Anomalous properties of hot dense nonequilibrium plasmas

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Abstract.
A concise overview of a number of anomalous properties of hot dense nonequilibrium plasmas is given. The possibility of quasistationary megagauss magnetic field generation due to Weibel instability is discussed for plasmas created in atom tunnel ionization. The collisionless absorption and reflection of a test electromagnetic wave normally impinging on the plasma with two-temperature bi-maxwellian electron velocity distribution function are studied. Due to the wave magnetic field influence on the electron kinetics in the skin layer the wave absorption and reflection significantly depend on the degree of the electron temperature anisotropy. The linearly polarized impinging wave during reflection transforms into an elliptically polarized one. The problem of transmission of an ultrashort laser pulse through a layer of dense plasma, formed as a result of ionization of a thin foil, is considered. It is shown that the strong photoelectron distribution anisotropy yields an anomalous penetration of the wave field through the foil.

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1. Introduction
Very intense, ultrashort laser pulses interacting with solid surfaces permit to investigate the properties of hot dense plasmas in strongly nonequilibrium states. Thanks to the shortness of the laser pulse, it is possible to create the conditions, when the newly formed plasma has relatively sharp boundary and the influence of electron-electron and electron-ion collisions is negligible to a large extent due to the high electron kinetic energy. Another important issue concerning plasmas created by a powerful ultrashort ionizing laser pulse is the following. As a result of ionization, a photoelectron distribution function with large anisotropy in the velocity space is created. Such plasmas are formed both in the regime of tunnel ionization and in the regime of barrier suppression [1, 2, 3, 4, 5] and may exhibit unusual new properties. Some of them are considered in the present communication. Anisotropic electron distribution functions (EDF) are responsible for the appearance of new features in an entire series of plasma phenomena, as well as the cause of other ones. Among the latter, one of the most interesting is the Weibel instability (see, for instance, [6, 7, 8]). The quasistationary magnetic field generation due to such instability is discussed in Section 2. The peculiar features of the collisionless absorption and
reflection are considered in Sections 3-5. In Section 3 we report the basic relations describing the components of the surface impedance of an anisotropic plasma, the absorption and reflection coefficients, and the phase-shift, measuring the degree of transformation of a linearly polarized incident wave into an elliptically polarized reflected wave. In Section 4 analytical and numerical results for the collisionless absorption coefficient are presented. The conditions are found when a significant decrease or increase of the collisionless absorption is expected. In Section 5 the phase-shift of the reflected wave is analyzed. Polarization modifications of the reflected wave are particularly important in the conditions of the anomalous skin effect and at high degree of temperature anisotropy. The radiation transmission through a thin foil is considered in Section 6. It is shown that for a plasma with EDF anisotropy the effective skin-depth becomes considerably larger than in the case of an isotropic plasma with the ensuing consequence that the thin foil becomes much more transparent.

2. Magnetic field generation
Let us investigate a plasma with the bi-maxwellian EDF

\[ F = \left( \frac{m}{2\pi} \right)^{3/2} \frac{N}{T_{\perp}\sqrt{T_x}} \exp \left[ -\frac{mv_x^2}{2T_x} - \frac{m}{2T_{\perp}} \left( v_y^2 + v_z^2 \right) \right] . \]  

(1)

with \( N \) the electron density, \( m \) the electron mass, and the effective temperatures \( T_x \) and \( T_{\perp} \) given in energy units. EDF like (1) is formed as a result of tunnel ionization of atoms [2].

As it is known, the distribution (1) is unstable with respect to Weibel instability development, resulting in a quasi-stationary magnetic field generation [9]. If \( T_x >> T_{\perp} \) then the instability growth rate has the following form [10]

\[ \gamma = q \sqrt{\frac{T_x}{m}} \left( 1 + \frac{q^2c^2}{\omega_L^2} \right)^{-1/2} >> q \sqrt{\frac{T_{\perp}}{m}} , \]

(2)

where \( \omega_L = (4\pi e^2 N/m)^{1/2} \) is the electron plasma frequency, \( e \) the electron charge, \( q \) the perturbation wave number, \( c \) the speed of light. The growth rate maximum is achieved in the domain of large wave numbers, where \( q > \omega_L/c \) and is given by

\[ \gamma_m = \frac{\omega_L}{c} \sqrt{\frac{T_x}{m}} << \omega_L . \]

(3)

The value of \( 1/2\gamma_m \) defines the characteristic time of a quasi-stationary magnetic field energy density increase with a non-uniformity scale less than \( c/\omega_L \). According to Eq.(3), the linear theory of instability gives a simple estimate of the ratio of the energy density of the quasi-stationary field \( B^2 \) to the energy density of the spontaneous magnetic field \( B_{sp}^2 \):

\[ \frac{B^2}{B_{sp}^2} \sim \exp \left( 2\gamma_m t \right) . \]

(4)

The linear theory gives the maximum possible increase in the initial noise, which is not realized when nonlinear effects become important. Numerical investigations of the nonlinear stage of Weibel instability show (see, for instance, [11]) that the energy density of the magnetic field in the saturated state is not greater than 10% of the average kinetic energy corresponding to the initial electron velocity distribution. In the light of this observation, we find for the energy of the quasi-stationary magnetic field the value

\[ \frac{B_m^2}{4\pi} \approx 0.1NT_x . \]

(5)
According to this estimate, in a plasma with electron density $N \approx 5 \cdot 10^{23} \text{cm}^{-3}$ and temperature $T_x \approx 500\text{eV}$ a magnetic field of the order of $20 \text{mG}$ is generated.

3. Optics of dense anisotropic plasma
Let us assume that an anisotropic plasma occupies the half-space $z > 0$, and consider the interaction of such a plasma with a weak linearly polarized electromagnetic wave, impinging normally on the plasma surface

$$\frac{1}{2} \overrightarrow{E} \exp(-i\omega t + ikz) + \text{c.c.}, \quad z < 0,$$

with $\overrightarrow{E} = (E_x, E_y, 0) = E(\cos \varphi, \sin \varphi, 0)$, $\varphi$ the angle between the electric field polarization direction $\overrightarrow{E}$ and the OX axis; $\omega$ and $k$ the frequency and wave number of the e.m. wave, $\omega = kc$. The wave frequency $\omega$ is assumed much smaller than the electron plasma frequency $\omega_L$. According to (1), the OX axis is also the symmetry axis of the EDF. Thus we have two new parameters characterizing the process as compared to the case of an isotropic plasma. One is the angle $\varphi$, the other is the degree of temperature anisotropy $\Delta = 1 - T_x/T_\perp$. We will show below that these new parameters effectively control the process new features. The wave (6) is reflected by the plasma surface and partially penetrates into the plasma, where it is absorbed as a result of interaction with the electrons. In the geometry of interaction chosen here, only two components of the electric field are different from zero. As the incident wave has low intensity, we can describe independently the interaction of each component with the plasma. First of all we consider the interaction of the $E_x$ component. Then for the reflected component we have

$$\frac{1}{2} E_x R_x e_x(\omega t - ikz) + \text{c.c.}, \quad z < 0,$$

where $R_x$ is the complex reflection coefficient. In its turn, the field inside the plasma has the form

$$\frac{1}{2} E_x(z) e_x(\omega t) + \text{c.c.}, \quad z > 0.$$

To define the field $E_x(z)$, from the Maxwell equations we have

$$\frac{d^2}{dz^2} E_x(z) + \frac{\omega^2}{c^2} E_x(z) = -\frac{4\pi ie\omega}{c^2} \int d\overrightarrow{v}v_x \delta f.$$

The perturbation $\delta f$ to the EDF (1) is found from the kinetic equation

$$-i\omega \delta f + v_z \frac{d}{dz} \delta f = -\frac{e}{m} \left\{ E_x(z) \frac{\partial F}{\partial v_x} + \frac{1}{c} B_y(z) \left[ v_x \frac{\partial F}{\partial v_z} - v_z \frac{\partial F}{\partial v_x} \right] \right\} \equiv -S_x(z, v_z),$$

where $B_y(z)$ is the magnetic field component in the plasma, created by the impinging wave

$$B_y(z) = -\frac{c}{\omega} \frac{d}{dz} E_x(z).$$

In plasmas with isotropic EDFs, the term in (10) containing $B_y(z)$ goes to zero. In other words, in isotropic plasmas, the magnetic field does not affect the electron kinetics in the skin layer. At the contrary, in plasmas with anisotropic EDFs like (1), the term containing $B_y(z)$ is not zero. Physically, it implies that the anisotropic distribution over velocities creates the conditions for the magnetic field $B_y(z)$ to rotate the electrons from one degree of freedom to the other.
Because of this the magnetic field contributes to determine $\delta f$, i.e. to influence significantly the electron kinetics. Besides, in the skin-effect conditions, the magnetic field in absolute value considerably exceeds the electric field, according to the inequality $c/\omega >> |d \ln E_x(z)/dz|^{-1}$. The joint manifestation of both these causes, as it will be shown below, is responsible for the appearance of new optical properties in plasmas with anisotropy in the EDF. Further, we consider the simplest boundary conditions on the plasma surface. Namely, we assume that electrons are specularly reflected by the plasma boundary. Provided these conditions are fulfilled, Eq.(10) gives

$$\delta f = \frac{1}{v_z} \int_z^{\infty} dz' S_x(z', v_z) \exp \left[ i \frac{\omega}{v_z} (z - z') \right], \quad v_z < 0;$$

$$\delta f = -\frac{1}{v_z} \int_0^z dz' S_x(z', v_z) \exp \left[ i \frac{\omega}{v_z} (z - z') \right] - \frac{1}{v_z} \int_0^{\infty} dz' S_x(z', -v_z) \exp \left[ i \frac{\omega}{v_z} (z + z') \right], \quad v_z > 0. \quad (12)$$

Substituting the perturbation $\delta f$ (12) into the r.h.s. of Eq.(9) and performing the Fourier transform over $z$, as it is usually done when the specular reflection conditions are assumed (see, for instance, [12]), from (9) we obtain

$$E_x(z) = -\frac{i k}{\pi} \frac{E_x(+0)}{Z_x} \int_{-\infty}^{\infty} \frac{dq}{q^2 - k^2 \varepsilon_x(q)} \exp (iqz), \quad z > 0, \quad (13)$$

where $E_x(+0)$ is the electric field on the plasma boundary. In writing the expression (13), we have used the notation $Z_x$ to indicate the component of the surface impedance

$$Z_x = \frac{E_x(+0)}{B_y(+0)}, \quad (14)$$

giving the ratio between the electric and magnetic fields on the plasma surface; besides, $\varepsilon_x(q)$ represents the x-component of the plasma dielectric susceptibility. If the electrons have a velocity distribution like (1), the function $\varepsilon_x(q)$ is given by

$$\varepsilon_x(q) = 1 - \frac{\omega^2}{\omega^2} J_+ \left( \frac{\omega}{q v_T} \right) - (1 - \frac{T_\perp}{T_x} \frac{\omega^2}{\omega^2} \left[ 1 - J_+ \left( \frac{\omega}{q v_T} \right) \right], \quad (15)$$

where $v_T = \sqrt{T_\perp/m}$, and the function $J_+$ has the form [13]

$$J_+(x) = J'_+(x) + iJ''_+(x), \quad (16)$$

$$J'_+(x) = x \exp \left( -\frac{x^2}{2} \right) \int_0^x dt \exp \left( \frac{t^2}{2} \right), \quad (17)$$

$$J''_+(x) = -\sqrt{\frac{\pi}{2}} x q |q| \exp \left( -\frac{x^2}{2} \right). \quad (18)$$

Further, using the continuity condition of the electric and magnetic fields on the plasma surface, from (6),(7) and (11), we find $R_x$ and $Z_x$. The interaction of the component $E_y$ with the plasma is considered in a completely similar way. The plasma response to this component is described by the functions $R_y$ and $Z_y$, which formally are similar to $R_x$ and $Z_x$, but $Z_y$ is determined
through $\varepsilon_y(q)$, whose expression follows from (15) letting $T_x = T_\perp$. As a final result, for the surface impedance components $Z_\alpha$ and the complex reflection coefficients $R_\alpha$ we find

$$Z_\alpha = -\frac{2i}{\pi} k \int_0^\infty \frac{dq}{q^2 - k^2 \varepsilon_\alpha(q)}, \quad \alpha = x, y,$$

$$R_\alpha = \frac{(Z_\alpha - 1)}{(Z_\alpha + 1)} = |R_\alpha| \exp(i\Psi_\alpha),$$

where the functions $|R_\alpha|$ and $\Psi_\alpha$ are expressed through the real and imaginary parts $Z'_\alpha$ and $Z''_\alpha$ of the surface impedance components $Z_\alpha = Z'_\alpha + iZ''_\alpha$. As a consequence of the Landau damping, the intensity of the reflected wave is smaller than the intensity of the incident one. The intensity decrease of the reflected wave is measured by the decrease of the function $|R_\alpha|$ with respect to unity. Another important effect is related to the circumstance that different components of the incident wave are reflected by the anisotropic plasma with different phase-shifts, $\Psi_x \neq \Psi_y$.

Due to the difference between $\Psi_x$ and $\Psi_y$, the reflected wave polarization differs as compared to that of the incident wave. Namely, the linearly polarized incident wave (6) is reflected by the anisotropic plasma as an elliptically polarized wave (7). Under these conditions, the ellipticity degree of the reflected wave is characterized by the phase-shift difference

$$\Psi = \Psi_x - \Psi_y.$$

In the physical conditions under consideration here, the real and imaginary parts of the impedance components in absolute value show small departures from unity. It allows to write the following approximate expressions for the absorption coefficient $A$,

$$A = 1 - R, \quad R = |R_x|^2 \cos^2 \varphi + |R_y|^2 \sin^2 \varphi,$$

$$A \simeq 4Z_x' \cos^2 \varphi + 4Z_y'' \sin^2 \varphi,$$

and the phase-shift $\Psi$

$$\Psi \simeq 2Z_y'' - 2Z_x''.$$

According to (22) the absorption coefficient is basically determined by the real parts of the impedance components. At the contrary, the phase-shift is determined by the difference of the imaginary parts. The following two Sections are devoted to the investigation of the functions $A$ (22) and $\Psi$ (23).

4. Collisionless absorption

The absorption coefficient (22) has been obtained under the assumption that the electron collision frequency, including both electron-electron and electron-ion collisions, is negligibly small. The departure of $A$ from zero is due to Landau collisionless absorption. From (22) it can be seen that to describe how the absorption coefficient depends on the plasma and radiation field parameters it is sufficient to give the corresponding description of the real parts of the impedance components (19). The real parts of the impedance components have been analyzed in [14] for different limiting cases. Taking advantage of the results of [14], we can write the following approximate formulae for the absorption coefficient for different physical limiting cases of interest:

a) strongly "high-frequency" skin-effect and arbitrary shape of the EDF, except $T_x > T_\perp \delta^{-2}$,

$$A = 8\sqrt{\frac{\pi}{2}} \Omega^3 \left( \frac{T_x}{T_\perp} \cos^2 \varphi + \sin^2 \varphi \right), \quad \delta << 1, \quad \delta^2 T_x << T_\perp;$$

$$A \simeq 4Z_x' \cos^2 \varphi + 4Z_y'' \sin^2 \varphi,$$
b) strongly "high-frequency" skin-effect and EDF strongly elongated along X,

$$A = 4\Omega \left( \frac{T_\perp}{T_x - T_\perp} \cos^2 \varphi + \sqrt{\frac{8}{\pi}} \delta^3 \sin^2 \varphi \right), \quad \delta << 1, \quad \delta^2(T_x - T_\perp) >> T_\perp; \quad (25)$$

c) strongly "anomalous" skin-effect and EDF strongly compressed along X,

$$A = 4\sqrt{\frac{7\Omega}{\pi}} \delta \left( \frac{T_x}{T_\perp} \ln \delta - 1 + \frac{1}{2} \left( \ln 2 - 0.577 \right) \cos^2 \varphi + \frac{(4\pi)^{1/3}}{3\sqrt{3}} \delta^{4/3} \sin^2 \varphi \right), \quad \delta >> 1, \quad T_x << T_\perp; \quad (26)$$

d) strongly "anomalous" skin-effect, EDF close to isotropic but weakly compressed along X,

$$A = 4\sqrt{\frac{7\Omega}{\pi}} \delta \left( \frac{T_x^2}{(T_x - T_\perp)^2} \left[ \ln \left( \frac{T_x}{T_\perp} \right) + 3 \ln \left( 1 - \frac{T_x}{T_\perp} \right) - 1 \right] \cos^2 \varphi + \frac{(4\pi)^{1/3}}{3\sqrt{3}} \delta^{4/3} \sin^2 \varphi \right), \quad T_\perp >> T_x >> T_\perp \delta^{-2/3}; \quad (27)$$

e) strongly "anomalous" skin-effect and EDF very close to isotropic

$$A = \frac{8}{3\sqrt{3}} \left( \frac{2}{\pi} \right)^{1/6} \Omega \delta^{1/3} \left[ 1 - \sqrt{\frac{3}{\pi}} \delta^{2/3} \left( 1 - \frac{T_x}{T_\perp} \right) \cos^2 \varphi \right], \quad T_\perp >> T_\perp \delta^{-2/3} >> |T_\perp - T_x|; \quad (28)$$

f) strongly "anomalous" skin-effect and EDF elongated along X

$$A = 4\Omega \left( \frac{T_x}{T_x - T_\perp} \cos^2 \varphi + \frac{2}{3\sqrt{3}} \left( \frac{2}{\pi} \right)^{1/6} \delta^{1/3} \sin^2 \varphi \right), \quad \delta >> 1, \quad (T_x - T_\perp) >> T_\perp \delta^{-2/3}. \quad (29)$$

In (24)-(29) $\Omega = \omega_c/\omega_l$ and $\delta = \varepsilon_{\varphi} \omega_{\varphi} / \omega_c$ is the parameter characterizing the skin-effect. As far as the parameter $\delta$ is concerned, we point out that it characterizes the degree of anomaly of the skin-effect in an isotropic plasma, when $T_x = T_\perp$. Altogether the asymptotic expressions (24)-(29) allow to understand the behavior of the absorption coefficient in all the limiting situations of physical interest. They are also useful for comparison with the results of the numerical calculations to be reported below. In Figs. 1 and 2 we report the results of calculations of the collisionless absorption coefficient $A$ (22), as a function of the angle $\varphi$, formed by the polarization vector of the absorbed wave with the EDF symmetry axis. Fig.1 shows the curves corresponding to $\delta = 9$ and three values of the temperature anisotropy $T_x = 5T_\perp; T_x = T_\perp; \text{and } T_\perp = 4T_\perp$. According to Fig.1, for the chosen plasma and laser parameters, a relative decrease of absorption takes place as compared to the case of an isotropic plasma. Besides, the absorption coefficient in a plasma with $T_x \neq T_\perp$ is found to grow monotonously with the angle $\varphi$ increase. The dependencies of $A$ vs $\varphi$, for $\delta = 0.3$ and the same values of temperatures as in Fig.1, are shown in Fig.2. We note that the curve with $T_\perp = 4T_\perp$ is qualitatively similar to the corresponding curve of Fig. 1. A different behavior, instead, is shown by the curve with $T_x = 5T_\perp$. In fact, a significant increase of absorption in an anisotropic plasma is observed. The dependency of $A$ on the angle $\varphi$ changes as well. For $\delta = 0.3$ and $T_x = 5T_\perp$ the absorption maximum occurs at $\varphi = 0$, when the wave field is polarized along the EDF symmetry axis.
Figure 1. Collisionless absorption coefficient $A$ versus $\varphi$ (in radians) the angle between the polarization vector of the e.m. wave and the EDF symmetry axis, in a plasma with $\delta = v_T \omega_L/\omega_c = 9$ and for three values of the temperature anisotropy: $T_x = 5T_\perp$; $T_\perp = T_x$; $T_\perp = 4T_x$.

Figure 2. The same function as in Fig.2, but for $\delta = 0.3$. 
5. Phase-shift at reflection

In this Section we discuss how the phase-shift of the reflected wave depends on the plasma parameters. First of all, we give a summary of approximate formulae of the phase-shift $\Psi$ (23) for different domains of the plasma parameters [15]:

$$\Psi \simeq 0, \quad |T_x - T_\perp| << T_\perp max(\delta^{-2}, \delta^{-2/3});$$  
(30)

$$\Psi \simeq -2\Omega + \frac{2\Omega}{\delta} \sqrt{\frac{T_\perp}{|T_x - T_\perp|}}, \quad \delta << 1, \quad \delta^2 |T_x - T_\perp| >> T_\perp;$$  
(31)

$$\Psi \simeq 2\Omega - \Psi_m, \quad \delta >> 1, \quad T_x << T_\perp;$$  
(32)

$$\Psi \simeq 2\Omega \sqrt{\frac{T_\perp}{|T_x - T_\perp|}} - \Psi_m, \quad T_\perp >> T_x - T_\perp >> T_\perp \delta^{-2/3};$$  
(33)

$$\Psi \simeq \frac{\pi}{2} \frac{\Omega}{\delta} \frac{T_\perp^2}{(T_x - T_\perp)^2} - \Psi_m, \quad T_\perp >> T_x - T_\perp >> T_\perp \delta^{-2/3};$$  
(34)

$$\Psi \simeq 2\Omega \frac{\delta}{\delta} \sqrt{\frac{T_\perp}{|T_x - T_\perp|}} - \Psi_m, \quad \delta >> 1, \quad |T_x - T_\perp| >> T_\perp;$$  
(35)

where $\Psi_m$ stands for the maximum absolute value of the phase-shift

$$\Psi_m = 4 \left(\frac{2}{\pi}\right)^{1/6} \Omega \delta^{1/3} = \frac{4}{3} \left(\frac{2}{\pi}\right)^{1/6} \left(\frac{v_T \omega_L^2}{\omega_p^2 c}\right)^{1/3}.\quad (36)$$

According to (30) the phase-shift is close to zero when the electron temperature anisotropy

![Figure 3. The phase-shift of the reflected wave $\Psi$ versus $\delta = v_T \omega_L / \omega c$ for three values of the electron temperature anisotropy: $T_\perp = 4T_x$; $T_\perp = 2T_x$; $T_\perp = 5T_x$.](image-url)
is relatively small. As the degree of anisotropy is increased, the absolute value of the phase-shift increases too. For $\delta \ll 1$ the phase-shift reaches the value $-2\Omega$ (31) for rather large temperature anisotropy, when $T_x > T_\perp$ and $\delta^{-2} >> T_\perp$. When, instead, $\delta >> 1$, one has that: i) the phase-shift absolute value maximum $\Psi_m$ becomes significantly larger than $2\Omega$; and ii) $\Psi_m$ is reached for relatively small temperature anisotropy, $|T_x - T_\perp| >> T_\perp \delta^{-2/3}$. That the phase-shift is larger in the case of large $\delta$ values is traced back to the fact that the larger the skin-effect degree of anomaly, the stronger the wave magnetic field influence on the electron kinetics in the skin-layer. Fig.3 reports plots of the phase-shift $\Psi$ versus $\delta$ for three values of the temperature anisotropy: $T_x = 2T_\perp; T_x = 5T_\perp; T_\perp = 4T_x$. For small $\delta$ all the curves are close to zero, in agreement with formula (30). Far from $\delta = 0$, the behavior of the function $\Psi$ depends on the temperature anisotropy. When $T_\perp = 4T_x$, the function $\Psi$ depends weakly on the parameter $\delta$, if $\delta \leq 9$. If $T_x = 5T_\perp$ the behavior of the curve in Fig.3 corresponds to formulae (30), (31) and (35). Finally, the curve of Fig.3 with $T_x = 2T_\perp$ is qualitatively described by formulae (30), (34) and (35). As a whole, Fig.3 shows that the larger $\delta$ and the ratio $T_x/T_\perp$, the larger the absolute value of the function $\Psi$ too.

6. Transmission through a thin foil
With the aim to define the reflection and the transmission coefficients by a plasma layer let us consider the electromagnetic field in the plasma layer. According to Maxwell’s equations for the electric and magnetic components of the radiation field inside the plasma, we have equations (9)-(11). In solving Eq.(10), we take advantage of the circumstance that the average electron kinetic energy along the $z$ axis $T_z = \frac{m}{N} \int dv \cdot v^2_z F/N$ is a small quantity. Namely, we assume that

$$l_s \omega >> \sqrt{\frac{T_z}{m}},$$

where $l_s$ is the effective skin-layer depth. The inequality (37) allows one to treat the spatial derivative in the left-hand side of Eq.(10) as a perturbation.

We consider the case, when $F$ does not depend on the test wave electric field. This evidently takes place for a prepared plasma. Then, taking into account Eq.(37) and keeping only the corrections containing the field derivative up to the second order, from Eq.(10) we obtain

$$\delta f = -\frac{i e}{m \omega} \left\{ E_x(z) \frac{\partial F}{\partial v_x} - \frac{i}{\omega} \frac{dE_x}{dz} v_x \frac{\partial F}{\partial v_z} - \frac{1}{\omega^2} \frac{d^2 E_x(z)}{dz^2} v_x v_z \frac{\partial F}{\partial v_z} \right\}. \quad (38)$$

This expression permits one to find the current density in the rhs of Eq.(9). Further, when $\omega \ll \omega_p$ and

$$T_x = m \int dv \cdot v^2_z F/N >> \frac{\omega^2}{\omega_p^2} mc^2,$$

Eq.(9) becomes

$$\frac{d^2}{dz^2} E_x(z) - \frac{E_x(z)}{l_s^2} = 0. \quad (40)$$

In Eq.(40), the length $l_s$ characterizes the field penetration depth inside a highly anisotropic plasma and is found as

$$l_s = \left( \frac{T_x}{m \omega^2} \right)^{1/2}. \quad (41)$$
This result implies that the inequality (37) amounts to $T_x >> T_z$. We note also that, thanks to inequality (39), the length $l_s$, given by Eq.(41), is much larger than the skin-layer dimension resulting when the EDF is isotropic. It is true both in the case of high-frequency and anomalous skin effects. In fact, in the first case, one has $l_s \sim c/\omega_p$, while in the second case $l_s \sim \left( c^2/T_x/\sqrt{m/\omega_p^2} \right)^{1/3}$. Equation (40) has the solution

$$E_x(z) = \frac{1}{\sinh(d/l_s)} \left\{ E(0) \sinh \left( \frac{d-z}{l_s} \right) + E(d) \sinh \left( \frac{z}{l_s} \right) \right\}, \quad (42)$$

where $E(0)$ and $E(d)$ are the fields in the plasma at the boundaries. The magnetic field is described by Eqs.(11) and (42). In its turn, the electric fields of the incident, reflected and transmitted waves have, respectively, the forms

$$E_i = \frac{1}{2} E_0 \exp (-i\omega t + ikz) + c.c.,$$
$$E_r = \frac{1}{2} E_0 R \exp (-i\omega t - ikz) + c.c.,$$
$$E_t = \frac{1}{2} E_0 T \exp (-i\omega t + ikz - ikd) + c.c., \quad (43)$$

where $R$ and $T$ are the complex reflection and transmission coefficients. We assume that all electric fields are directed along the OX axis. Using Eqs. (8), (42) and (43) and the continuity requirement of the magnetic and electric fields at the foil boundaries $z = 0$ and $z = d$, we find for $R$ and $T$:

$$R = \left( 1 + k^2 l_s^2 \right) \sinh \left( \frac{d}{l_s} \right) \left[ 2ikl_s \cosh \left( \frac{d}{l_s} \right) + (k^2 l_s^2 - 1) \sinh \left( \frac{d}{l_s} \right) \right]^{-1}, \quad (44)$$

$$T = 2ikl_s \left[ 2ikl_s \cosh \left( \frac{d}{l_s} \right) + (k^2 l_s^2 - 1) \sinh \left( \frac{d}{l_s} \right) \right]^{-1}. \quad (45)$$

From Eqs.(44) and (45) follows the relation

$$|R|^2 + |T|^2 = 1, \quad (46)$$

meaning that the incident wave energy is partly reflected and partly transmitted. In Eq.(46) the absorption coefficient is absent, as a result of the adopted approximate account of the spatial dispersion in Eq.(10). In other words, we have disregarded the small energy losses due to the field Cherenkov absorption. According to Eqs.(39) and (45), the ratio of the transmitted energy density flux to the incident one is given by

$$|T|^2 = \left\{ 1 + \frac{1}{4} \left( \sqrt{\frac{T_x}{mc^2}} + \sqrt{\frac{mc^2}{T_x}} \right)^2 \sinh \left( \frac{2\pi d}{\lambda} \sqrt{\frac{mc^2}{T_x}} \right) \right\}^{-1}, \quad (47)$$

where $\lambda = 2\pi/\omega$ is the wavelength of the incident radiation. According to Eq.(47) for $mc^2 >> T_x$, and relatively thin foil, such that

$$d < \frac{\lambda}{\pi mc^2}, \quad (48)$$

the transmission coefficient is close to unity. The foil transparency increase is due to the relatively large depth of field penetration inside the plasma, which possesses a highly anisotropic EDF. In its turn, the large penetration depth is a consequence of the wave magnetic field action on the electron motion in the skin-layer. It is just thanks to the magnetic field that a fast electron motion appears also in the direction perpendicular to the electric field yielding the penetration depth increase.
7. Conclusions
We have reported on investigations, addressing analytically and numerically the basic features of the collisionless absorption, reflection and transmission of linearly polarized radiation exhibited by a plasma possessing an anisotropic bi-maxwellian EDF. The reported results show significant differences as compared to the case of a plasma with an isotropic EDF. The main physical reason, to which the differences are to be traced back, is that in the conditions of an anisotropic EDF the wave magnetic field influences the electron kinetics inside the skin-layer. The reported new optical properties require, to become true, that electron collisions and Weibel instability have no influence on the electron distribution over velocities. It is possible at enough high effective temperatures and for reasonably short time intervals.

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