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A full wave theory of high-harmonic fast wave absorption in high-beta plasmas

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A theory of fast wave absorption in a high-beta plasma is given. A reduced, second-order, ordinary differential equation has been used which includes all collisionless electron dissipation mechanisms and ion cyclotron damping over many harmonics. This is relevant to the high-harmonic fast wave heating scheme proposed by Ono [Phys. Plasmas **2**, 4075 (1995)]. The parameters appropriate to the National Spherical Tokamak Experiment NSTX [J. Spitzer *et al.*, Fusion Technol. **30**, 1337 (1996)] have been used to investigate the absorption characteristics of electrons and ions under high-beta, high-harmonic conditions. [S1070-664X(98)00206-7]

I. INTRODUCTION

Small aspect ratio, spherical tokamaks represent a promising alternative route to magnetic fusion.[1–4] One of the advantages of spherical tokamaks is the high-beta condition, which is readily obtained[5] due to the low value of the magnetic field. In order to achieve fusion conditions in these devices additional heating schemes must be employed. Neutral beam injection is presently in use on the Small, Tight Aspect Ratio Tokamak (START[1]) and will also be utilized on the Mega Amp Spherical Tokamak (MAST[6]), the successor to START. On the other hand, a fast wave heating scheme will be installed on the National Spherical Tokamak Experiment (NSTX)[4] in the United States.

High-beta, low-magnetic-field plasmas provide a very different environment for radio frequency heating compared with conventional tokamaks. In order to take account of these very different conditions and to be able to use existing radio frequency systems, Ono[7] has proposed the use of high-cyclotron-harmonic fast wave heating for NSTX. This scheme will be required to heat the NSTX plasma starting from electron and ion temperatures which are of the order of 0.1–0.2 keV. The principal aim of the fast wave heating system is to heat the electrons both on and off axis and to drive the plasma current. Such a noninductive current drive scheme would have the highly desirable consequence of allowing the ohmic solenoid to be dispensed with from the central column of a spherical tokamak. Because of the low-field, high-beta conditions, the electron dissipation at these low temperatures is sufficient to commence heating, and further increase of the electron temperature causes still stronger electron absorption.

Ono[7] has carried out calculations which indicate that high-harmonic fast wave heating is a promising scheme for spherical tokamaks. In these calculations a parametric study of the absorption was made by solving the local dispersion relation for Im $k_\perp$. The conditions for strong electron heating were obtained assuming cold ions. It was also noted that when hot ions were included, ion absorption could become significant and would eventually provide an accessibility limit. This occurred for bulk ion temperatures between 1 and 2 keV.

The propagation of the fast wave at high cyclotron harmonics ($n \sim 20$) in a plasma containing hot ions poses a very different problem in a spherical tokamak from that in a conventional tokamak. The fast wave is usually thought of as a cold plasma wave in which the properties of the wave are dominated by the perpendicular motion of the ions. In the kinetic description, this corresponds to the $n = \pm 1$ harmonics which are the only ones to survive in the limit of a cold plasma. Other harmonic terms are only important when the wave frequency is very close to the particular ion cyclotron harmonic.

An important quantity for ion cyclotron interactions of the fast wave is the ratio $\rho_i/\lambda_\perp$, where ρ_i is the ion Larmor radius and $\lambda_\perp$ is the perpendicular wavelength of the fast wave. This is characterized by the parameter $\lambda_i \equiv k_\perp^2 v_{Ti}^2/2\Omega_i^2$ where the symbols have their usual meaning. For low cyclotron harmonics in conventional tokamaks, $\lambda_i \ll 1$. For the proposed NSTX with $B_0 = 0.25$ T, a plasma density of 5×10^{19} m^{-3}, and a harmonic number of $n = 20$, $\lambda_D \simeq 64 T_D$ (keV) where a deuterium plasma has been assumed and $k_\perp$ has been approximated by ω/c_A. Hence, for modest ion temperatures (e.g., $T_i \gtrsim 0.1$ keV) the fast wave will propagate under large Larmor radius conditions.

In these circumstances, the contribution of the $n = \pm 1$ harmonics to the propagation of the fast wave becomes negligible, as do a number of the other lower harmonics. Evidently, the fast wave behavior is governed by a few harmonics in the vicinity of the nearest ion cyclotron resonance. Nevertheless, the hot plasma contributions produce a response which is close to the cold value, namely $k_\perp \sim \omega/c_A$. In contrast to the situation in a low beta tokamak where the interaction of the fast wave with high ion cyclotron harmonics is negligible for a thermal plasma, this interaction becomes important, especially when the fast wave crosses many cyclotron harmonics, as it will in the field of a spherical tokamak.

As emphasized already,[7] the interaction of the ion cyclotron harmonics with the fast wave becomes crucial in assessing the effectiveness of fast wave heating in this regime. When the fast wave crosses each ion cyclotron harmonic resonance it will either undergo cyclotron damping, mode conversion to an ion Bernstein wave, or a combination of both effects. Incident fast wave energy which is not cyclotron damped or mode converted will be transmitted. For low-field side incidence the mode-converted ion Bernstein wave will give rise to reflection of the fast wave.

In order to quantify the relative significance of these effects we must use a full wave treatment. In general, this requires the solution of a coupled system of integrodifferential equations.[8,9] However, we shall attempt a much simpler approach in order to obtain the main features of fast wave heating under these unusual conditions and hence provide an efficient technique for assessing the heating performance.

The approach is based on the fast wave approximation[10,11] that has been used successfully in modeling ion cyclotron heating in conventional tokamaks. In the most recent version,[12,13] a second-order differential equation was constructed from the full electromagnetic, hot plasma dispersion relation. This allowed the electron damping to be calculated in addition to the ion absorption. This is important for modeling high-harmonic fast wave heating since its purpose is electron heating and current drive. However, the major problem for the fast wave approximation in high-beta conditions is to obtain a reliable description of the ion dissipation as the wave successively crosses many ion cyclotron harmonics. Since we use a second-order equation that only describes the propagation of the fast wave, mode conversion is not explicitly included. However, its effect is indicated by the occurrence of reflection from an ion cyclotron resonance region. Ion cyclotron damping is, of course, included explicitly and the aim is to provide an accurate calculation of this. The fast wave equation is formulated for a one-dimensional model in which the toroidal magnetic field, the equilibrium density, and temperatures vary as a function of x which simulates the radial coordinate. The amplitude of the fast wave is described by E_y and the other electric field components can be obtained from this with the aid of the local polarization relations.

II. THE FAST WAVE APPROXIMATION FOR HIGH HARMONICS

The problem of fast wave heating at high-ion-cyclotron harmonics in a spherical tokamak plasma contains a number of novel features. Because of the large magnetic shear in these plasmas, particularly in the edge region, where the poloidal magnetic field is comparable to or even larger than the toroidal magnetic field, the $k_\parallel$-spectrum is much more difficult to model and solutions are only now beginning to emerge. Consequently, the $k_\parallel$-spectrum is not known reliably at present. Another novel feature is that the magnetic field profile can be nonmonotonic. In this paper we shall not address either of these issues but will concentrate on the high-beta aspect of spherical tokamaks. As already discussed, this means that the fast wave propagates under conditions where the ion Larmor radius of the thermal ions can be much larger than the perpendicular wavelength. Under these circumstances the propagation and absorption of the fast wave in a nonuniform plasma gives rise to three coupled integrodifferential equations. Sauter and Vaclavik[8] and Brambilla[9] have developed codes to solve this system of equations for absorption of fast wave energy by alpha particles in conventional tokamaks. Smithe *et al.*[14] have recently produced a similar code to deal with the high-harmonic case. These codes are extremely complicated and very time consuming.

In this paper we have extended the fast wave approximation to deal with the high-harmonic case. The fast wave approximation was originally developed for low-beta conventional tokamaks, to deal with second harmonic,[10] minority, and ion hybrid resonance.[11] It was later extended to include electron dissipation.[12,13] In the present context we employ the fast wave approximation to provide a full wave theory of all electron dissipation mechanisms and ion cyclotron absorption at many cyclotron harmonics.

We begin with the full, hot plasma electromagnetic dispersion relation for a locally uniform plasma. This can be written in many different ways. However, it is convenient to write it in the following form for reasons that will become apparent shortly. Thus,

$$n_\perp^2 = \epsilon_{yy} - n_\parallel^2 - \frac{\epsilon_{xy}^2}{(\epsilon_{xx} - n_\parallel^2)} - \frac{\epsilon_{yz}^2}{(\epsilon_{zz} - n_\perp^2)} + \frac{2\epsilon_{xy}\epsilon_{yz}(\epsilon_{xz} + n_\perp n_\parallel)}{(\epsilon_{zz} - n_\perp^2)(\epsilon_{xx} - n_\parallel^2)} - \frac{(\epsilon_{yy} - n^2)(\epsilon_{xz} + n_\perp n_\parallel)^2}{(\epsilon_{zz} - n_\perp^2)(\epsilon_{xx} - n_\parallel^2)}. \tag{1}$$

For a cold plasma, the fast wave refractive index is given very accurately by the first three terms on the right-hand side of Eq. (1). This is because the remaining terms are either zero for the cold plasma limit or very small because ϵ_{zz} is much larger than the other elements of the dielectric tensor for the frequency range of interest. The latter condition is the reason that $E_\parallel \ll E_\perp$ for the fast wave which is still true for a hot plasma.

The fast wave approximation in a low-beta plasma is obtained by substituting the cold plasma refractive index into the hot plasma corrections to the dielectric tensor elements. This enables a reduced, second-order fast wave equation to be obtained which includes ion cyclotron damping, wave resonance (when it occurs), and all electron dissipation effects. Wave resonance arises when $\epsilon_{xx} - n_\parallel^2 \to 0$ in the third term on the right-hand side of Eq. (1). It is clear that this term contains information on the effect of ion Bernstein waves on the fast wave.

We follow a similar procedure for the high harmonic case. Here,

$$\epsilon_{xx} \simeq 1+\sum_j \frac{\omega_{pj}^2}{\omega k_\parallel v_{Tj}} \sum_{n=-50}^{50} \frac{n^2 I_n(\lambda_j) Z(\zeta_{nj})}{\lambda_j}, \tag{2}$$

where $\zeta_{nj} \equiv (\omega - n\Omega_j)/k_\parallel v_{Tj}$. In Eq. (2) only the $n=0$ term is important for the electrons, but for the ions the summation over the first 50 cyclotron harmonics is included. For the conditions relevant to NSTX the first 20 or so harmonics will be resonant across the machine. We have therefore carried out the sum over the ion harmonics to $n=50$ for greater accuracy. The same procedure is adopted for the other elements of the dielectric tensor where the notation of Stix[15] has been used.

In contrast to the low-beta case, if the cold plasma root is substituted into Eq. (1) it does not give a good approximation to the fast wave root for the high-beta, high-harmonic case. Hence, in order to obtain a reliable description of the propagation, reflection, and damping of the fast wave a different procedure has been used to generate the fast wave equation. This is derived with the aid of the full transcendental, electromagnetic dispersion relation which is obtained by substituting Eq. (2), and similar expressions for the other elements of the dielectric tensor, into Eq. (1). The exact, fast wave root $k_{\perp F}^2$ is extracted from this equation and this is used as the potential, $V(x)$, in the fast wave equation for the high-beta, high-harmonic case, which is written as

$$\frac{d^2 E_y}{dx^2} + V(x) E_y = 0. \tag{3}$$

The exact, fast wave root is a local value which is spatially varying due to the equilibrium profiles which specify the inhomogeneous plasma. The spatial variation for $k_\perp^2(x)$, which must be solved for at each position, gives the spatial dependence of the fast wave potential. For the present one-dimensional model we assume a straight magnetic field with a linear gradient, i.e.,

$$\mathbf{B}_0(x) = \hat{i}_z B_0 \left(1+\frac{x}{R}\right). \tag{4}$$

The density and temperatures were assumed to have profiles of the form

$$A = A_0 (1 - x^{\gamma_1})^{\gamma_2}. \tag{5}$$

We have taken $\gamma_1 = \gamma_2 = 2$ for the temperature profile and $\gamma_1 = 2$, $\gamma_2 = 0.5$ for the density profile.

The second-order fast wave equation given by Eq. (3) was integrated across the interaction region to obtain the transmission, reflection, and absorption coefficients for the fast wave. Since, by definition, Eq. (3) only describes the propagation of the fast wave, it does not include explicitly the phenomenon of mode conversion. However, the solution of the fast wave equation will contain a ''signature'' of mode conversion which may occur in the vicinity of an ion cyclotron harmonic resonance. If mode conversion is significant, the fast wave will be partially reflected at the corresponding ion cyclotron harmonic for low-field side incidence.

The power absorbed can be calculated by computing $\frac{1}{2}\mathrm{Re}\,\mathbf{J}\cdot\mathbf{E}^*$, giving

$$\begin{aligned} P = {} & \frac{\omega\epsilon_0}{2}|E_x|^2 \,\mathrm{Im}\,\epsilon_{xx} + \frac{\omega\epsilon_0}{2}|E_y|^2\,\mathrm{Im}\,\epsilon_{yy} \\ & + \frac{\omega\epsilon_0}{2}|E_z|^2\,\mathrm{Im}\,\epsilon_{zz} + \omega\epsilon_0\,\mathrm{Re}(E_x E_z^*)\mathrm{Im}\,\epsilon_{xz} \\ & - \omega\epsilon_0\,\mathrm{Im}(E_x E_y^*)\mathrm{Re}\,\epsilon_{xy} + \omega\epsilon_0\,\mathrm{Im}(E_z E_y^*)\mathrm{Re}\,\epsilon_{yz}. \end{aligned} \tag{6}$$

The power absorbed by the ions is dominated by the first, second, and fifth terms whereas the electron power absorption is given by the second, third, and sixth terms. For the ions it is the various harmonic terms which contribute whereas for the electrons only the $n=0$ term is important since $\omega \ll \Omega_e$.

Because the exact fast wave root of the full electromagnetic dispersion relation is used we must take account of the fact that the root is complex in order to calculate the power absorbed. Thus

$$k_{\perp F} = k_1 + i k_2. \tag{7}$$

Hence, the first term in Eq. (6) gives

$$\begin{aligned} & \frac{\omega\epsilon_0}{2}|E_x|^2\,\mathrm{Im}\,\epsilon_{xx}(k_{\perp F}) \\ & \quad \simeq \frac{\omega\epsilon_0}{2}|E_x|^2\,\mathrm{Im}\,\epsilon_{xx}(k_1) + \frac{\omega\epsilon_0}{2}|E_x|^2 k_2 \frac{\partial}{\partial k_1}\mathrm{Re}\,\epsilon_{xx}(k_1), \end{aligned} \tag{8}$$

where the first term on the right-hand side of Eq. (8) contributes to the power dissipated and the second term represents the contribution of the particles to the power flow. This is usually referred to as the kinetic power flow and can be negative in some regions and positive in others. It can therefore influence the amplitude of the electromagnetic field. The power dissipated is obtained from Eq. (6) by expanding all the terms in a similar manner, thus separating those responsible for absorption from those responsible for kinetic power flow.

III. NUMERICAL RESULTS

The NSTX experiment, which is presently under construction, will include a high-harmonic fast wave heating system. We have therefore used the fast wave equation to investigate the properties of high-harmonic heating for parameters relevant to NSTX. We have assumed a deuterium plasma with a central density, $n_{e0} = 5\times10^{19}\ \mathrm{m}^{-3}$. The major and minor radii have been taken as $R = 0.8$ m and $a = 0.64$ m. The frequency of the fast wave has been taken as $f = 41$ MHz. These are also the parameters used by Ono[7] when solving the hot plasma dispersion relation.

The first case we consider corresponds to the initial conditions in NSTX before the application of the high-harmonic heating. We therefore assume a central electron temperature, $T_{e0} = 0.2$ keV, and a central ion temperature, $T_{i0} = 0.1$ keV.

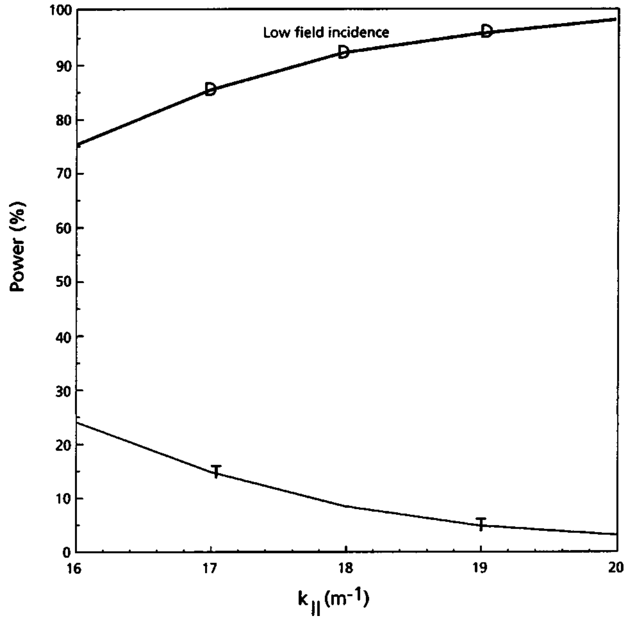

FIG. 1. Power transmission (T) and absorption (D) coefficients as a function of $k_{\|}(\mathrm{m}^{-1})$ for a deuterium plasma with NSTX parameters ($n_{e0}=5\times10^{19}\ \mathrm{m}^{-3}$, $R=0.8$ m, $a=0.64$ m, $B_0=0.25$ T, $f=41$ MHz, $T_e=0.2$ keV, and $T_D=0.1$ keV).

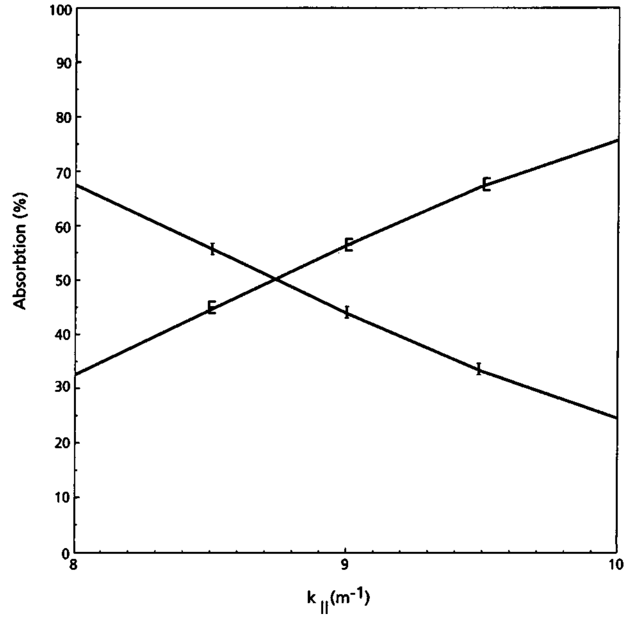

FIG. 3. Power absorbed by ions (I) and electrons (E) as a function of $k_{\|}(\mathrm{m}^{-1})$ for the parameters of Fig. 1 except $T_{e0}=1$ keV and $T_D=0.51$ keV.

The fast wave is assumed to be incident from the low-field side and for $k_{\|}$ in the range $8-10\ \mathrm{m}^{-1}$ the absorption is negligible. For the same plasma conditions but a higher range of parallel wave numbers with $k_{\|}$ between 16 and $20\ \mathrm{m}^{-1}$, Fig. 1 shows that the absorption is close to total for the larger values of $k_{\|}$ (i.e., for a single transit). The individual contributions to the power absorbed are shown in Fig. 2. The spatial location of the absorbed power is also given in Fig. 2, which is calculated for a value $k_{\|}=17\ \mathrm{m}^{-1}$, and shows that the power is deposited fairly close to the center of the discharge. The ion cyclotron absorption is seen to be negligible and the largest contribution comes from the electron Landau damping labelled by the number 3. The transit time damping, labelled by 2, is about a factor 3 smaller and the cross terms (labelled x) give their usual negative contribution.

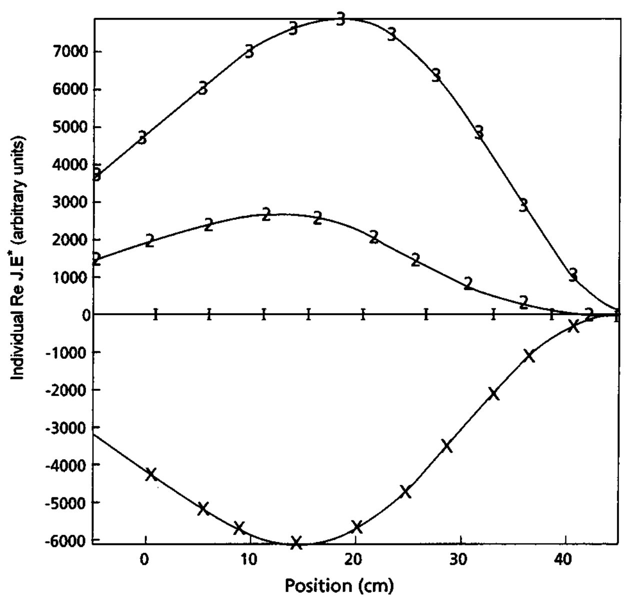

FIG. 2. Individual contributions to the power absorbed as a function of position for $k_{\|}=17\ \mathrm{m}^{-1}$ ($I=$ions, $3=$electron Landau damping, $2=$electron transit time damping, and $x=$electron cross terms).

Next, we consider a case where the central electron temperature $T_{e0}=1$ keV with a central ion temperature, $T_{i0}=0.51$ keV. For this case, the range of $k_{\|}$ is $8-10\ \mathrm{m}^{-1}$ and the power absorbed by the electrons, labelled (E), and the ions (I) is shown in Fig. 3. We see that this represents a critical case since for the lower values of $k_{\|}$, $P_i>P_e$, while for the larger values, $P_e>P_i$. This information is also exhibited in Fig. 4, which gives the various contributions to the electron absorption indicating that the transit time damping is larger than the electron Landau damping for these conditions. The total power absorbed as a function of position for $k_{\|}=9.5\ \mathrm{m}^{-1}$ is given in Fig. 5 which shows that for these hotter conditions the absorption occurs further from the center.

It is instructive to illustrate the separation of the kinetic power contributions from the power absorbed as shown in Eq. (8). The individual contributions to Re $\mathbf{J}\cdot\mathbf{E}^*$ are shown in Fig. 6, where the dielectric tensor elements have been evaluated with the exact complex value of $k_{\perp F}$. The power absorbed over the range from $x=-5$ to $x=5$ cm is shown and it can be seen that there are substantial regions where the ion contribution is negative. In Fig. 7, the power absorbed is shown over the same spatial region, calculated using only k_1, the real part of $k_{\perp F}$, as indicated in Eq. (8). It can be seen that the power absorbed by the ions is now positive definite. The ion and electron contributions to the kinetic power calculated according to Eq. (8) are shown in Fig. 8. Both are negative for these conditions.

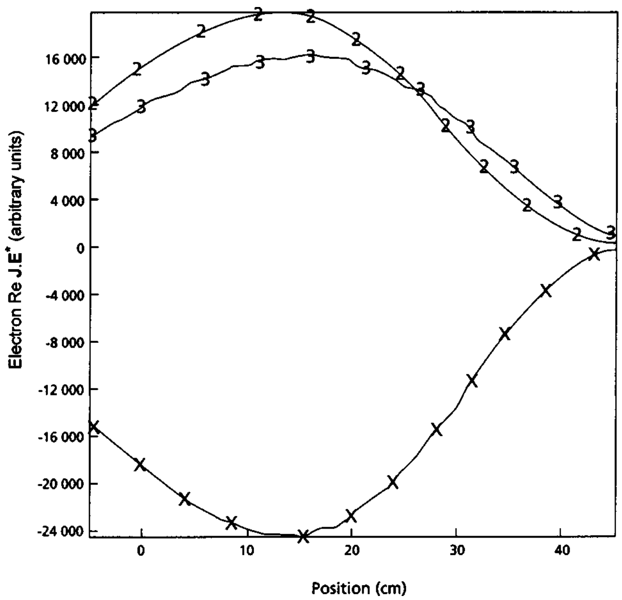

FIG. 4. Individual contributions to the electron power absorbed as a function of position with the notation of Fig. 2.

When both the electron and ion temperatures have the value of 1 keV and the parallel wave number is in the range $8-10\ \mathrm{m}^{-1}$, the power is almost totally absorbed by the ions. Over this range of parallel wave numbers the ion absorption falls from 100% for the lowest value of $k_\parallel$ to about 93% for the largest value. Furthermore, the power is absorbed close to the plasma edge. These conditions are clearly unfavorable for electron heating and current drive, either on or off axis. For these conditions, Fig. 9 shows that the kinetic power flow of the ions oscillates between negative and smaller positive values as a function of position.

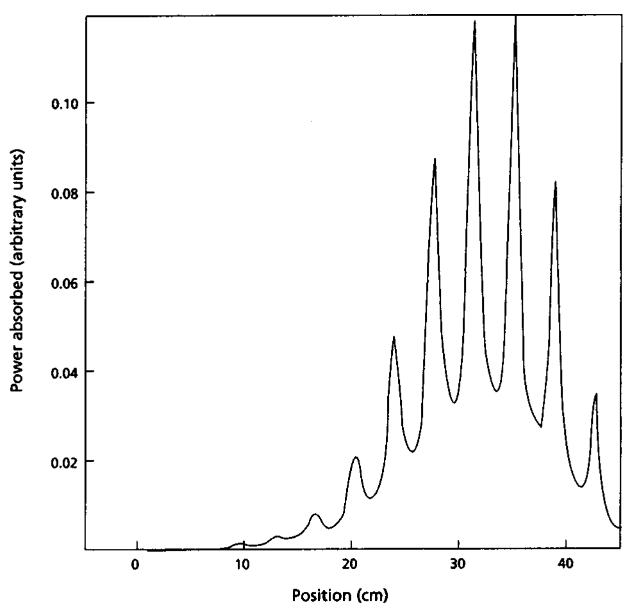

FIG. 5. Total power absorbed as a function of position for the conditions of Fig. 3 and $k_\parallel = 9.5\ \mathrm{m}^{-1}$.

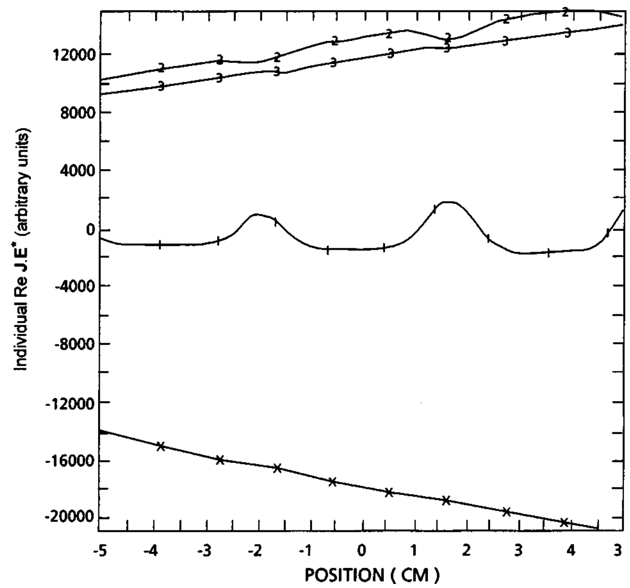

FIG. 6. Individual contributions to the power absorbed as a function of position for the parameters of Fig. 3 with $k_\parallel = 9.5\ \mathrm{m}^{-1}$. The power absorbed has not been separated from the kinetic power in this figure.

The previous case is put into perspective by the final examples. For these, $k_\parallel$ is again in the range $8-10\ \mathrm{m}^{-1}$ and the central ion temperature is still 1 keV, but now a central electron temperature of 2 keV is considered. For the lower end of the range of $k_\parallel$ the ions absorb about 70% of the power and the electrons about 30%. Equal sharing of the power occurs at $k_\parallel \simeq 9\ \mathrm{m}^{-1}$ and at the upper end of the range of $k_\parallel$ the electrons absorb about 65% of the power compared to 35% by the ions.

The sensitivity of the ratio P_e/P_i to the value of $k_\parallel$ is illustrated in Figs. 10 and 11. For these two cases $T_{e0} = 1.5$ keV and $T_{i0} = 0.8$ keV. In Fig. 10 the power absorbed

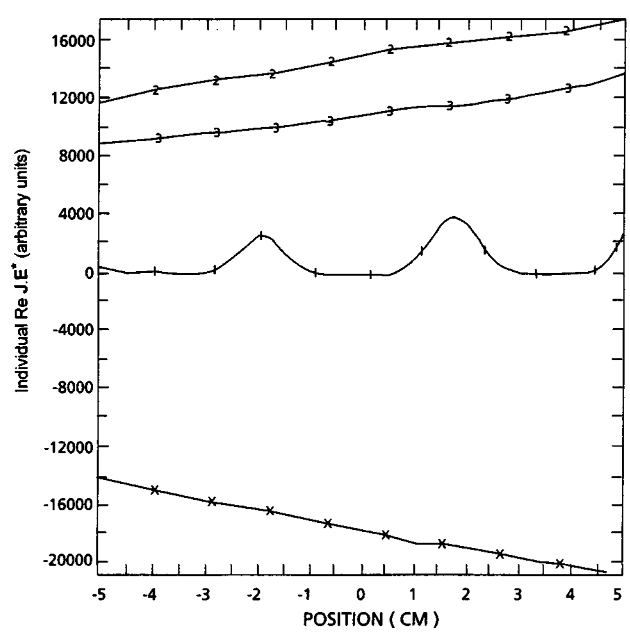

FIG. 7. Same as Fig. 6 but with the kinetic power subtracted.

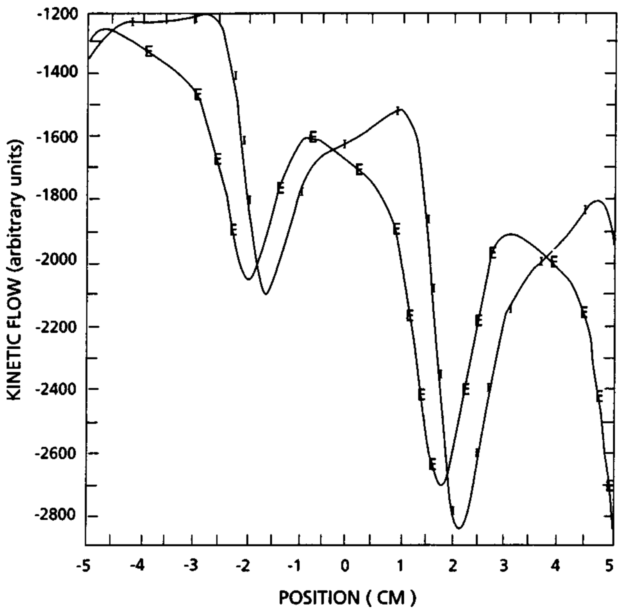

FIG. 8. Ion and electron contributions to the kinetic power for the parameters of Fig. 6, plotted as a function of position.

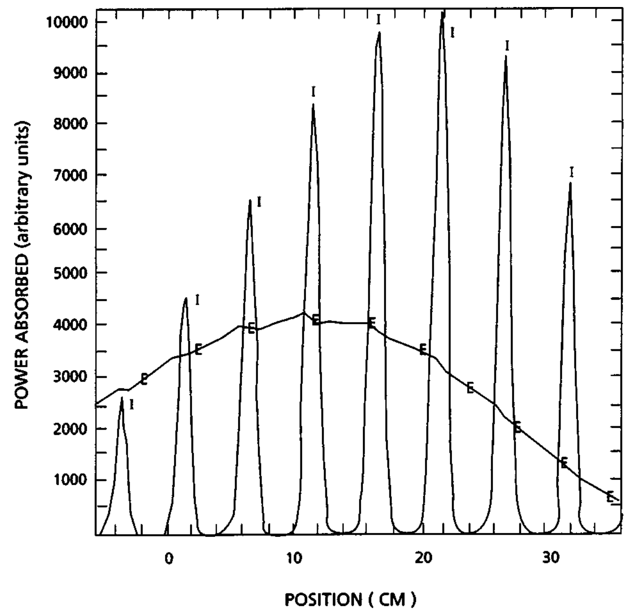

FIG. 10. Power absorbed by ions (I) and electrons (E) as a function of position for the parameters of Fig. 1 except $T_{e0}=1.5$ keV, $T_D=0.8$ keV, and $k_{\|}=6\ \mathrm{m}^{-1}$.

by the electrons and the ions is plotted as a function of position for $k_{\|}=6\ \mathrm{m}^{-1}$. For this case the powers are again approximately equal. In Fig. 11 the value of $k_{\|}$ is $8\ \mathrm{m}^{-1}$ and now the electron power is much larger than the ion power.

Mode conversion of the fast wave to the ion Bernstein wave does not appear to be significant for the conditions relevant to high-harmonic heating in high-beta, spherical tokamaks. When there is strong coupling of the fast wave to the ions it is through harmonic ion cyclotron damping which evidently suppresses mode conversion. In the calculations that we have performed we have found almost no reflection of the fast wave as it crosses the ion cyclotron harmonic resonances.

IV. SUMMARY AND CONCLUSIONS

We have extended the fast wave approximation to provide a reduced, full wave model for high-harmonic fast wave heating in high-beta conditions relevant to spherical tokamaks. The conditions expected in the future NSTX have been considered for illustrative purposes. Our results confirm Ono's[7] conclusion that electron absorption can be strong, even for the initial conditions, $T_{i0}\sim 0.1$ keV and T_{e0}

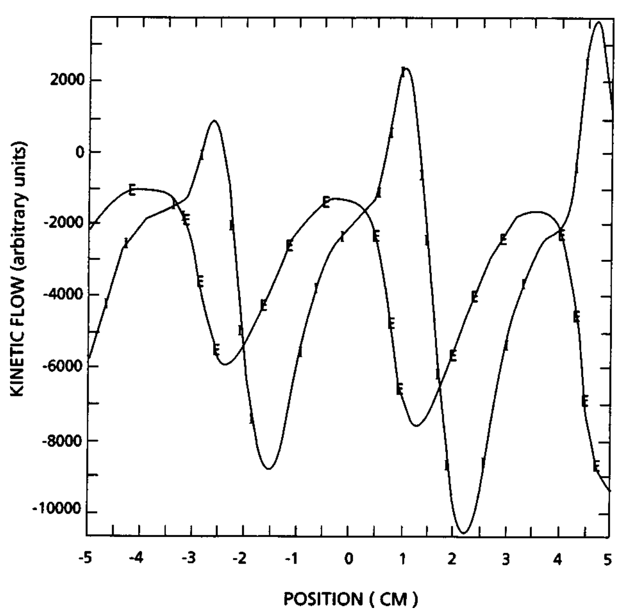

FIG. 9. Ion and electron contributions to the kinetic power as a function of position for the parameters of Fig. 1 except $T_{e0}=1$ keV, $T_D=1$ keV, and $k_{\|}=8.5\ \mathrm{m}^{-1}$.

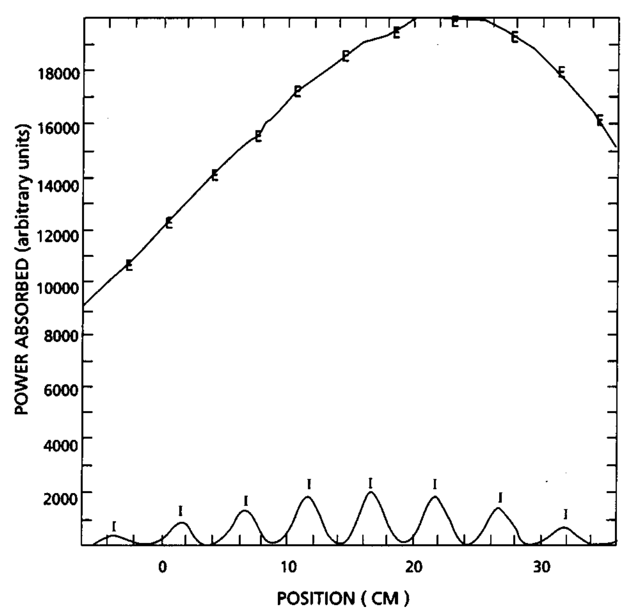

FIG. 11. Same as Fig. 10 except $k_{\|}=8\ \mathrm{m}^{-1}$.

$\sim$0.2 keV, although large values of $k_\parallel \sim 16-20\ \mathrm{m}^{-1}$ are evidently required for significant electron absorption. For these conditions, ion absorption is negligible.

For hotter electrons, $T_{e0} \sim 1$ keV, the electron absorption becomes still stronger and hence the absorption region moves outward, making off-axis current drive feasible but central current drive difficult. For higher ion temperatures, ion damping becomes significant. When T_{i0} is in the range 0.5–1 keV, the ion absorption can become comparable or even much larger than the electron absorption. Nevertheless, conditions can still be achieved where $P_e \gg P_i$ if $T_{e0} > T_{i0}$. The ratio P_e/P_i is very sensitive to the value of $k_\parallel$, making the role of the, inevitable, strong magnetic shear crucial.

We have also shown that it is important to distinguish the kinetic power in order to obtain positive definite values for the power absorbed. For low electron temperatures, electron Landau damping makes the largest contribution to electron dissipation, whereas, for higher temperatures, transit time damping is larger, although electron Landau damping is still comparable. The cross terms always cancel part of the transit time damping and the electron Landau damping.

Clearly, it is important to carry out further calculations using a sheared magnetic field model in order to provide a more realistic assessment of high-harmonic fast wave heating for spherical tokamaks. Such a model will be investigated in future work. The present, very simple model is too crude to provide a definitive conclusion but it does indicate that ion absorption can become important at the rather modest ion temperatures of 0.5–1 keV.

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