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TOKAMAK PLASMA DIAGNOSIS BY
SURFACE PHYSICS TECHNIQUES

BY

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ABSTRACT

The utilization of elementally-sensitive surface techniques as plasma diagnostics is discussed with emphasis on measuring impurity fluxes, charge states, and energy distributions in the plasma edge. A model of plasma flow to the probe is presented and applied to the interpretation of data. Limits on time and energy resolution, and sensitivity are given. The overlap of these techniques with conventional plasma diagnostics is described.

I. Introduction

Surface physics techniques are being routinely applied in tokamaks to characterize wall conditions. In this mode of operation, samples are inserted to the wall position in a tokamak, exposed to high power pulses or discharge cleaning, and then retracted into an analysis chamber. To date, AES [1],[2],[3], SIMS [1], RIBS [1],[4],[5], ESCA [4], SXAPS [6], thermal desorption [7], and nuclear reactions [4] have been used in such studies. The main goal of these efforts is to correlate wall conditions with plasma behavior. By this approach it should be possible to learn what state of the walls is associated with minimum plasma contamination caused by global sputtering and desorption.

The question naturally arises, can the same surface physics techniques also be utilized to diagnose important plasma properties? To perform this function a sample must be inserted into the plasma to act as a collector for ions and neutrals which impact on it. Several groups [8],[9],[10],[11] are currently using probes in this manner. Surface analysis techniques applied to the sample would identify and quantify the deposited elements. A model for probe behavior is then needed to relate the data to plasma parameters.

In this paper attention is restricted to the use of elementally-sensitive surface techniques in the determination of fluxes, energies and charge states of impurities present in tokamak plasmas. In section II are presented discussions of sample configurations, heat and plasma flow to the probe,

and the effect of varying probe potentials. Use of the probe as an ion thermometer and charge state discriminator is discussed in section III. The overlap of this approach with conventional plasma diagnostics is described in section IV.

II. The Probe

A. General Considerations

If a cube-shaped probe is floated in an infinite, homogeneous, magnetized, collisionless plasma with the magnetic field, $\vec{B}$, normal to two faces, then the two faces experience a greater ion flux than the other four. This is due to the fact that motion along the magnetic field, $V_{\parallel}$, is free streaming, while motion perpendicular to $\vec{B}$, $V_{\perp}$, is diffusive, i.e., for the ions

$$iV_{\parallel} = 2_i \Gamma_{\parallel} / n_i \sim (2 k T_i / m_i)^{1/2}, \quad (1)$$

$$iV_{\perp} = i \Gamma_{\perp} / n_i \sim D_{\perp} \nabla_{\perp} n_i / n_i, \quad (2)$$

where $D_{\perp}$ is the diffusion coefficient perpendicular to $\vec{B}$ and T_i , m_i and n_i are the ion temperature, mass and density. In the edge plasma of tokamaks $V_{\parallel} \approx 1 \times 10^6$ cm/s and $V_{\perp} \approx 5 \times 10^3$ cm/s [12,13] for 10 eV oxygen ions. Assuming, for the moment, that all ions which impact on the sample surface stick to it, and that there is no erosion or diffusion, then the probe collects 200 times more ions/per unit area by parallel flux $i \Gamma_{\parallel}$, than by perpendicular flux, $i \Gamma_{\perp}$. Of course, the ion

flux to a probe can vary without there being a change in the plasma density. Three obvious ways for this to occur are: variations in the potential of the probe relative to the plasma; variations of the ion or electron temperatures; and variation of the collisionality of the plasma causing the free streaming motion parallel to $\vec{B}$ to become viscous. These will be discussed later in more detail.

To determine the perpendicular and parallel energies, $E_{\perp}$ and $E_{\parallel}$, of the impacting ions, use can be made of the fact that their gyro-radii are macroscopic. Consider the cube face which is perpendicular to B , and a shield with an aperture smaller than the ion gyro-radius, λ_i , in front of this face. With the shield sufficiently far from the cube face, the ions that pass through the aperture (see Fig. 1) can land on the cube face as far as $2\lambda_i$ from the geometric shadow of the aperture on the cube face. In contrast (Fig. 1), with a shield in front of a $\vec{B}$ -parallel face, most ions would not impact on that face at all because of the small $V_{\parallel}$ required. Those ions with $V_{\parallel} < 5 \times 10^4$ cm/s represent only $(mV_{\parallel}^2/nkT)^{1/2}$ of the total n_i which, in this case, is about $2 \times 10^{-3} n_i$.

In the above configuration, time resolution of the ion flux may be obtained by moving the probe face parallel to the shield. In this manner ions arriving at the probe at different times will land at different locations. Probes of this type have used spinning "movie" film [9] and rotating discs [5] or cylinders [11]. Another approach has been to keep the probe stationary and open and close a shutter [14]. In the case

where the probe is moved behind a hole, the time resolution, Δt , can be as short as W/S , where W is the diameter of the hole and S the speed at which the probe is moved. This value of Δt may be obtained only if the probe is less than $\sim .2 \%$ from the shield. As shown in Fig. 2, when the distance x , between the probe and shield is increased to 1.5% , the time resolution degrades to $(W + 2.5 \%)s$ and then asymptotes to $(W + 2.1 \%)s$ as $x \rightarrow \infty$. The criterion used is that 75% of the deposited ions be within $\pm SA/2$ of the area being analyzed. The details of the deposition profiles leading to the above are shown in Sec. III.

At this juncture it is again emphasized that any potential difference between the plasma and the probe will change the flux and energy spectrum of the ions which reach the probe. Thus for the probe to be used as a flux meter or ion thermometer, seemingly it should be maintained at the space potential. In fact, this is not the case, as will be shown in section IIC. Another effect of operating the probe at space potential is that the power flux to the probe, q , due to the nearly saturated electron current flowing to it, should be about 10 times greater than if the probe were maintained at the floating potential, i.e., $q_E = n_e V_e (2 T_e)/2$ at space potential vs. $q_P = n_i V_i (6 T_e + 2 T_i)/2$ at floating potential.

The power flux to and the temperature rise of the probe surface place a limit on how far inside the tokamak plasma it may be useful. The temperature of the probe surface should not be allowed to exceed 300°C because of possible diffusion of the implanted ions into the bulk, or evaporation of them from the surface. The surface temperature that would occur if the

probe were positioned at a certain minor radius and held at a fixed potential is given by a standard formula [15]

$$T_s = \frac{2q}{K} \left(\frac{Kt}{\rho C \pi} \right)^{1/2}$$

where T_s = surface temperature ($^{\circ}\text{K}$) 300 $^{\circ}\text{C}$
 q = heat flux (W/cm^2)
 K = thermal conductivity ($\text{W}/\text{cm } ^{\circ}\text{K}$) 1
 t = duration of heat flux (s)
 C = specific heat ($\text{J}/\text{gm } ^{\circ}\text{K}$) .15
 ρ = density (gm/cm^3) .

This is valid for $t \leq d^2 \rho C / 4K$. The values given are those for a tungsten probe. For a 1 sec duration discharge, the upper limit on q is about 300 W/cm^2 . Figure 3 is a plot of electron temperature and density profiles for the edge plasma in PLT [16] as well as the power flux for the two cases - the probe at floating potential and at space potential. As can be seen, a probe at floating potential may be useful 3 cm further in, to $r = 44$ cm, than one at space potential. Time resolved probes could tolerate a $(\Delta t)^{-1/2}$ greater power flux, as long as the shield could stand the increased heat load. For $\Delta t = 100$ ms a floating probe could operate at $r = 42$ cm.

The hierarchy of scale lengths important in this problem are shown in the following, with the smallest length at the bottom of the list.

- (a) R = tokamak major radius = 10^2 cm
- (b) $L_c = V_t / q_0 n$ = collisional scale length = 1 to 10^3 cm
- (c) $L_{\parallel} = n / \nabla_{\parallel} n$ = parallel scale length = 5 to 50 cm
- (d) a = tokamak minor radius = 30 cm
- (e) d = probe size = 1 cm
- (f) $L_r = n / \nabla_r n$ = radial scale length = 0.5 to 5 cm
- (g) ρ_i = ion gyroradius = .91 to 0.5 cm
- (h) λ_D = Debye length = 10^{-3} cm
- (i) ρ_e = electron gyroradius = 10^{-4} cm .

(3)

The use of this probe to measure absolute fluxes depends on detailed knowledge of surface physics phenomena and techniques. The energy dependent sticking coefficient of ions must be measured for the systems of interest. Among the impurity ions that will impinge on the collector are C, O, Cl, Fe, Mo, and W. In the past, collector surfaces have been C, Al, Al_2O_3 , Si, SiO_2 , Fe, In, Ag and Be. Some sticking coefficients are known for these systems from sputter deposition technology, and are close to unity for massive projectiles landing on light atom substrates. However much work remains.

Diffusion of the implanted plasma ions may also be a problem, particularly for light atoms. High fluxes of protons to the surface may also knock the implanted heavy ions beyond the detection range of some surface techniques. For this reason complementary analysis by AES and RBS [14] would be useful.

Other mechanisms exist by which deposited atoms may be removed from the probe surface. Erosion has been cited by Staib and Staudenmaier [10] as a possible explanation for changes they have observed in ion deposition in TFR.

Among several erosion mechanisms are sputtering, arcs, and evaporation. All three would become more serious as the probe is inserted further into the plasma, thus experiencing higher heat fluxes and more energetic and multiply charged ions. Under collisional conditions, massive ions will be dragged along at the proton speed and hence would have superthermal energy and enhanced sputtering. Negative probe potentials would enhance sputtering by multiply charged ions. Sputtering, arcing, and evaporation could be detected a posteriori by visual inspection.

Much has been written about the surface analysis techniques that can be used to analyze the deposited elements. Scanning AES does have sufficient sensitivity and spatial resolution to be used in all the applications discussed in this paper. As an example, consider a plasma with an average oxygen content of $0.1\% n_e$. For a probe located at the edge where the plasma density is typically $1/10$ its central value of $10^{14}/\text{cm}^3$, the amount of oxygen deposited on the probe in 1 ms would be $\sim n_i V_i t \sim (10^{13} 10^{-3}) 10^{+6} 10^{-3} = 10^{13}/\text{cm}^2$, which is within the sensitivity limits of AES.

B. Flow to the Probe

In this section two points are discussed: where in the plasma do the ions which strike the probe originate, and how does the plasma collisionality affect ion flow to the probe.

As previously noted, most of the flow to the probe occurs along field lines. In steady state, the rapid loss of plasma to the probe from the flux tube attached to it (see Fig. 4) must be compensated by new ionizations and by perpendicular transport of plasma into the flux tube. The continuity

equation for this case is.

$$\Gamma_i + S_i = \frac{dn_i}{dt} = 0, \quad (4)$$

where Γ_i = flux of ions into the flux tube, and S_i = source of ions, e.g. ionization, in the flux tube. In Fig. 5 are plotted radial distributions of first ionization for 0.1 and 1.0 eV O atoms and 1.0 eV H atoms coming from the direction of the wall. These curves are calculated using the data in Fig. 3 and the ionization rates of Lotz [17]. It is evident that little ionization occurs in most of the region accessible to the probe. If the electron temperature at the edge were raised, more ionization would occur at larger r . However the power flux to the probe would also increase and that region would then become inaccessible. Hence the source term in Eq. (4) can be set to zero.

(High ionization states are pertinent to the current discussion due to the fact that a large fraction of impurity ions should leave the plasma multiply charged [18],[19],[20],[21]. This, of course, occurs because the ionization rate is much more rapid than recombination, even in the plasma edge.)

With the above considerations, Eq. (4) becomes

$$\nabla_{\perp} \Gamma_{\perp} + \nabla_{\parallel} \Gamma_{\parallel} = 0. \quad (5)$$

By applying Gauss' theorem to the flux tube with end faces at $L_{\parallel}$ (where the plasma is isotropic) and λ_D from the probe, Eq. (5) becomes

$$\sum_i \Gamma_i A_i + \sum_i \Gamma_i A_i = 0 \quad , \quad (6)$$

where $A_i = \bar{B}$ -normal cross sectional area of the probe $= d^2$
and $A_i =$ surface area of the flux tube $\approx 4 d L_i$. Again
the perpendicular flow is diffusive, that is, Eq. (2) holds.
However, depending on the collisionality of the plasma, the
parallel flow may be either free streaming or viscous. This
is determined by comparing L_i , calculated using Eq. 6, with L_c ,
the Spitzer [22] ion-ion 90° scattering length,

$$\begin{aligned} p-p L_c &= V_p \tau_{p-p} \quad (a) \\ p-i L_c &= V_p \tau_{p-i} \quad (b) \\ i-i L_c &= V_i \tau_{i-i} \quad (c) \quad , \quad (7) \end{aligned}$$

where $\tau_{p-p} = 90^\circ$ proton-proton scattering

$$\begin{aligned} &\approx \tau_{p-i} (2m_p/m_i) Z_i^2 \quad , \\ &\approx \tau_{i-i} (m_p/m_i)^{1/2} (n_i/n_p) Z_i^4 \quad , \end{aligned}$$

and Z_i is the charge state of the ions.

if $L_c \geq L_i$ the flow is necessarily free streaming. And,
because $d \sim L_r$, when $L_c < L_i$, the flow will be viscous over
sizeable fractions of A_i and L_i . To find L_i when $V_i = (2kT/m)^{1/2}$
we combine Eqs. (1), (2), and (6) using the values of
 $5 \times 10^3 \text{ cm}^2/\text{s}$ for D_i , 10^6 cm/s for the impurity ion parallel
velocity, $5 \times 10^6 \text{ cm/s}$ for the proton parallel velocity,
 $2 \times 10^8 \text{ cm/s}$ for the electron thermal velocity, and $d/2$ for
the perpendicular scale length, L_r , giving

$$\begin{aligned}
 L_{\parallel} &= \frac{v_{\parallel} d^2}{16 D_{\perp}} \approx 12 \text{ cm (for impurity ions)} \\
 &\approx 60 \text{ cm (for protons)} \\
 &\approx 2.4 \times 10^3 \text{ cm (for electrons) ,}
 \end{aligned}
 \tag{8}$$

The results are shown in Fig. 6. Near the wall ($r = 51$ cm) the plasma density is so low that the free streaming approximation is good for H^+ , O^+ and even O^{+6} ions. Inside $r = 44$ cm the high ion temperature again causes the plasma to be collisionless over distances greater than 10 cm. However between $r = 47$ cm and $r = 44$ cm the plasma is sufficiently collisional to allow viscosity to affect the parallel motion of O^{+6} and H^+ .

There are two main effects of viscosity. The first is to slow down the flow of ions to the collector because of collisions of the ions streaming inside the flux tube with ions diffusing across it. For the case of $d \gtrsim L_r$, these collisions would reduce the parallel flux by $\sim L_c/L_{\parallel}$. As can be seen in Fig. 6, this only occurs for the highly ionized states: O^{+6} - O^{+6} if $n_o/n_e \gtrsim .03$, and for O^{+6} - H .

However, the second effect is an increase in the velocity of the massive impurities because of the drag of the protons on them. For a probe with $d \sim L_r$ the net result of the two effects will approximately cancel. The details must be worked out by a kinetic treatment for varying d/L_r ratios. Non-classical viscous effects may also be important.

C. Effects of Probe Potential, ϕ_p , on Ion Collection and Ion Energy.

The presence of an ideal probe in a plasma should not change the trajectories of ions from what they would be if the probe were absent. However, the plasma is necessarily perturbed since the probe also must act as a sink to the ions and electrons. Thus the plasma just outside the probe is not isotropic; there are only particles streaming towards the probe. With this in mind one must examine how to operate the probe so that the collected ions can be related to the distribution of ions in the plasma, $f_i(\vec{v})$. For simplicity we treat only two cases: space potential and floating potential.

(i) Space Potential, ϕ_s

With the probe at space potential, $\phi_p = \phi_s$, the electron flow at its surface will be at thermal speed, $v_{te} = (T_e/2m_e)^{1/2}$. n_e will vary along B. As shown schematically in Fig. 7, this saturated electron current, I_e , to the probe will cause an $I_e R$ "rise" between the plasma and the probe, where R is the classical resistance of the plasma in the flux tube. Because the electron collision length, L_c (Fig. 6), is much shorter than the electron free streaming length, Eq. 8, several collision lengths distant from the probe, the electron flow will be at a drift velocity less than the electron thermal velocity but still faster than the ion thermal velocity. Still further away (Fig. 7) an electron rich region must be built up to be a source for the electron flow towards the probe, i.e., the

electrons lost to the probe from the plasma must be balanced by other electrons entering the tokamak volume. Moving along B towards the probe ϕ increases by $\Delta\phi = IR$. A kinetic theory derivation, performed by SanMartin [23] for probes with $r_i > d$, show this $\Delta\phi$ hill. At a distance of eL_c the electrons free fall towards the probe. Their density decreases and an ion rich region is left behind. Within a few Debye lengths of the probe an electron sheath forms as the potential returns to ϕ_s .

$\Delta\phi$ is proportional to the electron temperature, as can be seen from the following: Because of collisions with ions the electron flow is diffusive,

$$e \dot{r}_n = e D_n \nabla_n n_e = n_e e v_n, \quad (9)$$

so

$$e v_n = \frac{e v^2 \tau_{e-i}}{L_n}. \quad (10)$$

Now

$$R = \rho L_n / A_n, \quad (11)$$

where

$$\rho = \frac{1}{.58} \frac{\pi^{3/2} m_e^{1/2} z^2 e^2 \ln \Lambda}{2 (2 kT)^{3/2}},$$

and

$$I = e n_e v_n A_n. \quad (12)$$

Multiplying (11) by 12) and using (10) and the Spitzer collision time, we get

$$\Delta\phi = IR = e n_e \left(\frac{e v_t^2 \tau_{ei}}{L_n} \right) A_n \left(\rho \frac{L_n}{A_n} \right) \quad (13)$$

$$\approx 1.9 \frac{kT_e}{e}$$

Because this model is crude, the accuracy of the factor 1.9 in (13) is open to question.

As $\Delta\phi$ increases, the ion flow to the probe is reduced by the factor $\exp(-\Delta\phi'/kT_i)$, where $\Delta\phi'$ is the potential difference between the last position in the flux tube where the ion is thermalized and the top of the $\Delta\phi$ hill. The resulting charge separation puts an upper limit of $\sim T_i$ on $\Delta\phi$. The impurity ion flux to the probe would be as shown in Fig. 8, where we have assumed: that $\Delta\phi' = \Delta\phi$; that the impurity ion content is a constant fraction of the electron density; and that the rise in ion temperature due to I^2R heating or instabilities is small. The dependence of the flux on minor radius, i.e., a peak occurring where $T_e \sim T_i$, is a qualitative prediction of this model. Detailed comparison with experiment will require measurements of T_e , T_i and n_e because of the strong dependence of the ion flux on these quantities. The ion temperature profile shown in Fig. 3 is an estimate based on numerical simulations using the Baldur code [24].

Two final points should be made about operating the probe at the space potential. First, the $\Delta\phi$ hill will be more effective in reducing the flux to the probe of multiply ionized ions than singly charged ions. Secondly, space potential

measurements that do not disturb the plasma are possible with heavy ion beam probes [25]. Results from ST show the space potential on the plasma axis to be ~ 100 Volts negative with respect to the vacuum vessel. The potential becomes less negative at larger minor radii. Thus if a probe in ST were electrically attached to the vacuum vessel the ion flux would be reduced even more than if the probe were at space potential.

(ii) Floating potential, ϕ_f

As the probe potential is lowered from ϕ_s , the potential along B will vary as shown in Fig. 7(d). The $\Delta\phi$ will decrease as less electron current is drawn. Finally, at $\phi_p = \phi_f$, the electron and ion fluxes balance and there is no change in ϕ along B until a few Debye lengths from the probe. There the usual ion-electron double sheath occurs. The electron distribution is essentially isotropic along B until that location because the loss of electron to the probe is small compared with the random "thermal current", $n_e e v_t$. Under these conditions the ion flux, Eq. (14), to the probe should be at the ion acoustic speed [23],[27], as long as viscous [28] effects are small.

$$\Gamma_i \approx n_i C_s \approx n_i [(T_i + Z_i T_e)/m_i]^{1/2}, \quad (14)$$

where C_s is the ion acoustic speed. Measurements on FM-1 [26] have shown the plasma flow velocity into a poloidal divertor to be at $.4 C_s$ for $T_i \ll T_e$.

(iii) Effects of Probe Potential on Ion Energy

In section III the use of the probe to measure E_i and E_e are described. Here we address the question, how will

the operating voltage of the probe affect the energy distribution of ions reaching it? We limit our attention to a probe as shown in Fig. 1, with its face perpendicular to B , having a shield with a small aperture in front of it. In order that all potential changes occur at the shield and that the region between the shield and probe is free of perturbing electric fields, two conditions must be met. The first condition is that the aperture be smaller than the Debye length. This severely reduces the flux of ions to the probe. One way to avoid this is to place a grid with spacing $< \lambda_D$ over a larger hole, say with diameter $\sim .2 \lambda_i$. The second condition is an upper limit on the impurity concentrations in the plasma. As the plasma electrons stream in a tight pencil-like beam between the aperture and the probe, the massive impurity ions make excursions of $2 \lambda_i$ from the beam axis. So that the ion trajectories not be affected, the potential caused by this charge separation must be $< 0.1 T_i$. Integrating Poisson's equation shows that the upper limit on n_i is $\sim 4 \times 10^{10} \text{ cm}^{-3}$ for $10 \text{ eV } O^+$ and a $.1 \lambda_i$ radius aperture.

From Fig. 7(d) the change in $E_{\parallel}$ is obvious for probe operation at either ϕ_s or ϕ_f . For $\phi_p = \phi_s$ there is no change in $E_{\parallel}$ for those ions which reach the probe. However the $\Delta\phi$ hill acts as a retarding field energy analyzer passing only those ions with $E_{\parallel} \geq 1.9 Z_i T_e$. At $\phi_p = \phi_f$ the sheath drop accelerates the ions into the probe, and the Maxwellian energy distribution is shifted $\sim 4T_e$ upward.

Changes in $E_{\perp}$ are caused by radial electric fields, ϵ_r . For $x \gg \lambda_d$ and $\phi_p = \phi_f$, $\epsilon_r = 0$. For $\phi_p = \phi_s$ or ϕ_f and $x \lesssim \lambda_d$, $\epsilon_{\perp}$ may be sizeable. An upper limit on $\epsilon_{\perp}$

may be obtained by assuming $(n_i - n_e) = 0$ in the sheath region, but that the full electric field is still present. Using Laplace's equation,

$$\nabla \cdot \epsilon = 0, \quad (15)$$

we find the radial field evaluated at the aperture edge is

$$\epsilon_r \approx - \frac{2 kT_e W}{e \lambda_d^2}, \quad (16)$$

where

$$\begin{aligned} \alpha &= 2 & \text{for } \phi_p &= \phi_s \\ &= 4 & \phi_p &= \phi_f \end{aligned}$$

and W = aperture width or grid spacing.

Using the configuration of a grid over a larger hole with $W \sim \lambda_d$, yields

$$\epsilon_r = - \frac{\alpha kT_e}{e \lambda_d}, \quad (17)$$

The change in perpendicular energy caused by an ion passing through this region can be found using the impulse approximation.

$$\Delta P_{\perp} \approx Z_i \epsilon_r e \Delta t \sim Z_i e \epsilon_r \frac{\lambda_d}{i v_t} = Z_i \alpha \frac{kT_e}{i v_t}, \text{ and} \quad (18)$$

$\Delta E_{\perp} = v_t \Delta P_{\perp} = Z_i \alpha kT_e$. Thus, unless $W \ll \lambda_d$, $\Delta E_{\perp}/E_{\perp}$ will be positive and of $O(1)$.

III. An Ion Thermometer and Charge State Discriminator

The radial distribution of ions passing through a pinhole in shield and impacting on the probe has been calculated in Ref. 29. For the geometry shown in Fig. 1 (B-normal face), they find that areal density of ions on the probe is given by:

$$F(r) = \frac{n_i A \beta}{\gamma^4} \int_{r/2}^{\infty} dr' \exp\left(\frac{-\alpha r'^2}{\gamma^2}\right) \sum_m \frac{x^2}{g_m^3(r', r)} \quad (19)$$

$$\exp\left[\frac{-\alpha x^2}{\gamma^2 g_m^2(r', r)}\right] / \left[\left(1 - \frac{r^2}{4r'^2}\right)^{1/2} r\right]$$

where $\alpha = m_i/2kT_i$,

$\beta = (\alpha/\pi)^{3/2}$,

$\gamma = m_i/Z_i eB$,

A = area of hole,

$g_m(r', r) = 2\pi m \pm \sin^{-1}(r/2r')$, and

x = distance from the aperture to the probe.

For $x \leq \ell_i/2$, Eq. (19) can be put into the form of effusive flow of unmagnetized particles through a hole.

For $x \gg \lambda_i$ the density distribution asymptotically approaches,

$$F(r) \propto \exp(-r^2/\lambda_i^2)/r. \quad (20)$$

So for a single charge state of an ion, measuring $F(r)$ at large x will give λ_i , and measuring $F(r)$ vs. x would give E_n . If we make the reasonable assumption that the ions are isothermal, then the different charge states may be unfolded from $F(r)$ by an iterative fitting technique. A normalized plot of $F(r)$ vs. x is shown in Fig. 9.

In the near field region, $x < .5 \lambda_i$, the $x^2/(x^2 + r^2)^2$ behavior of effusive flow is evident. At $x \approx .7 \lambda_i$ a small "tit" begins to grow on axis as ions are bent back in by the magnetic field. At $x \approx 2\lambda_i$ the "tit" attains its largest height. By $x \approx 12 \lambda_i$ the deposition is essentially of the form described by Eq. (20). A sheath drop would stretch out $F(r)$ along x so that its asymptotic form would only be reached at larger x . A small radial electric field would cause a broadening of $F(r)$ at large x . Large radial fields would cause $F(r)$ to assume a volcano-type shape.

When operating at $\phi_p = \phi_s$ the multiply ionized particles will get rejected compared with singly charged ions. Thus the E_1 distribution could be assumed to be due to $Z_i = 1$ ions. In principle, unfolding ion charge states from the deposition profiles is accomplished by systematically varying x and ϕ_p .

IV. Comparison with Standard Plasma Diagnostics

The primary diagnostics on tokamaks for measuring impurity fluxes, concentrations and energies are visible, U.V. and x-ray spectroscopies [30]. The emissivity of a plasma is proportional to $\exp(-\Delta E/kT_e)$ where ΔE is the energy for excitation. The minimum ΔE for energy resolved x-ray techniques is ~ 200 eV, thus few x-ray transitions will occur outside the limiter radius. Runaway electrons with keV to MeV energies are too infrequent on occurrence to be useful for exciting x-ray emission in the plasma edge. Visible and U.V. spectroscopies extend to ΔE down to ~ 1 eV, thus the emissivity of the plasma edge can be high. It is, in fact, high enough that lines of Cr I to $\sim$ Fe XVII have been observed across the plasma diameter. However, only the low ionization states have been observed at large minor radii for the same reason as noted above, i.e., highly ionized atoms have large ΔE . Recombination of electrons with highly ionized atoms can provide cascades of photons. But each recombination event will only result in 1 photon for each transition. For example, the radiative recombination time of O^{+6} to O^{+5} at $T_e = 5$ eV and $n_e = 1 \times 10^{13} \text{ cm}^{-3}$ is ~ 30 ms. In contrast, the excitation time for $\Delta n = 0$ O^{+1} line under the same conditions is about 1 μs , and thus is about 3×10^4 times brighter.

Spectroscopic techniques do provide measurements of perpendicular fluxes.¹² In this sense the surface physics approach complements the optical.

Spectroscopic techniques, in particular line broadening, are used to measure ion temperatures. Most emphasis has been for measurements in the plasma core, but only minor changes would be required for measurements at the edge.

Detailed descriptions of the application of mass spectroscopy to plasma diagnosis can also be found in Ref. [30].

V. Summary

The flow of plasma to a probe immersed in the edge plasma of a tokamak will lead to the deposition of impurity ions on its surfaces. Most ions that reach the probe will have originated further inside the tokamak and diffused into the flux tube attached to the probe. Flow along the field lines, being more rapid than flow across them, will result in greater deposition rates on the B-normal surfaces. The flow is sensitive to potential and density gradients. No complete theory for the probe exists. However, maintaining the probe at the floating potential should result in ion flow to it at the acoustic speed. Possible erosion of the implanted ions, diffusion into the bulk, or non-unity sticking coefficients will reduce the actual surface concentration from that predicted from the acoustic speed model.

Measurements of impurity energy distribution of charge states can be made using an aperture in front of the probe and measuring the deposition profile resulting from the

macroscopic gyroradii. These measurements of field parallel impurity ion fluxes, charge states and energy distributions are complementary to most standard plasma diagnostics.

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FIGURE CAPTIONS

Fig. 1. A cube-shaped probe immersed in a magnetized plasma. The probe is inside a shield which has a pinhole in each face. The flux of ions along the field is much greater than that across it.

Fig. 2. Time resolution, Δt , of a probe moved behind an exposure slit at a speed S . The slit width is W . The peaking of Δt at $1.5 \lambda_i$ is due to the magnetic field. At large x Δt asymptotes to $(W + 2.1 \lambda_i)/S$.

Fig. 3. Electron temperature T_e , ion temperature T_i , electron density n_e , and heat flux q , in the PLT plasma. T_e and n_e were measured by Langmuir probes [16] for $44 < r < 51$ cm. T_i is taken from the Baldur code [24]. q is defined in the text.

Fig. 4. Flow of plasma ions into and along the flux tube connected to the probe. Most ions which reach the probe originate further inside the tokamak and diffuse in the flux tube. Few first ionizations take place in the region accessible to the probe.

Fig. 5. Radial distribution of ionization of 0.1 and 1.0 eV O atoms and 1.0 eV H atoms entering PLT from the direction of the wall. The T_e and n_e profiles in Fig. 3 were used to calculate these curves.

Fig. 6. O^+ free streaming length, $O^+ L_{\infty}$, and various collision lengths $L_C(O^+ - H^+)$, $L_C(O^{+6} - H^+)$, $L_C(O^{+6} - O^{+6})$ and $L_C(e - H^+)$. These are defined in equations (3b), (3c), (7) and (8). For $43 < r < 48$ cm the plasma is essentially collisionless.

Fig. 7. (a) Potential distribution along B for a probe at space potential, ϕ_s (b) and (c) The electric field, E , and charge distribution, $(n_i - n_e)$, along B for a probe at ϕ_s (d) The variation of potential as the probe potential is reduced from ϕ_s to ϕ_f , the floating potential.

Fig. 8. Radial dependence of impurity flux to, or deposition on, a probe at space potential. A probe inserted further and further into the plasma would at first be bombarded by more then fewer ions because of a potential rise along its flux tube. A similar behavior could occur if sputtering or evaporation eroded the deposited atoms.

Fig. 9. (a) and (b) Normalized deposition profiles, $F(r)/F(o)$, of ions incident through a pinhole onto a target, as a function of target distance, x . (a) clearly shows effusive flow for $x \leq .5 \lambda_i$. (b) shows the asymptotic behavior of $F(r)$ at large x .

Fig. 1. 783309

Fig. 2. 783329

Fig. 3. 783313

Fig. 5. 783308

Fig. 6. 783312

Fig. 7. 783314

Fig. 8. 783310

Fig. 9. 791247