



NRL/MR/6707--00-8451

Use of RF Bias in LAPPS (Large Area Plasma Processing System)

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June 19, 2000

20000710 068

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REPORT DOCUMENTATION PAGE

Form Approved
OMB No. 0704-0188

Public reporting burden for this collection of information is estimated to average 1 hour per response, including the time for reviewing instructions, searching existing data sources, gathering and maintaining the data needed, and completing and reviewing the collection of information. Send comments regarding this burden estimate or any other aspect of this collection of information, including suggestions for reducing this burden, to Washington Headquarters Services, Directorate for Information Operations and Reports, 1215 Jefferson Davis Highway, Suite 1204, Arlington, VA 22202-4302, and to the Office of Management and Budget, Paperwork Reduction Project (0704-0188), Washington, DC 20503.

1. AGENCY USE ONLY (<i>Leave Blank</i>)	2. REPORT DATE <p style="text-align: center;">June 19, 2000</p>	3. REPORT TYPE AND DATES COVERED <p style="text-align: center;">Interim</p>	
4. TITLE AND SUBTITLE <p style="text-align: center;">Use of RF Bias in LAPPS (Large Area Plasma Processing System)</p>		5. FUNDING NUMBERS	
6. AUTHOR(S) <p style="text-align: center;">Wallace M. Manheimer, Martin Lampe, and Richard F. Fernsler</p>		8. PERFORMING ORGANIZATION REPORT NUMBER <p style="text-align: center;">NRL/MR/6707--00-8451</p>	
7. PERFORMING ORGANIZATION NAME(S) AND ADDRESS(ES) <p style="text-align: center;">Naval Research Laboratory Washington, DC 20375-5320</p>		10. SPONSORING/MONITORING AGENCY REPORT NUMBER	
9. SPONSORING/MONITORING AGENCY NAME(S) AND ADDRESS(ES) <p style="text-align: center;">Office of Naval Research 800 N. Quincy Street Arlington, VA 22217</p>		11. SUPPLEMENTARY NOTES	
12a. DISTRIBUTION/AVAILABILITY STATEMENT <p style="text-align: center;">Approved for public release; distribution is unlimited.</p>		12b. DISTRIBUTION CODE	
13. ABSTRACT (<i>Maximum 200 words</i>) <p>In the LAPPS processing scheme, RF bias can be used for either of two purposes: to provide energetic ion bombardment of a substrate, or to controllably raise the electron temperature (which is intrinsically very cool) to the desired value. The physics of RF bias in LAPPS differs from the situation in conventional processing reactors, for several reasons: (1) The plasma density adjacent to the substrate can be so high that the ion plasma frequency exceeds the microwave frequency. (2) Plasma transport to the substrate is across a magnetic field. (3) Ionization occurs only in a thin, well-defined planar sheet, and thus the volume occupied by plasma is very broad in two dimensions but thin in the third dimension. (4) The surface area of the substrate is comparable to that of the containment vessel. We discuss the modifications to the theory of RF bias that are needed to account for these factors. We examine the partition of RF power into various plasma channels and show that, for a given RF current, the presence of the magnetic field does not substantially change the ion bombardment energy or power, but does significantly increase the electron heating power. We address the issue of return current path for the RF current, and show that the presence of the magnetic field does not influence the return current flow. In conventional reactors, the vessel area is large compared to the substrate, and return current can be transported through the plasma and dissipated over the vessel walls. But in LAPPS, it may be preferable to provide a return current collector immediately adjacent to the plasma sheet.</p>			
14. SUBJECT TERMS		15. NUMBER OF PAGES <p style="text-align: center;">21</p>	16. PRICE CODE
17. SECURITY CLASSIFICATION OF REPORT <p style="text-align: center;">UNCLASSIFIED</p>	18. SECURITY CLASSIFICATION OF THIS PAGE <p style="text-align: center;">UNCLASSIFIED</p>	19. SECURITY CLASSIFICATION OF ABSTRACT <p style="text-align: center;">UNCLASSIFIED</p>	20. LIMITATION OF ABSTRACT <p style="text-align: center;">UL</p>

1. Introduction

LAPPS (“Large-Area Plasma Processing System”) is the acronym for an approach developed at the Naval Research Laboratory in which a sheet electron beam of energy several keV, with cross-sectional dimensions typically of the order of one or two cm thick by 1 m wide, is used to ionize a process gas and create a plasma sheet with dimensions as large as 1m×1m. A substrate whose dimensions are nearly as large can be placed adjacent to the plasma sheet. A magnetic field parallel to the substrate with a B of order 10^{-2} tesla, is needed to guide the electron beam. This approach is of interest for a variety of processing applications, and has a number of notable potential advantages, including uniformity over very large areas, high plasma density, separately controllable ion flux and neutral radical flux to the substrate, electron temperature T_e which can be low ($< 1\text{eV}$) or can be controllably increased, and large area access for gas flow and pumping. In a recent paper [1] we have outlined the basic theory, parameter requirements and operating regimes for LAPPS, and discussed and analyzed some initial experiments on ion flow to a grounded or dc-biased collector. We pointed out that the physics of LAPPS differs in a number of important respects from the more familiar high-density plasma sources: ionization occurs only within a sharply-defined planar region, the plasma is very broad but thin, the area of the substrate is comparable to the exposed area of the containment vessel, and the presence of the magnetic field has significant effects on plasma properties.

In the present paper, we discuss the use of rf bias in LAPPS to controllably increase the energy of ions bombarding the substrate. Initial experiments with rf bias are in progress now, and our primary objective here is to set the context for these experiments. In Sec. 2 we review the standard theory of ion acceleration in rf sheaths. In Sec. 3 we discuss modifications to the theory which are necessary for high-density plasma. In Secs. 4 and 5 we concentrate on issues that are specific to LAPPS, in particular the effects of the magnetic field and of the geometry in the context of rf bias. We shall be particularly interested in the partition of rf energy among ion acceleration (usually the preferred channel), ohmic heating and stochastic electron heating, discussed in Sec. 4, and in the issue of controlling the rf return current so as to optimize processing conditions, discussed in Sec. 5. To emphasize the practical consequences of the theory, many of the results are written in units scaled to the typical operating regime for LAPPS, and oxygen plasma [1] is used as an example to illustrate the theory. Throughout this paper MKS units are used except for temperature, which is in volts. This is converted to joules by the Boltzman constant $k_B = 1.6 \times 10^{-19} \text{J/eV}$ (k_B is the electron charge in coulombs). In MKS units the pressure is given in pascals ($1\text{Pa} = 7.6 \times 10^{-3} \text{Torr}$).

2. RF Sheaths in Low-Density Collisionless Plasmas: Godyak-Lieberman Theory

We shall begin with a brief review of the conventional rf sheath model [2-5] for unmagnetized plasma with density such that

$$\omega_{pi} \ll \omega, \tag{1a}$$

$$\omega \ll \omega_{pe}. \tag{1b}$$

Here $\omega \equiv 2\pi f$ is the angular frequency of the rf bias, and ω_{pi} , ω_{pe} are the ion and electron plasma frequencies. Consider a 1-D sheath model as shown in Fig. 1, with x the distance from the rf-biased surface into the plasma. An rf potential $\phi(x,t)$, which is assumed to be of large amplitude compared to T_e and oscillating at frequency ω , is imposed on the plasma by the electrode bias. At nearly all times in the rf cycle, the potential $\phi(0,t)$ at the electrode is lower than the potential in the central plasma by $\gg T_e$. Equation (1) implies that the electrons react instantaneously to the field $E_x(x,t) = -\partial\phi(x,t)/\partial t$, and thus electrons are fully excluded at time t from a region $0 < x < s(t)$ which we shall refer to as the electron sheath. In the region $x \geq s(t)$, the electron density $n_e(x,t)$ is equal to the ion density n_i , and the plasma is quasineutral. The location of $s(t)$ oscillates non-sinusoidally at frequency ω , from 0 to a maximum value s_0 . Within the quasineutral plasma, gradients of the instantaneous potential are small compared to the scale length of sheaths, and we may assume that $\phi(x,t) \equiv \phi_p(t)$ is uniform for $x \geq s(t)$. When Eq. (1) is satisfied, the ions cannot respond on the ω time scale, but rather react only to the time-averaged potential $\bar{\phi}(x)$, which is a constant $\bar{\phi}_p$ for $x \geq s_0$ and decreases for $x < s_0$. Thus the ion density $n_i(x)$ is time-independent and is essentially uniform for $x \geq s_0$ but falls off in the region $0 \leq x < s_0$, which we refer to as the ion sheath. The ion sheath is thus a time-independent structure, while the electron sheath occupies a region $0 \leq x < s(t)$ which oscillates from 0 to s_0 . The region $x \geq s_0$ is quasineutral at all times, while any point in the region $0 < x < s_0$ is quasineutral during part of the rf cycle, but has a time-averaged charge density which is positive. This is illustrated in Fig. 1.

The ion velocity $u_i(x)$ and density $n_i(x)$ are determined as follows. Since ion collisions can be neglected in the sheath, and (according to the Bohm condition [1,6]) ions enter the sheath at the ion acoustic speed $c_s \equiv (k_B T_e / M)^{1/2}$, energy conservation gives

$$M u_i^2(x) = k_B T_e - 2e [\bar{\phi}(x) - \bar{\phi}_p] \quad (2)$$

where M is the ion mass. Continuity of ion current then relates $n_i(s)$ to the density n_0 in the bulk plasma, $x \geq s_0$:

$$n_i(s) = \frac{n_0}{u_i(s)} \sqrt{\frac{k_B T_e}{M}} \equiv n_0 \sqrt{\frac{T_e}{T_e - 2[\bar{\phi}(s) - \bar{\phi}_p]}} \quad (3)$$

The oscillating current carried by the electrons in the plasma, denoted as $J_{rf} \sin \omega t$, is related to $s(t)$ by

$$J_{rf} \sin \omega t = e n_i(s) \frac{ds}{dt}, \quad (4)$$

since $n_e(s,t) = n_i(s)$ at the instantaneous edge of the electron sheath, where quasineutrality still applies. Combining Eqs. (3) and (4) and estimating $|ds/dt|$ as order ωs_0 , we arrive at an

approximate relation between the time-averaged sheath potential $\bar{\phi}_0 \equiv \bar{\phi}_p \equiv -\bar{\phi}(0)$ and the sheath width s_0 ,

$$s_0 = \frac{\alpha J_{rf}}{n_0 e \omega} \sqrt{\frac{\bar{\phi}_0}{T_e}}, \quad (5)$$

where α is a coefficient of order unity, whose exact evaluation requires a more elaborate theory, or an experimental determination. The Langmuir-Child law

$$J_i = \frac{4\epsilon_0}{9} \sqrt{\frac{2e}{M}} \frac{\bar{\phi}_0^{3/2}}{s_0^2} \quad (6)$$

relates $\bar{\phi}_0$ to the sheath thickness s_0 and the ion current density J_i flowing into the sheath. Using Eq. (5) for s_0 and the Bohm relation to specify $J_i = n_0 e c_s$, we find the dc bias of the substrate to be

$$\bar{\phi}_0 = \frac{1.9 \times 10^{-4} \alpha^4 J_{rf}^4}{\left(\frac{n_0}{10^{16}}\right)^2 \left(\frac{f}{13.56 \text{ MHz}}\right)^4 T_e}, \quad (7)$$

where we have scaled the units to typical values in plasma processing. Since $\bar{\phi}_0$ depends on the fourth power of our estimate of sheath width, Eq. (7) as a simple estimate of voltage is much less accurate than (5) as an estimate of sheath width. It may be best to regard Eq. (7) as a scaling law with a coefficient to be determined by experiment or a more complete theory. Sobolewski [7] has performed an extensive series of experiments on inductive discharges in argon plasma at three different frequencies, 100 kHz, 1 MHz and 10 MHz. Only at 10 MHz is Eq. (1) reasonably satisfied. At this frequency, Sobolewski found that the rf impedance scaled as $\bar{\phi}_0^{3/4}$, from which one may infer that $\bar{\phi}_0 \propto J_{rf}^4$ in agreement with Eq. (7). Using his data at density $n_0 = 10^{17} \text{ m}^{-3}$, we find that α^4 in Eq. (7) should be set to 0.46. Equation (7) then indicates that an rf current $\sim 100 \text{ A/m}^2$ at 13 MHz will produce 100 eV ions in a LAPPS plasma with $n_0 = 10^{17} \text{ m}^{-3}$ and $T_e = 1 \text{ eV}$. According to Eq. (5), the sheath width is $s_0 \approx 5 \times 10^{-4} \text{ m}$. For typical ion species with charge-exchange cross-section $4 \times 10^{-19} \text{ m}^2$ in gas at pressure 6.6 Pa, the ion mean free path is $1.6 \times 10^{-3} \text{ m}$, so the assumption of a collisionless sheath is well justified. The power per unit area of the ion bombardment, in W/m^2 , is

$$P_{\text{ion}} = J_i \bar{\phi}_0 = \frac{5.6 \times 10^{-4} J_{rf}^4}{\left(\frac{n_0}{10^{16}}\right) \left(\frac{f}{13.56 \text{ MHz}}\right)^4 T_e^{1/2}} \left(\frac{M_{\text{O}_2}}{M}\right)^{1/2}, \quad (8)$$

where we have chosen to scale the ion mass to that of O_2 .

The rf-powered electrode also supplies power to the plasma through ohmic heating and stochastic heating of the electrons in the vicinity of the sheath. We shall discuss these processes in Sec. 4. In LAPPS, it is preferable that the power into these channels be small compared to power in the ion stream, and we shall see that this is normally the case.

We note in passing that the scaling $\bar{\phi}_0 \propto J_{rf}^4$ applies only to discharges where the plasma is sustained by a power source separate from the rf bias of the electrode, as is the case for inductive and ECR discharges as well as LAPPS. In capacitively coupled rf discharges, which are powered entirely by the rf biased electrode, n_0 is itself proportional to J_{rf} , and thus the overall scaling is $\bar{\phi}_0 \propto J_{rf}^2$, as is well known.

3. The Effect of High Plasma Density

The Godyak-Lieberman (GL) rf sheath model is based on the low-density assumption (1a), which ensures that the ion space charge is determined by the response to the time-averaged fields. In addition, it is assumed that the time τ for an individual ion to cross the sheath is long compared to the rf period,

$$\omega\tau \gg 2\pi, \quad (9a)$$

and therefore the individual ions react to the time averaged (dc) fields. This operating regime has the attractive property that all ions arrive at the substrate with very nearly the same energy, independent of the phase of the rf at the time the ions arrive at the sheath. (There is a thermal spread of energy of ions entering the sheath, but it is always $<T_e$.) When (1a) holds, (9a) is equivalent to a requirement on the electron rf oscillation velocity in the bulk plasma $v_{os} \equiv J_{rf}/n_0e$,

$$v_{os} \gg c_s, \quad (9b)$$

and both (9a) and (9b) are satisfied provided only that

$$\bar{\phi}_0 \gg T_e. \quad (9c)$$

RF bias is only useful when (9c) is satisfied and ions are accelerated to superthermal energies, so (1a) is the primary requirement for validity of the GL model. In the GL limit, the sheath impedance (discussed in more detail in Sec. 5) is predominantly capacitive, with a resistive component. But (1a) may not be satisfied in high density plasmas. For example, for O_2 plasma with rf bias frequency 13.56 MHz, (1a) holds if the plasma density n_0 near the sheath is $<1.2 \times 10^{17} \text{ m}^{-3}$. LAPPS can operate at higher density, where the GL model is not appropriate.

Various theories and experiments have explored these higher density regimes [7-10]. In Ref. 8, we show that the characteristics of a collisionless rf sheath can be specified in terms of two dimensionless parameters, ω_{pi}/ω and v_{os}/c_s . As shown in that paper, and confirmed dramatically in Sobolewski's experiments, the high density case $\omega_{pi}/\omega \gg 1$ with the strong bias

specification (9c) reduces essentially to a time-varying dc-biased sheath. The sheath impedance is almost entirely resistive. (Reference 8 shows that a capacitive high density sheath is generally not possible because it would demand a larger particle flux than what the plasma can supply.) In this high density limit of a resistive sheath, $\omega\tau \ll 2\pi$, so ions cross the sheath in a time short compared to the rf period, and each ion sees only the sheath fields at the particular rf phase when it enters the sheath. During half the rf period, the sheath potential $\phi(t)$ is negative and accelerates ions. During this time, the ion current is very nearly constant at n_0ec_s , the usual ion saturation current. During the other half-period, no ion current crosses the sheath, but electrons flow freely to the substrate in such quantity as to neutralize the ion charge. Since $\phi(t)$ varies sinusoidally over the rf period, the energy gained by an individual ion can be anywhere between 0 and a maximum value $\bar{\phi}_{\max}$, depending on the rf phase when the ion enters the sheath. This type of sheath is different in that the "Ohm's Law" specifies a constant current density nec_s . The independent variable then becomes the rf voltage and the ion stream power dissipated per unit area is the average of $nec_sV(t)$ over the half cycle with the wall negative with respect to the plasma.

It is generally advantageous to have the ion bombardment of the substrate occur at a specified ion energy of choice, but in etching practice a spread in ion energies is often tolerated in order to operate in the high density regime. If it is important that the ion bombardment be monoenergetic, one may choose to use an rf frequency that is higher than the usual 13.56 MHz, in order to insure that $\omega > \omega_{pi}$ and the parameters fall in the GL regime.

4. Magnetic Effects and rf Electron Heating

In LAPPS, ions and electrons transporting to the substrate must cross a magnetic field **B** of order 0.01 tesla which is parallel to the substrate. The magnetic field is sufficiently weak that the ions may be considered unmagnetized, but the plasma energetics and the rf sheath properties are altered in several ways by magnetic effects on the electrons. Some of these effects have been discussed by Lieberman and Lichtenberg [11].

A. Plasma resistivity and ohmic heating of the quasineutral bulk plasma

In the quasineutral plasma, the rf current carried by the magnetized electrons is given by a tensor Ohm's law, $\underline{J}_{rf} = \underline{\sigma}(\omega) \cdot \underline{E}_{rf}$, where the rf conductivity tensor, in the two dimensions (x,y) perpendicular to the magnetic field, is

$$\underline{\sigma}(\omega) = -\frac{\epsilon_0 \omega_{pe}^2}{\left[\Omega_e^2 - (\omega + iv_e)^2 \right]} \begin{bmatrix} i\omega - v_e & \Omega_e \\ -\Omega_e & i\omega - v_e \end{bmatrix} \quad (10)$$

Here Ω_e is the electron gyrofrequency and v_e is the electron collision frequency. The substrate to be processed may or may not be a conductor, but it is mounted on an rf-biased platen which is a conductor, and which in the case of LAPPS is very large (typically 1 m). The electric field parallel to the surface, $E_{rf,y}$, must be zero within the platen, and one may assume that it is zero as

well within the LAPPS plasma, which lies within a few cm of the platen. In the typical situation in which the LAPPS plasma lies between *two* large conducting plates, the entire area between the plates will have $E_{rf,y}$ close to zero. It follows that the rf electron current $J_{rf,x}$ flowing toward the platen is given by the simple Ohm's law

$$J_{rf,x} = \sigma_{\perp}(\omega)E_{rf,x} = -\frac{\epsilon_0 \omega_{pe}^2 (i\omega - \nu_e)}{[\Omega_e^2 - (\omega + i\nu_e)^2]} E_{rf,x} \quad (11)$$

and that there is a large rf current $J_{rf,y}$ flowing adjacent to the plate.

The Ohm's law (11) has the same form as for an unmagnetized plasma, but the conductivity element $\sigma_{\perp}(\omega)$ is much smaller in the magnetized plasma. Therefore, for a given value of $J_{rf,x}$, the electric field $E_{rf,x}$ in the plasma is much larger than it would be in the unmagnetized plasma. Ohmic dissipation is thus much larger in the magnetized plasma than in unmagnetized plasma at the same current. In essence, the power is deposited in the plasma primarily by dissipation of the electron stream energy in the y direction, which is much larger than the streaming energy in the x direction. The Ohmic power P_{Ω} deposited in the plasma, per unit substrate area, is $(J_{rf,x}^2 / \sigma_{\perp})w$ where w is the width of the plasma. For the case of molecular oxygen plasma, ν_e is given by Hake and Phelps [12] as $\nu_e = 4 \times 10^{-14} T_e^{2/3} n_{mol}$ where n_{mol} is the number density of molecules. Assuming $\Omega_e \gg \omega, \nu_e$, we find P_{Ω} in W/m^2 to be

$$P_{\Omega} = \frac{1.7 \times 10^2 B^2 J_{rf}^2 w}{\left(\frac{n}{10^{16}}\right) P T_e^{2/3} \left[1 + \frac{63}{P^2 T_e^{4/3}} \left(\frac{f}{13.56 \text{MHz}}\right)^2\right]} \quad (12)$$

B. Electric field in the electron sheath

We now consider whether the rf electric field in the electron-free region $0 \leq x < s(t)$ is also modified by the presence of the magnetic field. From Poisson's equation, the electric field in the sheath is given by

$$E_x(x) = -\epsilon_0^{-1} e \int_x^{s(t)} n_i(x') dx' + E_p(t), \quad (13)$$

where $E_p(t)$ is the spatially uniform field in the region $x > s(t)$ where the plasma is quasineutral. In an unmagnetized discharge, it is always well justified to neglect E_p in Eq. (13), since the plasma conductivity is very large. However, as we have seen σ_{\perp} is much smaller in a magnetized plasma, and one may wonder if last term on the right of Eq. (13) represents a significant modification to $E_{rf,x}$. Setting the displacement current in the sheath equal to the conduction current in the plasma gives

$$4\pi J_{rf} \cos \omega t = -4\pi n_i(s) e \frac{ds}{dt} + \epsilon_0 \frac{dE_p}{dt}. \quad (14)$$

Using (10) to specify E_p in terms of $J_{rf,x}$, we find that

$$\left(1 - \frac{\omega \Omega_e^2 (\omega - i\nu_e)}{\omega_{pe}^2 (\omega^2 + \nu_e^2)} \right) J_{rf} = -n_i(s) e \frac{ds}{dt}. \quad (15)$$

For LAPPS parameters, where normally $\omega_{pe} \gg \Omega_e \gg \nu_e$, ω , ν_i , the second term in brackets is very small. Therefore, the presence of the magnetic field does not significantly modify the standard relation between $J_{rf,x}$ and ds/dt , and it follows that Eqs. (7) and (8), which gives $\bar{\Phi}_0$, the ion bombardment energy, and the ion bombardment power in terms of J_{rf} , are unchanged.

C. Stochastic heating

As the sheath edge $s(t)$ oscillates back and forth, electrons arriving at the sheath edge bounce off the potential barrier at $x=s(t)$ and back into the plasma. If the incident electron speed is v_x , an electron which strikes the sheath at time t bounces off with speed $v_x + 2\dot{s}(t)$. In the conventional case of an unmagnetized plasma, the difference between the energy flux of electrons arriving at the sheath and the energy flux of electrons reflected off the sheath is [13]

$$H = \mu u_0^2 n_0 \sqrt{\frac{k_B T_e}{8\pi m}} \quad (16a)$$

where u_0 is the maximum sheath velocity which is roughly $s_0 \omega / 2$, where s_0 is given in Eq. (5).

We now consider how stochastic heating is modified by the magnetic field. In LAPPS we normally have $\Omega > \omega$. This being the case, an electron reflected off the moving sheath does not continue into the bulk plasma; it is turned around by the magnetic field and returns for repeated encounters with the sheath edge, at time intervals $\sim \pi / \Omega_e$. This is illustrated in Fig. 2. The repeated coherent bouncing of an electron off the sheath edge is eventually terminated by collisional decorrelation, as pointed out by Lieberman and Lichtenberg[11]. Summing up the repeated energy gain until the bouncing electrons are collisionally decorrelated, they find that the stochastic energy flux in a magnetized plasma is given by

$$H = n_0 \mu u_0^2 \sqrt{\frac{k_B T_e}{2\pi m}} \frac{\Omega_e}{\pi (\nu_e^2 + \omega^2)} \left(\nu_e + \frac{\Omega_e}{\pi} \right). \quad (16b)$$

Since $\Omega_e > \omega$ and $\Omega_e \gg \nu_e$ in LAPPS, stochastic heating is greatly enhanced by the magnetic field. In W/m^2 , the stochastic power into the plasma, per unit area, is

$$P_{sto} = \frac{2.8 \times 10^{-2} J_{rf}^2 T_e^{1/2} B}{\left(\frac{n}{10^{16}} \right) \left[1.1 \left(\frac{f}{13.56 \text{MHz}} \right)^2 + 5.4 \times 10^6 T_e^{4/3} P^2 \right]} \left(\frac{\bar{\phi}_0}{T_e} \right). \quad (17)$$

In Eq. (17) we have left the factor $\bar{\phi}_0/T_e$ as a dimensionless factor. Typically this factor varies between about 1 and 100 for processing plasmas.

D. Power partition in LAPPS

In LAPPS, the primary purpose of rf bias is to accelerate ions into the substrate, and we will normally wish to have

$$P_{ion} > P_{\Omega}, P_{sto}. \quad (18)$$

so that the rf power is used efficiently. The formulas (8), (12), (17) for P_{ion} , P_{Ω} , and P_{sto} are all rather complicated, and it is not easy to state general rules for the satisfaction of Eq. (18). We note that (17) is indeed satisfied for $B=0.01$, $T_e=1$ eV, $P=13$ Pa, plasma width $w = \text{few} \times 10^{-2}$ m, $\bar{\phi}_0/T_e \leq 100$, and $J_{rf} \geq 100 \text{A/m}^2$. In general, increasing B results in a greater fraction of the rf power going into ohmic and stochastic electron heating. Indeed, this is the physical basis of the MERIE [11, 14] (magnetically-enhanced reactive ion etching) scheme, in which the magnetic field is used simply to increase the coupling of rf energy into the electrons and therefore achieve higher plasma density in a capacitive rf discharge. In LAPPS, the discharge is sustained by beam ionization, but rf power can be used as an additional control over the electron temperature. Beam-produced plasma is intrinsically very cool [1], typically with $T_e < 1$ eV. For some applications, it is advantageous to operate at higher T_e , and rf power can be applied in a controlled way to increase T_e . The total power deposited in the plasma by the beam, per unit area, is

$$P_b = J_b \frac{dE_b}{dz} w, \quad (19)$$

where J_b is the beam current density and E_b is the beam electron energy. In Ref. [1] we discuss the parameter dependence of beam heating at some length. In W/m^2 , for an oxygen plasma, we find

$$P_b = 9.5 \times 10^5 J_b w P \frac{\ln(40E_b)}{E_b(\text{keV})} \quad (20)$$

For typical LAPPS parameters, where E_b is several keV and J_b is 100A/m^2 or more, the beam power P_{beam} substantially exceeds the rf electron heating powers P_{Ω} and P_{sto} and is the primary energy source sustaining the discharge.

5. Return Path for the rf Current

When a substrate is rf-biased to induce ion bombardment, the rf circuit is completed by current flow through the plasma to other bounding surfaces. This leads to the possibility of ion bombardment of these other surfaces, which is generally undesirable. Thus it is important to determine how and where the rf return current is distributed. In conventional reactors, e.g. reactive ion etchers, a fairly uniform plasma fills a grounded vessel whose surface area is large compared to the substrate. The rf return current is spread over a large area of the vessel, so that the return current density J_{ret} is much lower than the rf current density J_{rf} at the driven substrate. The ion bombardment voltage $\bar{\phi}_0$ and power flux P_{ion} scale as J_{rf}^4 , in accordance with Eqs. (7) and (8), so very little of the ion power is dissipated in the vessel. But in LAPPS the substrate is very large, and all of the plasma ionization occurs within a broad, thin sheet parallel to the substrate. The high-density plasma is localized within a slab-shaped region which is similarly broad but only a few cm thick. It is natural to use a vessel which of similar shape; the area of the substrate will then be comparable to the interior area of the vessel. Thus ion bombardment of the vessel can be significant. In this section we shall discuss the strategies that can be used to control and optimize the effects of rf return current in LAPPS.

We consider first the possible effect of the magnetic field in guiding the return current flow. The mobility of the magnetized electrons is much greater along than across the magnetic field \mathbf{B} , so there might be a concern that the rf return current will be directed along \mathbf{B} to a small area at the end of the vessel, as shown in Fig. 3a. However, under normal circumstances potential variations within the plasma are small compared to the potential drop through the sheath from the plasma to ground. Hence $\bar{\phi}_0$ is essentially the same at all points where the plasma makes contact with a grounded surface, and according to Eqs. (7) and (8) J_{rf} and the ion power flux are uniform to all such points. The magnetic field plays little or no role in determining the path taken by the return current, because the impedance of the current flow path through the plasma (either along or across the field) is much less than the capacitive impedance of the GL sheath at the grounded return current collector, or the resistive impedance of the high-density sheath. A more detailed discussion of the return current circuit properties is given in the Appendix.

Next, one might consider whether the rf current can be deposited on floating insulator surfaces, e.g. an insulator containment vessel. In effect, the connection to ground would then be through stray capacitance from the vessel to ground. We show in the Appendix that the impedance associated with this capacitive path far exceeds the sheath impedances discussed in the previous paragraph. Thus the return current will always take the path to any available grounded conductor, rather than depositing on an insulating surface, and it is in fact essential that a suitable grounded return current collector be provided.

In LAPPS, it would be natural to provide a grounded return-current collector on the opposite side of the plasma sheet from the substrate, as shown in Fig. 3b. If the collector is placed symmetrically at the same stand-off from the plasma sheet, the ion bombardment of this electrode will be equal to that of the substrate. To avoid wasting ion energy in bombardment of a passive surface, the "return-current collector" could be a second substrate placed on a grounded

platen, to be processed simultaneously. Alternatively, it may be preferable to leave the vessel wall opposite the substrate uncovered, for gas inflow and/or pumping ports. The return-current collector need not be the vessel wall itself; it could be a grounded mesh placed at a location of choice adjacent to the plasma sheet. Equation (7) indicates that the ion bombardment voltage $\bar{\phi}_0$ is proportional to n_0^{-2} , where n_0 is the plasma density near the surface, and therefore $P_{\text{ion}} = n_0 e c_s \bar{\phi}_0$ is proportional to n_0^{-1} . To maximize the fraction of the rf power that goes into ion bombardment of the substrate, and minimize ion bombardment of the collector, one might choose to locate the return current collector as close as possible to the plasma source, where the plasma density is highest.

An alternative strategy would be to allow the return current to flow to an area on the grounded vessel which is much larger than the substrate, as in conventional RIE. This is illustrated in Fig. 3c. It would be necessary to use a large-volume containment vessel, filled with low-density plasma which surrounds the main LAPPS plasma sheet. This plasma could be the result of diffusion from the LAPPS source, or it could be actively generated in various ways. The return current would then flow through this low-density plasma to the walls of the chamber. Since $\bar{\phi}_0 \propto n_0^{-2}$ and $P_i \propto n_0^{-1}$, and the plasma density n_0 is lower at the remote walls than at the substrate, there is a tendency for the bulk of the ion power to be deposited in the walls rather than the substrate. On the other hand, $\bar{\phi}_0 \propto J_{\text{ret}}^4$, and the return current density J_{ret} is inversely proportional to the vessel area A_{wall} . To insure that the total ion power dissipated by the return current is less than the ion power to the substrate, it is necessary that

$$\left[n_0 A^3 \right]_{\text{substrate}} < \left[n_0 A^3 \right]_{\text{wall}}. \quad (20)$$

It is not likely that this strategy will be used in LAPPS, since it would mean using a vacuum vessel that is much larger than would otherwise be necessary, and filling it with plasma.

A fourth possibility, which is unique to LAPPS, is that the plasma falls off to essentially zero density (more precisely, to very low conductivity) before making contact with any wall. One might then envision this as a free surface of the plasma, as shown in Fig. 3d, supporting an rf-varying surface charge density (and hence a displacement current in the region outside the plasma). In that case, there need not be an rf conduction return current to any part of the vessel. However, we show in the Appendix that this scenario is untenable, because the rf electric field outside the plasma would easily break down the adjacent un-ionized background gas. The highly conducting LAPPS plasma sheet would, in effect, function as an electrode coupling the rf power into the adjacent gas, and thus driving a capacitive discharge which of necessity fills the vessel with plasma.

We conclude that in all cases the rf circuit must be completed by current flowing from the biased substrate, through the plasma sheet, to a grounded collector on the other side. In a LAPPS reactor, this collector is expected to be either a screen located at a suitably chosen distance from the plasma sheet, or a second substrate symmetrically located on a grounded platen.

6. Summary

Etching and certain other processes require directional ion bombardment of the substrate at controlled energies well above the floating potential. In LAPPS as in other types of plasma reactor, this is accomplished by rf-biasing the platen on which the substrate sits. However the phenomena associated with rf bias in LAPPS differ in some ways from the usual plasma reactors. In LAPPS, it is possible to operate with high plasma density adjacent to the substrate, so that $\omega < \omega_{pi}$. In this case, the usual Godyak-Lieberman theory becomes inadequate, and one must use the high-density model outlined in Sec. 3. In this regime, the ions arrive at the substrate with a broad energy distribution. If it is desirable to operate at high plasma density with a narrow ion spectrum, one could choose to use an rf bias frequency that is higher than the usual 13.56 MHz, so as to remain in the regime $\omega > \omega_{pi}$. Because the plasma electrons are magnetized, the fraction of the rf energy that goes into ohmic and stochastic electron heating is greatly increased. In LAPPS it is not necessary to use rf heating to maintain the plasma density, but the rf electron heating can be used to increase and control the electron temperature. A LAPPS reactor must be designed so that the rf return current does not dissipate an excessive power load in areas other than the substrate. The usual design solution is to carry the rf return current to the reactor walls through fairly high-density plasma, and spread it out over a large area of the walls. In LAPPS, it may be preferable to direct the return current to a large-area collector immediately adjacent to the plasma sheet, on the side opposite the rf-biased substrate. In various situations, this may be the reactor wall, or a collector screen, or another substrate symmetrically located for processing.

Appendix: Circuit analysis of the rf return current

A. Plasma and sheath impedances

For cross-field rf current flow to a collector located a distance w from the substrate, the plasma impedance is $Z_{\perp} = w/A\sigma_{\perp}(\omega)$, where A is the area of the substrate and $\sigma_{\perp}(\omega)$ is the cross-field rf conductivity. Taking the large- Ω limit of Eq. (11), we see that Z_{\perp} is the sum of a resistive and an inductive impedance,

$$Z_{\perp} = R + i\omega L, \quad (A1)$$

where

$$R = \frac{v\Omega^2}{\omega_{pe}^2 \epsilon_0 (\omega^2 + v^2)} \frac{w}{A}, \quad (A2a)$$

$$L = -\frac{\Omega^2}{\omega_{pe}^2 \epsilon_0 (\omega^2 + v^2)} \frac{w}{A}. \quad (A2b)$$

Notice that the bulk magnetized plasma in the LAPPS configuration has a *negative* inductance.

For rf current flow to the end walls, the flow path would be mostly parallel to \mathbf{B} , and the $\Omega_e=0$ limit of Eq. (10) applies. The resistance and inductance is then

$$R = \frac{v}{\omega_{pe}^2 \epsilon_0} \frac{w_{end}}{A_{end}}, \quad (A3a)$$

$$L = \frac{1}{\omega_{pe}^2 \epsilon_0} \frac{w_{end}}{A_{end}}, \quad (A3b)$$

where A_{end} is the mean cross-sectional area of the current path and w_{end} is a mean distance from the edge of the substrate to the end plate. For the unmagnetized case, the plasma has a positive inductance. On the other hand, the sheaths are capacitive with capacitance give by

$$C_{sh} = \frac{\epsilon_0 A}{s_0}, \quad (A4)$$

where s_0 is the sheath thickness and A is the area of the sheath. Taking typical LAPPS parameters $\omega \sim 8.5 \times 10^7 \text{ sec}^{-1}$, $v \sim 10^8 \text{ sec}^{-1}$, $\Omega \sim 1.7 \times 10^9 \text{ sec}^{-1}$, $\omega_p \sim 5.6 \times 10^{10} \text{ sec}^{-1}$, $w \sim 0.04 \text{ m}$, $s_0 \sim 10^{-3} \text{ m}$, it is clear that $Z_{||} \ll Z_{\perp} \ll Z_{sh}$, as noted in Sec. 5. Thus the sheath impedance is dominant, and the anisotropic plasma impedance plays little role in determining the rf current flow pattern through the plasma.

The rf current densities to the collector plate and to the end walls arrange themselves so that the sheath voltage drop is the same at both locations. From Eq. (7), this implies that

$$\frac{J_{rf,end}}{J_{rf,plate}} = \sqrt{\frac{n_{0,end}}{n_{0,plate}}}. \quad (A5)$$

Since the dc ion current density is $n_0 e c_s$, the ion bombardment powers P_i per unit area, at the two locations, are related by

$$\frac{P_{i,end}}{P_{i,plate}} = \frac{n_{0,end}}{n_{0,plate}}. \quad (A6)$$

The density n_0 may be larger at the end wall, which is located directly in the path of the electron beam source. Thus the ion power per unit area may be larger at the small-area end walls. This is probably not a major concern, since the end wall must also serve as a dump for the keV electron beam. However, locating the collector plate as close as possible to the plasma source serves to minimize P_i at the end wall. If it is desirable to further decrease $J_{rf,end}$ and $P_{i,end}$, this can be done by connecting the end wall to ground through a large inductor, to increase the rf impedance.

B. RF current flow to an insulating surface

RF current flow often terminates on an insulating substrate, or on an insulating film layer on the inside of a conducting vessel. This is possible because the insulating surface is separated from a grounded or driven conductor by a thin layer of width δw , so that the capacitive impedance of the insulating layer, $-i\epsilon_0\delta w/\omega A$, is small. However, if rf current terminates on an insulating part of a vessel which separated from ground by a substantial distance w_{gr} , the capacitive impedance to ground will be of the order of $-i\epsilon_0\delta w_{gr}/\omega A$ (with some dependence on the details of the geometry), and is typically large compared to other impedances in the system. Thus the return current will always seek out a grounded terminus, and it is important to provide an appropriate grounded path for the return current.

C. RF breakdown of any region of un-ionized gas

Assume that within the containment vessel there is an interface between plasma and un-ionized gas, and let E_g be the rf electric field on the gas side, and E be the rf electric field on the plasma side. Using the continuity equation and the time derivative of Poisson's equation, and integrating once, we get

$$\epsilon_0 \frac{\partial E}{\partial t} + J_{rf} = \epsilon_0 \frac{dE_g}{dt}, \quad (A7)$$

Fourier transforming Eq. (A7) in time, and using Ohm's law in the form (9), we find

$$E_g(\omega) = \left(1 + \frac{i\sigma_{\perp}(x, \omega)}{\epsilon_0\omega} \right) E(x, \omega). \quad (A8)$$

Using σ_{\perp} from Eq. (9), and assuming $\Omega_e \gg \omega, v_e$, we find

$$E_g(\omega) = \frac{\omega_{pe}^2}{\Omega_e^2} \left(1 + \frac{iv_e}{\omega} \right) E(x, \omega) \gg E(x, \omega) \quad (A9)$$

Taking a typical magnitude $J_{rf} \sim 100 \text{ A/m}^2$ for the rf current density, $E \sim 10 \text{ V/m}$ for typical LAPPS parameters, and thus $E_g \sim 10^4 \text{ V/m}$. E_g is many orders of magnitude in excess of the breakdown field for a gas at pressure 5 to 10 Pa. Thus the use of rf bias at 13.56 MHz guarantees that the plasma-filled region will be spatially extended until it makes electrical contact with the vessel.

Up to this point, we have tacitly assumed that ion conductivity is negligible compared to electron conductivity. For rf operation at the usual frequency of 13.56 MHz, this is a very good

assumption, but at lower rf frequencies, ion conductivity can become important. If we include the ion contribution to σ_1 , Eq. (A9) becomes

$$E_g(\omega) \equiv \left[\frac{\omega_{pe}^2}{\Omega_e^2} \left(1 + \frac{iv_e}{\omega} \right) - \frac{\omega_{pi}^2}{\omega^2 + iv_i \omega} \right] E(x, \omega). \quad (\text{A10})$$

Notice that if collisions are neglected E_g goes to zero at the lower hybrid resonance frequency $\omega = \omega_{LH} \equiv (m/M)^{1/2} \Omega_e$. At resonance the rf space charge is sustained by the lower hybrid waves within the plasma. One might think that gas breakdown could be avoided by choosing f_{LH} (which is 1.2 MHz for an O_2 plasma with $B = 10^{-2}$, well below the usual 13.56 MHz) as the rf bias frequency. But electron collisionality is *not* negligible at $\omega = \omega_{LH}$, it is in fact dominant, and Eq. (A10) thus reduces to

$$E_g(\omega) = \frac{iv_e}{\omega} \frac{\omega_{pe}^2}{\Omega_e^2} E(x, \omega), \quad (\text{A11})$$

Since $v_e \gg \omega$ when $\omega = \omega_{LH}$, the surface electric field is even larger at f_{LH} than at 13.56 MHz, so gas breakdown still occurs.

Acknowledgements: The authors would like to thank Robert Meger and Darren Leonhardt of NRL, and Michael Lieberman of Berkeley for a number of useful discussions. This work was supported by the Office of Naval Research.

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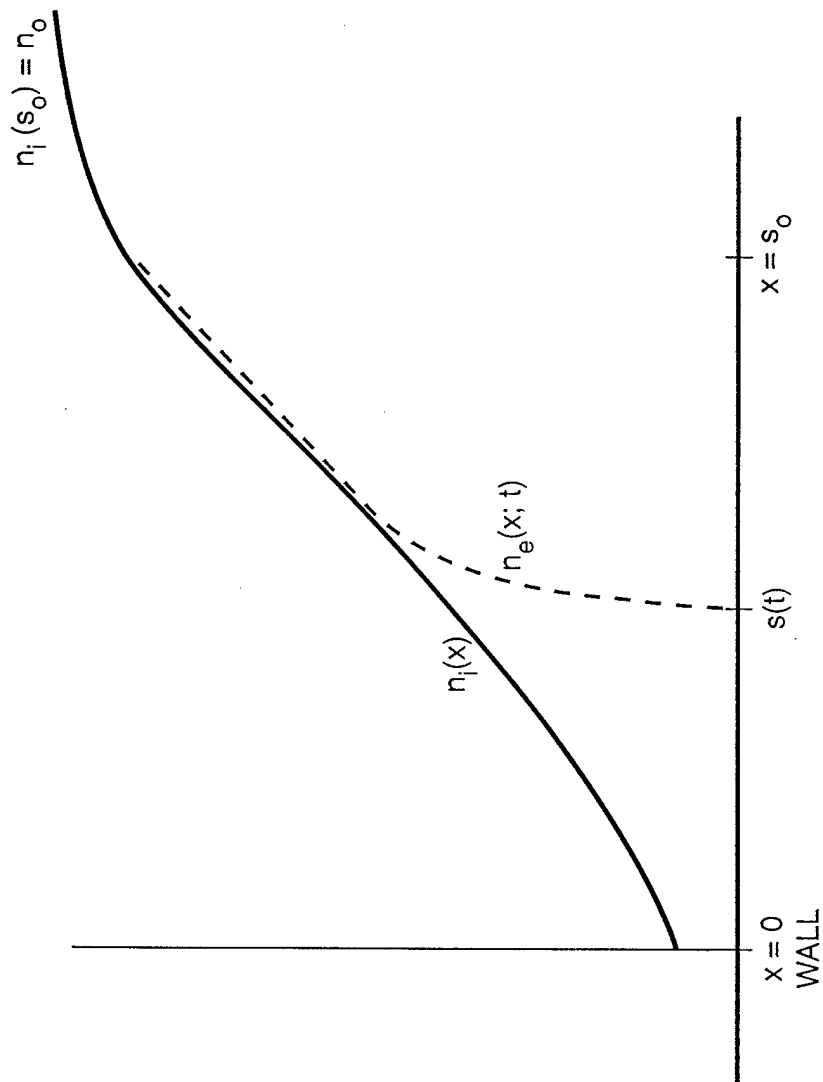


Fig. 1. A schematic of the electron and ion density in an rf sheath where $\omega_{pi} \ll \omega \ll \omega_{pe}$

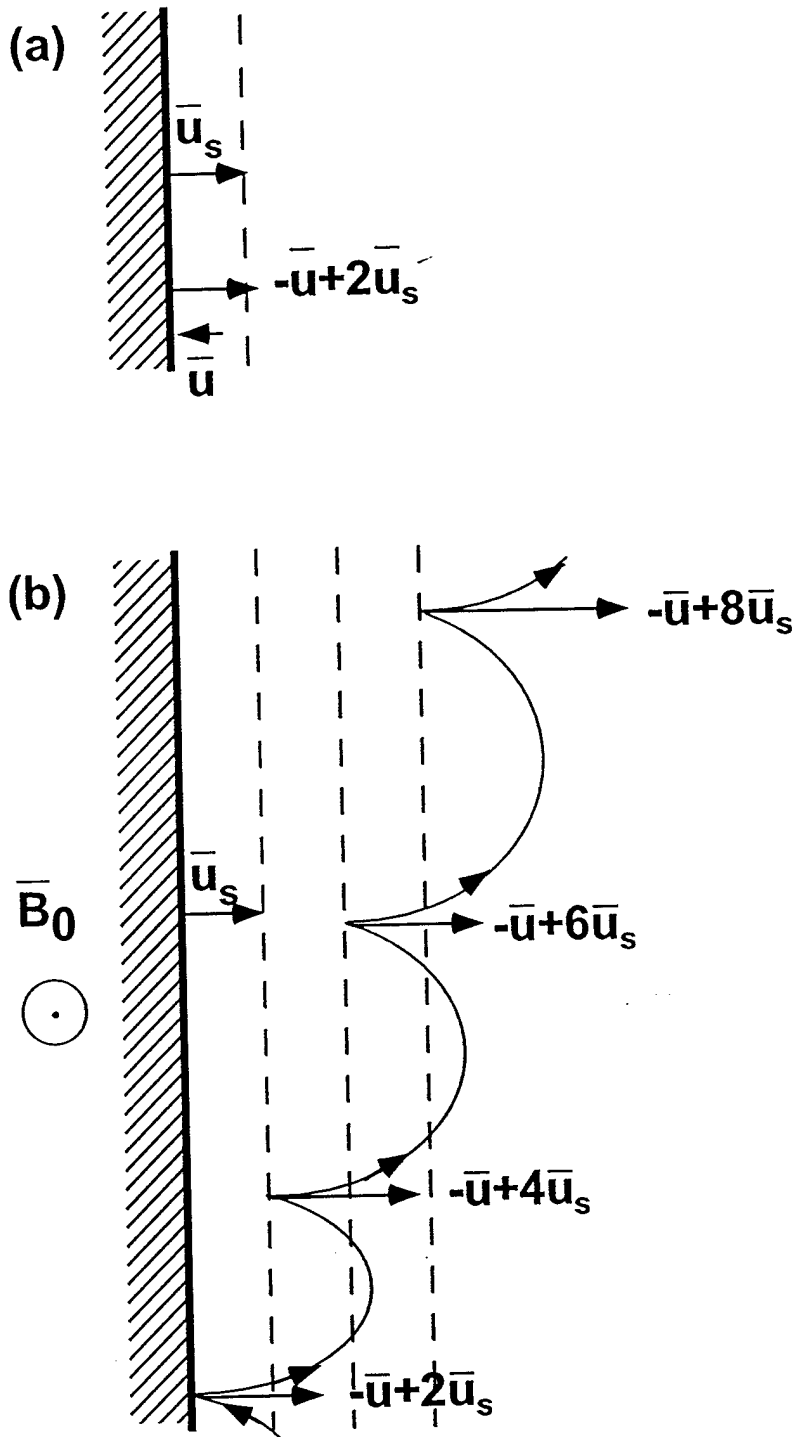


Fig. 2. Repeated bounces of stochastically heated electrons in a magnetized sheath.

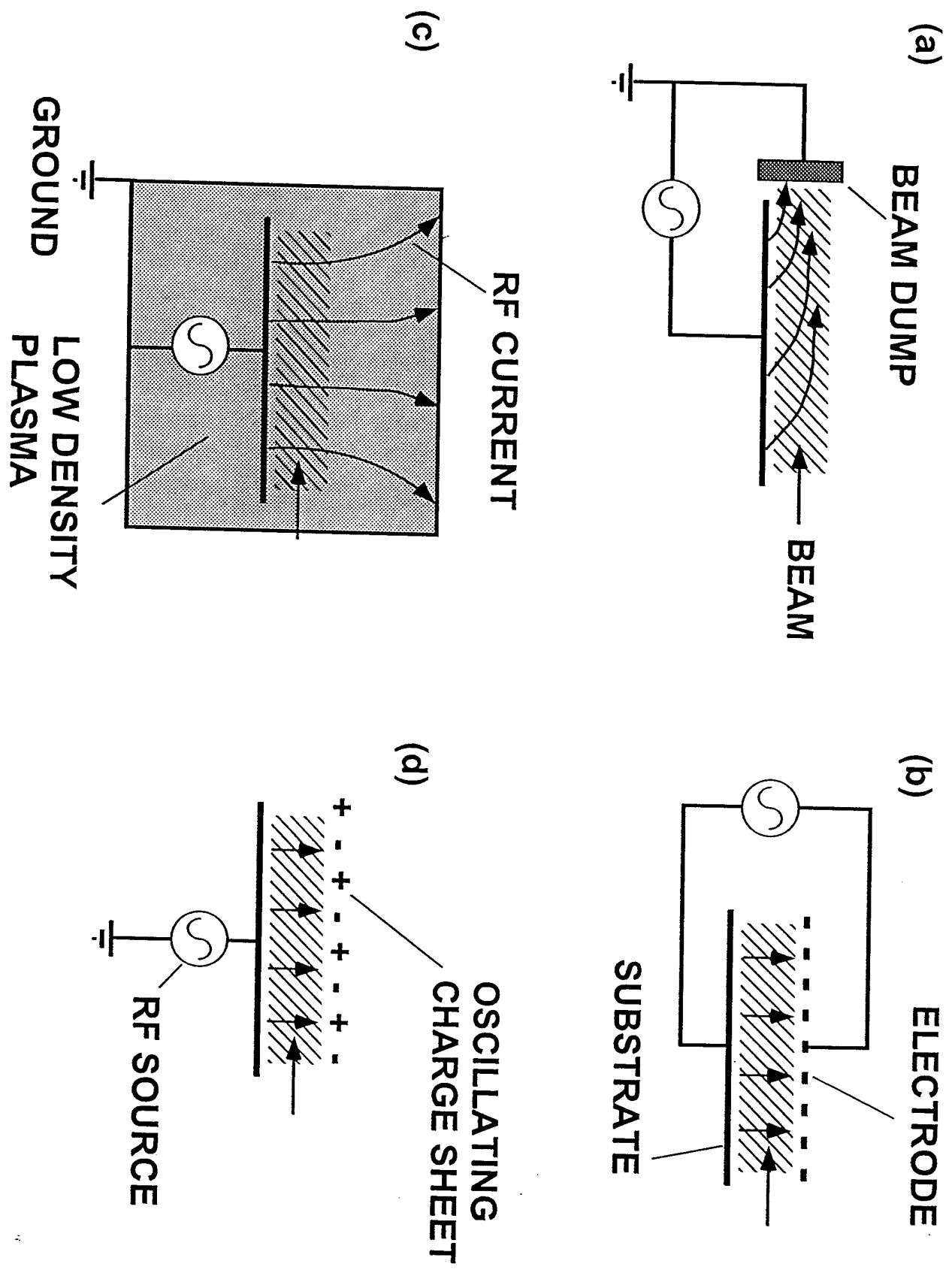


Fig. 3. Possible return current paths in an rf sheath in a LAPPS plasma.