CONF-860522--5

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PARTICLE FLUX MEASUREMENTS WITH PUMP LIMITERS*

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CONF-860522--5

DE86 011069

ABSTRACT

The pressure buildup in pump limiters is a function of the incident particle flux, the pump limiter geometry, particle-surface interactions, and, under certain conditions, non-linear effects due to interactions of the neutral gas with the incoming plasma. At sufficiently low density and for a short pump limiter throat, the pressure observed in the pump limiter chamber is a direct measure of the particle flux at the limiter slot. This paper discusses a simple model for the correlation between particle flux and pump limiter pressure. In addition, some estimates are presented on the thermalization of fast neutrals in the pump limiter duct and its implications for "enhanced particle collection."

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*Research sponsored by the Office of Fusion Energy, U.S. Department of Energy, under Contract No.

DE-AC05-84OR21400 with Martin Marietta Energy Systems, Inc.

1. INTRODUCTION

The pressure buildup in pump limiters is a function of the particle flux incident on the limiter and deflected into the pumping chamber [1]. Modeling this process is complicated because of the complexity of the involved atomic and surface processes. The plasma particles incident on the neutralizer plate are subject to a variety of surface phenomena such as reflection, implantation, desorption, and sheath effects, which all depend on particle energy, material, temperature, surface state, etc. Neutrals re-emitted from the neutralizer plate encounter a large number of possible interactions with the plasma. Detailed modeling of all processes involved is complicated, and sophisticated numerical codes like DEGAS [2] have been developed to handle these complex interactions. However, the necessary complexity of these codes often obscures the governing physics processes that lead to a particular result. Therefore, it is in many cases useful to study the basic physics processes with global models and simple approximations. Although simple models lack the completeness of sophisticated codes, they can contribute to the understanding of the basic processes.

2. PARTICLE FLUX AND PRESSURE BUILDUP

We consider a pump limiter system in equilibrium; that is, particle fluxes and pressures are constant in time. The equilibrium is considered to be flux equilibrium rather than pressure equilibrium, since the plasma pressure outside the limiter is one to two orders of magnitude higher than the measured neutral pressure. We assume the densities to be low enough so that gas-plasma interactions in the limiter throat can be neglected. Furthermore, the system is not being pumped. In this case the incoming particle flux is neutralized and fills the pump limiter chamber with gas up to a pressure at which the backflow of neutrals equals the incident ion flux. This situation is depicted in Fig. 1. The incident particle flux is $\phi = \phi_0 \exp[-(x - h)/\lambda]$, with h and λ the height of the pumping slot and the decay length of the particle flux, respectively. The fraction $\eta\phi$ of the incident flux is collected in

the limiter chamber and is balanced by the backflow of neutrals $n_s \bar{v}_s/4$ at the slot. The particle balance is then

$$\eta w \phi_0 \int_0^h \exp[-(x-h)/\lambda] dx = 2hw \frac{n_s \bar{v}_s}{4}, \quad (1)$$

where the factor of 2 accounts for molecular particles and w is the width of the pump limiter slot.

The neutral flux of molecules at the slot $n_s \bar{v}_s/4$ is unknown, but it can be inferred from the measured neutral pressure in the pump limiter chamber. A pressure measurement (with an ion gauge) is usually a measurement of the neutral density combined with the assumption that the gas is at room temperature using the relation $p = nkT$. This is a prerequisite for our further considerations. For this case, the neutral flux can be determined: $n \bar{v}/4 = p/\sqrt{2\pi mkT}$ where m is the molecular mass and T is room temperature. Since there is flux equilibrium inside the pump limiter, the neutral backflow at the throat must be equal to the flux at the gauge $n_s \bar{v}_s = n \bar{v}$, even though n_s and $\bar{v}_s$ are not known independently. The incident flux ϕ_0 in Eq. (1) can then be expressed as

$$\phi_0 = \frac{2hp}{\eta \int_0^h \exp[-(h-x)/\lambda] dx \sqrt{2\pi mkT}}. \quad (2)$$

To determine ϕ_0 , the parameters λ and η must be known in addition to the measured pressure p . If λ is unknown, only the integrated ion flux into the slot can be determined. The remaining unknown is the collection probability η .

3. COLLECTION PROBABILITY

The collection probability η is that fraction of the incident flux ϕ that is scattered into the pump limiter chamber and contributes to the pressure buildup. The fraction $(1 - \eta)$

escapes from the chamber after the collision at the neutralizer plate. The collection probability is not to be confused with the trapping fraction γ , which is that part of the incident flux that is removed by pumping and does not return to the system. Without pumping, the trapping fraction is zero.

To estimate the collection probability, we must discuss the particle-surface interaction. Depending on the particle energy, which is determined by the sheath potential, a certain fraction of the incident ion flux is reflected as neutral atoms and the remaining fraction is implanted in the surface. For incident energies of 100 eV, typical particle reflection is $R_n \approx 0.5$ and energy reflection $R_E \approx 0.3$ [3]. The implanted particles eventually recombine and desorb as thermal molecules from the surface. Upon saturation of the surface ($\sim 4 \times 10^{17} \text{ cm}^{-2}$), each incident particle produces one returning particle. We assume that reflected particles as well as desorbed particles are emitted with a cosine distribution. With increasing energies, the reflected fraction decreases and the process becomes dominated by desorption. We assume that the plasma density is low enough that the particles reflected or desorbed from the plate do not interact with the plasma.

From Fig. 1 it is obvious that the collection probability is a function of x and the base b_1 or b_2 , respectively. Elementary geometry yields for the collected flux $\phi_{\text{coll}} = \eta\phi$

$$\phi_{\text{coll}} = \frac{1}{2} \phi_0 \int_0^h \exp[-(h-x)/\lambda] \left[1 + \sin \left[\frac{\pi}{4} - \arctan \frac{x}{b-x} \right] \right] dx . \quad (3)$$

The integral can be solved by linearization or numerical integration. It can be seen from Fig. 1 that increasing the base b does not enhance the collection probability proportionally because (1) the increased base is the projection of the incremental solid angle and (2) the intensity of the incremental flux decreases proportional to the cosine of the emitted flux distribution.

The collection probability η as a function of the normalized pump limiter base b/h is shown in Fig. 2 for the pump limiters of the Impurity Study Experiment (ISX-B) with $h = 2.5 \text{ cm}$ and $\lambda = 5 \text{ cm}$. In ISX-B, two different pump limiters were installed: one with

one-sided collection and $b = 11$ cm and one with two-sided collection and $b = 6$ cm. According to Fig. 2, the collection probability is only slightly enhanced, from 0.73 to 0.83, for the longer base. As a design criterion, we find that for b/k larger than 4, the collection improvement is minute. The actual pump limiter in ISX-B had a mushroom configuration, and Fig. 1 represents the central cut only. Integration over the cylinder geometry yields $\eta = 0.63$ for two-sided collection and $\eta = 0.76$ for one-sided collection. According to these considerations, the collection probability is fairly high, which is a consequence of (1) the 45° angle of the neutralizer plate and (2) the cosine distribution of the emitted particle flux.

4. EXPERIMENTAL RESULTS

The simple model described here should be best suited for low-density ohmic discharges in which plasma effects in the pump limiter throat are expected to be negligible. Figure 3 shows the pressures in the two pump limiter modules in equilibrium at line-averaged densities of $1 \times 10^{13} \text{ cm}^{-3}$. The one-sided limiter shows higher pressure than the two-sided one. This is mostly a consequence of toroidal asymmetries of the particle flux; that is, the flux on the electron-drift side has been observed to be lower than that on the ion-drift side. Thus, the slot on the electron-drift side represents only a small increase in particle collection, but doubles the loss. The net effect is higher pressure with one-sided collection. This effect is even more pronounced in beam-heated discharges. For the particle flux analysis, we consider only the one-sided pump limiter module.

Figure 3 shows how the limiter pressure follows the plasma density with about a 20-ms delay, which is the response time of the pump limiter system. At constant density, the pressure is in equilibrium. The H_α signal at the limiter, which is not shown here, has the same shape as the density. This indicates that pressure and particle flux are proportional to each other. By applying Eq. (2) with $\eta = 0.76$ and $\lambda = 7$ cm (Langmuir probe measurement), we obtain the particle flux at the upper edge of the pump limiter, $\phi_0 = 1.75 \times 10^{18} \text{ cm}^{-2} \cdot \text{s}^{-1}$. We can compare this particle flux deduced from the pressure

measurement with the flux measured using a Langmuir probe at the pump limiter slot. From Langmuir probe data, one usually assumes $n_e v_i = n_e \alpha c_s$, with c_s being the ion sound speed and $\alpha \approx 0.1, \dots, 1$. Since we do not know α , we plot ϕ_0 from Eq. (2) over $n_e c_s$ with $c_s = \sqrt{2kT_e/m_i}$ and n_e, T_e from Langmuir probe measurements. Figure 4 shows this plot for ohmic and for beam-heated discharges. The highest fluxes correspond to $\bar{n}_e = 5 \times 10^{13} \text{ cm}^{-3}$ in OH discharges and $\bar{n}_e = 3.5 \times 10^{13} \text{ cm}^{-3}$ in NBI discharges. Increasing the density beyond 2.5 in beam-heated discharges showed nonlinear effects for which this model is not applicable. As Fig. 4 shows, fluxes deduced from the pressure are consistent with Langmuir probe data for $\alpha \approx 0.2$. For ohmic discharges, α might be slightly higher. In either case, α seems to be a little lower than the values of 0.3 or 0.4 most commonly used. This could reflect too high an estimate for the collection probability η , which determines ϕ_0 according to Eq. (2). If $\alpha = 0.3$ is assumed, the data are consistent with collection probabilities between 0.5 and 0.63. There is uncertainty in the assumptions made here about the backscattering at the neutralizer plate, although the maximum possible value of η appears to be 0.8 and our best estimate of $\eta = 0.65 \pm 0.15$ cannot be too far off. Thus, a simple pressure measurement yields the particle flux in the scrapeoff layer, as was shown for OH discharges with densities up to $5 \times 10^{13} \text{ cm}^{-3}$ and for NBI discharges up to $3.5 \times 10^{13} \text{ cm}^{-3}$.

The data shown in Figs. 3 and 4 are all without active pumping in the limiter. When the Zr/Al getter pump inside the pump limiter chamber was activated, the pressure dropped by a factor of 3.2 at high densities to 4 at low densities. A pressure drop of 3.2 requires the particle removal through pump and slot to be 3.2 times higher than through the slot alone, i.e., $\phi_p + \phi_s = 3.2\phi_s$. Hence, the pump removes 68% of the collected flux. The trapping fraction γ is then $\gamma = 0.68\eta \approx 0.5$; that is, 50% of the incident flux ϕ is removed from the system by pumping.

5. PARTICLE TRANSPORT IN THE PUMP LIMITER

For the global particle balance in the pump limiter, it was assumed that the gas at the pressure gauge was at room temperature. To verify this assumption, we try to estimate the time after which a fast particle is thermalized and the corresponding distance it has traveled. Although an analytical approach which includes all the involved physics processes would be too complicated, reasonable assumptions and simplifications can lead to useful results and provide valuable insight.

The diffusion length is $x = \sqrt{Dt}$ with a diffusion coefficient $D = (1/4)\lambda\bar{v}$. Assuming a cosine distribution of the emitted particles, the mean free path in a tube of diameter d can be computed to be $\lambda = (\pi d)/2$, where collisions between particles are negligible. The resulting diffusion coefficient $D = (1/8)\pi d\bar{v}$ can be shown to be compatible with the usual formula for the conductance of molecular flow. The average particle velocity $\bar{v}$ changes as a function of the wall collisions, and the diffusion coefficient and the time between collisions change correspondingly. In the interval between collisions k and $(k + 1)$, we have $D_k = (1/4)\lambda\bar{v}_k$ and $t_k = \lambda/\bar{v}_k$. The diffusion length is then

$$\bar{x} = \left(\sum_{k=1}^N D_k t_k \right)^{1/2} = \frac{1}{4} \pi d \sqrt{N} , \quad (4)$$

where N is the number of collisions; that is, the diffusion length is determined by the tube diameter and the number of collisions.

We now estimate the number of collisions and the elapsed time that a fast particle needs for thermalization with the wall. In the first approach, we consider the slowing-down process of monoenergetic particles. The energy reflection coefficient R_E for energies between zero and 100 eV can conveniently be represented as $R_E = 1 - 0.082 \ln(E/E_w)$ (for deuterium on iron) where E is the incident energy and E_w is the energy corresponding to wall temperature. If a wider energy range is to be covered, the analytical expression for R_E becomes more complicated [4]. When the incident particle energy corresponds to the

wall temperature, the energy reflection coefficient becomes unity and the energy transfer to the wall is zero.

The energy of the reflected particle is $E_r = R_E E$. After N collisions it is

$$E_N = E_0 \prod_{k=1}^N [1 - 0.082 \ln(E_{k-1}/E_w)] , \quad (5)$$

where Π indicates the product for $k = 1, \dots, N$ and E_0 is the initial particle energy. After the k th collision, the particle has a velocity $v_k = \sqrt{2E_k/m}$ and the time between collisions is $t_k = \lambda/v_k$. After the n th collision, the elapsed time is

$$t_N = (\pi d/2) \sum_{k=1}^N 1/\sqrt{2E_k/m} . \quad (6)$$

The particle energy as a function of the number of collisions is plotted in Fig. 5 with initial particle energies of $E_0 = 25$ eV and 100 eV. The graph shows that the slowing-down process is fast at the beginning when R_E is low and becomes very slow when the particle energy approaches wall temperature and R_E approaches unity. The corresponding time required for the slowing-down process is shown in Fig. 6. For 25-eV or 100-eV particles in a 3-cm-diam tube, thermalization time is on the order of 10^{-3} s.

In the second approach, we consider a particle ensemble with a Maxwell-Boltzmann velocity distribution. The average velocity of the distribution is $\bar{v} = \sqrt{8kT/\pi m}$, and the average energy that a particle carries to the wall is $2kT$ [5]. The energy transferred to the wall in a collision is $2kT(1 - R_E)$. The energy transferred to the wall, per unit time, must be equal to the energy lost by the particle ensemble:

$$\dot{E} = -\frac{3}{2} nk \dot{T} \cdot V = A \frac{n\bar{v}}{4} 2kT(1 - R_E) . \quad (7)$$

Substituting $R_E = 1 - 0.082 \ln(T/T_w)$, we obtain the differential equation

$$\frac{dT}{dt} = -CT^{3/2} \ln \left(\frac{T}{T_w} \right), \quad (8)$$

where C contains the constants of Eq. (7). Integration of Eq. (8) yields the elapsed time for the thermalization from the initial temperature T_0 to the temperature T_N :

$$t_N = \frac{1}{C} \int_{T_0}^{T_N} \frac{dT}{T^{3/2} \ln(T/T_w)}. \quad (9)$$

With the substitution $2z = \ln(T/T_w)$, the integral can be transformed into an exponential integral, which can be computed by its infinite series or by numerical integration. For deuterium gas in a 3-cm-diam circular tube, the constant is $C = 2.8 \times 10^4$. The temperature as a function of elapsed time is shown in Fig. 6 for the initial temperatures of 100 eV and 25 eV. Comparison with the monoenergetic particle solution indicates that the slowing-down functions are similar, but that the thermalization of the Maxwell-Boltzmann distribution is somewhat faster at the end. This is a consequence of the higher average energy $2kT$ per particle transported to the wall by the particle distribution. Neither the number of collisions nor the thermalization time changes much when we compare 25-eV and 100-eV particles. Slowing down from 100 eV to 25 eV takes only two collisions because the energy reflection is low at the high-energy end, and the additional time is negligible since the collision times are short.

Once the number of collisions for thermalization is known, the diffusion length can be estimated with Eq. (4). For particles with the initial kinetic energy of 100 eV, the diffusion length for thermalization in a 3-cm-diam tube is $x = 18$ cm and the elapsed time is on the order of 1 ms. The thermalization point is somewhat arbitrary because the wall temperature is approached asymptotically.

In the ISX-B pump limiter module, the ion gauge was connected to the pump limiter chamber with a 50-cm-long, 3-cm-diam tube; that is, the tube was about 3 times longer than the thermalization length. Hence, we assume that the gas at the ion gauge was in thermal equilibrium with the wall. The space within the pump limiter chamber was to a large extent filled by the Zr/Al getter cartridge such that the mean free path for particles was only on the order of 1–3 cm. Therefore, we assume that the particles were thermalized fairly rapidly even within the pump limiter chamber. Consequently, enhanced pumping due to energetic particles could not be expected.

The estimates of the thermalization length can also provide some insight into the effects of “enhanced conductance.” This enhancement can be expected if the thermalization length is longer than the pumping duct. As an example, we consider a 10-cm-diam, 100-cm-long duct. The thermalization length according to Eq. (4) is about 50 cm. Hence, the first half of the duct has enhanced conductance, while the second half has normal thermal conductance. Since the resulting conductance is determined by the smaller value, enhancement cannot be expected for particle energies typical for the plasma edge of fusion devices.

6. CONCLUSIONS

Some simple modeling has been presented of the global particle balance and the thermalization in pump limiters. The modeling fits the ISX-B pump limiter data reasonably well. It allows us to deduce the particle flux in the scrapeoff layer from simple pressure measurements. If particle control is not the objective, a smaller probe can be designed for scrapeoff layer studies. Estimates of the thermalization length of fast particles in the pump limiter can be obtained from an energy-dependent reflection coefficient and the slowdown process of either monoenergetic particles or a Maxwell-Boltzmann distribution. The thermalization length is almost independent of the initial particle energy as it is observed in the scrapeoff layer of fusion devices. Improved particle collection due to “enhanced conduction” can only occur in short ducts of sufficiently large diameter.

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FIGURE CAPTIONS

Fig. 1. Geometry and particle fluxes in the ISX-B pump limiter module.

Fig. 2. Collection probability $\eta = \phi_{\text{coll}}/\phi$ as function of b/h , the pump limiter base b normalized to the height h .

Fig. 3. Density and gas puff of low-density OH discharge and corresponding limiter pressures due to (a) single-sided collection and (b) two-sided collection.

Fig. 4. Particle flux ϕ_0 deduced from pressure measurement plotted over $n_e c_s = (1/\alpha)n_e v_{\parallel}$ from Langmuir probe measurements.

Fig. 5. Particle energy as a function of the number of collisions with the wall for $E_0 = 100$ eV and 25 eV.

Fig. 6. Thermalization of a Maxwell-Boltzmann distribution of particles and of monoenergetic particles as function of time.