Nonlocal electron kinetics in a low-pressure inductively coupled radio-frequency discharge

V. I. Kolobov, D. F. Beale, L. J. Mahoney, and A. E. Wendt

Engineering Research Center for Plasma-Aided Manufacturing, University of Wisconsin-Madison, Madison, Wisconsin 53706

(Received 17 March 1994; accepted for publication 18 May 1994)

The Boltzmann equation for electrons is analyzed for a low-pressure inductively coupled rf discharge in argon driven by a planar coil. Spatially resolved probe measurements of the electron distribution function (EDF) indicate that the total energy of electrons is an argument of the EDF. Pressure dependence of the light emission distribution is explained on the basis of nonlocal electron kinetics.

Inductively coupled rf discharges are currently being investigated as plasma sources for material processing.¹ For the pressure range of interest, 1–100 mTorr, the electron energy relaxation length exceeds the characteristic discharge dimensions (of the order of 10 cm) and the electron distribution function² (EDF) is not in equilibrium with the local electric field (see for example Ref. 3 and references therein). In such a nonlocal regime, a kinetic treatment of electrons is necessary to calculate the ionization rate, plasma density, electron current density, etc. In this letter we analyze the main consequences of nonlocal electron kinetics by examining the Boltzmann equation for a given spatial distribution of electric fields (both electrostatic and rf fields) to explain experimental data for a planar inductively coupled discharge in argon. The discharge's most distinguishing feature appears in the spatial separation of electron heating and ionization. Though the electron heating occurs in the region of a high rf field, a substantial portion of ionization is produced near the maximum of the electrostatic potential where the rf field is absent.

The discharge is driven by the electric field from a spiral coil placed on the dielectric roof of a squat cylinder with metallic walls and bottom (Fig. 1). The azimuthal rf field E_θ is localized in the vicinity of the antenna and peaks in an annular region at roughly half of the chamber radius. The electron heating mechanism depends on the relative magnitude of the three characteristic dimensions: field penetration depth δ , electron mean-free path λ , and the length v/ω which electron passes during the field period.⁴ When λ is the smallest of them, the ohmic heating takes place, when the smallest is δ , the collisionless heating is expected to occur.⁵

Since the electron energy relaxation length exceeds the spatial scale, the total energy $\epsilon = w - e\phi(z, r)$ is the argument of the EDF.^{6–8} The observed profile of the electrostatic potential $\phi(z, r)$ is shown schematically in Fig. 2. Curves 1 and 2 correspond to axial distribution of $\phi(z, r)$ at $r=0$ and at the point r_0 , where the maximal amplitude of the rf field is reached. The potential of the metallic wall ϵ_1 with respect to the plasma exceeds the excitation potential ϵ^* . The dashed area indicates the region in (ϵ, z, r) space where inelastic collisions occur. The electrons with $\epsilon > \epsilon^*$ have a kinetic energy w sufficient for excitation only in a region slightly above curve 3, where their potential energy is approximately zero.

We will consider here the case of small anisotropy of the

EDF when the conventional two-term approximation is applicable $f = f_0(\epsilon, z, r) + \mathbf{v}/v \cdot \mathbf{f}_1(\epsilon, z, r, t)$, where $f_1 \ll f_0$. For the rf field $E_\theta(z, r, t) = E_\theta(z, r)e^{i\omega t}$, neglecting the energy loss in elastic collisions and electron-electron collisions, the kinetic equation can be written in the form^{6–8}

$$\text{div} \left(\frac{v^3}{3\nu} \right) \text{grad } f_0 + \frac{\partial}{\partial \epsilon} \nu D_\epsilon(z, r, \epsilon) \frac{\partial f_0}{\partial \epsilon} = \nu \nu^*(w) f_0 - \sqrt{v^2 + 2\epsilon^*/m} (w + \epsilon^*) f_0(\epsilon + \epsilon^*), \quad (1)$$

$$\mathbf{f}_1 = -\frac{v}{\nu + i\omega} \left(\text{grad } f_0 + e \mathbf{E}_\theta \frac{\partial f_0}{\partial \epsilon} \right), \quad (2)$$

where ν and ν^* are the momentum transfer and inelastic collision frequencies ($\nu \gg \nu^*$), and $D_\epsilon = e^2 E_\theta^2(z, r) v^2 \nu / 3(\nu^2 + \omega^2)$. For $\omega \ll \nu^*$ the function f_0 is modulated only slightly and can be assumed constant. The last term in Eq. (1) representing the electrons which have suffered inelastic collisions is written for the simplest case of the excitation of a single effective level with energy ϵ^* . The rf heating results in diffusion in ϵ . The energy diffusion coefficient D_ϵ vanishes at $z > \delta$ and depends on r for $z < \delta$.

FIG. 1. Schematic of the discharge. The contour lines show calculated strength of the azimuthal rf field.

FIG. 2. Axial distributions of the electrostatic potential $\phi(r,z)$: (1) on the axis, $\epsilon = -e\phi(0,z)$, (2) in the point r_0 of maximal rf field strength, $\epsilon = -e\phi(r_0,z)$. Curves 3 and 4 shifted by ϵ^* with respect to curves 1 and 2 indicate the region in (ϵ, z, r) space where inelastic collisions take place. The heating occurs mainly at $z < \delta$.

The electrons with $\epsilon < \epsilon_1$ are trapped. Since their energy relaxation is slow compared to spatial diffusion, the EDF for $\epsilon < \epsilon_1$ depends only on ϵ .⁶⁻⁸ The low-energy electrons originate from ionization and inelastic collisions. They cannot overcome the ambipolar potential barrier and reach the region where the heating takes place. Their only relaxation mechanism is electron-electron collisions. So an EDF slope change may be observed at $\epsilon = \epsilon_b$ (Fig. 2) when the diffusion in ϵ switches on. At energies $\epsilon > \epsilon_b$, when the source of slow electrons and e - e collisions can be neglected in Eq. (1), the EDF has the form

$$f_0 = A \int_{\epsilon}^{\epsilon_0} \frac{d\epsilon'}{\langle v(\epsilon') D_{\epsilon}(\epsilon') \rangle}. \quad (3)$$

The spatial averaging $\langle \rangle$ designates the integration over the region bounded by $\epsilon = -e\phi(r,z)$ which is available for electrons with energy ϵ . The normalization constant A is governed by the plasma density and the upper limit ϵ_0 is determined by matching the EDF tail (in the absorbing-wall approximation $\epsilon_0 = \epsilon^*$). For electrons with small energy ϵ , the averaged diffusion coefficient in Eq. (3) is small. The more energetic electrons penetrate deeper into the rf field region and their energy diffusion is large. Consequently, the EDF slope is large at low energies and diminishes with the energy growth.

All these features are observed in the experiment. Spatially resolved probe measurements were carried out in a planar 20-cm-diam by a 14-cm-long inductive coupled rf discharge in Ar driven by a spiral coil at 13.56 MHz with a pressure range 10–100 mTorr and power levels 100–200 W.⁹ In Fig. 3, the EDFs measured at different points across the radius (a) and along the axis (b) are given as functions of the total electron energy. The EDFs coincide over the measured range of total energy and are shifted by the local ambipolar potential.

Regardless of the ionization mechanism, the wall potential in the nonlocal case must exceed the excitation potential

FIG. 3. The axial and radial evolution of the EDF. Curve labels correspond to radial distance (in cm) from the center at a fixed axial distance (4.4 cm) from the dielectric window (a) or axial distance from the dielectric on axis (b); 50 mTorr, 100 W.

$\epsilon_1 > \epsilon^*$. The energy relaxation of the fast electrons $\epsilon > \epsilon^*$ is characterized by the length $\lambda^* = \lambda \sqrt{\nu/\nu^*}$. If λ^* exceeds the characteristic discharge dimension L , the EDF in the region $\epsilon^* < \epsilon < \epsilon_1$ depends only on ϵ and falls off with the energy scale $T_1 \sim e\langle E_0 \rangle \lambda^* \ll \epsilon^*$, where $\langle E_0 \rangle$ is the spatially averaged value of the rf field. Since the total number of ionization events (and thus the total number of events involving loss to the wall) falls off exponentially with increasing of $(\epsilon_1 - \epsilon^*)$, the value of $(\epsilon_1 - \epsilon^*)$ cannot exceed a few T_1 in the case of stepwise ionization (in the case of direct ionization, the value of ϵ_1 should correspondingly exceed the ionization potential). As the probability for loss to the wall for an electron with $\epsilon > \epsilon_1$ is much larger than the probability for an inelastic collision in the case $\lambda^* > L$, the wall potential can be estimated as

$$(\epsilon_1 - \epsilon^*) \sim \sqrt{\frac{e\langle E_0 \rangle \lambda^*}{\epsilon^*}} \ln(1 + 1/\gamma), \quad (4)$$

where $\gamma < 1$ is the probability of ionization relative to excitation events.

At energies $\epsilon > \epsilon_1$ (see Fig. 2) electrons can escape to the metallic wall and the EDF becomes a function of the spatial coordinates. As trapped electrons give zero contribution to the current, only free electrons with $\epsilon > \epsilon_1$ are responsible for the electron current transport. For the metallic surface the ion and electron fluxes should be equal integrated over the entire surface. As the energy diffusion occurs in the vicinity of the

FIG. 4. Normalized Abel transformed optical emission profiles at 200 W for two pressures: (a) 20 mTorr, (b) 50 mTorr. Measurements of the EDF at 200 W demonstrate the same features as those shown in Fig. 3.

dielectric wall, $z=0$ the potential of the wall $\epsilon_2(r)$ must exceed ϵ_1 in such a way to equilibrate the electron and ion fluxes locally. The fast electrons with $\epsilon_2 > \epsilon > \epsilon_1$ are able to climb the hill where $\epsilon_2(r) < \epsilon$ but are trapped with respect to the dielectric surface, where $\epsilon_2(r) > \epsilon$. Thus the boundary condition at $z=0$ can be written in the form

$$\frac{\partial f_0}{\partial z} = 0 \quad \text{at } \epsilon_1 < \epsilon < \epsilon_2(r), \quad f_0 = 0 \quad \text{at } \epsilon_2(r) < \epsilon. \quad (5)$$

Since the characteristic time for the wall loss is shorter than that between inelastic collisions, $\lambda^* > L$, only the first two

terms in Eq. (1) can be kept in the energy range $\epsilon_2 > \epsilon > \epsilon_1$. At $L \sim R$, the calculation of the EDF in this energy range constitutes a two-dimensional problem. The EDF at $\epsilon > \max[\epsilon_2(r)]$ falls off with the scale $T_2 \sim eE_0 \delta \ll T_1$. The difference $[\epsilon_2(r) - \epsilon_1]$ is defined by the local ion flux into the dielectric wall at a point r .

Fast electrons are responsible for the ionization and excitation of atoms. The light emission maximum observed experimentally¹⁰ near the heating region (Fig. 4) is mainly due to the free electrons (localized above curve 4 in Fig. 2). The trapped fast electrons (which have kinetic energy sufficient for excitation in the region above curve 3 in Fig. 2) are responsible for the second maximum of light emission observed near the potential maximum at low pressures [Fig. 4(a)].

If $L > \lambda^*$, the EDF becomes a function of the spatial coordinates even at $\epsilon > \epsilon^*$ and is enriched by the fast electrons at distances less than λ^* from the heating zone. For this case, the emitted light intensity decays outside the heating zone with the spatial scale λ^* ($\lambda^* \sim 4$ cm at 50 mTorr for 15 eV electrons). If the energy loss in inelastic collisions is faster than energy gain by the ambipolar field, the emission maximum in the center disappears [Fig. 4(b)]. Self-consistent discharge modeling involving the proposed electron kinetics is in progress. A quantitative comparison with the experimental data^{9,10} will be published elsewhere.

The authors would like to thank W. N. G. Hitchon for useful discussions. This work was supported in part by NSF.

¹M. A. Lieberman and R. A. Gottscho, *Physics of Thin Films*, edited by M. Francombe and J. Vossen (Academic, New York, 1994), Vol. 18, p. 1–119.

²By EDF we denote the density in the phase space that coincides with the second derivative of probe current with respect to probe potential.

³V. A. Godyak and R. B. Piejak, *Appl. Phys. Lett.* **63**, 3137 (1993).

⁴E. M. Lifshitz and L. P. Pitaevskii, *Physical Kinetics, Course of Theoretical Physics* (Pergamon Press, Oxford, 1981), Vol. 10.

⁵M. M. Turner, *Phys. Rev. Lett.* **71**, 1844 (1993).

⁶I. B. Bernstein and T. Holstein, *Phys. Rev.* **94**, 1475 (1954).

⁷L. D. Tsensin, *Zh. Eksp. Teor. Fiz.* **66**, 1638 (1974) [*Sov. Phys. JETP* **39**, 805 (1974)].

⁸I. D. Kaganovich and L. D. Tsensin, *IEEE Trans. Plasma Sci.* **20**, 66 (1992).

⁹L. J. Mahoney, A. E. Wendt, C. J. Richards, E. Barrios, and J. L. Shohet (unpublished).

¹⁰D. A. Beale, A. E. Wendt, and L. J. Mahoney, *J. Vac. Sci. Technol. A* (to be published).