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Damping of electron cyclotron waves in dense plasmas of a Compact Ignition Tokamak

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Abstract

Absorption of electromagnetic waves by hot and dense plasmas is investigated in the electron cyclotron range of frequency. It is shown that the strong reduction of the damping of the extraordinary mode, caused by finite Larmor radius effects on waves propagating perpendicularly to the magnetic field, becomes insignificant at large values of the parallel component of the refractive index.

With an appropriate form of the relativistic dispersion relation which includes high order Larmor radius terms, heating of dense plasmas in a Compact Ignition Tokamak is investigated. It is shown that by using the extraordinary mode with oblique propagation and frequency of 190 GHz it is possible to bring to thermonuclear ignition a dense ohmic plasma with a toroidal magnetic field of 105 kG and a central density of $1 \times 10^{15} \text{ cm}^{-3}$.

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28

I. Introduction

Recently,¹⁻⁴ the linear theory of electron cyclotron waves, approximated to the lowest significant order in the Larmor radius, has been used for the study of plasma absorption of extraordinary waves propagating obliquely to the magnetic field with frequency ω significantly smaller than the electron cyclotron frequency ω_c . This type of wave absorption offers several advantages: It minimizes the wave frequency for a given value of magnetic field, which makes possible the use of existing high power sources in present large tokamaks.^{2,4} It maximizes the plasma density for a given value of ω , which makes it suitable for heating high β plasmas. It minimizes trapped particle effects in tokamaks, since the wave is absorbed by passing electrons. Finally, it allows control of the power deposition in both the physical and the velocity space, which can be used for profile control or current drive in tokamaks.

The damping of electron cyclotron waves in hot and dense plasmas, as those considered in Refs. 1-4, can be affected by finite Larmor radius (FLR) effects. It has been shown previously⁵ that plasma absorption of electromagnetic waves propagating perpendicularly to the magnetic field with the extraordinary mode is strongly reduced by FLR effects when the plasma frequency $\omega_p \approx \omega_c \approx \omega$. In this article we address the same issue for arbitrary values of the parallel refractive index $N_{\parallel}$. Section II presents the relativistic dispersion relation of electron cyclotron waves with arbitrary values of $N_{\parallel}$ in hot Maxwellian plasmas. In Sec. III we study the solution of the dispersion equation for the extraordinary mode with frequency around the electron cyclotron frequency. Section IV discusses the issue of auxiliary heating in the Compact Ignition Tokamak⁶ (CIT), and Sec. V presents numerical results on the absorption of electron cyclotron waves by dense CIT plasmas with a toroidal magnetic field of 105 kG and a wave frequency of 190 GHz. Finally, our conclusions are summarized in Sec. VI.

II. The dispersion relation

The relativistic dispersion relation of electron cyclotron waves in a Maxwellian plasma has been discussed in several articles. Here we use a form derived recently,⁷ which is particularly useful for investigating the absorption of waves with oblique propagation.

A power series expansion of the dielectric tensor with terms of up to second order in Larmor radius was derived in Ref. 7. In this paper we add the next significant order terms, which we have found important for

a correct description of the damping of electron cyclotron waves in dense plasmas.

In the orthogonal frame of reference with basis $(\mathbf{e}_1, \mathbf{e}_2, \mathbf{e}_3)$, where the magnetic field $\mathbf{B} = B\mathbf{e}_3$ and the complex index of refraction $\mathbf{N} = N_\perp\mathbf{e}_1 + N_\parallel\mathbf{e}_3$, the general form of the dispersion relation is given by the equation

$$\begin{aligned} D \equiv & \epsilon_{11}N_\perp^4 + 2N_\parallel\epsilon_{13}N_\perp^3 \\ & + N_\perp^2(\epsilon_{13}^2 + N_\parallel^2\epsilon_{11} - \epsilon_{11}\epsilon_{22} - \epsilon_{12}^2 + N_\parallel^2\epsilon_{33} - \epsilon_{11}\epsilon_{33}) \\ & + 2N_\parallel N_\perp(N_\parallel^2\epsilon_{13} - \epsilon_{22}\epsilon_{13} + \epsilon_{12}\epsilon_{23}) + \epsilon_{11}\epsilon_{23}^2 - \epsilon_{22}\epsilon_{13}^2 + 2\epsilon_{12}\epsilon_{13}\epsilon_{23} \\ & + N_\parallel^2(\epsilon_{13}^2 - \epsilon_{23}^2) + N_\parallel^2\epsilon_{33}(N_\parallel^2 - \epsilon_{11} - \epsilon_{22}) + \epsilon_{23}(\epsilon_{12}^2 + \epsilon_{11}\epsilon_{22}) = 0, \quad (1) \end{aligned}$$

where $\epsilon_{ij} = \epsilon'_{ij} + i\epsilon''_{ij}$ are the components of the hot plasma dielectric tensor. The Hermitian part ϵ'_{ij} was given in Ref. 7 and is not presented here. For the anti-Hermitian part one obtains

$$\begin{aligned} \epsilon''_{11} &= a_{11} + N_\perp^2(b_{11} + c_{11}) + N_\perp^4(d_{11} + f_{11} + g_{11}), \\ \epsilon''_{12} &= -i[a_{11} + N_\perp^2(2b_{11} + c_{11}) + N_\perp^4(3d_{11} + 3f_{11}/2 + g_{11})], \\ \epsilon''_{22} &= a_{11} + N_\perp^2(3b_{11} + c_{11}) + N_\perp^4(37d_{11}/5 + 2f_{11} + g_{11}), \\ \epsilon''_{13} &= N_\perp[a_{13} + N_\perp^2(b_{13} + c_{13}) + N_\perp^4(d_{13} + f_{13} + g_{13})], \\ \epsilon''_{23} &= iN_\perp[a_{13} + N_\perp^2(2b_{13} + c_{13}) + N_\perp^4(3d_{13} + 3f_{13}/2 + g_{13})], \\ \epsilon''_{33} &= N_\perp^2[a_{33} + N_\perp^2(b_{33} + c_{33}) + N_\perp^4(d_{33} + f_{33} + g_{33})]. \end{aligned} \quad (2)$$

These expressions contain the contribution of the first (a_{ij}, b_{ij}, d_{ij}) , the second (c_{ij}, f_{ij}) , and the third (g_{ij}) harmonic of the electron cyclotron frequency, and are therefore valid whenever the effect of the fourth harmonic is negligible, i.e., for thermal plasmas with $T_e \leq 50 \text{ keV}$ when $\omega < \omega_c$.

The various terms of Eq. (2) are given in the Appendix. Here we note that every coefficient in these equations contains a factor of type $F(p) = \exp[-\mu(\gamma - 1)]$, with $\mu = mc^2/T_e$, $\gamma = (1 + p^2)^{1/2}$, and p a solution of the resonant equation

$$(1 + p^2)^{1/2} - n\omega_c/\omega = N_\parallel p, \quad (3)$$

where n is the harmonic number. It is this factor that makes the use of the relativistic theory essential for a correct description of electron cyclotron wave absorption even in plasmas with $T_e \ll mc^2$.

III. Numerical results

Our previous results¹⁻⁴ were obtained with a reduced dispersion relation, where only the lowest order terms (a_{ij}) were used in the dielectric tensor. In this section we shall investigate the validity of this approximation.

The imaginary parts of $N_{\perp}$ given by the reduced and the full dispersion relation are shown in Fig. 1 and 2 as a function of $(\omega_p/\omega)^2$ for the extraordinary mode with $\omega_c/\omega = 1.3$, $T_e = 20$ keV, and with $N_{\parallel} = 0$ and $N_{\parallel} = 0.5$, respectively. In agreement with Ref. 5, we find that high order FLR terms cause a reduction of the wave damping.

The ratio of the two results (full/reduced) is shown in Fig. 3 as a function of $(\omega_p/\omega)^2$ for the same plasma parameters and for several values of $N_{\parallel}$. It appears that the contribution of high order FLR terms, which for perpendicular propagation causes a strong reduction of wave damping at large values of $(\omega_p/\omega)^2$, becomes rapidly insignificant as the value of $N_{\parallel}$ increases. This is also shown by Fig. 4, where the ratio is displayed as a function of $N_{\parallel}$ for several values of T_e . From these results it appears that for the case of perpendicular propagation the FLR effect is not very sensitive to the value of electron temperature. It is also a weak function of ω_c/ω , as shown by Fig. 5 where the ratio is plotted for several values of this parameter.

Finally, in agreement with Ref. 5 we have found that in the same range of plasma parameters the contribution of high order FLR terms to the dispersion of ordinary waves is always negligible for $\omega < \omega_c$.

In conclusion, the results of Refs. 1-4, where the plasma absorption of electron cyclotron waves with large frequency downshift and strong oblique propagation was investigated, are only slightly modified by the inclusion of high order FLR terms in the wave dispersion relation. Nevertheless, in applying this heating scheme to reactorlike plasmas, one is forced to consider wave propagation close to perpendicular for the sake of achieving accessibility to the highest values of plasma density which are allowed by a given value of wave frequency. In this case the full dispersion relation, Eq. (2), must be used for a correct description of wave absorption. This is the topic of the remainder of this paper where we investigate the absorption of electron cyclotron waves at downshifted frequency in the dense plasmas of a Compact Ignition Tokamak⁶ (CIT).

IV. Auxiliary Heating in CIT

CIT is a small, high field tokamak experiment designed to study the

physics of α -particle heating.⁶ It is expected that a large toroidal magnetic field (≈ 100 kG) will allow the values of plasma density (10^{15} cm⁻³) and current (10 MA) which are needed for reaching thermonuclear ignition in a small device. Auxiliary heating is expected to be necessary as well, but transport studies⁸ indicate that it will also be necessary either to avoid the deterioration in confinement caused by auxiliary heating (H-mode), or to control very carefully plasma density and temperature profiles.

ICRF has been selected for CIT.⁶ Unfortunately, the use of this type of auxiliary heating under CIT plasma conditions faces two major obstacles:

- (a) the plasma coupling of the fast magnetosonic wave is extremely sensitive to plasma edge conditions which are very difficult to control;
- (b) with ICRF it is difficult to control the spatial distribution of the energy absorbed by the plasma, especially in the presence of a large α -particle component.

In this article we suggest that Electron Cyclotron Heating (ECH) at down shifted frequency be also considered a viable candidate for CIT. The plasma coupling of electron cyclotron waves is not sensitive to plasma edge conditions or to α particles, and, more importantly, with this type of plasma waves it is possible to control the electronic power deposition profile in both the physical and the velocity space, a crucial requirement for profile control in tokamaks.

V. ECH in CIT

The use of a self-consistent magnetic configuration is important for a correct description of ECH in Tokamaks, since the damping of electron cyclotron waves is very sensitive to the amplitude and direction of the magnetic field which is modified by plasma current and pressure.

The CIT magnetic configuration has been modelled by solving the Grad-Shafranov equation for the poloidal magnetic flux ($2\pi\psi$)

$$\nabla \cdot (\nabla\psi/R^2) = -dp/d\psi - (F/R^2)dF/d\psi,$$

with the plasma pressure (p) and the poloidal current flux ($2\pi F$) given by $p = p_1\psi^2 + p_0$, and $F^2 = f_1\psi^2 + f_0$, respectively. This choice of flux surface functions reduces the problem to the solution of an ordinary differential equation.

Figure 6 shows an example of a CIT magnetic configuration with a vacuum toroidal magnetic field $B_o = 105$ kG at $R = 125$ cm, and a volume average beta $\langle \beta \rangle = 0.05$. With a plasma current $I = 8.0$ MA, the value

of the safety factor q is 1.0 on axis and 3.5 on the outmost magnetic surface shown in Fig. 6. If the central electron density $n_e(0) = 1 \times 10^{15} \text{ cm}^{-3}$, we get $T_e(0) = T_i(0) = 24 \text{ keV}$.

The wave trajectory was obtained from the familiar ray equations⁹

$$\begin{aligned} \frac{d\mathbf{r}}{ds} &= - \frac{\partial H / \partial \mathbf{k}}{|\partial H / \partial \mathbf{k}|}, \\ \frac{d\mathbf{k}}{ds} &= \frac{\partial H / \partial \mathbf{r}}{|\partial H / \partial \mathbf{k}|}, \end{aligned} \quad (4)$$

where $\mathbf{r}$ is the spatial coordinate, s is the arc length along the ray, and $\mathbf{k} = N\omega/c$ is the wave vector. The function H in Eq. (4) is the real part of D of Eq. (1), which in the range of electron temperature considered in this paper is well approximated by the cold plasma theory.

The fraction of power which is absorbed by the plasma is given by

$$P(s) = 1 - \exp\left(-2 \int_0^s \mathbf{k}^i \cdot \frac{d\mathbf{r}}{ds}\right),$$

where the integral is along the ray, and $\mathbf{k}^i$ is obtained from the imaginary part of the solution of Eq. (1).

Some examples of ray trajectories are given in Fig. 6. Here waves with the extraordinary mode and a frequency of 190 GHz are launched into a CIT discharge, with the plasma parameters listed above, from two different antennas: one located on top of the torus and the other near the equatorial plane. The rays of Fig. 6, which are labelled with the initial toroidal ($N_t = \sin\alpha$) and poloidal ($N_p = \sin\beta$) component of $\mathbf{N}$, have $\alpha = 12^\circ$.

The fraction of wave energy absorbed inside a magnetic surface, for the case of a Gaussian beam with a full half-intensity width of 5° , is shown in Fig. 7 as a function of the surface volume for the two launching geometries of Fig. 6. In both cases the wave is almost completely absorbed in a single pass. The absorbed power density profile, averaged over the magnetic surface, is shown in Fig. 8 for the case of top launching. The absorbing region is near the center of the torus, even though the central value of the electron cyclotron frequency is more than 50% larger than the wave frequency. Its position can be easily controlled by adjusting the angle β , since by increasing its value the absorbing layer is shifted towards the outside of the torus.

To obtain strong absorption at lower values of T_e one must either increase the value of α (i.e., $N_{||}$) or lower that of B_0 . The first method is

described in Fig. 9 for the case of horizontal launching. Here we assume that the rf power is turned on with a value of $\alpha \approx 65^\circ$, when $n_e(0) = 4.5 \times 10^{14} \text{ cm}^{-3}$ and $T_e(0) = 5 \text{ keV}$. These are plasma conditions that one can safely predict for CIT with only ohmic heating. Then, as T_e increases, the value of α is decreased and that of n_e is increased to the values of Fig. 6 in such a manner that the conditions for total single pass absorption are satisfied, as shown in Fig. 9.

The second method is described in Fig. 10 for the case of top launching. Here we assume that the rf power is turned on with a value of $\alpha \approx 12^\circ$ when $B_0 = 80 \text{ kG}$, $T_e(0) = 4 \text{ keV}$, and $n_e(0) = 7 \times 10^{14} \text{ cm}^{-3}$. Again, these are plasma conditions easily attainable in CIT with only ohmic heating. The extraordinary electron cyclotron wave is completely absorbed by such a plasma, with a well-localized power deposition profile, as shown in Fig. 11. In this heating scenario we assume that, as T_e increases, the values of B_0 and n_e are risen to the final values of Fig. 6 in such a manner that total single pass absorption is achieved with constant value of α , as shown in Fig. 10. The power deposition remains well localized with a profile that changes from the shape of Fig. 11 to that of Fig. 8.

In both heating schemes the wave is absorbed by suprathermal electrons. Their energy, given by Eq. (3), is $\approx 5 T_e$ and, in the range of plasma parameters considered for CIT, it is much smaller than the Dreicer critical energy. Since the relaxation time of these resonant electrons is also much smaller than the plasma confinement time in CIT, the absorption of wave energy results in plasma bulk heating, which is not much different from ohmic heating. Note also that, since the wave energy is released to suprathermal electrons, this type of electron heating is also suited for current profile control.

Of the two heating schemes described above, the second (Fig. 10) has two major advantages. The first is that, by starting the discharge during a rising B_0 , one extends the burn time, inevitably very short in any compact device. The second advantage is more important and has to do with the initial polarization of the wave which, to couple to the extraordinary mode, must possess the right degree of ellipticity. While in the first heating scenario (Fig. 9) the initial polarization changes drastically during the heating pulse, from close to circular to almost linear, in the second scenario it remains nearly constant and, more importantly, close to linear. In this case, if a linearly polarized wave is launched with the electric field perpendicular to the local magnetic field, less than 5% of the wave energy is coupled to the ordinary mode, which is not completely absorbed

in a single pass. This simplifies enormously the launching structure. We can envision a simple antenna, where by bouncing the wave on a plane mirror whose position can be changed during the heating pulse, one can control the electron distribution by simply adjusting the angles α and β .

In summary, the plasma absorption of extraordinary waves with a frequency of 190 GHz can be used for auxiliary heating of CIT with a toroidal magnetic field of 105 kG. This type of plasma heating leads to a very simple method of profile control, where by adjusting the wave's initial direction one can control the electron power deposition in both the physical and the velocity space.

VI. Conclusion

In conclusion, we have investigated the damping of electron cyclotron waves in dense and hot plasmas by taking into account high order Larmor radius terms in the wave dispersion relation. We have shown that the strong reduction caused by FLR effects on the damping of extraordinary waves with perpendicular propagation and frequency around the electron cyclotron frequency becomes insignificant at large values of the parallel component of the refractive index.

We have shown that ECH at downshifted frequency can be successfully used for auxiliary heating of CIT. A simple heating scenario has been described where a linearly polarized wave is launched from the top of the plasma torus. It has been shown that with a wave frequency of 190 GHz it is possible to bring to thermonuclear ignition a dense, ohmic CIT plasma with a toroidal magnetic field of 105 kG and a central density of $1 \times 10^{15} \text{ cm}^{-3}$. One of the many advantages offered by this heating scheme is the control of the power deposition profile which is achieved by simply adjusting the wave's launching direction.

These results prove that ECH at downshifted frequency constitutes a feasible alternative to ICRF for auxiliary heating in CIT. Wave sources with frequency of 190 GHz are not at present commercially available. Nevertheless, recent impressive advances in the development of both gyrotrons¹⁰ and free-electron lasers¹¹ give confidence on a favorable solution of this technological issue.

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Appendix: dielectric tensor coefficients

The coefficients on the r.h.s. of Eq. (2) are given by:

$$a_{11} = aRS \sum_{s=\pm 1} (1 + s/\xi) F(p_s) ,$$

$$a_{13} = aRS(\omega/\omega_c) \sum_{s=\pm 1} [p_s + s(p_s + 2sq)/\xi + 3sq/\xi^2] F(p_s) ,$$

$$a_{33} = aRS(\omega/\omega_c)^2 \sum_{s=\pm 1} [p_s^2 + sp_s(p_s + 4sq)/\xi + 6sq(p_s + sq)/\xi^2 + 12sq^2/\xi^3] F(p_s) ,$$

$$b_{11} = aRS(R/N_{||}\mu)(\omega/\omega_c)^2 \sum_{s=\pm 1} s(1 + 3s/\xi + 3/\xi^2) F(p_s) ,$$

$$b_{13} = aRS(R/N_{||}\mu)(\omega/\omega_c)^3 \sum_{s=\pm 1} s[p_s + 3s(p_s + sq)/\xi + 3(p_s + 4sq)/\xi^2 + 15q/\xi^3] F(p_s) ,$$

$$b_{33} = aRS(R/N_{||}\mu)(\omega/\omega_c)^4 \sum_{s=\pm 1} s[p_s^2 + 3sp_s(p_s + 2sq)/\xi + 3(p_s^2 + 8sqp_s + 4q^2)/\xi^2 + 30q(p_s + 2sq)/\xi^3 + 90q^2/\xi^4] F(p_s) ,$$

$$d_{11} = aRS(R/N_{||}\mu)^2(5/8)(\omega/\omega_c)^4 \sum_{s=\pm 1} (1 + 6s/\xi + 15/\xi^2 + 15s/\xi^3) F(p_s) ,$$

$$d_{13} = aRS(R/N_{||}\mu)^2(5/8)(\omega/\omega_c)^5 \sum_{s=\pm 1} [p_s(1 + 6s/\xi + 15/\xi^2 + 15s/\xi^3) + 2(q/\xi)(2 + 15s/\xi + 45/\xi^2 + 105s/2\xi^3)] F(p_s) ,$$

$$d_{33} = aRS(R/N_{||}\mu)^2(5/8)(\omega/\omega_c)^6 \sum_{s=\pm 1} [p_s^2(1 + 6s/\xi + 15/\xi^2 + 15s/\xi^3) + 4(qp_s/\xi)(2 + 15s/\xi + 45/\xi^2 + 105s/2\xi^3) + 10(q/\xi)^2(2 + 18s/\xi + 63/\xi^2 + 84s/\xi^3)] F(p_s) ,$$

$$c_{ij} = -4b_{ij}(\omega_c \Rightarrow 2\omega_c) ,$$

$$f_{ij} = -(128/5)d_{ij}(\omega_c \Rightarrow 2\omega_c) ,$$

$$g_{ij} = (243/5)d_{ij}(\omega_c \Rightarrow 3\omega_c) ,$$

where the symbol $b_{ij}(\omega_c \Rightarrow 2\omega_c)$ means that ω_c must be replaced by $2\omega_c$ in b_{ij} . The various quantities in these expressions are:

$a = \pi(\omega_p/\omega)^2/4N_{||}^2 K_2(\mu) \exp(\mu)$, $\mu = mc^2/T_e$, $K_2(\mu)$ is the McDonald function, $q = R/(1 - N_{||}^2)$, $\xi = \mu N_{||} R/(1 - N_{||}^2)$, $R = [(\omega_c/\omega)^2 - 1 + N_{||}^2]^{1/2}$,

S is a step function, i.e., $S = 1$ for $R^2 \geq 0$ and $S = 0$ for $R^2 < 0$,
 $F(p_s) = \exp[-\mu(\gamma_s - 1)]$, $\gamma_s = (1 + p_s^2)^{1/2}$, and p_s is a solution of

$$(1 + p_s^2)^{1/2} - \omega_c/\omega = N_{||}p_s ,$$

i.e., $p_s = [N_{||}(\omega_c/\omega) + sR]/(1 - N_{||}^2)$.

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Fig. 1 Imaginary part of $N_{\perp}$ with (solid) and without (dash) high order FLR terms as a function of $(\omega_p/\omega)^2$ for $\omega_c/\omega = 1.3$, $T_e = 20 \text{ keV}$, and $N_{\parallel} = 0$.

Fig. 2 Same as Fig. 1 with $N_{\parallel} = 0.5$.

Fig. 3 Ratio of the imaginary parts of $N_{\perp}$ obtained with and without the high order FLR terms as a function of $(\omega_p/\omega)^2$ for $\omega_c/\omega = 1.3$, $T_e = 20$ keV, and several values of $N_{||}$.

Fig. 4 Same as Fig. 3 as a function of N_{II} for $\omega_c/\omega = 1.2$, $(\omega_p/\omega)^2 = 0.5$, and several values of T_e .

Fig. 5 Same as Fig. 3 for $\omega_c/\omega = 1.1$, $T_e = 20 \text{ keV}$, $N_{||} = 0$, and several values of ω_c/ω .

Fig. 6 CIT magnetic configuration with $B_o = 105 \text{ kG}$, $I = 8 \text{ MA}$, $n_e(0) = 1 \times 10^{15} \text{ cm}^{-3}$, $T_e(0) = 24 \text{ keV}$, $\langle \beta \rangle = 0.05$. Shown are also some ray trajectories for the extraordinary mode with $f = 190 \text{ GHz}$, and $\alpha = 12^\circ$.

Fig. 7 Fraction of wave energy absorbed inside a magnetic surface as a function of the normalized surface volume for the two launching cases of Fig. 6; solid-top launching, dash-side launching.

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Fig. 8 Absorbed power density profile (surface average in arbitrary units) versus the magnetic surface volume for the case of top launching of Fig. 6.

Fig. 9 Values of α (solid) and $T_e(0)$ (dash) as a function of $n_e(0)$ for complete single path absorption in the case of side launching with $B_0 = 105 \text{ kG}$.

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Fig.10 Values of $n_e(0)$ (solid) and $T_e(0)$ (dash) as a function of B_0 for complete single path absorption in the case of top launching with $\alpha \approx 12^\circ$.

Fig.11 Same as Fig. 8 for $\alpha \approx 12^\circ$, $n_e(0) = 7 \times 10^{14} \text{ cm}^{-3}$, $T_e(0) = 4 \text{ keV}$, and $B_0 = 80 \text{ kG}$.

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