{14} \text{cm}^{-3}$ and

$$n_4 \approx \left(1.46 \times 10^6 \frac{j^2}{(j+1)^6} \right)^{1/3.5} \quad (\text{B13})$$

Functions θ and ψ are tabulated in Ref. [39].

APPENDIX C: ADDITIONAL IONIZATION AND EXCITATION RATE COEFFICIENTS

Drawin [41] ionization rate coefficient was expressed as:

$$S_j = 5.75 \times 10^{-8} \left(\frac{I_H}{I_j} \right)^{3/2} \xi_j^1 (u_j^1)^{1/2} \psi_1(u_j^1, b_j) \exp(-u_j^1) \quad (C1)$$

where

$$\psi_1(u, b) = \exp(u) \int_u^\infty \left(1 - \frac{u}{x}\right) e^{-x} \ln(1.25 \frac{bx}{u}) dx \quad (C2)$$

and

$$b_j = 1 + \frac{j-1}{j+2} \quad (C3)$$

All of the other notations are those in Appendix B.

Van Regemorter excitation rate coefficients [42] using Belly [57] Gaunt factor $\bar{g}$ are given by:

$$S_j^{mn} = 7.2 \times 10^{-15} A_j^{nm} \frac{g_j^n}{g_j^m} \frac{I_H}{(E_j^{nm})^3} \left(\frac{I_j^n}{kT_e} \right)^{1/2} \bar{g} \exp(-u_j^{nm}) \quad (C4)$$

with $\bar{g} = A + (Bu + C) \exp(u) E_1(u)$.

For oxygen, for $j = 1, 3, 5$, and 7 , the coefficients are, respectively: $A = 0.3$, $B = 0.25$, $C = 0.25$, and for $j = 2, 4, 6$, and 8 , we have: $A = 0.5$, $B = 0.25$, and $C = 0.25$.

Seaton excitation rate coefficients [43] are given by

$$S_j^{mn} = 1.36 \times 10^{-13} \frac{\xi_j^{m, nm, n} A_j^{nm, n} g_j^n}{(E_j^{nm})^3 (kT_e)^{1/2} g_j^m} \exp(-u_j^{nm}) \quad (C5)$$

where the spontaneous transition probability, A_j^{nm} , was taken from [60].

APPENDIX D: TRANSITIONS STUDIED

The carbon and oxygen transitions (Tables VI and VII, respectively) which were considered in evaluating the power radiated were obtained using [61,62,63].

INTERNAL DISTRIBUTION IN ADDITION TO TIC UC-20

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