Absorption of High-frequency Electromagnetic Energy in a High-temperature Plasma

By R. S. Sagdeyev and V. D. Shafranov

In a plasma of high temperature, the collision frequency is very low and there is a decrease in the effectiveness of known heating mechanisms which are based on collisions. The problem then arises of noncollisional plasma heating with high-frequency electromagnetic fields. By the term high-frequency field, we mean electromagnetic oscillations whose period is much smaller than the characteristic confinement times of the hot plasma.

The interaction of a high-frequency field with a plasma can also be used as a probe for studying its properties.

The solution of these two theoretical problems is associated with an investigation of the penetration of electromagnetic waves into a plasma, i.e., with an analysis of the electromagnetic absorption spectrum caused by various mechanisms and with the investigation of the role which these mechanisms play in heating of plasma.

In this paper an analysis of the cyclotron and Cherenkov mechanisms is given. These are two fundamental mechanisms for noncollisional absorption of electromagnetic radiation by a plasma in a magnetic field. The expressions for the dielectric permeability tensor, \( \varepsilon_{ab} \), for a plasma with a nonisotropic temperature distribution in a magnetic field, are obtained by integrating the kinetic equation with Lagrangian particle co-ordinates in a form suitable to allow a comprehensive physical interpretation of the absorption mechanisms.

Using the tensor \( \varepsilon_{ab} \) and taking into account the thermal motion, the oscillations of a plasma column stabilized by a longitudinal field have been analyzed. For a uniform plasma, the frequency spectrum has been obtained together with the direction of electromagnetic wave propagation when both the cyclotron and Cherenkov absorption mechanisms take place. The influence of nonlinear effects on the electromagnetic wave absorption and the part which cyclotron and Cherenkov absorption play in plasma heating have also been investigated.

Any mechanism responsible for the absorption of electromagnetic energy involves the accumulation of energy by a charge in an electric field and the subsequent redistribution of this energy among the various degrees of freedom in the random motion. Absorption mechanisms can be physically classified as noncollisional or resonant processes and collisional or nonresonant processes.

A continuous accumulation of energy by a charge is possible in a static or quasi-stationary field if there is no dissipation. But if the charge is in a high-frequency field, it is continually accumulating and losing energy alternately. If plasma heating is to be effective, it is necessary that the characteristic collision time required for energy dissipation be comparable with the oscillation period. The corresponding absorption mechanism, which is associated with frequent collisions, will be called here the nonresonant or collisional mechanism. It includes ordinary Joule heating as well as betatron or gyrorelaxational heating proposed by Budker in 1951 and independently by Schlüter.

If the collision time is much longer than the oscillation period, the continuous effective energy accumulation by a charged particle will be possible only under resonance conditions when the direction of the electric field coincides with the direction of increasing velocity. This mechanism will be called resonant or noncollisional absorption.

In principle, resonances are possible in the following cases.

(A) If a charge performs a periodic motion there will be a resonance when the field frequency coincides with the motional frequency.

(B) For a charge in a magnetic field there is an ion or electron cyclotron resonance. If the particle velocity along a line of force is \( v_z \), the frequency spectrum of the absorbed radiation is determined by the Doppler effect according to the relation

\[
\omega - k_z v_z = \omega_{ce},
\]

where \( k_z \) is the \( z \) component of the wave vector, \( \omega_{ce} = eH/mc \) and the other symbols have their usual meaning.

(C) Resonance can also take place in the absence of a periodic motion provided that the charged particle

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velocity is equal to the phase velocity of the electromagnetic wave. This is, of course, possible only for slow waves.

(D) If there is no magnetic field, the resonant accumulation of energy by charged particles takes place only for longitudinal electromagnetic plasma oscillations. The corresponding absorption caused by electron oscillations is known as Landau damping. Damping of a similar type takes place when there are ion (sonic) oscillations.

(E) For a sufficiently dense plasma in a magnetic field, resonance for the transverse branch of the electromagnetic oscillations is also possible. The index of refraction is of the order \( n^2 \approx e^2 / m \), where \( c_\text{e} \). The ratio of the phase velocity to the thermal velocity are of the same order. Under these conditions, the inverse Cherenkov absorption mechanism becomes substantial. A charge moving along a magnetic field is in resonance with the wave field where the following condition is satisfied

\[ \omega - k z v_z = 0. \]  

If \( \omega/k < c \), this condition, as well as condition (1), determines a velocity of the charges which will absorb energy.

Because of the spread of thermal velocities, the frequency range for which absorption takes place may become very broad for both cyclotron and Cherenkov absorption.

There is one fundamental difference between collisional and non-collisional heating. In the first case, when collisions are important, the velocity distribution during the heating process is nearly Maxwellian while the temperature is rising. In the second case, when only the charges moving in resonance accumulate energy, the velocity distribution during heating will differ greatly from the Maxwellian distribution. In this case, the term heating will mean an increase in the average random energy of the electrons and ions.

The interaction of an electromagnetic field with a plasma is determined, in principle, by the Maxwell self-consistent equations together with the kinetic equations for the electron and ion plasma. Under the assumption that the velocity associated with the regular motion of the charges in a wave field is smaller than the average thermal velocity, these equations are linear in the field quantities. In this case, many problems concerning plasma oscillations can be suitably solved by using the dielectric permeability tensor \( \varepsilon \), whose antherimetric part \( (\varepsilon_z + \varepsilon_x - \varepsilon_y)/2 \) determines the absorption. When there are no collisions, the general form of the dielectric tensor can be expressed explicitly provided that the kinetic equations are integrated in Lagrangian coordinates. The correction to the distribution function which one obtains in this way is given by

\[ f_i = -e \int_{-\infty}^{\infty} \left\{ \mathbf{E}(x(t'), t') + c e \mathbf{v}(t') \times \mathbf{H}(x(t'), t') \right\} \times \partial_0^i \mathbf{p}(t') \, dt'. \]  

If the fields vary harmonically in space and time, i.e., as \( e^{i(k \cdot x - \omega t)} \), the corresponding tensor \( \varepsilon_{a\beta} \) will be

\[ \varepsilon_{a\beta} = \delta_{a\beta} + (i/2) \int_{-\infty}^{\infty} K_{a\beta}(t) \, dt, \]  

with real and imaginary parts

\[ \varepsilon_{a\beta}' = \delta_{a\beta} + \left( \frac{i}{2} \right) \int_{-\infty}^{\infty} K_{a\beta}(t) \, dt, \]  

\[ \varepsilon_{a\beta}'' = \frac{i}{2} \int_{-\infty}^{\infty} K_{a\beta}(t) \, dt. \]

The summation is carried out with respect to the different species of charged particles. The function \( K_{a\beta}(t) \) for each species, is

\[ K_{a\beta} = \frac{4\pi e^2}{m} \int v_a(t) \left( 1 - \frac{\mathbf{v} \cdot \mathbf{v}}{\beta} \right) \times \mathbf{E} \mathbf{v} \frac{\partial E}{\partial \mathbf{v}} \]  

\[ \times \left[ \exp \left( i \omega t - \mathbf{k} \cdot \mathbf{v}(t') \right) \right], \]  

where \( \mathbf{v}(t) \) is the particle velocity in the unperturbed plasma expressed in terms of the initial velocity \( \mathbf{v} \) and the time \( t \). The integration is carried out with respect to the initial momenta. For the particular case of the Maxwellian distribution, the expression in brackets reduces to the simple expression \( \mathbf{v} \times \mathbf{E} \).

For a plasma with no collisions in a uniform magnetic field directed along the \( z \) axis, we have

\[ v_{x}(t) = v_{x} \cos \omega_0 t - v_{y} \sin \omega_0 t \]
\[ v_{y}(t) = v_{x} \sin \omega_0 t + v_{y} \cos \omega_0 t \]
\[ v_{z}(t) = v_{z}. \]

The quantity in the exponent of (5) can then be expressed in terms of (6)

\[ \omega t - \mathbf{k} \cdot \mathbf{v} \, dt' = (\omega - k z v_z) t - (k z v_z / \omega_0) \sin \omega_0 t \]
\[ + (k z v_z / \omega_0) (1 - \cos \omega_0 t), \]  

where \( \mathbf{k} = (k_z, 0, k_z). \)

If the unperturbed distribution function, \( f_0 \), is a product of Maxwellian distributions with different temperatures corresponding to particle motions parallel and perpendicular to \( \mathbf{H}_0 (\mathbf{T}_1 \text{ and } T_2) \), the x, y components of the expression in the brackets of (5) will be

\[ v_{x,y} \left[ \frac{1}{T_\perp} - \frac{k z v_z}{\omega_0 (T_\perp - 1)} \frac{1}{T_1} \right] f_0 \]

and the z component is

\[ v_z \left[ \frac{1}{T_\parallel} + \frac{k z v_z}{\omega_0 (T_\parallel - 1)} \frac{1}{T_1} \right] f_0. \]
By substituting these expressions, together with (6) and (7), into (5) and (4), closed formulae are readily obtained for the tensor $e_{a\beta}$. In a strong magnetic field and for relatively low frequencies, $\omega \lesssim \omega_{ce}$, the expression for $e_{\alpha\beta}$ is greatly simplified by expanding the exponent in powers of $k v_\perp / \omega_{ce}$. In the zero-order approximation, we have

$$
\frac{\varepsilon}{\omega} = 1 + i \sum 2 \left( \frac{\omega_0}{\omega} \right)^2 \left[ 1 - \frac{T_{\perp}}{T_1} - \frac{\omega H}{\omega} \right] \frac{\delta_+ (\omega - \omega_0 - k v_\perp)}{k v_\perp} - \left( \frac{\omega_0}{\omega} \right)^2 \left[ 1 - \frac{T_{\perp}}{T_1} - \frac{\omega H}{\omega} \right] \frac{\delta_+ (\omega - \omega_0 + k v_\perp)}{k v_\perp},
$$

$$
\frac{g}{\omega} = 1 + i \sum 2 \left( \frac{\omega_0}{\omega} \right)^2 \left[ 1 - \frac{T_{\perp}}{T_1} - \frac{\omega H}{\omega} \right] \frac{\delta_+ (\omega - \omega_0 - k v_\perp)}{k v_\perp} - \left( \frac{\omega_0}{\omega} \right)^2 \left[ 1 - \frac{T_{\perp}}{T_1} - \frac{\omega H}{\omega} \right] \frac{\delta_+ (\omega - \omega_0 + k v_\perp)}{k v_\perp},
$$

$$
\eta = 1 + \sum 2 \left( \frac{\omega_0}{k v_\perp} \right)^2 \frac{i \omega H}{k v_\perp} W \left( \frac{\omega - \omega_0}{k v_\perp} \right) \frac{\delta_+ (\omega - \omega_0 - k v_\perp)}{k v_\perp} + 1,
$$

where $v_\perp = (2T_1/m)^{1/2}$ and

$$
W(z) = e^{-z^2} (1 + 2iz - z^2 dz)
$$

is the probability integral with a complex argument. The approximation $k v_\perp / \omega_{ce} \ll 1$ made in (10) is valid provided that the Larmor radius is small compared with the characteristic transverse length in the problem. In this approximation it is easy to obtain a solution to several problems associated with plasma oscillations in a uniform magnetic field even when the extent of the plasma is limited in the direction perpendicular to the field.

Let us consider the oscillations in a pinched plasma stabilized by a longitudinal magnetic field, taking into account the thermal motion of the charged particles. The plasma pressure can be counterbalanced by the magnetic pressure of the longitudinal and azimuthal external fields. The azimuthal field is produced by a longitudinal current flowing on the surface of the plasma. The oscillations are taken to be of the form $\exp(i k z + m \phi - \omega t)$. The basic equations are the Maxwell equations

$$
\nabla (\nabla \cdot \mathbf{D}) = -\nabla (\omega / c)^2 \mathbf{B}
$$

where $D_e = e E_z + i g E_A$, $D_\phi = -i g E_r + e E_A$, and $D_e = \eta E_z$. The boundary conditions are:

1. The tangential components of the electric field on the surface of the pinch are continuous and vanish on the surface of an infinite conductivity (metal) cylinder which surrounds the plasma column.

2. The normal component of the pressure tensor is continuous at the surface of the perturbed plasma cylinder.

Solutions of the equations for the electric field are expressed in terms of Bessel functions whose argument is $k v_\perp r$, where
\[ x_1 z^2 = \left( \eta \left( \omega^2 \right) - k^2 \right) + i g (i g - \lambda_2) (\omega/c)^2 / \epsilon \]
and the variables \( \lambda_1 \) and \( \lambda_2 \) are roots of the equation
\[ k^2 + i g \eta^2 + [ (\eta/c)^2 - \epsilon] (\eta - \epsilon)] / (\theta \epsilon)^{-1} + [ (\eta/c)^2 + \eta - \epsilon] = 0. \]

Neglecting the thermal motion with \( k = 0 \), a dispersion equation first deduced by Kôrper\(^4\) can be obtained.

Let us now investigate cyclotron absorption by ions. For frequencies satisfying the condition
\[ \frac{\omega}{v_A}, \frac{\omega}{c} \gg \omega \gg \frac{\omega}{v_A}, \frac{\omega}{c}, \]
(where \( v_A \) is the Alfvén velocity), an expansion in \( \eta^{-1} \) may be made. If the plasma pressure at the pinch boundary is also neglected in comparison with the magnetic pressure, the dispersion equation can be written as follows for the \( m = 0 \) mode
\[ \frac{\omega \alpha^2 I_0(\alpha a)}{I_1(\alpha a)} = 1 - \frac{\alpha^2 \alpha^2}{k_0^2 a^2} k_0^2 (k_0 a) I_1 (k_0 a) - I_1 (k_0 a) I_1 (k_0 a) \]
where
\[ \alpha \equiv (H_2)/H_0(a), \quad k_0 \equiv (H_2)/H_0(a), \]
and \( \alpha^2 = \{ (a^2 - (\omega/c)^2)^2 - (\omega/c)^2 \} / ((\omega/c)^2 e - k^2 \}, \]
and \( a, b \) are the radii of unperturbed pinch and external conductor, respectively. With \( H_2 = 0, b = \infty, (H_2)_{\infty} = (H_2)_{\infty} \), and neglecting the thermal motion in \( e \) and \( g \), this formula agrees with a result deduced by Stix\(^5\) for the oscillations of a cold plasma cylinder in a longitudinal magnetic field.

In contrast to other works,\(^4,5\) the influence of the thermal motion on noncollisional cyclotron damping is taken into account. The dispersion equation permits us to find the imaginary part of the frequency which determines the rate of cyclotron energy absorption by the plasma. Then, knowing the expressions for the electromagnetic field, one can determine the amount of energy absorbed by the plasma.

There is no Cherenkov absorption in Eq. (12). To take this into account, the higher-order terms in the \( \eta^{-1} \) and \( k_i v_i/\omega H \) expansions should be evaluated. The Cherenkov absorption of low-frequency waves in a pinched plasma is discussed elsewhere.\(^6\)

The analysis of the complete dispersion equation for a plasma cylinder is a rather complicated procedure. It is of interest to investigate first the absorption spectrum for the simpler case of a uniform plasma. In this case, the wave damping is determined by the imaginary part of the index of refraction, \( N = p(1 + i\epsilon) \). With the above assumption, namely \( k_i v_i/\omega_{pe} \ll 1 \), the indices of refraction of both the extraordinary and ordinary waves, \( N_e \) and \( N_p \), are expressed in terms of \( g_{ep} \) by the same formulae which were employed when thermal motion was not taken into account. Therefore, for the propagation across the magnetic field of a linearly polarized wave with its electric vector perpendicular to the constant magnetic field, one finds \( N_e^2 = (\epsilon^2 - g_{ep}^2) / \epsilon \); while \( N_p^2 = \eta \) for waves with the electric vector along the direction of the magnetic field.

In the case of transverse wave propagation, \( k_p = 0 \), there is no absorption and the formulae for \( g_{ep} \) take the same form as when the thermal motion is not taken into account. However, even for purely transverse wave propagation, there exist absorption regions in the plasma that are beyond the range of validity of the assumptions made in order to obtain the expression (10) for the dielectric tensor. For example, the calculation of the relativistic Doppler effect shows that there is cyclotron absorption with higher harmonics.

For frequencies substantially greater than \( \omega_{pe} \) and in a frequency range where \( Re N \gg 1 \), Cherenkov absorption will also occur when the thermal motion of charged particles across the magnetic field is taken into account. Formally, this result follows from the fact that with \( \omega \gg \omega_{pe} \) it can be assumed in Eq. (7) that \( \sin \omega_{00} = \omega_{00} \), and, as a consequence, one obtains \( \omega = k_{H}^2 \), for the resonant frequency. This is just the condition for Cherenkov radiation and the corresponding Cherenkov absorption. For waves propagating in the direction of the magnetic field the refraction index squared is \( N_e^2 = \epsilon + g \) and, therefore, has no component \( \eta \) which can give rise to Cherenkov absorption. The absence of Cherenkov absorption for both longitudinal and transverse waves when \( \omega \leq \omega_{pe} \) is connected with the fact that according to the theory of Cherenkov radiation, there can be no (Cherenkov) radiation in these directions from the motion of a charged particle along a line of force. For frequencies \( \omega \gg \omega_{pe} \) the conception that the charge is "attached" to a line of force is not valid and, therefore, transverse Cherenkov radiation is also possible.

Let us now determine the absorption edges where \( \Im N \ll \Re N \) and \( \Im \epsilon \ll \Re \epsilon \); then one can make an expansion in terms of the imaginary parts of \( \epsilon, \eta \) and \( \omega \) and obtain the index of refraction in the form \( N = p(1 + i\epsilon) \). For longitudinal wave propagation, when \( \alpha \ll 1 \), the cyclotron absorption by electrons is described by the following expressions\(^2\)
\[ p^2 = \frac{2 \epsilon x}{v_{1 p}^2} \left( \frac{\omega_{pe}}{\omega} \right)^2 e - \left( \theta \omega c \right) / (\theta \omega c)^2 \]
\[ x = \frac{2 \pi \epsilon}{2 f^2 v_{1 p}^2} \left( \frac{\omega_{pe}}{\omega} \right)^2 e - \left( \theta \omega c \right) / (\theta \omega c)^2, \]
where
\[ \tilde{\omega} \equiv (\omega - \omega_{pe}) / \omega. \quad (13) \]

For oblique wave propagation, the absorption edge is caused by Cherenkov absorption by electrons with \( \omega \ll \omega_{pe} \), where the cyclotron resonance absorption is no longer significant. With the assumption that
\[ \omega_{pe} / \omega_{ce}^2 \gg 1 \]
and
\[ \sin^2 \theta / \cos^2 \theta \ll 4 \omega_{pe}^4 / \omega_{ce}^2 \omega^5, \]
which have been studied by Braginsky and Kasantzev.* Unity can be dropped from the expression for $N^2$ in sufficiently dense plasmas when $(\omega_0/\omega_p)^2 \gg 1$; the absorption then depends only on the ratio of the electron or ion pressure to the magnetic pressure.

$$N^2 = \frac{\omega_0^2}{\omega_0^2 - \omega^2} \times \left[ 1 + \frac{\pi^4}{2} \left( \frac{c}{v_1} \right)^3 \frac{1}{\omega_0^2} \frac{\omega^2}{\omega_0^2 - \omega^2} \times \frac{\sin^2 \theta}{\cos \theta} \exp \left( - \frac{\omega_0 \cos \theta}{\omega_0^2 \cos \theta - \omega} \right) \left( \frac{c}{v_1} \right)^3 \right].$$

For ion cyclotron resonance, when $|\theta - \pi/2| \gg (n/M)$, we have the relations

$$\beta_i^2 = \frac{\omega_i^2}{\omega_0^2 - \omega^2} = 2\omega_0 \cos \theta$$

$$\nu = \frac{\pi^2}{4} \left( \frac{\omega_i^2}{\omega} \right)^2 \frac{c}{v_1} \beta_i^2 \cos \theta \left[ \exp \left( - \frac{\omega_i \cos \theta}{\beta_i \omega} \right) \right]$$

$$\times \left( 1 - \cos \theta \right) \left[ \omega_i^2 \sin^2 \theta + 4\omega_i^2 \cos^2 \theta \right]^{1/2} \left[ \omega_i \sin^4 \theta + 4\omega_i \cos^2 \theta \right]^{1/2}. \quad (15)$$

For oblique wave propagation in the region $\omega \ll (\omega_0/\omega_p)$, there are also noncollisional absorptions which have been studied by Braginsky and Kasantzev.* Unity can be dropped from the expression for $N^2$ in sufficiently dense plasmas when $(\omega_0/\omega_p)^2 \gg 1$; the absorption then depends only on the ratio of the electron or ion pressure to the magnetic pressure.

$$\beta_0 \approx 8\pi n T_e^2 / \hbar^2$$

$$\beta_i \approx 8\pi n T_i^2 / \hbar^2,$$

for $n_e = n_i = n$. \quad (16)

If we now introduce the variables $x = \omega/\omega_H$ and $\alpha = [(\omega - \omega_H)/(\omega - x^2)]/c/v_1$, then Eq. (13) has the form

$$\frac{1}{\pi} \left| 1 - x^2 \right|^2 \int_0^\infty e^{-i\alpha \omega} \left( 1 - x^2 \right) x e^{-\alpha \omega} d\omega / \alpha \omega.$$
frequency, \( \omega_{He} = eH/m_e c \), when the ions can be regarded as being at rest. In the kinetic equation for electrons we shall neglect the term containing the magnetic field. This approximation can be partially justified by the fact that in the ordinary linear approximation, with an unperturbed isotropic Maxwellian distribution, a corresponding term is absent; namely, \( \mathbf{v} \times \mathbf{B} \partial f_0/\partial \mathbf{v} = 0 \) with \( f_0(v) = f_0(v^2) \). If all the variables depend only on one spatial coordinate directed along the constant magnetic field, then a solution to the problem with the initial conditions \( f(\mathbf{v}, z, 0) = f_0(v_z^2 + v_y^2, v_x^2), E_x(z, 0) = E^0_x(z), E_y(z, 0) = E^0_y(z), E_z(z, 0) = 0 \) can be obtained for transverse waves by “integrating along trajectories”. The appropriate kinetic equation and wave equation are

\[
\frac{\partial f}{\partial t} + \mathbf{v} \cdot \frac{\partial f}{\partial \mathbf{v}} + \left( \frac{eE}{mc} \right) \frac{\partial f}{\partial \mathbf{v}} = 0
\]

and

\[
\nabla \times (\nabla \times \mathbf{E}) = \left( \frac{1}{c^2} \right) \frac{\partial^2 \mathbf{E}}{\partial t^2} = \left( 4\pi e^2 \right) \delta (\mathbf{v} / \partial \mathbf{v}) / \partial t. \tag{18}
\]

The following result for the distribution function is obtained from these equations

\[
f(v_x, v_y, v_z, z, t) = f_0([v_x - v_1(t, z)]^2 + [v_y - v_2(t, z)]^2, v_z^2)
\]

where

\[
v_1(t, z) = \frac{e}{m} \cos \omega_0 t \int_0^t \left[ E_x(t', z') \cos \omega_0 t' - E_y(t', z') \sin \omega_0 t' \right] dt' - \frac{e}{m} \sin \omega_0 t \int_0^t \left[ E_y(t', z') \cos \omega_0 t' + E_x(t', z') \sin \omega_0 t' \right] dt',
\]

\[
z' = z - v_z(t - t')
\]

and \( v_2(t, z) \) has a similar form.

Equation (18) is linear in the field quantities and can be easily solved by the Fourier-Laplace method when the longitudinal velocities have a Maxwell distribution initially. From (19) and (20) one can also see that the average random thermal energy will not increase during the wave damping, i.e., the field energy will appear in the kinetic energy of the ordered motion of the plasma electrons. This is because the electrons which receive energy from the high-frequency field are always in phase with the field and, therefore, with one another. As a result, the thermal distribution of electron velocities associated with the distribution of the phases of their Larmor rotation will remain unchanged. To transform the energy of regular electron rotation into energy of random motion, it is necessary that the phases be “mixed”. In addition to phase mixing due to the extremely rare collisions in a high-temperature plasma, randomization of phases can be caused by the nonuniformity of the magnetic field along a line of force. This can occur because the electrons rotating at different places and at different frequencies, \( \omega_{He} \), have their phases mixed due to the thermal motion along the line of force.

In the above case of cyclotron resonance the energy is absorbed by charged particles moving along a constant field with a velocity \( v_z = (\omega_{He} - \omega)/k_z \). Here, the energy in the transverse motion increases and, therefore, a change takes place mainly in the distribution of transverse velocities. Since the number of particles which are in resonance is determined by the longitudinal velocity distribution (which is almost unchanged), the absorption takes place continuously.

The situation is different with Cherenkov absorption since during the process of absorbing energy there is a change in the longitudinal velocity distribution which determines the resonance of the charged particles. Here the particles are brought out of resonance and a balance may occur between the absorption and emission of radiation for certain particles. Mathematically, this is expressed by the vanishing of the derivative \( \partial f_0/\partial v_z \) in the velocity range for which the resonance \( v_z = \omega/k_z \) takes place.

The problem of noncollisional absorption influenced by nonlinear effects is a subject under further investigation.

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