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Quasilinear collisionless electron heating in a weakly magnetized inductively coupled plasma

F. SOHBATZADEH[1,2] and H. LATIFI[1,2]

[1]Laser Research Institute, Shahid Beheshti University, Evin, Tehran, Iran
[2]Department of Physics, Shahid Beheshti University, Evin, Tehran, Iran
(f-sohbatzadeh@cc.sbu.ac.ir)

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Abstract. Collisionless electron heating in a low-pressure weakly magnetized inductively coupled plasma is considered by means of quasilinear theory. We use a one-dimensional slab model of an inductively coupled plasma with a penetrating radio frequency (RF) electric field given by an exponential function. Diffusion coefficients in velocity space are obtained regarding the induced RF electric field and induced polarization electric field (DC electric field). It is shown that the induced polarization electric field has a significant contribution to the collisionless electron heating. We consider the influence of the quasilinear collision integral on the evolution of the symmetric part of the electron energy distribution function.

1. Introduction

Low-pressure inductively coupled plasmas (ICPs) are used in various fields of technology from material processing to fusion [1, 2]. In such plasmas, the electron–neutral collision frequency ν is much smaller than the radio frequency (RF) frequency ω and the electron mean free path l is comparable or larger than the plasma dimensions. Under such conditions, the mechanism of electron heating is collisionless (Landau damping) rather than the collisional Joule heating which is dominant at high pressures. Collisionless electron heating is essential in sustaining such plasmas and has attracted a great deal of experimental and theoretical investigations [3–6].

The application of a weak and steady-state magnetic field on an inductively coupled plasma can effect the plasma parameters [7–9], and it gives rise to a polarization electric field (DC field) and an enhanced penetration of the electromagnetic field (skin depth) in the plasma [10,11]. The skin depth increases for the $k_\parallel$-spectrum of the RF field, where $k_\parallel$ is the parallel component of the wavevector $\mathbf{k}$ with respect to the applied magnetic field $\mathbf{B}_0$. The induced polarization electric field (IPEF) and the skin depth could effect plasma heating and should be taken into account in the investigation.

In this paper, we focus our attention on the electron heating in a weakly magnetized low-pressure ICP such that $\omega_c \ll \omega_p$, where ω_c is the electron cyclotron frequency and ω_p is the electron plasma frequency. We consider the electron velocity diffusion coefficient using quasilinear theory. We consider two infinite parallel plane

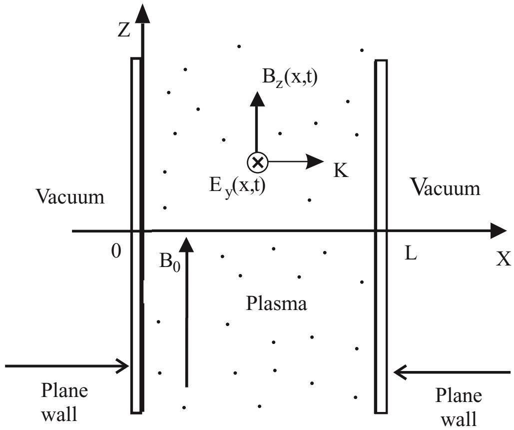

Figure 1. Schematic picture of a magnetized ICP in slab geometry.

walls located at $x = 0$ and $x = L$. The plasma is confined between the walls and we assume that the plasma density and electromagnetic waves may only vary in the x-direction, i.e. $\nabla_z = \nabla_y = 0$. A RF coil for generating an electromagnetic field is assumed to exist to the left of the $x = 0$ boundary. We assume that a transverse electromagnetic wave penetrates into the plasma and has the form $\mathbf{E} = E_y(x,t)\mathbf{e}_y$ and $\mathbf{B} = B_z(x,t)\mathbf{e}_z$, where $\mathbf{E}$ and $\mathbf{B}$ are related to each other via Maxwell equations. The weak external DC magnetic field B_0 is applied parallel to the walls in the positive z-direction. Figure 1 illustrates a schematic picture of a simplified one-dimensional magnetized ICP in slab geometry.

It has already been shown that the weak external DC magnetic field does not change the penetration of the RF field into the one-dimensional magnetized plasma, because of the mutual compensation of Pederson and Hall effects. However, the polarization field will be generated and must be taken into account in the investigation [10, 11].

To facilitate further calculations, we do not solve the self-consistent equation for the electric field $\mathbf{E}(x,t)$ and instead assume an exponential decaying spatial dependence in the form of $\mathbf{E}(x,t) = E_0 \exp(-x/\delta) \exp[i(\kappa x - \omega t)]\mathbf{e}_y$. In this form of $\mathbf{E}(x,t)$, $\delta = (c^2 m_e v_{th}/4\pi e^2 n_0 \omega)^{1/3}$ is the collisionless expression of skin depth, E_0 is a real parameter, v_{th} is electron thermal velocity and κ is the wave number, which in the frequency range 100 KHz–10 MHz is much less than $1/\delta$ [5, 12]. In addition, we neglect nonlinear effects (such as the nonlinear induced polarization field) related to the induced RF magnetic field which have been discussed in [10, 13].

This paper is organized as follows. In Sec. 2, we obtain a qualitative expression for the IPEF in a magnetoactive ICP in slab geometry. In Sec. 3, we discuss the quasilinear theory of collisionless electron heating in a magnetoactive plasma. In this section, we derive the diffusion coefficient in velocity space for a boundless magnetoactive plasma. In Sec. 4, we apply the results of Sec. 2 to a bounded magnetoactive ICP, and we compare the contribution of the IPEF to the collisionless electron heating with that of the induced RF electric field. In Sec. 5, we consider quasilinear collision integral influence on the evolution of the symmetric part of the electron energy distribution function (EEDF). Finally, Sec. 6 gives a summary and a discussion of our results.

2. IPEF

In this section, we obtain a qualitative expression for the IPEF in a magnetoactive ICP in slab geometry. The IPEF is closely related to the Hall effect. We consider the linear electron equation:

$$\frac{\partial \mathbf{V}}{\partial t} = -\frac{e}{m}\left[\mathbf{E} + \frac{1}{C}\mathbf{V}\times\mathbf{B}\right] - \nu\mathbf{V}, \tag{1}$$

where ν is the electron–neutral collision frequency. The induced RF magnetic field leads to the nonlinear polarization field in an ICP, therefore it is neglected in the linear electron equation. One can solve (1) for $\mathbf{V}$ to obtain

$$V_x = \left(\frac{-ie}{m\hat{\omega}}\right)\frac{E_x(x) - i(\omega_c/\hat{\omega})E_y(x)}{1-(\omega_c/\hat{\omega})^2}, \tag{2}$$

$$V_y = \left(\frac{-ie}{m\hat{\omega}}\right)\frac{E_y(x) + i(\omega_c/\hat{\omega})E_x(x)}{1-(\omega_c/\hat{\omega})^2}, \tag{3}$$

where $\hat{\omega} = \omega + i\nu$ and $\omega_c = eB_0/mc > 0$ is the electron cyclotron frequency. Next, making use of Maxwell equations, one obtains

$$\nabla(\nabla\cdot\mathbf{E}) - \nabla^2\mathbf{E} = -\frac{4\pi}{c^2}\frac{\partial\mathbf{J}}{\partial t}, \tag{4}$$

where the displacement current is neglected. This assumption is valid in a typical ICP. It is easy to see that the x-component of the left-hand side of this equation vanishes. Therefore, $\mathbf{J}$ can have only y-components. One can now find the polarization field from (2) by setting $J_x = 0$ and obtaining

$$E_x(x) = i\frac{\omega_c}{\hat{\omega}}E_y(x). \tag{5}$$

This is a DC electric field which varies in space on the characteristic length scale of the skin depth δ. By inserting (5) into (3), it is seen that J_y does not change in the one-dimensional magnetized ICP. This means that the DC external magnetic field does not change the skin depth.

3. Quasilinear theory of collisionless electron heating in a magnetoactive plasma

We solve kinetic equation for electrons and obtain electron velocity diffusion coefficient in a magnetized low pressure ICP. At first we consider a boundless plasma, then we extend the model to a bounded ICP. The oscillatory motion of ions in the presence of RF waves can be ignored due to their relatively small mobility. The EEDF $F(\mathbf{r},\mathbf{v},t)$ in an arbitrary electric field $\mathbf{E}(\mathbf{r},t)$ and magnetic field $\mathbf{B}(\mathbf{r},t)$ is

$$\frac{\partial F(\mathbf{r},\mathbf{v},t)}{\partial t} + \mathbf{v}\cdot\nabla_r F(\mathbf{r},\mathbf{v},t) - \frac{e}{m}\left[\mathbf{E} + \frac{\mathbf{v}\times\mathbf{B}}{c}\right]\cdot\nabla_v F(\mathbf{r},\mathbf{v},t) = S(F), \tag{6}$$

where $S(F)$ is the collision integral. In the low-pressure plasma, the electron heating characteristic length scale δ becomes smaller than the electron MFP (Mean free path) l; thus, one can decompose the electric field $\mathbf{E}(\mathbf{r},t)$ and magnetic field $\mathbf{B}(\mathbf{r},t)$ into two parts:

$$\mathbf{E}(\mathbf{r},t) = \bar{\mathbf{E}}(\mathbf{r},t) + \tilde{\mathbf{E}}(\mathbf{r},t), \quad \mathbf{B}(\mathbf{r},t) = \bar{\mathbf{B}}(\mathbf{r},t) + \tilde{\mathbf{B}}(\mathbf{r},t), \tag{7}$$

where the characteristic length scale of $\tilde{\mathbf{E}}(\mathbf{r},t)$ and $\tilde{\mathbf{B}}(\mathbf{r},t)$ is the skin depth δ and that of $\bar{\mathbf{E}}$ and $\bar{\mathbf{B}}$ is electron collision MFP l. To estimate the numerical values of δ and l, a typical ICP of helium gas at a pressure of 1 mTorr is considered. With $n_e = 10^{11}$ cm^{-3} and $T_e = 5$ eV, one obtains $\delta = 3.16$ cm and $l = 50.0$ cm [14]. We can also separate the distribution function F into two parts. The first part is $F_0(\mathbf{r}, \mathbf{v}, t) = \langle F(\mathbf{r}, \mathbf{v}, t)\rangle_l$, averaged over the scale of the electron collision MFP l. The second part is $\tilde{F}(\mathbf{r}, \mathbf{v}, t)$, which describes deviations of the EEDF F from the first part with space scale smaller than MFP l, and $\tilde{F} \ll F$,

$$F(\mathbf{r}, \mathbf{v}, t) = F_0(\mathbf{r}, \mathbf{v}, t) + \tilde{F}(\mathbf{r}, \mathbf{v}, t). \tag{8}$$

We can then separate (6) into slow- and fast-varying parts in the form of quasilinear equations [15]:

$$\frac{\partial F_0}{\partial t} + \mathbf{v} \cdot \nabla_r F_0 - \frac{e}{m}\left(\bar{\mathbf{E}} + \frac{\mathbf{v} \times \bar{\mathbf{B}}}{c}\right) \cdot \nabla_v F_0 - \frac{e}{m}\left\langle \tilde{\mathbf{E}} + \frac{\mathbf{v} \times \tilde{\mathbf{B}}}{c} \cdot \nabla_v \tilde{F}\right\rangle_\delta = S(F_0), \tag{9}$$

$$\frac{\partial \tilde{F}}{\partial t} + \mathbf{v} \cdot \nabla_r \tilde{F} - \frac{e}{m}\left(\bar{\mathbf{E}} + \frac{\mathbf{v} \times \bar{\mathbf{B}}}{c}\right) \cdot \nabla_v \tilde{F} - \frac{e}{m}\left(\tilde{\mathbf{E}} + \frac{\mathbf{v} \times \tilde{\mathbf{B}}}{c} \cdot \nabla_v F_0\right) = S(\tilde{F}), \tag{10}$$

where the last term in the left-hand side of (9) is the nonlinear wave particle interaction term which is averaged over the plasma heating characteristic length scale of δ. It expresses collisionless plasma heating in the quasilinear theory.

The equation for the oscillating part of the electron distribution function $\tilde{F}$ can be written as

$$\frac{\partial \tilde{F}}{\partial t} + \mathbf{v} \cdot \nabla_r \tilde{F} - \frac{e}{mc}\mathbf{v} \times \mathbf{B}_0 \cdot \nabla_v \tilde{F} = \frac{e}{m}\left(\tilde{\mathbf{E}} + \frac{\mathbf{v} \times \tilde{\mathbf{B}}}{c}\right) \cdot \nabla_v F_0 - \nu \tilde{F}, \tag{11}$$

where ν is the electron–neutral collision frequency, and the collision integral $S(\tilde{F})$ was approximated by $-\nu\tilde{F}$. The solution of (11) can be found using the integral along the unperturbed orbit of electrons [16, 17],

$$\tilde{F}(x, v, t) = \frac{e}{m}\int_{-\infty}^{t}\left[\tilde{\mathbf{E}}(x', t') + \frac{\mathbf{v} \times \tilde{\mathbf{B}}(x', t')}{c}\right] \cdot \nabla_{v'} F_0 \exp[-\nu(t - t')]\, dt', \tag{12}$$

where $\tilde{\mathbf{E}}(x', t')$ is the sum of the induced RF electric field $\tilde{\mathbf{E}}_y(x', t')$ and generated polarization field $\tilde{\mathbf{E}}_x(x')$, $x' = x - (v_\perp/\omega_c)\{\sin\phi - \sin[\phi - \omega_c(t - t')]\}$ is the electron trajectory in the wave and in the presence of a static external magnetic field $\mathbf{B}_0$, and $v_\perp$ is the perpendicular component of the electron velocity with respect to $\mathbf{B}_0$. In order to evaluate (12), it is convenient to use the Fourier transform of the electric fields $\tilde{\mathbf{E}}(x') = \int_{-\infty}^{+\infty} \mathbf{E}(k)\exp(ikx')\, dk$, one can then obtain

$$\begin{aligned}\tilde{F}(x, v, t) = \frac{e}{m}\int_{-\infty}^{\infty} dk \sum_{n,n'=-\infty}^{\infty} \\ \times \left\{\frac{E_y(k)J_{n'}(\xi)J'_n(\xi)}{\hat{\omega} - n\omega_c}\exp(-i\omega t) + \frac{inJ_{n'}(\xi)J_n(\xi)E_x(k)}{\xi(i\nu - n\omega_c)}\right\} \\ \times \frac{\partial F_0}{\partial v_\perp}\exp[i(n - n')\phi]\exp(ikx),\end{aligned} \tag{13}$$

where $\xi = kv_\perp/\omega_c$, $J_n(\xi)$ is the nth order of the Bessel function, $J'_n(\xi)$ is a derivative of $J_n(\xi)$ with respect to ξ, and $E_y(k)$ and $E_x(k)$ are the Fourier components of the induced RF electric field and the IPEF, respectively.

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We now obtain the quasilinear collision integral and diffusion coefficient. In the low-pressure ICP, the RF frequency ω is larger than the inelastic electron collision frequency which permits us to assume that F_0 does not depend on time. By the temporal averaging of (9) over the wave period T, we get

$$\mathbf{v} \cdot \nabla_r F_0 = S(F_0) + S_{\mathrm{QL}}(F_0), \tag{14}$$

where

$$S_{\mathrm{QL}} = \frac{e}{m}\left\langle \left(\tilde{\mathbf{E}} + \mathbf{v} \times \frac{\tilde{\mathbf{B}}}{c}\right) \cdot \nabla_v \tilde{F} \right\rangle_{\delta,T}, \tag{15}$$

is the quasilinear collision integral, and the brackets denote averaging over the wave period T and space scale δ, such that $\delta \ll l$. Making the averages, it gives

$$S_{\mathrm{QL}} = \frac{e}{2m}\delta(x-x_0)\,\mathrm{Re}\left\{\nabla_v \cdot \left[\left(\tilde{\mathbf{E}}^*(k) + \mathbf{v} \times \frac{\tilde{\mathbf{B}}^*(k)}{c}\right)\tilde{F}\right]\right\}, \tag{16}$$

where the delta function indicates the localization of collisionless electron heating in the range $x \approx x_0$. Averaging (16) over the gyroangle ϕ in velocity space and using (13), one can arrange the quasilinear collision integral in the form

$$S_{\mathrm{QL}} = \delta(x-x_0)\frac{1}{v_\perp}\frac{\partial}{\partial v_\perp}v_\perp\left(D_{\perp\perp}\frac{\partial F_0}{\partial v_\perp}\right), \tag{17}$$

with diffusion coefficient

$$D_{\perp\perp} = \pi\left(\frac{e}{m}\right)^2 \sum_{n=-\infty}^{\infty}\int_{-\infty}^{\infty} dk\left(d_n^{E_y} + d_n^{E_x}\right), \tag{18}$$

with the terms

$$d_n^{E_x} = \frac{\nu|E_x(k)|^2 n^2 J_n^2(\xi)}{\xi^2\left(n^2\omega_{\mathrm{c}}^2 + \nu^2\right)}, \tag{19}$$

$$d_n^{E_y} = \frac{\nu|E_y(k)|^2 J_n'^2(\xi)}{(\omega - n\omega_{\mathrm{c}})^2 + \nu^2}, \tag{20}$$

where (19) and (20) express diffusion coefficient components related to the IPEF $\mathbf{E}_x(x)$ and the induced RF electric field $\mathbf{E}_y(x)$, respectively. In Fig. 2 the variation of $d_1^{E_x}$ and $d_1^{E_y}$ are shown as functions of $\xi = kv_\perp/\omega_{\mathrm{c}}$. $d_1^{E_x}$ and $d_1^{E_y}$ represent electron heating in collisionless power absorption. Although the collision frequency is small ($\nu \ll \omega$), collisions are necessary for collisionless power dissipation. As can be seen from (19) and (20), if $\nu \to 0$, collisionless heating also tends to zero.

4. Diffusion coefficients in a bounded magnetoactive ICP in slab geometry

To use (19) and (20) in a bounded magnetoactive ICP, it is necessary to set the boundary conditions for electromagnetic waves and EEDF $\tilde{F}$ on the plasma surfaces. It can be assumed that charged particles undergo mirror reflection from the plasma surfaces. This implies that (11) can be extended in the complete x range ($-\infty < x < +\infty$) if the following relations are taken into account:

$$\tilde{F}(v_x, v_y, v_z, x<0) = \tilde{F}(-v_x, v_y, v_z, x>0), \tag{21}$$

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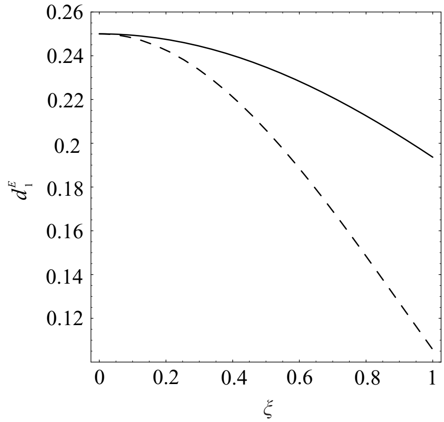

Figure 2. Variation of $d_1^{E_x}$ (full curve) and $d_1^{E_y}$ (dashed curve) in arbitrary units as a function of $\xi = kv_\perp/\omega_c$.

for EEDF $\tilde{F}$ and

$$E_y(x<0) = E_y(x>0), \quad E_x(x<0) = -E_x(x>0), \tag{22}$$

$$B_y(x<0) = -B_y(x>0), \tag{23}$$

for electromagnetic waves on $x=0$. The problem now is reducing to a boundless plasma. We now use the Fourier transform of a periodically continued electric field with period L as

$$E_y(k) = \frac{1}{L}\int_0^L E_y(x)\cos(kx)\,dx, \tag{24}$$

and

$$E_x(k) = \frac{1}{L}\int_0^L E_x(x)\sin(kx)\,dx, \tag{25}$$

where $k = n\pi/L$, $E_y(x) = E_0\exp(-x/\delta)\cos(\kappa x)$ and $E_x(x) = (-\omega_c\omega E_0)/(\omega^2+\nu^2)$ $\exp(-x/\delta)\sin(\kappa x)$. Figures 3 and 4 show the profile of the normalized RF electric field and the IPEF as a function of x/δ, respectively.

To demonstrate the polarization field effect on the electron heating, we consider the ratio of $d_n^{E_x}$ to $d_n^{E_y}$ in the limit $\nu \ll \omega < \omega_c$. To illustrate the numerical values of collision frequency, electron cyclotron frequency, electron plasma frequency and electron gyroradius in an ICP, let us assume that $B_0 = 10$ G, $n_e = 10^{11}$ cm^{-3}, $T_e = 5$ eV and $\omega/2\pi = 13.56$ MHz. One then obtains $\nu = 2\times10^6$ s^{-1}, $\omega_c = 1.76\times10^8$ s^{-1}, $\omega_p = 1.79\times10^{10}$ s^{-1} and $r_c = 0.85$ cm [14]. Since $J_n(\xi)/J_n'(\xi) \approx 1$ and $|E_x(k)|^2/|E_y(k)|^2 \approx (\omega_c/\omega)^2$, one obtains

$$\frac{d_n^{E_x}}{d_n^{E_y}} \approx \begin{cases} \left(\dfrac{L\omega_c}{\pi r_c\omega}\right)^2 & n \neq 0, \\ \left(\dfrac{L\omega_c}{\pi r_c\nu}\right)^2 & n = 0, \end{cases} \tag{26}$$

where $r_c = v_\perp/\omega_c$ is the electron gyroradius and L is gap length of the plasma. Hence, the gap length L does not exceed the energy relaxation length λ_ϵ, i.e. $r_c < L < \lambda_\epsilon$, and one gets $d_n^{E_x}/d_n^{E_y} \gg 1$. For $\nu \ll \omega < \omega_c$, it can be concluded that

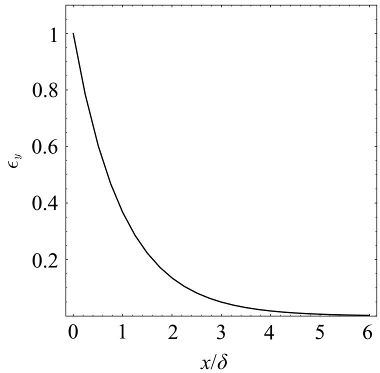

Figure 3. Evolution of the normalized RF electric field $\epsilon_y = E_y(x)/E_0$ as a function of x/δ for $\delta = 3.16$ cm and $\kappa = 0.01$ cm^{-1}.

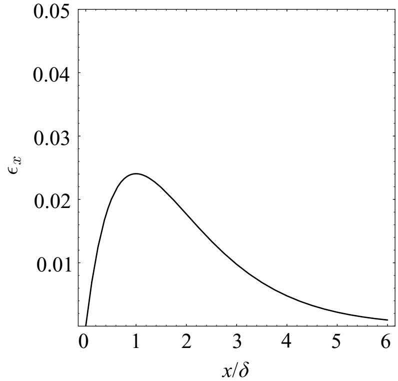

Figure 4. Evolution of the normalized polarization electric field $\epsilon_x = E_x(x)/E_0$ as a function of x/δ for $\delta = 3.16$ cm, $\kappa = 0.01$ cm^{-1}, $\omega/2\pi = 13.56$ MHz, $\omega_c = 1.76 \times 10^8$ s^{-1} and $\nu = 2 \times 10^6$ s^{-1}.

the IPEF has a significant contribution to the plasma heating in a one-dimensional magnetized ICP. Thus, the application of a weak steady-state magnetic field ($\omega_c \ll \omega_p$) on an ICP which leads to a DC electric field (polarization field) also introduces additional collisionless electron heating. The corresponding diffusion coefficient is larger than that of the penetrating RF electric field in the investigated model.

5. Symmetric distribution function F_0

In this section we consider the effect of the quasilinear collision integral on the slowly varying EEDF F_0 using quasilinear theory. It is assumed that the EEDF F is essentially isotropic, which requires that $\nu \gg \nu^*$, where ν^* is the inelastic collision frequency. In this case, the EEDF F can be expanded into spherical harmonics (Lorentz approximation) and terms beyond first order may be neglected,

$$F(\mathbf{r}, \mathbf{v}, t) = F_0(\mathbf{r}, \mathbf{v}, t) + \frac{\mathbf{v}}{v} \cdot \mathbf{F}_1(\mathbf{r}, \mathbf{v}, t), \tag{27}$$

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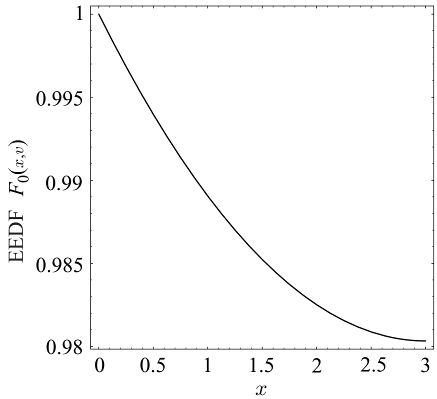

Figure 5. Normalized electron energy distribution function $F_0(x,v)$ as a function of x at a magnetic field of 10 G.

where F_0 is the symmetric part of the EEDF F and the second term on the right-hand side of (27) is the anisotropy part, with $F_1 \ll F_0$. One can obtain an equation for symmetric part of the EEDF F as follows [18]:

$$-\frac{v^2\nu}{3\left(\nu^2+\omega_c^2\right)}\nabla_\perp^2 F_0 - \frac{v^2}{3\nu}\nabla_\parallel^2 F_0 = S_{\mathrm{QL}}(F_0) + S_0^*, \tag{28}$$

where $\nabla_\parallel$ and $\nabla_\perp$ are parallel and perpendicular components of gradient with respect to $\mathbf{B}_0$, S_0^* is the inelastic collision integral, and $S_{\mathrm{QL}}(F_0)$ is defined by (17). For the case of the infinite dimension of two parallel plasma walls as indicated in Sec. 1, one can suppose that only perpendicular gradients survive (28), and $\nabla_\perp^2$ can be replaced by $\partial^2/\partial x^2$. Therefore, (28) with an assumption of $S_0^* = -\nu^* F_0$ becomes

$$\frac{v^2\nu}{3\left(\nu^2+\omega_c^2\right)}\frac{\partial^2 F_0}{\partial x^2} - \nu^* F_0 = -S_{\mathrm{QL}}(F_0). \tag{29}$$

Next, integration of (29) over the interval $x = x_0 - \epsilon$ to $x = x_0 + \epsilon$, taking into account the plasma heating symmetry with respect to the position x_0, gives

$$\frac{2}{3}\frac{v^2\nu}{\nu^2+\omega_c^2}\frac{\partial F_0}{\partial x}\bigg|_{x=x_0+\epsilon} = -S_{\mathrm{QL}}(F_0), \tag{30}$$

where ϵ is very small. The above equation can be used as the boundary condition for

$$\frac{v^2\nu}{3\left(\nu^2+\omega_c^2\right)}\frac{\partial^2 F_0}{\partial x^2} - \nu^* F_0 = 0. \tag{31}$$

Using the boundary condition (30) for a planar slab plasma, the EEDF F_0 can be written as

$$F_0(v,x) = F_0(v)\frac{\cosh((x-L/2)/\lambda_\varepsilon)}{\cosh(L/2\lambda_\varepsilon)}, \tag{32}$$

where $\lambda_\varepsilon = \sqrt{v^2\nu/3(\nu^2+\omega_c^2)\nu^*}$ is the energy relaxation length. In Fig. 5 the variation of $F_0(v,x)/F_0(v)$ as a function of x at a magnetic field of 10 G is shown. It is seen that the energy relaxation length decreases due to the applied external magnetic field. Thus, the theory is valid for a weak magnetic field. The kinetic equation for the EEDF F_0 can be deduced from (30) taking $x_0 = 0$ and $\epsilon \to 0$,

$$F_0(v) = \frac{-1}{2\nu^* \lambda_\varepsilon} S_{\mathrm{QL}}(F_0) \coth\left(\frac{L}{2\lambda_\varepsilon}\right). \tag{33}$$

The effect of the nonlinear wave–particle interaction term is included in this kinetic equation via the quasilinear collision integral $S_{\mathrm{QL}}(F_0)$. Although the collisionless electron heating takes place at the boundaries, inelastic collisions transfer the heating into the bulk plasma since $\lambda_\varepsilon \gg L$. In this region, one can get the following continuity equation

$$\nu^* F_0 + \frac{S_{\mathrm{QL}}}{L} = 0, \tag{34}$$

which indicates that the collisionless electron heating is balanced by electron inelastic collisions. Furthermore, the isotropic EEDF F_0 is spatially homogeneous in this limit.

6. Summary and conclusions

We have investigated collisionless electron heating in a low-pressure weakly magnetized ICP via quasilinear theory. We have used a one-dimensional slab model of magnetized ICP with a prescribed penetrating electric field profile given by the exponential function. An expression for the induced polarization electric field due to the Hall effect was obtained qualitatively. The diffusion coefficients in velocity space were investigated in a bounded magnetized ICP. In the simplified model of magnetized ICP, It was shown that the IPEF leads to the additional electron heating which has a significant contribution to the collisionless electron heating compared with that of the induced RF electric field. The kinetic equation of the isotropic part of EEDF F has been obtained regarding the quasilinear collision integral.

References

[1] Ferreira, C. M. and Moisan, M. 1993 *Microwave Discharges: Fundamentals and Applications*. Nato ASI Series B **302**: Physics. New York: Plenum Press.

[2] Lieberman, M. A. and Lichtenberg, A. J. 1994 *Principles of Plasma Discharges and Materials Processing*. New York: Wiley.

[3] Seo, S.-H., Chung, C. W. and Chang, H.-Y. 2000 *Surf. Coating Technol.* **131**, 1.

[4] Aliev, Yu. M., Kaganovich, I. D. and Schlüter, H. 1997 *Phys. Plasmas* **4**, 2413.

[5] Tyshetskiy, Yu. O., Smolyakov, A. I. and Godyak, V. A. 2002 *Plasma Sources Sci. Technol.* **11**, 203.

[6] Kaganovich, I. D. 1999 *Phys. Rev. Lett.* **82**, 327.

[7] Chung, C. W., Kim, S. S. and Chang, H. Y. 2002 *Phys. Rev. Lett.* **88**, 095002.

[8] Lee, H. J., Yang, I. D. and Wang, K. W. 1996 *Plasma Sources Sci. Technol.* **5**, 383.

[9] Lee, H. J., Kim, J. H., Wang, K. W. and Joo, J. H. 1996 *J. Vac. Sci. Technol.* A **14**, 1007.

[10] Smolyakov, A., Godyak, V. and Duffy, A. 2000 *Phys. Plasmas* **7**, 4755.

[11] Chen, F. F. 2001 *Phys. Plasmas* **8**, 3008.

[12] Godyak, V. A. and Kolobov, V. A. 1997 *Phys. Rev. Lett.* **79**, 4589.

[13] Cohen, R. H. and Rognlien, T. D. 1996 *Plasma Sources Sci. Technol.* **5**, 442.

[14] Raizer, Y. P. 1997 *Gas Discharge Physics*. Springer, Berlin.

[15] Akhiezer, A. I., Akhiezer, I. A., Polovin, R. V., Sitenko, A. G. and Stepanov, K. N. 1975 *Plasma Electrodynamics, Volume 2, Non-linear Theory and Fluctuations*. New York: Pergamon Press.

[16] Brambilla, M. 1998 *Kinetic Theory of Plasma Waves*. Oxford: Clarendon Press.

[17] Stix, T. H. 1992 *Waves in Plasmas*. New York: American Institute of Physics/Springer.

[18] Shkarofski, I. P., Johnston, T. W. and Bachynski, M. P. 1966 *The Particle Kinetics of Plasma*. Reading, MA: Addison-Wesley.