Experimental Evidence of Collisionless Power Absorption in Inductively Coupled Plasmas

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Electromagnetic power absorbed in inductively coupled plasma driven by a planar coil has been found directly from axial distributions of the rf electric field and current density measured with magnetic probes. It is shown that at gas pressure around 1 mTorr the absorbed power is much larger than that found from the cold plasma theory, and is in reasonable agreement with the one calculated in the framework of a theory for the anomalous skin effect in a two- dimensional system accounting for spatial dispersion of plasma conductivity. [S0031-9007(98)05767-6]

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Nonlocal electron kinetics and nonlocal plasma electrodynamics are now well recognized as playing important roles in inductively coupled plasmas (ICP) at low gas pressure [1–5]. Both the electron energy distribution and the rf current density in such discharges are not local functions of the rf electric field. The nonlocality of the current is due to thermal electron motion which results in spatial dispersion of the plasma conductivity in the weakly collisional regime.

The spatial dispersion of conductivity underlies the anomalous skin effect and results in nonmonotonic distributions of the rf field and current density [6] and a local negative power absorption in ICP [7]. It has been shown long ago that the spatial dispersion effect may lead to collisionless power dissipation in the skin layer [8]. Recently, collisionless power absorption in inductive plasma has been a subject of many theoretical and modeling efforts [2–5,9–11]. Although collisionless electron heating is now widely accepted, to our knowledge no direct experimental proof of its existence has been given.

In this Letter we report on experimental observation of collisionless electron heating in ICP by comparing the measured power absorption with that calculated using the cold plasma theory and that calculated using the theory of the anomalous skin effect in a plasma slab [12] extended to two-dimensional systems as described in Refs. [4,7].

The experiments were carried out in a cylindrical ICP in argon gas driven with a planar induction coil in a stainless steel chamber with a Pyrex glass bottom [6,7]. The chamber inside diameter $2R = 19.8$ cm, its length $L = 10.5$ cm, and the glass thickness was 1.27 cm. A five turn planar induction coil was mounted 1.9 cm below the bottom surface of the discharge chamber. An electrostatic shield between the glass and the coil has practically eliminated capacitive coupling between rf coil and the plasma to the extent that the rf plasma potential referenced to the grounded chamber wall was less than 0.1 V.

The measurements of the azimuthal rf electric field and current density and their phase distribution along the chamber axis at a fixed radial position of 4 cm (near the maximum in radial distribution of the rf electric field) were performed with a miniature magnetic probe. Detailed probe description, signal processing, and validation of results obtained with this probe are given in Ref. [13]. The axial distributions of the azimuthal rf field and current density [their rms magnitudes, $E(z)$ and $J(z)$ and phases, $\phi_E(z)$ and $\phi_J(z)$] measured in this experimental setup are reported in Ref. [6].

Langmuir probe measurements were performed both in the discharge center ($r = 0, z = 5$ cm) [14] and along the axis at a fixed radial position of $r = 4$ cm [7] to determine the electron energy distribution function (EEDF). The integrals of the EEDF such as electron density n_e, effective electron temperature T_e, and electron-atom collision frequency in the rf field, ν_{en}, calculated from these measurements are given in Table I.

The axial distribution of the absorbed power density was found directly from the measured quantities as: $P(z) = E(z)J(z)\cos[\phi_E(z) - \phi_J(z)]$. It has been shown in Ref. [7] that under conditions of the anomalous skin effect $P(z)$ can be nonmonotonic and can even become negative beyond the skin layer. The power absorption by the plasma can be characterized by the absorbed power flux $S(z)$

$$S(z) = \int_0^z P(x)\,dx\,, \tag{1}$$

which is equal to the loss of the Poynting flux $(c/4\pi)EH$. The absorbed power flux is shown in Fig. 1 for 1 and 0.3 mTorr. For both pressures, the driving frequency was 6.78 MHz and the total absorbed discharge power $P_d = 100$ W. For comparison, the power flux $S_{\mathrm{col}}(z)$ calculated from a cold plasma theory accounting only for the collisional heating is also shown in Fig. 1. It is seen

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TABLE I. The discharge parameters measured in the skin layer ($z = 1$ cm and $r = 4$ cm) for $\omega/2\pi = 6.78$ MHz.

p mTorr	P W	n 10^{10} cm^{-3}	T_e eV	$\omega_{\rm eff}/\omega$	ν_{en} 10^7 s^{-1}	$\nu_{\rm eff}$ 10^7 s^{-1}	T_g K
0.3	100	2.7	10	1.06	0.15	5.6	520
1.0	100	3.9	5.8	1.04	0.46	3.8	480
10	100	10	3.6	1.19	4.0	3.9	400
100	50	33	1.7	1.83	10	16	380
300	50	100	1.3	2.23	17	27	420

that the measured power flux $S_{\rm exp}(z)$ is significantly larger than the collisional one.

A similar divergence between $S_{\rm exp}(z)$ and $S_{\rm col}(z)$ was observed at 3.39 and 13.56 MHz for different discharge powers. The results of measurements for a wide range of argon pressure are summarized in Fig. 2. It is seen that the ratio $S_{\rm exp}/S_{\rm col}$ at $z = L$ is close to unity at pressures above 10 mTorr and can exceed an order of magnitude at the lowest gas pressure.

In distinguishing between the total and collisional parts of the absorbed power it is important to correctly evaluate the collisional part, which is defined by electron-atom transport collision frequency in the rf field ν_{en}. Two points must be taken into account. First, in a high density plasma typical for materials processing and lighting applications, the gas temperature in the discharge chamber T_g can be considerably higher than the ambient temperature T_0. This leads to a reduction of gas density in a discharge with a controlled gas pressure p, that should be taken into account in the evaluation of ν_{en}. Analysis of gas heating performed for the conditions of our experiment has shown that the major mechanism of gas heating at low pressure ($p < 50$ mTorr) is due to ion energy transfer in charge exchange collisions with atoms; at higher pressure gas heating due to electron collisions prevails. Gas temperatures in the skin layer calculated with measured T_e and $n_e(z)$ are given in Table I.

Second, due to the strong Ramsauer effect in argon gas, ν_{en} depends on both the rf frequency ω and the EEDF form. The latter is always non-Maxwellian in gas discharge plasmas. Therefore, the plasma conductivity and absorbed power evaluated with ν_{en} found in textbooks [where it is usually calculated for the limiting cases of dc ($\omega \ll \nu_{en}$) or microwave ($\omega \gg \nu_{en}$) fields assuming a Maxwellian EEDF] can differ significantly from the actual values [15]. For calculation of the collisional power absorption we used the conductivity of a cold plasma in the form:

$$\sigma = \frac{e^2 n_e}{m(\nu_{en} + j\omega_{\rm eff})}, \qquad (2)$$

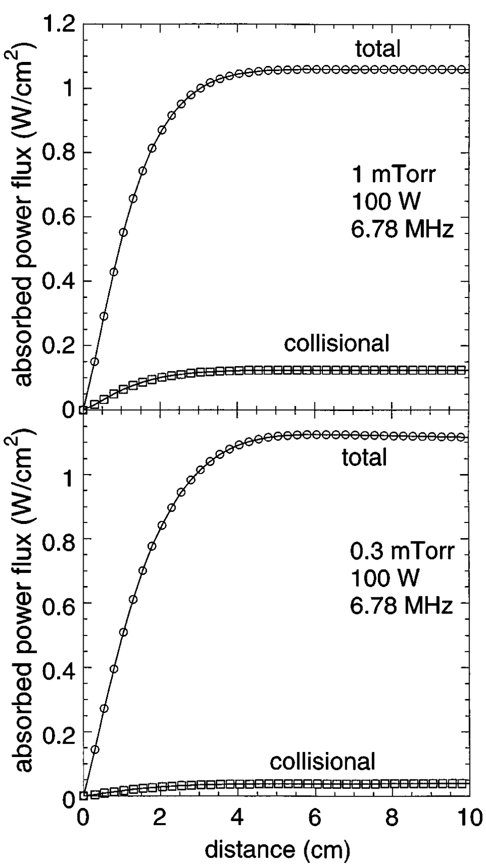

FIG. 1. Absorbed power flux for 0.3 and 1.0 mTorr.

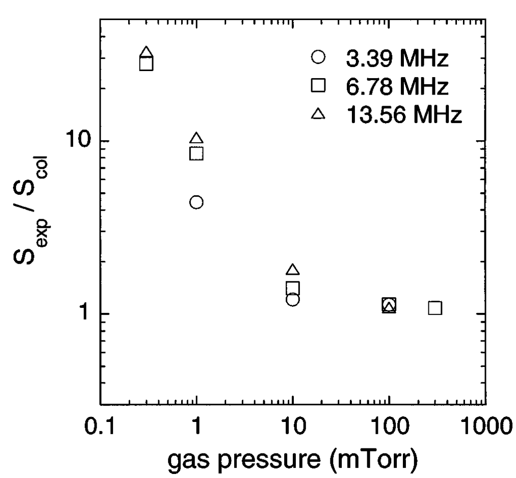

FIG. 2. The ratio of the total measured to collisional power flux as a function of argon pressure.

where ν_{en} and $\omega_{\rm eff}$ were found according to kinetic theory (see Ref. [15]):

$$\sigma = -\frac{2e^2 n_e}{3m} \int_0^\infty \frac{\varepsilon^{3/2}}{\nu_c(\varepsilon) + j\omega} \frac{df(\varepsilon)}{d\varepsilon}\, d\varepsilon\,. \quad (3)$$

Here e is the electron charge, m is the electron mass, $\nu_c(\varepsilon)$ is the transport electron collision frequency, $f(\varepsilon)$ is the EEDF, and ε is the electron energy. The values of ν_{en} and $\omega_{\rm eff}$ calculated with the EEDFs measured in the middle of the skin layer ($z = 1$ cm and $r = 4$ cm) are given in Table I for $\omega/2\pi = 6.78$ MHz together with other relevant parameters.

There are several ways to evaluate the collisional power flux $S_{\rm col}$. The first one follows from the theory of the normal skin effect (see, for example, Ref. [2]):

$$S_{\rm col}(L) = \frac{e^2 E^2(0) n_e \delta_n \nu_{en}}{2m(\nu_{en}^2 + \omega_{\rm eff}^2)}\,. \quad (4)$$

Here $E(0)$ is the rms rf field at the plasma boundary ($z = 0$), and δ_n is the normal skin depth. This formula assumes a uniform plasma and an exponential profile of the rf field and current density; neither of these assumptions is valid in this experiment [6]. Another way to evaluate $S_{\rm col}$ consists of using the measured distributions of the rf electric field or current density:

$$S_{\rm col}(L) = \int_0^L {\rm Re}(\sigma) E^2(x)\, dx = \int_0^L \frac{e^2 E^2(x) n_e(x) \nu_{en}}{m(\nu_{en}^2 + \omega_{\rm eff}^2)}\, dx \quad (5)$$

and

$$S_{\rm col}(L) = \int_0^L {\rm Re}(\sigma^{-1}) J^2(x)\, dx = \int_0^L \frac{J^2(x) m \nu_{en}}{n_e(x) e^2}\, dx\,. \quad (6)$$

Finally, $S_{\rm col}$ can be found with $E(z)$ and $J(z)$ measured in the experiment, using the collisional power factor $\cos\psi_{\rm col} = (1 + \omega_{\rm eff}^2/\nu_{en}^2)^{-1/2}$ for the conductivity of cold plasma:

$$S_{\rm col}(L) = \int_0^L E(x) J(x) \cos\psi_{\rm col}\, dx\,. \quad (7)$$

All of the last three expressions for $S_{\rm col}$ use the experimental distributions of $E(z)$ and $J(z)$. The last one, which has the advantage of being independent of $n_e(z)$ (thus, having less error), has been used for the calculation of $S_{\rm col}$ shown in Figs. 1. All four expressions for $S_{\rm col}$ give results which agree with each other within 10%–30%.

The collisionless power absorption is frequently accounted for by introducing an effective collision frequency $\nu_{\rm eff} > \nu_{en}$ in the classical expression for the plasma conductivity. In this way, one obtains $P = {\rm Re}(\sigma) E^2$ with $\sigma = e^2 n_e / m(\nu_{\rm eff} + j\omega_{\rm eff})$. Here $\omega_{\rm eff}$ is a function of ω and ν_{en}, and $\omega_{\rm eff} \approx \omega$ in the limit $\nu_{\rm eff} \le \omega$, similar to that in a collisional plasma [15]. The effective frequency $\nu_{\rm eff}$ accounts for total heating, collisionless as well as collisional. The local value of $\nu_{\rm eff}$ can be found using $E(z)$, $J(z)$, and $\psi(z)$ measured in the experiment: $\nu_{\rm eff} = (e^2 n_e E \cos\psi)/mJ$ [6]. The values of $\nu_{\rm eff}$ at $z = 1$ cm and $r = 4$ cm found this way are given in Table I. Comparing the magnitudes of $\nu_{\rm eff}$ and ν_{en} one can see that for $p < 10$ mTorr, $\nu_{\rm eff} > \nu_{en}$, which well correlates with the pressure dependence of the ratio $S_{\rm exp}/S_{\rm col}$ found in this experiment.

The calculated ratio $\eta = S_{\rm tot}(L)/S_{\rm col}(L)$ is shown in Fig. 3 as a function of the rf frequency for 0.3 and 1 mTorr. The total power density $P_{\rm tot} \propto {\rm Re}(JE^*)$ and the collisional power density $P_{\rm col} \propto EJ(1 + \omega^2/\nu_{en}^2)^{-1/2}$ used for calculations of $S_{\rm tot}$ and $S_{\rm col}$ are obtained from the solution of a coupled set of Maxwell and Boltzmann equations by the Fourier method [7]. This solution rigorously accounts for effects of thermal electron motion in the axial direction and the resulting spatial dispersion of the plasma conductivity. The spatial distributions of $E(z)$ and $J(z)$ are found neglecting thermal electron motion in the planes orthogonal to the axis, assuming a spatially homogeneous plasma with a Maxwellian EEDF, and an energy independent collision frequency ν_c. It is shown in [7] that in spite of these assumptions, the calculated profiles of $E(z)$, $J(z)$, and $P(z)$ are close to the experiment. In Fig. 3, the calculated values of η are compared to experimental data. Plasma parameters used for the calculations are taken from Table I. Reasonable agreement of theory and experiment is seen. The 2 times lower η obtained in the calculations can be due to neglect of radial electron motion and use of an energy independent collision frequency in the model. The strong energy dependence of $\nu(\varepsilon)$ in argon due to Ramsauer effect might be the main reason for the discrepancies. It is worth noting that due to the integral character of the absorbed power flux $S(L)$, our calculations based on a

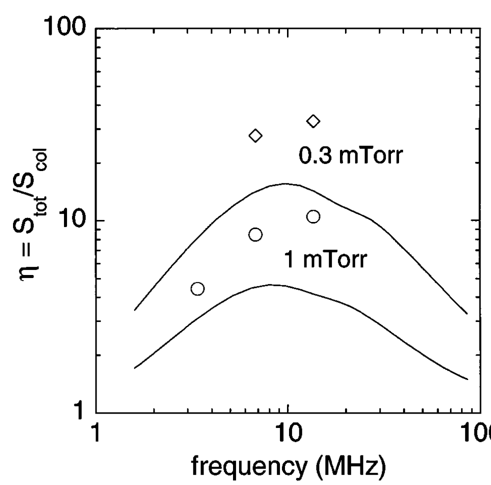

FIG. 3. The ratio of the total to collisional power flux as a function of frequency. The experiment is presented by symbols and the theory is presented by lines.

simple model of the collisionless heating [2] also give results close to those shown in Fig. 3.

As seen in Fig. 3 there is a maximum of $\eta(\omega)$ with respect to frequency; a similar effect has been also observed in calculations [2] and [3]. The appearance of this maximum could be attributed to resonant electron interaction with the rf field, which takes place when the half period of the field is close to the electron transit time through the skin layer [16]. The resonance condition, $\omega = v_T/\delta$, where v_T is the electron thermal velocity and $\delta \approx c/\omega_p$ is the skin depth, corresponds to the boundary of nonlocality $\Lambda = (\omega_p v_T/\omega c)^2 \approx 1$. Here ω_p is the electron plasma frequency and c is the speed of light. This boundary separates the domain of local electrodynamics ($\Lambda \ll 1$) from the domain of nonlocal electrodynamics ($\Lambda \gg 1$). The former corresponds to normal skin effect and the absence of spatial dispersion, and the latter corresponds to the anomalous skin effect with strong spatial dispersion and nonlocality of the rf current. Our calculations of the total power absorption as well as the calculation in Ref. [2] do not show a maximum in frequency dependence of $S_{\mathrm{tot}}(L)$. The maximum of $\eta(\omega)$ appears due to the different frequency dependencies of the collisionless ($\nu_{en} = 0$), S_{st}, and collisional, S_{col}, power absorption. For $\nu_{en} \ll \omega$, $S_{\mathrm{col}} \propto \omega^{-2}$ whereas S_{st} is almost frequency independent at $\omega \ll v_T/\delta$ and decays rapidly $S_{\mathrm{st}} \propto \omega^{-4}$ at $\omega \gg v_T/\delta$ [2].

For moderate nonlocality ($\Lambda \geq 1$) typical to experiments at 0.3 and 1 mTorr, $\eta(\omega)$ increases with frequency whereas the nonlocality parameter Λ decreases with frequency. This counterintuitive observation suggests that a larger degree of nonlocality does not necessarily mean a larger collisionless power absorption, although the latter is originated by the spatial dispersion of conductivity (nonlocality of electron current). In fact, η increases with Λ only for $\omega > v_T/\delta$. It is interesting to note that collisionless power absorption for evanescent low frequency waves (which are typical for ICP) was first demonstrated theoretically for the condition of the normal skin effect ($\omega \gg v_T/\delta$) where nonlocal effects are small [8].

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