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By

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ELECTRON CYCLOTRON HEATING OF A TOKAMAK REACTOR

AT DOWN-SHIFTED FREQUENCIES

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Abstract

The absorption of electron cyclotron waves in a hot and dense tokamak plasma is investigated for the case of the extraordinary mode for outside launching. It is shown that, for electron temperatures $T_e > 5$ keV, strong absorption occurs for oblique propagation at frequencies significantly below the electron gyrofrequency at the plasma center. A new density dependence of the wave absorption is found which is more favorable for plasma heating than the familiar $n_e^{-1}$ scaling.



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Recent experimental results obtained in the T-10[1] and Doublet III[2] tokamaks have shown that electron cyclotron resonance heating (ECRH) is a very efficient method of plasma heating. So far, the most attractive type of ECRH for a tokamak reactor has been considered that with an outside launching of waves in the ordinary mode. However, this method is very questionable in the case of dense plasmas in large magnetic fields because it requires the use of high frequency microwave sources. For instance, for parameters similar to those of an INTOR-type tokamak, i.e., toroidal magnetic field $B(o) = 55$ kG, minor radius $a = 120$ cm, major radius $R = 500$ cm, central electron density $n_e(o) = 3 \times 10^{14}$ cm$^{-3}$, the use of the ordinary mode requires wave sources at frequency $f$ close to the central electron gyrofrequency $f_c(o) = 154$ GHz. Moreover, since the central plasma frequency $f_p(o) = f_c(o)$, a large part of the plasma column is not accessible.

In order to overcome these difficulties, ECRH at frequencies $f < f_c(a)$ with the extraordinary mode launched from the low field side has been proposed.[3] In this method, bulk heating is achieved via wave absorption by the fast electron tail generated during rf current drive in a low temperature plasma. In this note, we point out that the method is also applicable to a Maxwellian plasma for sufficiently high values of the electron temperature $T_e$. At first glance, ECRH by the extraordinary mode in dense plasmas also seems questionable because of the well-known $n_e^{-1}$ dependence of the damping rate. As shown below, this undesirable density scaling is generally valid for wave absorption at $f = f_c$. For large down-shift of the resonant frequency, i.e., for $f_c^2 - f^2 \gtrsim f_p^2$, wave damping for moderate values of the plasma density is nearly independent of $n_e$. For the case considered below it is only for densities $n_e > 10^{14}$ cm$^{-3}$ that the $n_e^{-1}$ law holds. The mildly relativistic dispersion relation[4] for oblique propagation is discussed elsewhere[5] and we only present the numerical results.

We consider the following spatial profiles of the plasma parameters: $n_e(r) = n_e(o)(1-r^2/a^2)$, $T_e(r) = T_e(o)(1-r^2/a^2)^{1.5}$, $B(x) = B(o)/(1+x/R)$. In order to couple the extraordinary mode to the plasma for wave launching from the low-B side of the torus, the wave frequency must fulfill the accessibility condition $f_-(o) < f < f_+(a)$, where

$$f_\pm(r) = [f_c(r)/2]\ [(1 + \frac{4f_p^{\,2}(r)/f_c^{\,2}(r)}{1 - N_\parallel^{\,2}})^{1/2} \pm 1] \quad , \tag{1}$$

and $N_\parallel$ is the parallel refractive index. For $f_p < f_c$ and $N_\parallel^2 \ll 1$, one finds $f_-(o) \lesssim 0.62\ f_c(o)$ and $f_+(a) = f_c(o)/(1 + a\cos\theta/R)$, where $\theta$ is the poloidal coordinate. For $f_c(o) = 154$ GHz, $f_-(o) \lesssim 95.5$ GHz and $f_+(a) > 123.2$ GHz, thus, for $f = 115$ GHz the accessibility condition is fulfilled.

We now discuss the case of the extraordinary mode for propagation in the equatorial plane ($\theta = 0$). We first consider the INTOR case for $T_e(o) = 5$ keV. In Fig. 1, we present $\eta(x) = 1 - \exp(-2\int_{x_o}^{x} k''_e\, dx')$ versus x for $f = 115$ GHz, $n_e(o) = 10^{14}$ cm$^{-3}$, and $\psi = 25°$ and $30°$, where $k''_e$ is the imaginary part of the x component of the propagation vector, $N_\parallel \approx \sin\psi$, and $x_o$ is the initial position at the plasma edge. It appears that appreciable absorption in the first transit occurs for $\psi \gtrsim 25°$. The fraction of the wave energy absorbed between $x_o$ and x is very sensitive to the value of $N_\parallel$ and the maximum of the power deposition occurs near $x \approx 50$ cm. As $T_e$ increases above the initial value, the plasma density can be raised further. In Fig. 2, we show $\eta(x)$ versus x for $\psi = 25°$, $n_e(o) = 3 \times 10^{14}$ cm$^{-3}$, and $T_e(o) = 8$ and 10 keV. Most of the wave energy is deposited near $r \approx a/2$. Note that for wave absorption midway between the plasma axis and the plasma edge, the relevant accessibility condition is $f_-(x_1) < f < f_+(a)$, where $x_1$ is defined by the region of maximum power deposition. In the case of $T_e(o) = 8 - 10$ keV, $x_1 = 60$ cm. This allows a raise in the central density. For $T_e(o) = 10$ keV, $f = 115$ GHz, $\psi = 25°$, and $n_e(o) = 5 \times 10^{14}$ cm$^{-3}$, the value of $\eta$ before the cut-off point ($x_- = 70$ cm) is $\eta = 0.4$ in the first transit. The reflected part of the incident wave will cross again the absorption region located near 90 cm and it is almost totally absorbed in the second transit.

We now investigate the density dependence of the wave damping. As known, several authors[5-8] have shown that the damping of the extraordinary mode scales as $n_e^{-1}$. This result holds for $f \approx f_c$ for which the anti-Hermitian part of the dielectric tensor is predominant in the computation of the wave polarization. For $N_\parallel^2 \ll 1$, the wave polarization is approximately given by $(1 - i\ \varepsilon_{12}/\varepsilon_{11})$, where $\varepsilon_{11}$ and $\varepsilon_{12}$ are the components of the plasma dielectric tensor. For $f \approx f_c$, it is found that $|1 - i\ \varepsilon_{12}/\varepsilon_{11}| \propto n_e^{-1}$, hence, $k''_e \propto |1 - i\ \varepsilon_{12}/\varepsilon_{11}|^2 \mathrm{Im}\ \varepsilon_{11} \propto n_e^{-1}$. For large values of $(f_c - f)$, however, the wave polarization tends to that of a cold plasma, i.e.,

$$1 - i\ \varepsilon_{12}/\varepsilon_{11} = 1 - (f_c/f) f_p^2/(f_c^2 - f^2 + f_p^2) \quad .$$

For $f_c^2 - f^2 \gtrsim f_p^2$, $(1 - i\ \varepsilon_{12}/\varepsilon_{11})$ decreases slower than $n_e^{-1}$ for increasing values of $n_e$. Since $\mathrm{Im}\ \varepsilon_{11} \propto n_e$, we find that $k''_e$ is nearly independent of $n_e$ for moderate values of $n_e$. For $f_c^2 - f^2 \ll f_p^2$, $(1 - i\ \varepsilon_{12}/\varepsilon_{11})$ is similar to the case $f \approx f_c$ and thermal effects are important. In this case, $k''_e \propto n_e^{-1}$. This qualitative argument is illustrated in Fig. 3 where we show $k''_e$ versus $x$ for $T_e(o) = 5$ keV, $\psi = 25°$, $f = 115$ GHz, and $n_e(o) = 3 \times 10^{13}$, $10^{14}$, and $3 \times 10^{14}$ cm$^{-3}$. The same pattern is obtained for $T_e(o) > 5$ keV. It appears that $k''_e$ is nearly independent of $n_e$ for $3 \times 10^{13}$ cm$^{-3}$ $< n_e(o) < 10^{14}$ cm$^{-3}$ where $f_c^2 - f^2 \gtrsim f_p^2$ and decreases approximately as $n_e^{-1}$ for $n_e(o) >$

$10^{14}$ $cm^{-3}$. The maximum value of $k''_e$ for $n_e(o) = 3 \times 10^{14}$ $cm^{-3}$ is approximately 1/3 of the value for $n_e(o) = 3 \times 10^{13}$ $cm^{-3}$ rather than 1/10 as predicted by the $n_e^{-1}$ dependence.

It is of interest to determine the range of velocities of the absorbing electrons. This is obtained using the relativistic resonance condition

$$p_\parallel/mc = [N_\parallel(f_c/f) \pm (N_\parallel^2 - 1 + f_c^2/f^2)^{1/2}]/(1 - N_\parallel^2) \quad . \tag{2}$$

where m is the electron rest mass, c is the speed of light, and $p_\parallel$ is the electron parallel momentum. For the case of $T_e(o) = 8$ keV, $n_e(o) = 3 \times 10^{14}$ $cm^{-3}$, and $\psi = 25°$ (Fig. 2), the maximum power deposition occurs near $x = 70$ cm where $f_c/f = 1.17$ and $N_\parallel = 0.42$. Using Eq. (2) with minus sign (the other root is irrelevant), we obtain $p_\parallel/mc \approx -0.30$, i.e., $|p_\parallel| \approx 3.3\ (mT_e)^{1/2}$. This value of $p_\parallel$ lies in the low velocity part of the electron tail for which the density of the absorbing electrons is high enough for appreciable wave damping and Coulomb collisions with the electron body are sufficiently frequent to achieve thermalization in a time less than the energy confinement time. As shown in Figs. 1 and 2 and by Eq. (2), electron heating takes place via superthermal electrons in the momentum space and away from the plasma axis. This calls for two comments. First of all, for large wave powers the relatively weakly collisional electrons may experience quasilinear effects which in principle tend to reduce the damping rate. A quantitative answer to this problem requires a Fokker-Planck kinetic code. However, we expect that the power dependence of $k''_e$ is rather weak since the wave energy is deposited over magnetic surfaces of large radius. Secondly, as shown by experiments[1] for noncentral ECRH, the global energy confinement time is not worse than that obtained for power deposition near the plasma axis. Furthermore, for electron heating far away from the plasma axis, the possibility of suppressing the m = 2 mode has been demonstrated. Note also that numerical simulations[9] of the ignition experiment in the INTOR tokamak point out the advantage of noncentral heating. The results presented in Figs. 1, 2, and 3 are obtained using a slab model and are valid for propagation in the equatorial plane. We have checked that in this case a ray tracing code yields the same results. Plasma heating for lower values of B(o) and f is possible in present day large tokamak devices. For instance, using the TFTR parameters, a = 80 cm, R = 250 cm, B(o) = 30 kG, f = 60 GHz, $\psi$ = 45°, $n_e(o) = 5 \times 10^{13}$ cm$^{-3}$, and $T_e(o)$ = 4 keV, we find that more than 90% of the extraordinary mode is absorbed in the first transit, and that half of the energy is absorbed in the regions with r < 40 cm and r < 55 cm, when the wave is launched from the high field side and the low field side, respectively. Most of the wave energy is deposited in the vicinity of 45 cm and 55 cm, respectively. The investigation of wave absorption for large down-shifted frequencies for $T_e(o)$ = 1-2 keV is also possible if the electron tail is activated by an external pusher. This occurs in a low density ohmic plasma or in the lower hybrid current drive regime. In both cases, the electron distribution possesses a mildly relativistic tail which can be used to investigate the properties of electron cyclotron wave absorption by energetic electrons.[3] In addition, an interesting by-product of wave absorption for down-shifted frequencies is the generation of an additional current and, therefore, the method can be used for optimization of current drive.[10]

In conclusion, we have shown that the two major difficulties of ECRH, namely, the development of high frequency gyrotrons and the cut-off density of the ordinary mode can in principle be overcome by using the extraordinary mode for external launching at down-shifted frequencies. We have also shown that for wave absorption far away from the region where $f \approx f_c$, the damping of the extraordinary mode is weakly dependent on the plasma density. This contrasts with the density dependence for wave absorption near $f \approx f_c$. Since, in high temperature plasmas, most of the wave energy is absorbed far away from the region where $f \approx f_c$, the limitation due to density raising is less severe than generally expected. Finally, in most cases considered in this work, the ordinary mode damping is found to be negligible.

ACKNOWLEDGMENT

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FIGURE CAPTIONS

FIG. 1 $\eta$ vs x for $n_e(o) = 10^{14}$ cm$^{-3}$, $T_e(o)$ = 5 keV, f = 115 GHz, and $\psi$ = 25° and 30°.

FIG. 2 $\eta$ vs x for $n_e(o) = 3 \times 10^{14}$ cm$^{-3}$, f = 115 GHz, $\psi$ = 25° and $T_e(o)$ = 8 keV and 10 keV.

FIG. 3 $k''_e$ in cm$^{-1}$ vs x for f = 115 GHz, $T_e(o)$ = 5 keV, $\psi$ = 25°, and $n_e(o)$ = $3 \times 10^{13}$ cm$^{-3}$, $10^{14}$ cm$^{-3}$, and $3 \times 10^{14}$ cm$^{-3}$.

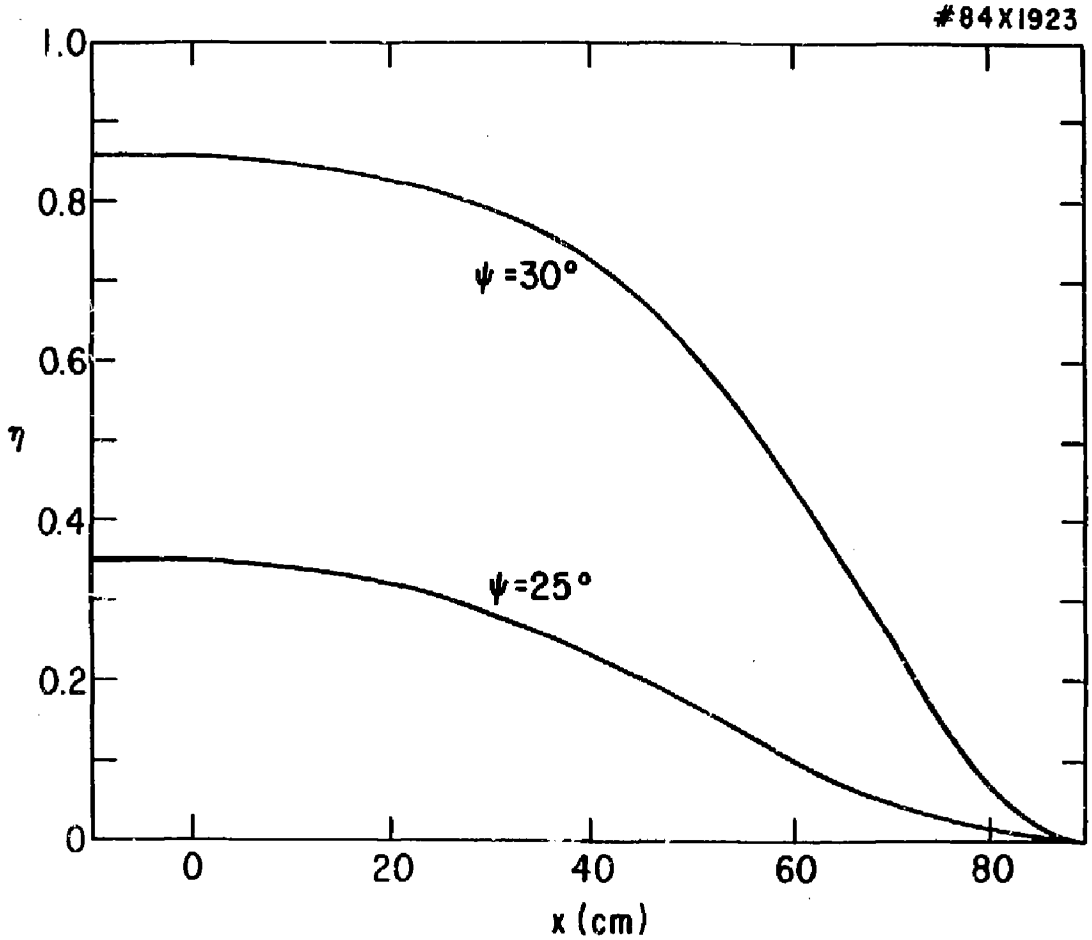



Fig. 1

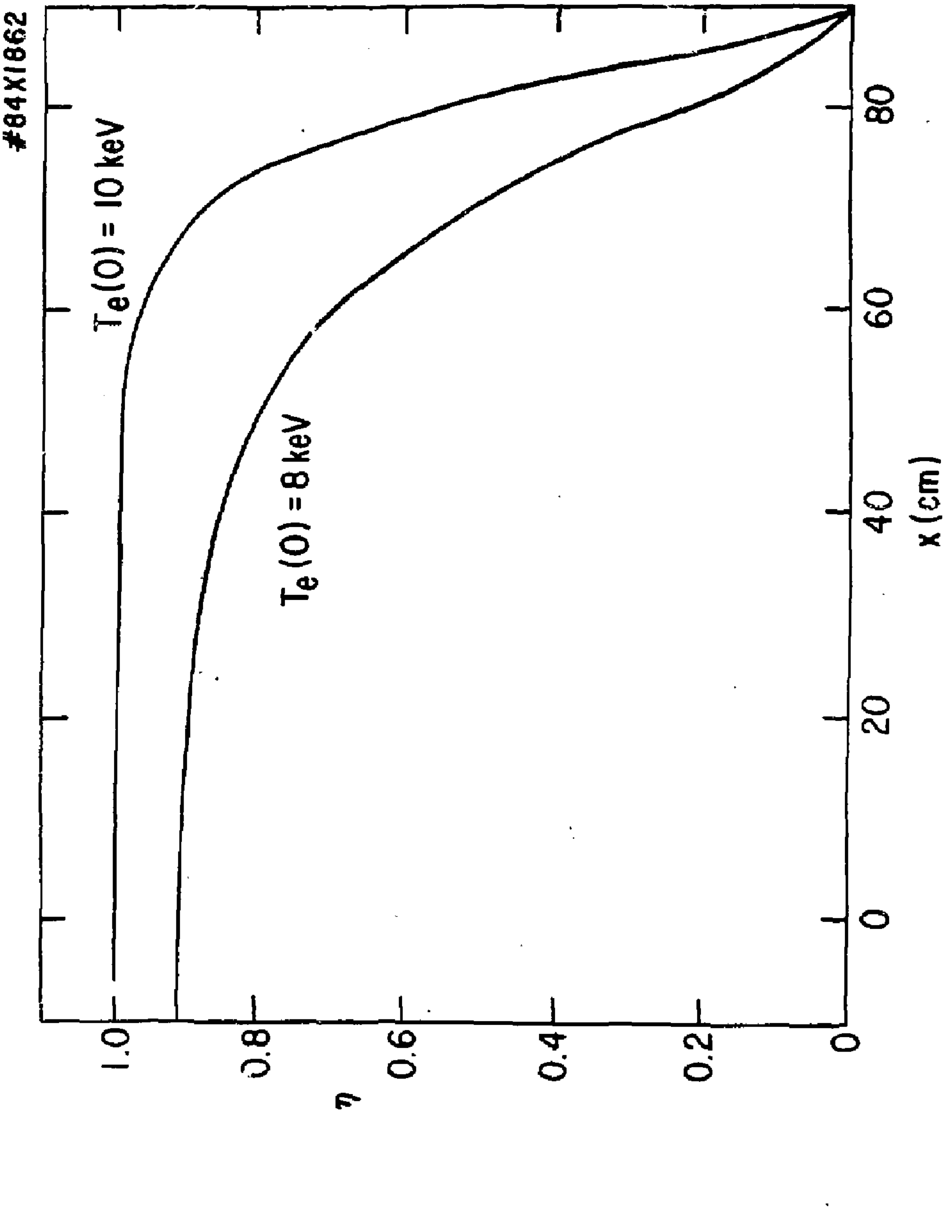
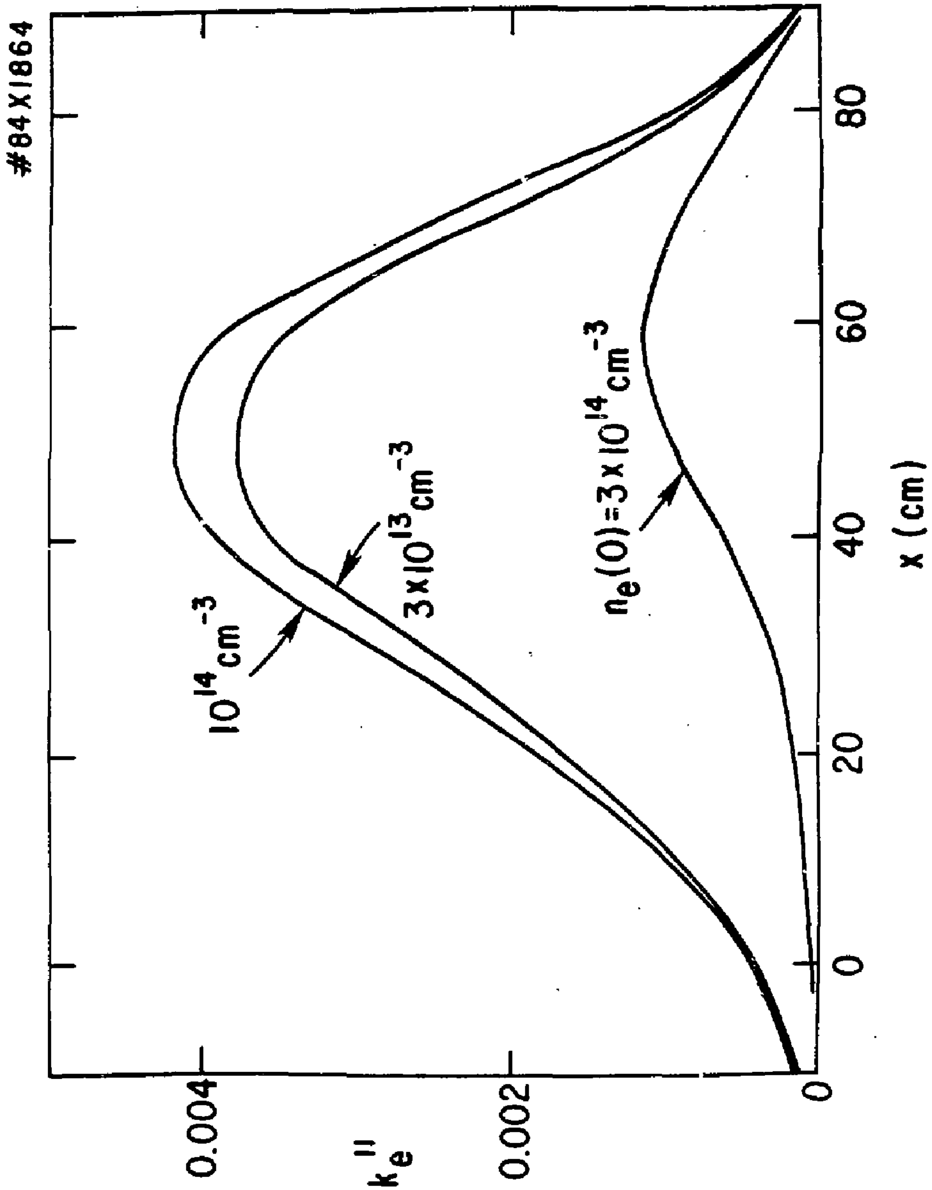



Fig. 2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0.002
0
$k_e''$
0
20
40
60
80
x (cm)
$10^{14}cm^{-3}$
$3\times10^{13}cm^{-3}$
$n_e(0)=3\times10^{14}cm^{-3}$

Fig. 3

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