Quasilinear damping of ion Bernstein waves on the harmonic resonant layer

A. Cardinalli, a) C. Castaldo, b) R. Cesario, and F. De Marco

Associazione EURATOM-ENEA sulla Fusione, Centro Ricerche Energia Frascati, 00044 Frascati, Rome, Italy

(Received 7 November 1997; accepted 11 May 1998)

The damping of an ion Bernstein wave (IBW) across a resonant layer of the minority species of the plasma is calculated within the framework of quasilinear theory. The model uses the analytical solution of the one-dimensional Fokker–Planck equation for the ion distribution function. A greater reduction of the damping has been found in this case compared with the one obtained with the linear model. This allows the wave to penetrate further towards the plasma core. In addition, a particle flux induced by the sharp spatial gradient of the distribution function across the resonant layer, due to the IBW field, has also been evaluated. This shows that modification of the distribution function generates a flux of particles from the outside which tends to increase the particle density in the vicinity of the resonance itself. © 1998 American Institute of Physics. [S1070-664X(98)02208-3]

I. INTRODUCTION

In past few years, the interest in the coupling of the ion Bernstein wave (IBW) power to a tokamak plasma has greatly increased especially since this heating and current-drive scheme can also provide a tool for active control of plasma transport by turbulence suppression.1,2

The coupling, propagation, and absorption of the IBWs in tokamak plasmas was originally studied by Ono 3 within the framework of linear theory. In this scheme, the IBW, mode converted from a lower hybrid wave (LHW) that was launched by a slow wave antenna, is damped on the plasma particles near the layer of an integer harmonic of the ion cyclotron frequency. In addition, Porkolab 4 studied the nonlinear absorption of the IBW (regulated by a power threshold mechanism) near the layer of a semi-integer ion cyclotron harmonic.

In a review paper, 5 Ono gave an extensive summary of the main characteristics of IBW absorption (linear and nonlinear) and of overall activities in the IBW experimental domain.

Evidence of quasilinear effects on the energy distribution function of the majority species plasma ions was observed during the IBW experiment at the Japanese Institute of Plasma Physics Tokamak II-Upgrade (JIPPT-II-U) 6 (130 MHz, 3 T). In fact, the calculated nonthermal component of the energy distribution function was found to be in reasonable agreement with the experimental measurements of the neutrals’ flux energy spectrum obtained using a fast particle analyzer. In that model, the calculation was performed in the framework of quasilinear theory taking into account only a high-velocity part of the spectrum of the plasma ions. In the present article, we focus our attention on the quasilinear diffusion induced by the IBW at the ion cyclotron resonant layers of a minority ion species, where the wave damping is expected to take place.

In this study, we are evaluating the steady state solution 6,7 of the one-dimensional (1D) Fokker–Planck equation (outlined in Sec. II) in the whole velocity spectrum of a minority ion species. We note that the theory presented here cannot be considered as a generalization of the model used in Ref. 6. In fact, we consider isotropic distribution functions and we evaluate the Coulomb diffusion operator in a test-particle approximation. The above assumptions are no longer valid in the case of a majority ion heating scheme considered in Ref. 6. Our calculations refer to the Alcator-C (183.6 MHz, 9.3 T) experiment 8,9 and to the Frascati Tokamak Upgrade (FTU) (433 MHz, 7.8 T) experiment that is planned.10 In both cases, we have observed very strong diffusion in the thermal range where the propagating IBW field flattens the distribution function. Moreover, in Sec. III, we have calculated the IBW power quasilinear damping showing that the IBW power damping across the resonance layer $\omega = 3\Omega_D$ and $\omega = 9\Omega_D$ (relevant for the above IBW experiments) is much weaker than the damping predicted by linear theory.

In addition, the IBW produces a sharp transition in the shape of the ion distribution function between the resonance layer and the inner Maxwellian plasma. This suggests that the ions can diffuse across the magnetic surfaces like in classical transport, where temperature and pressure gradients induce radial particle fluxes.11 We calculate this diffusive flux in Sec. IV and evaluate the consequent modification of the resonant ions’ density profile. In Sec. V, we summarize and discuss the effects of the quasilinear damping on wave absorption and on the fluxes induced.

II. VELOCITY DISTRIBUTION DURING IBW HEATING

In Sec. II, we outline the steady state solution of the 1D Fokker–Planck equation, describing the interaction of an IBW with a minority ion species at a cyclotron resonance $\omega = n\Omega_R$.
The electric field of the IBW in the resonant region is modeled as a narrow wave packet of amplitude $E$, with central frequency $\omega$, and wave vector $k = k_0 e_x$, that is perpendicular to the magnetic field $B = Be_z$ and directed towards the plasma center. The values of the electric field amplitude and the perpendicular wave number prior to the resonance layer being obtained assuming radio frequency (rf) power flux conservation.\textsuperscript{12} In the resonant region, the wave damping should be taken into account by solving the nonlinear equation: $dE/dr = \gamma(r,E^2)E$ where $\gamma$ is the quasilinear spatial damping rate. We also assume that the spectrum in the parallel wave vector is symmetrical with respect to the sign, and peaked around $k_z = \pm k_{ij}$, with $k_{ij} \approx k_{\perp}$.

The quasilinear kinetic equation for the distribution function of the ion species of mass $m$ and electric charge $q = Ze$ is

$$\frac{\partial f_0}{\partial t} = C(f_0) + Q(f_0),$$

(1)

where $f_0 = f_0(t,r,v)$ is the dominant slowly averaged part of the total distribution function (the average is taken over a number of wave periods in both time and space); $C(f_0)$ is the Coulomb collision term and $Q(f_0)$ is the quasilinear diffusion term.\textsuperscript{13} The remaining dependence on $t$ and $r$ of the distribution function is the variation occurring on longer time and space scales compared to the wave periods and wavelengths.

Choosing a spherical coordinate system in the velocity space $(v, \xi \equiv \cos \theta, \phi)$ (where $\theta_0$ is the pitch angle, i.e., the angle between $B$ and the particle velocity $v$) and $\phi$ is the gyro angle, we limit ourselves to time scales that are long compared to the gyro period $\Omega^{-1}$ and length scales that are long compared to the Larmor radius $\rho$. These inequalities hold:

$$\left| \frac{\partial \ln f_0}{\partial t} \right| \ll \Omega, \quad |\nabla \ln f_0| \ll \rho^{-1}.$$  

(2)

In this manner, $f_0$ does not depend on the gyro angle at the lowest order.

Following Ref. 14, where an extensive depiction of the meaning and limits of this method can be found, we average Eq. (1) over a magnetic flux surface of major and minor radii $R_0$ (tokamak major radius) and $r$, respectively, crossing the resonant layer at the poloidal angle:

$$\theta_0 = \arccos(r_0/r),$$

(3)

where $r_0$ is the distance of the resonance from the plasma center. In addition, we average Eq. (1) over $\xi$ in the interval $[-1,1]$, assuming that the distribution function is isotropic in the velocity space.

The assumption that we made neglecting the anisotropy, which develops in the ion distribution function due to the fact that the $E$ field involves mainly the perpendicular ion dynamic, can be justified for the minority heating. In fact, since the quasilinear diffusion coefficient decreases for high velocities, the collisions are much more efficient in limiting the distortion of the distribution function at high energies.\textsuperscript{7} Moreover, the isotropizing effect of the collisions are even effective at low and moderate energies where quasilinear diffusion distorts the unperturbed Maxwellian even in the presence of strong rf power absorption. The isotropic solution of the Fokker–Planck equation may be a good approximation of the pitch-angle averaged distribution function, as long as the assumption $\langle Q(f) \rangle = \langle Q(f) \rangle$ can still be considered valid\textsuperscript{15} (the bracket indicates the pitch-angle average). It is worth noting that the pitch-angle average is a relevant quantity in the evaluation of macroscopic physical quantities involving velocity space averages.

The averages indicated are performed assuming the ion-cyclotron frequency profile:

$$\Omega = \Omega_0 \left(1 + \frac{r}{R_0} \cos(\theta)\right)^{-1},$$

(4)

where $\Omega_0 = qB_0/me$, and $B_0$ is the magnetic field intensity at $r = 0$.

We also note that the Coulomb collision term is evaluated in the test particle approximation, considering the resonant ion species immersed in a Maxwellian background of ions and electrons with temperature $T$.

At this point, the 1D Fokker–Planck equation is obtained and its steady state solution, relevant for IBW experiment analyses, can easily be found:6,7

$$f(r,v) = N(r) \exp\left(-\int_{0}^{v_{th}} du R(u,r)\right),$$

(5)

where

$$R(u,r) = \frac{u \Psi(u)}{0.5 \Psi(u) + u \Gamma(u,r)},$$

(6)

$N(r)$ is the normalization factor of the distribution function according to

$$\int dv f = n_i(r),$$

(7)

$$n_i(r) = \frac{n_i(r)}{\int_{-\infty}^{\infty} dv \exp(-\int_{0}^{v_{th}} du R(u,r))}.$$  

(8)

$n_i(r)$ is the density of the resonant ions, and the adimensional functions $\Psi(u)$ and $\Gamma(u,r)$ account for collisional and quasilinear diffusion, respectively. The function $\Psi(u)$ is given by

$$\Psi(u) = \lambda \Sigma_i \mu_i G(l_iu),$$

(9)

where

$$G(x) = \text{erf}(x) - \frac{2x}{\sqrt{\pi}} \exp(-x^2);$$

(10)

$$l_i = (m_i/m_f)^{1/2}, \quad \mu_i = (m_i/m_f)(n_f/n_i)(Z_i/Z_f)^2, \quad n_f$$

(11)

is the particle density, $\lambda$ is the Coulomb logarithm, and the subscript $f$ designates the background plasma particles, both electrons ($f = e$) and ions ($f = i$), of mass $m_f$ and electric charge $Z_fe$ and density $n_f$.

The function $\Gamma(u,r)$ is expressed by

$$\Gamma(u,r) = \sigma(r) S(r) n \frac{F_p(bu)}{b}.$$  

(12)
The function $\sigma$ is defined by

$$\sigma = \frac{(ZE)^2}{16\pi n_0 k_i e^2}. \quad (11)$$

The function $S(r)$ is a geometrical factor given by

$$S(r) = \frac{R_0 + r \cos(\theta_r)}{r \sin(\theta_r)} \frac{\Omega_0}{\Omega_R}, \quad (12)$$

and $\Omega_R$ is the ion-cyclotron frequency at the resonance.

The factor $n[F_\perp(\mathbf{b}u)/b]$ in Eq. (10) with

$$b = \frac{k_i v}{\Omega_R},$$

and

$$\mathbf{b}u = \frac{k_i v}{\Omega_R}, \quad (13)$$

$$F_\perp(x) = \int_{-1}^{1} d\xi J_n^2(x(1-\xi^2)^{1/2}) \quad (14)$$

($J_n$ is the Bessel function of order $n$) includes the effects of the harmonic number and finite Larmor radius.

In the limit $b \gg 1$ we have $R(u, r) \rightarrow 2u$, and the distribution function tends to be a Maxwellian. When $b \ll 1$ (which corresponds to the LHW limit), the quasilinear diffusion affects the resonant particles with $u \approx O(b^{-1})$ and the bulk distribution function is practically unmodified.

For IBWs we have $b = O(1)$ and the quasilinear term $\Gamma(u, r)$ can dominate in the thermal range as is shown in Sec. III. In addition, we observe that in either case $\Gamma(u, r)$ tends to vanish, going to high energies ($u \gg 1$) and high harmonic orders ($n \gg 1$).

In conclusion we note that the steady-state solution of the 1D Fokker–Planck equation considered here is only valid in a limited radial range. A lower bound for $r$ follows from the condition that the equation in the unknown $\alpha$ (range 0, 2$\pi$):

$$\omega - \frac{n \Omega_0}{1 + \frac{r}{R_0} \cos \alpha} - k_i v_T = 0, \quad (15)$$

has a solution for almost all the particles. This implies, in fact, that $r > r_0 + \delta$, where

$$\delta = g \frac{k_i v_T}{\omega} (R_0 + r_0), \quad (16)$$

with $g \approx 2$ [to ensure that 99% of all the particles satisfy Eq. (15)].

An upper bound for $r$ is $r_1 = (r_0^2 + h^2)^{1/2}$, where $h$ is the vertical width of the area illuminated by the wave from the equatorial plane.

III. QUASILINEAR DAMPING OF THE IBW

Here in Sec. III, we discuss the numerical evaluations of the steady-state distribution function [see Eq. (5) in Sec. II] and the corresponding IBW quasilinear damping for the IBW experiment on Alcator-C (183.6 MHz, 9.3 T) and on FTU (433 MHz, 7.8 T). In both cases, a deuterium minority is present in a hydrogen plasma. Following linear theory, this minority would completely absorb the IBW across the resonant layer, hampering wave penetration at the plasma center.

To simplify our calculations, we neglected the radial dependence of the distribution function in the resonant region by using the following approximation of the geometrical factor $S(r)$:

$$S(r) \approx O(\epsilon^{-1}), \quad (17)$$

where $\epsilon$ is the inverse aspect ratio. We also assumed constant plasma and IBW parameters including the electric field amplitude. Indeed, for weak damping (which is relevant in the experimental scenario examined) we can take a constant average value of the squared electric field on the resonant layer: $E^2 = (E_1^2 + E_0^2)/2$, where $E_1^2$ is the incident squared field amplitude (defined below in terms of the injected power) and $E_0^2$ is the final squared amplitude.

We evaluated the perpendicular wave number from the simplified dispersion equation:

$$\epsilon_{Re}^2 = 1 - \frac{m_i}{m_e} \frac{k_0^2 \rho_e^2}{2} \frac{\omega}{\pi^2 k_0^2 \rho_e} \left(\omega - \frac{n_i}{\Omega_e} \right) = 0, \quad (18)$$

where $\epsilon_{Re}$ is the IBW dispersion function, $N_i$ is the resonant harmonic number of the majority ion.

The plasma density and electron temperature at the resonant layer were calculated from the following profiles:

$$n_e(r) = n_{edge} + (n_0 - n_{edge}) \left[1 - \left(\frac{r}{a}\right)^2\right]^p, \quad (19)$$

$$T_e(r) = T_{edge} + (T_0 - T_{edge}) \left[1 - \left(\frac{r}{a}\right)^2\right]^q, \quad (20)$$

where $p = 1.2$ for Alcator-C and $p = 1$ for FTU, and $a$ is the plasma radius.

The electric field amplitude $E_i$ of the incoming wave packet before the resonance corresponds to a coupled power $P$ and can be evaluated from the following expression:

$$P = \frac{\omega}{8\pi} \frac{E_i^2}{\delta k_{\perp}} \Sigma, \quad (21)$$

where $\Sigma$ is the surface area illuminated by the wave. In order to allow for wave spreading, we also assume that $\Sigma \approx 2 \Sigma_0$ (where $\Sigma_0$ is the antenna surface area).

We have plotted in Figs. 1(a) and 1(b) and in Figs. 2(a) and 2(b) the function

$$F = 4 \pi \left(\frac{v}{v_T}\right)^2 \left(\frac{v}{v_T}\right)^3 \left(\frac{v^3}{n_i}\right) f$$

vs $v/v_T$, numerically evaluated by using the plasma and IBW parameters reported in Table I. Figure 1 pertains to the Alcator-C experiment, while in Fig. 2 we used the parameters of the FTU tokamak.

In both Figs. 1 and 2, it is possible to notice a great distortion of the plotted distribution function with respect to the Maxwellian curve in the thermal range ($v/v_T < 3$). An increase in the coupled power results into movement of the
maximum of the distribution towards higher energies and the distribution shows a larger spread.

The quasilinear diffusion flattens the distribution function in the thermal range. Thus, a reduction of the damping of the wave is expected. The spatial damping rate $\gamma$ (defined in the direction set by the wave packet group velocity) is given by

$$\gamma \approx -\frac{\epsilon_{\text{Im}}}{|\partial \epsilon_{\text{Re}}/\partial k_\perp|},$$

where $\epsilon_{\text{Im}}$ and $\epsilon_{\text{Re}}$ are the imaginary and real parts of the dielectric function, respectively. In particular, we have

FIG. 1. The function $F = [4 \pi f (v/v_T)^2 (\nu_T^2/n_e)]$ vs $(v/v_T)$ for the Alcator-C IBW experiment: $B_0 = 9.3$ T, $f_{\text{IBW}} = 183.6$ MHz, deuterium harmonic number $m = 3$, at coupled power of (a) 5 and (b) 100 kW.

FIG. 2. The function $F = [4 \pi f (v/v_T)^2 (\nu_T^2/n_e)]$ vs $(v/v_T)$ for the IBW experiment on the FTU: $B_0 = 7.8$ T, $f_{\text{IBW}} = 433$ MHz, deuterium harmonic number $m = 9$, at coupled power of (a) 30 and (b) 100 kW.
TABLE I. Plasma and rf parameters for the Alcator-C and FTU experiments used in the calculation.

<table>
<thead>
<tr>
<th>Parameters</th>
<th>Alcator-C</th>
<th>FTU</th>
</tr>
</thead>
<tbody>
<tr>
<td>Major radius</td>
<td>64 cm</td>
<td>93.5 cm</td>
</tr>
<tr>
<td>Minor radius</td>
<td>12 cm</td>
<td>30 cm</td>
</tr>
<tr>
<td>Edge density</td>
<td>$7 \times 10^{12}$ cm$^{-3}$</td>
<td>$10^{13}$ cm$^{-3}$</td>
</tr>
<tr>
<td>Central density</td>
<td>$1 \times 10^{14}$ cm$^{-3}$</td>
<td>$2.5 \times 10^{14}$ cm$^{-3}$</td>
</tr>
<tr>
<td>Edge temperature</td>
<td>40 eV</td>
<td>100 eV</td>
</tr>
<tr>
<td>Central temperature</td>
<td>900 eV</td>
<td>3600 eV</td>
</tr>
<tr>
<td>Majority species</td>
<td>Hydrogen</td>
<td>Hydrogen</td>
</tr>
<tr>
<td>Minority species</td>
<td>Deuterium</td>
<td>Deuterium</td>
</tr>
<tr>
<td>Frequency</td>
<td>183.6 MHz</td>
<td>433 MHz</td>
</tr>
<tr>
<td>$n_i$</td>
<td>5</td>
<td>5</td>
</tr>
<tr>
<td>$n_i'$</td>
<td>2230</td>
<td>943</td>
</tr>
<tr>
<td>Harmonic number</td>
<td>3</td>
<td>9</td>
</tr>
<tr>
<td>Resonance position</td>
<td>10 cm</td>
<td>22 cm</td>
</tr>
<tr>
<td>Max. power</td>
<td>100 kW</td>
<td>300 kW</td>
</tr>
<tr>
<td>Magnetic field</td>
<td>9.3 T</td>
<td>7.8 T</td>
</tr>
</tbody>
</table>

$$e_{\text{Im}} = -2\pi^2 \frac{\omega_{pD} \omega}{k_B T} v_T^{-3} \times \int_{-\infty}^{\infty} du \int_{n(y)}^{\infty} \frac{b(u^2 - u^2(y))^{1/2}}{u} \frac{dG}{du} \, dt,$$  

where $y$ is the distance from the resonant layer, $\omega_{pD}$ is the deuterium plasma frequency, and

$$u(y) = \frac{c}{n_i(R + x_R)} \frac{y}{v_T},$$

$$b = \frac{n k_i v_T}{\omega},$$

$$G = \frac{v_T^3}{n_i} f.$$

It is important to note that the main contribution to the integral of Eq. (23) comes from the low energy ions of the bulk where the distribution function remains essentially isotropic.

The expression for $\partial e_{\text{Re}}/\partial k_\perp$ is not modified by the quasilinear diffusion because $e_{\text{Re}}$ depends mainly on the electrons and on the majority ion species. It can be obtained using the approximated form of $e_{\text{Re}}$ given by the left-hand side of Eq. (18).

In Figs. 3 and 4 the total damping (expressed as the rate of the absorbed power) is plotted versus the coupled power for the Alcator-C and the FTU IBW experiments.

In the IBW experiment on Alcator-C, the quasilinear damping allows wave penetration across the deuterium minority resonant layer $\omega = 3\Omega_D$ at an operating deuterium concentration in the range of 1%–4%.

IV. RADIAL PARTICLE FLUXES INDUCED BY THE IBW

The strong distortion of the Maxwellian distribution function in the resonant plasma layer and the thinness of the transition layer towards the inner nonresonant plasma are both characteristic of IBW quasilinear physics. These two conditions could be the cause of the generation of a significant radial flux of deutos (as in classical transport, where temperature and/or pressure gradients induce radial particle fluxes in Maxwellian plasmas$^{11}$). We evaluate this flux following in great detail the formalism described in Ref. 11.

The force acting on the deutos (perpendicular to the magnetic field) due to collisions with electrons and protons is

$$F_\perp = \int d^3v m_D v \frac{\partial f_D}{\partial t} \bigg|_e.$$  

In a stationary regime, the effect of that force is a deuton flux:

$$J = \frac{e_c \wedge F_\perp}{m_D\Omega_D}.$$  

Due to the rotation symmetry of the collision operator, the force $F_\perp$ is zero if the distribution function is not dependent on the gyro-phase angle $\phi$.

For this reason, in order to evaluate $F_\perp$, we have to restore the gyro-phase angle dependence. This dependence was already neglected when we calculated the steady-state
deuton distribution. To overcome this difficulty in the following analysis we consider the deuton distribution function depending explicitly on the radial coordinate. The spatial diffusion we are considering occurs on a longer time scale compared to the characteristic time of quasilinear and collisional diffusion. For this reason, the spatial diffusion is neglected in the above Fokker–Planck equation.

We rewrite the distribution function, averaged over a number of wave periods both in space and time (see Sec. II), as

$$f_D = f_D^0(r,v,t) + f_D^1(r,v,\phi,t).$$  \hspace{1cm} (27)

Assuming that the distribution is nearly constant on the gyro trajectory to satisfy the assumed ordering of Eq. (2), we have

$$f_D(r,v,\phi,t) = h_D[r - \rho(\phi),v,t],$$  \hspace{1cm} (28)

where $\rho = e_{\parallel} \Delta v_\perp / \Omega_D$ is the gyro radius vector. A first-order Taylor expansion of $h_D$ gives

$$f_D^0(r,v,t) = h_D(r,v,t),$$  \hspace{1cm} (29)

$$f_D^1(r,v,\phi,t) = - \rho(\phi) \nabla h_D(r,v,t).$$  \hspace{1cm} (29)

Replacing the function $h_D$ with the steady-state solution $g_D$ of the 1D Fokker–Planck equation we obtain the radial flux,

$$J = \frac{4\pi}{3\Omega_D} \epsilon_i \int_0^\infty dv \langle \Delta v_\perp \rangle v^3 \frac{\partial g}{\partial r},$$  \hspace{1cm} (30)

where $\langle \Delta v_\perp \rangle$ is the Chandrasekhar diffusion parameter defined by Spitzer.  \hspace{1cm} (30)

We approximate the radial gradient of the distribution in the transition region of width $2\delta$:

$$\frac{\partial g_D^0}{\partial r} \approx \frac{g_D^e - g_D^m}{2\delta},$$  \hspace{1cm} (31)

where $\delta$ is defined by Eq. (16), $g_D^e$ is the steady-state solution of the 1D Fokker–Planck equation, and $g_D^m$ is the Maxwellian distribution corresponding to the inner nonresonant layers.

In Fig. 5 is plotted the deuton radial flux towards the resonance layer vs the injected power for (a) Alcator-C and (b) FTU plasma parameters with a deuterium minority dilution of 1%.
ergy values in the range 0 eV–3 keV injected and propagates towards the plasma center with en-

existence of a superthermal component of the deuterium minority. In addition, we note that a radial particle flux induced by the transition between the resonant layer and the outer Maxwellian plasma can also be produced. In the cases of FTU and Alcator-C examined, this flux should be directed towards the resonance. We suggest that an effect of deuton accumulation at their ion-cyclotron resonance has to be expected. A way to experimentally distinguish between direct minority heating and heating followed by spatial diffusion is the use of a CO\(_2\) laser scattering technique as described in Ref. 8. If the heating is followed by spatial diffusion, the minority density consequently rises at the resonant layer. This should increase the wave damping and reduce the amplitude of the detected signal during the IBW pulse.

V. CONCLUSIONS

Based on the steady-state solution of the 1D Fokker–Planck equation, we have numerically evaluated the quasi-linear diffusion and calculated the quasilinear damping of an IBW propagating across an ion-cyclotron resonant layer of a minority ion species. In a typical tokamak plasma experiment, we have found that the IBW can propagate across a resonant layer with minority dilution being only a few percent. Under the same conditions, linear theory predicts, instead, complete wave absorption. This result is important for the IBW experiment on the FTU, which operates in a hydrogen plasma at the ninth ion-cyclotron harmonic of the deuterium minority. Moreover, we have found that the IBW penetrates across the resonant layer \(\omega = 3\Omega_D\) when injecting a rf power in the range of 20–100 kW. This result was found for the Alcator-C (\(\nu = 183.6\) MHz, \(B = 9.3\) T) hydrogen plasma in the presence of 1%–4% dilution of the deuterium minority. This result is in agreement with the central ion heating observed during that experiment.

We have obtained a large distortion of the ion distribution function in the thermal range \((\nu < 3\nu_{th})\) in the ion-cyclotron resonant layer. This distortion of the distribution function can generate a diffusive radial particle flux towards the nonresonant Maxwellian plasma due to the existence of spatial gradients. We have calculated this flux in the framework of classical transport theory. In the cases examined, the radial particle flux directed towards the resonance layer produces a local density rise of the minority ion species. This result suggests that the IBW could be used to collect impurity ion species at some peripheral ion-cyclotron resonant layers.

ACKNOWLEDGMENTS

The authors would like to thank Dr. E. Barbato and Dr. F. Romanelli for their valuable comments on the manuscript and acknowledge Dr. F. Paoletti and L. M. Carlucci for their scientific/editorial contribution.

7M. Brambilla, Nucl. Fusion 34, 1121 (1994).