Nonlinear absorption of ultrapower laser radiation by relativistic underdense plasma

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Abstract.
The nonlinear absorption of laser radiation of relativistic intensities in the underdense plasma by a mechanism of stimulated bremsstrahlung of electrons on the ions/nuclei is investigated in the low frequency approximation. Coefficient of nonlinear inverse-bremsstrahlung absorption is studied for relativistic Maxwellian plasma at asymptotically large values of laser fields and high temperatures of electrons.

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1. Introduction

In the last two decades, laser technologies have made a giant leap forward such that laser sources of ultrarelativistic intensity level can be attained [1]. Near-infrared laser beams are available up to intensities of $10^{22}$ W/cm$^2$ and much stronger lasers will be available in near future [2]. For such lasers the dimensionless relativistic invariant parameter of intensity $\xi_0 \equiv eE_0/mc\omega >> 1$ ($e$ - elementary charge, $m$ - electron mass, $E_0$, $\omega$ - electric field amplitude and frequency of a laser radiation, $c$ - light speed in vacuum). The latter represents the work of the field on the one wavelength in the units of the particle rest energy. Interaction of such lasers with the matter at extreme conditions in ultrashort space-time scales have attracted broad interest over the last few years conditioned by a number of important applications, such as generation and probing of highenergy-density plasma [3], ions acceleration and inertial confinement fusion [4], vacuum nonlinear optics [5], compact laser-plasma accelerators [6], etc. Generally, the interaction of such fields with the electrons in the presence of a third body makes available the revelation of many nonlinear relativistic electrodynamic phenomena. As a third body can serve ion and in the superintense laser fields one can observe relativistic above threshold ionization [7] and high order harmonic generation [8], electron-positron pairs production on nuclei [9], and multiphoton stimulated bremsstrahlung (SB) of electrons on the ions/nuclei [10]. The latter is one of the fundamental processes at the interaction of superstrong laser pulses with plasma and under the same circumstances inverse-bremsstrahlung absorption may become dominant mechanism of absorption of strong electromagnetic (EM) radiation in underdense plasma.

With the advent of lasers many pioneering papers have been devoted to the theoretical investigation of the electron-ion scattering processes in gas or plasma in the presence of a laser field using nonrelativistic [11–20] as well as relativistic [21–25] considerations. The appearance of superpower ultrashort laser pulses of relativistic intensities has initiated new interest in SB in relativistic domain [26–28], where investigations were carried out mainly in the Born approximation over the scattering potential. Meanwhile for ions with the large charge and for the clusters [29], when electron interaction with the entire dense cluster ion core that composed of a large number of ions is dominant, the Born approximation is not applicable. The theoretical description of SB in superstrong EM fields and scattering centers of large charges requires one to go beyond the scope of Born approximation over the scattering potential and the perturbation theory over laser field. In this context, when quantum effects are considerable, one can apply eikonal [23, 30] or generalized eikonal approximation [31]. For the infrared and optical lasers, in the multiphoton interaction regime, one can apply classical theory and the main approximation in the classical theory is low frequency (LF) or impact approximation [16, 17]. LF approximation have been generalized for relativistic case in Refs. [24, 25], where the effect of an intense EM wave on the dynamics of SB and non-linear absorption of intense laser radiation by a monochromatic electron beam due to SB have been carried out. Regarding the absorption of an EM of relativistic
intensities in plasma due to inverse SB there are not investigations which could help to clarify the behavior of nonlinear absorption of an intense EM radiation at very large \( \xi_0 \). Hence, it is of interest to clear up how nonlinear SB effect will proceeds in the plasma, taking also into account initial relativism of plasma electrons.

In the present paper inverse-bremsstrahlung absorption of an intense laser radiation in relativistic Maxwellian plasma is considered in the relativistic LF approximation \cite{24, 25}. In this approximation one can consider as superstrong laser fields as well as scattering centers with large charges. The radiation power absorption in such plasma is investigated as for circularly polarized wave (CPW), as well as for linearly polarized wave (LPW). We consider the dependence of the absorption coefficient on the intensity and polarization of the laser radiation, as well as on the temperature of the plasma electrons.

The organization of the paper is as follows: In Sec. II the relativistic absorption coefficient of the EM wave of arbitrary polarization and intensity due to the mechanism of SB process is presented. In Sec. III we consider the problem numerically along with derivation of asymptotic formulas for absorption coefficient. Conclusions are given in Sec. IV.

2. Nonlinear inverse-bremsstrahlung absorption coefficient

The absorption coefficient \( \alpha \) for an EM radiation field of arbitrary intensity and polarization, in general case of the homogeneous ensemble of electrons of concentration \( n_e \), with the arbitrary distribution function \( f(p) \) over momenta \( p \), at the inverse bremsstrahlung on the scattering centers with concentration \( n_i \), can be represented in the form:

\[
\alpha = \frac{n_e}{T} \int d\mathbf{p}_0 f(\mathbf{p}_0) W, \tag{1}
\]

where \( W \) is the classical energy absorbed by a single electron per unit time from the EM wave of intensity \( I \) due to SB process on the scattering centers. For the homogeneous scattering centers \( W \sim n_i \). For the generality, we assume Maxwellian plasma with the relativistic distribution function:

\[
f(\mathbf{p}_0) = \frac{\exp \left( -\frac{\mathcal{E}(\mathbf{p}_0)}{k_B T_e} \right)}{4\pi m^2 c k_B T_e K_2(mc^2/k_B T_e)}, \tag{2}
\]

where \( k_B \) is the Boltzmann’s constant, \( T_e \) is the temperature of electrons in plasma, \( \mathcal{E}(\mathbf{p}_0) \) is relativistic energy-momentum dispersion law of electrons, \( K_2(x) \) is the McDonald’s function; \( f(\mathbf{p}_0) \) is normalized as

\[
\int f(\mathbf{p}_0) d^3\mathbf{p}_0 = 1. \tag{3}
\]

To obtain \( W \) for the SB process, the electron interaction with the scattering potential and EM wave in the LF approximation can be considered as independently proceeding processes \cite{17, 25}, separated into the following three stages, schematically
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depicted in Fig. 1. Field free electron with energy $E_0$ and momentum $p_0$ interacts with the EM wave. The exact solution of the relativistic equation of motion of an electron in a plane EM wave is well known \[10\]. The energy and momentum in a plane EM wave field can be written as:

$$ p_\perp(\psi) = p_\perp(\psi_0) - e\frac{A(\psi_0) - A(\psi)}{c}, $$

$$ \nu p(\psi) = \nu p(\psi_0) + \frac{1}{2c(\mathcal{E}(\psi_0) - c
u p(\psi_0))} \times \left[ e^2 (A(\psi_0) - A(\psi))^2 - 2ecp(\psi_0) (A(\psi_0) - A(\psi)) \right], $$

$$ \mathcal{E}(\psi) = \mathcal{E}(\psi_0) + c\nu (p(\psi) - p(\psi_0)), $$

where

$$ A(\psi) = A_0(\psi)(\hat{e}_1 \cos \psi + \hat{e}_2 \zeta \sin \psi) $$

is the vector potential of the EM wave of currier frequency $\omega$ and slowly varying amplitude $A_0(\psi)$. Here $\psi = \omega \tau$ is the phase, $\tau = t - \nu r/c$, $\nu$ is an unit vector in the EM wave propagation direction, $\hat{e}_{1,2}$ are the unit polarization vectors, and $\arctan \zeta$ is the polarization angle. At the second stage the elastic scattering of the electron in the potential field takes place at the arbitrary, but certain phase $\psi_s$ of the EM wave. Thus, taking into account adiabatic turn on of the wave ($A(\psi_0) = 0$) from Eqs. (4)-(6) before the scattering one can write

$$ p_\perp(\psi_s) = p_\perp(\psi_0) + \frac{eA(\psi_s)}{c}, $$

$$ \nu p(\psi_s) = \nu p(\psi_0) + \frac{1}{2c\Lambda} \left[ e^2 A^2(\psi_s) + 2ecp_0A(\psi_s) \right], $$

$$ \mathcal{E}(\psi_s) = \mathcal{E}(\psi_0) + c\nu (p(\psi_s) - p_0), $$

where

$$ \Lambda = \mathcal{E}(\psi_s) - c\nu p(\psi_s) = E_0 - c\nu p_0 $$

is the integral of motion for a charged particle in the field of a plane EM wave. The mean energy of an electron in the wave-field before the scattering will be

$$ \langle \mathcal{E}(\psi) \rangle_i = \mathcal{E}_0 + \frac{e^2 \langle A^2 \rangle}{2\Lambda}. $$

Then it takes place elastic scattering on the scattering center $U(r)$. Due to the instantaneous interaction of the electron with the scattering potential the wave phase does not change its value during the scattering. The electron with initial momentum $p(\psi_s)$ is acquired the momentum $p'(\psi_s)$ ($\mathcal{E}'(\psi_s) = \mathcal{E}(\psi_s)$) after the scattering directly at the same phase $\psi_s$ of the wave, which can be defined from the generalized consideration
of the elastic scattering. Thus, measuring the scattering angle \( \vartheta \) from a direction \( p(\psi_s) \), with corresponding azimuthal angle \( \varphi \) for the scattered momentum one can write

\[
\begin{bmatrix}
  p_x'(\psi_s) \\
  p_y'(\psi_s) \\
  p_z'(\psi_s)
\end{bmatrix} = p(\psi_s) \hat{R} \begin{bmatrix}
  \sin \vartheta \cos \varphi \\
  \sin \vartheta \sin \varphi \\
  \cos \vartheta
\end{bmatrix},
\]

where \( \hat{R} = \hat{R}_z(\varphi_0) \hat{R}_y(\vartheta_0) \hat{R}_y(\varphi_0) \) and \( \hat{R}_z(\varphi_0) \) are basic rotation matrices about the \( y \) and \( z \) axes:

\[
\hat{R} = \begin{bmatrix}
  \cos \vartheta_0 \cos \varphi_0 & -\sin \varphi_0 & \sin \vartheta_0 \cos \varphi_0 \\
  \cos \vartheta_0 \sin \varphi_0 & \cos \varphi_0 & \sin \vartheta_0 \sin \varphi_0 \\
  -\sin \vartheta_0 & 0 & \cos \vartheta_0
\end{bmatrix}.
\]

As a \( Oz \) axis we take wave propagation direction \( \nu \), \( \vartheta_0 \) is the polar angle and \( \varphi_0 \) is the azimuthal angle in the wave-polarization plane.

At the third stage the electron again interacts only with the wave, moving in the wave field with the momentum and energy defined from Eqs. (4)-(6):

\[
p_{\perp f}(\psi) = p'(\psi_s) - eA(\psi_s) - A(\psi),
\]

\[
\nu p_f(\psi) = \nu p'(\psi_s) + \frac{1}{2c\Lambda'} \left[ e^2 (A(\psi_s) - A(\psi))^2 
- 2cep'(\psi_s) (A(\psi_s) - A(\psi)) \right],
\]

\[
\mathcal{E}_f(\psi) = \mathcal{E}(\psi_s) + c\nu (p(\psi) - p'(\psi_s)),
\]

where

\[
\Lambda' = \mathcal{E}_f(\psi) - c\nu p_f(\psi) = \mathcal{E}(\psi_s) - c\nu p'(\psi_s).
\]

The mean energy of an electron in the wave field after the scattering will be

\[
\langle \mathcal{E}_f(\psi) \rangle = \mathcal{E}(\psi_s) + \frac{1}{2\Lambda'} \times \left[ e^2 (A^2) + \langle A^2 \rangle \right] - 2cep'(\psi_s)A(\psi_s).
\]

The energy change due to SB can be calculated as a difference of mean energy in the field before and after the scattering:

\[
\Delta \mathcal{E}(\vartheta, \varphi, \psi_s, p_0) = \langle \mathcal{E}_f(\psi) \rangle - \langle \mathcal{E}(\psi) \rangle.
\]
Taking into account Eqs. (12) and (19) we obtain:

$$\Delta \mathcal{E} = \frac{e^2 \mathbf{A}^2(\psi_s)}{2} \left( \frac{1}{\Lambda'} - \frac{1}{\Lambda} \right)$$

$$- \frac{e c \mathbf{p}'(\psi_s) \mathbf{A}(\psi_s)}{\Lambda'} + \frac{e c \mathbf{p}(\psi_s) \mathbf{A}(\psi_s)}{\Lambda}.$$  \hfill (20)

For the energy absorbed by a single electron per unit time from the EM wave due to SB process on the scattering centers one can write

$$W = \frac{n_i}{2\pi} \int_0^{2\pi} d\psi_s \int v(\psi_s) \Delta \mathcal{E} d\sigma(\vartheta, p(\psi_s)),$$  \hfill (21)

where $$v(\psi_s) = \frac{e^2 p(\psi_s)}{\mathcal{E}(\psi_s)}$$ is the velocity of an electron in the wave-field, $$p(\psi_s) = \sqrt{\mathcal{E}^2(\psi_s) - m^2 c^4}/c$$ is the momentum, and $$d\sigma(\vartheta, p(\psi_s))$$ is the differential cross section of the elastic scattering in the potential field $$U(r)$$. Taking into account that the main contribution in the integral (21) comes from the small angle scatterings one can write

$$W = \frac{n_i}{2\pi} \int_0^{2\pi} d\psi_s \int v(\psi_s) \frac{\partial^2 \Delta \mathcal{E}}{\partial^2 \vartheta} d\sigma_{\text{tr}}(\vartheta, p(\psi_s)),$$  \hfill (22)

where

$$d\sigma_{\text{tr}}(\vartheta, p(\psi_s)) = (1 - \cos \vartheta) d\sigma(\vartheta, p(\psi_s))$$  \hfill (23)

is the transport differential cross section. For the Coulomb scattering centers with potential energy

$$U(r) = \frac{Ze^2}{r}$$

of electron interaction with ion of charge $$Ze$$, one can use relativistic cross section for elastic scattering at small angles [32] and make integration over $$\vartheta$$ and $$\varphi$$ to obtain:

$$W = n_i Z^2 e^4 \int_0^{2\pi} d\psi_s \frac{m^2 c^2}{\Lambda^3} \left[ \frac{e^2 \mathbf{A}^2(\psi_s) + \langle \mathbf{A}^2 \rangle}{2m^2 c^4} \times (\mathcal{E}(\psi_s) \Lambda - m^2 c^4) + e c \mathbf{p}(\psi_s) \mathbf{A}(\psi_s) \right] \frac{\mathcal{E}(\psi_s)}{p^3(\psi_s)} L_{\text{cb}},$$  \hfill (24)

where

$$L_{\text{cb}} = \ln \left( \frac{v^2(\psi_s) p(\psi_s)}{Ze^2 \omega} \right)^2$$  \hfill (25)

is the Coulomb logarithm. The latter has been obtained taking $$\rho_{\text{min}} = Ze^2/v p$$ as a lower limit of the impact parameter, while for the upper limit we assume $$\rho_{\text{max}} = v/\omega$$. Taking into account Eqs. (1), (7), and (24) for the absorption coefficient we obtain:

$$\alpha = \frac{n_i n_e Z^2 e^4}{I} \int d\mathbf{p}_0 f(\mathbf{p}_0) \int_0^{2\pi} d\psi_s \left[ \frac{e^2 \mathbf{A}^2(\psi_s) + \langle \mathbf{A}^2 \rangle}{2m^2 c^4} \times (\mathcal{E}(\psi_s) \Lambda - m^2 c^4) + e c \mathbf{p}(\psi_s) \mathbf{A}(\psi_s) \right] \frac{m^2 c^2 \mathcal{E}(\psi_s)}{\Lambda^3 p^3(\psi_s)} L_{\text{cb}},$$  \hfill (26)
where

\[ I = (1 + \zeta^2)\omega^2 A_0^2/8\pi c. \]

Thus, Eq. (26) represents nonlinear inverse-bremsstrahlung absorption coefficient \( \alpha \) for an EM radiation field of arbitrary intensity and polarization, for homogeneous ensemble of electrons of concentration \( n_e \), with the arbitrary distribution function \( f(p_0) \) over momenta \( p_0 \).

Note that the LF approximation in the intense laser field is applicable when

\[ \lambda \gg \lambda_D, \quad (27) \]

where \( \lambda \) is the laser radiation wavelength and \( \lambda_D = \sqrt{\kappa_B T_e/4\pi n_e e^2 Z} \) is the Debye screening length:

\[ \lambda_D[cm] = 7.43 \times 10^2 \times \sqrt{T_e[eV]/Zn_e[cm^{-3}].} \quad (28) \]

Besides, for an underdense plasma one should take into account condition \( \omega > \omega_p \), where \( \omega_p = \sqrt{4\pi n_e e^2/m_*} \) is the plasma frequency with "effective mass" \( m_* \) of the relativistic electron in the EM wave [33]:

\[ m^* = m\sqrt{1 + \langle \xi^2(\psi) \rangle}. \quad (29) \]

Thus for CPW and for large \( \xi_0^2 \) one can write

\[ n_e < m\xi_0\omega^2/4\pi e^2 = 1.1 \times 10^{21} \times \xi_0 \times \lambda^{-2}(\mu m). \quad (30) \]

There is also limitation on the pulse duration \( \tau \) of an EM wave. The SB should be the main mechanism, which is responsible for the absorption of the laser radiation in plasma. This condition is failed, if the influence of the strong EM wave lead to development of an instability. Hence, pulse duration \( \tau \) of an EM wave has to satisfy the condition \( \mu \tau \lesssim 1 \), where \( \mu \) is the maximal increment of the instability of the plasma in the strong laser field.

In general, analytical integration over momentum \( p_0 \) and scattering phase \( \psi_s \) is impossible, and one should make numerical integration. The latter along with derivation of asymptotic formulas for absorption coefficient \( \alpha \) will be done in the next section.

3. Numerical Treatment: Asymptotic Formulas

As we are interested in superintense laser pulses of relativistic intensities, then it is convenient to represent the absorption coefficient (26) in the form of dimensionless quantities:

\[
\frac{\alpha}{\alpha_0} = \frac{1}{2\pi(1 + \zeta^2)} \xi_0^2 \int d\bar{p}_0 \bar{f}(\bar{p}_0) \int_0^{2\pi} d\psi_s \frac{\gamma(\psi_s)}{\bar{p}^3(\psi_s)\Lambda^3} \\
\times \left( \xi^2(\psi_s) + \langle \xi^2(\psi_s) \rangle \right) \left( \gamma(\psi_s)\Lambda - 1 \right) + \bar{p}_0\xi(\psi_s) + \xi^2(\psi_s) \right) \Lambda_{cb}. \quad (31)
\]
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Here
\[
\alpha_0 = 4Z^2 r_e^3 \lambda^2 n_i n_e,
\]
and \(r_e\) is the classical electron radius. In Eq. (31) the dimensionless momentum, energy, and temperature were introduced as follows:
\[
\bar{P} = \frac{P}{mc}, \quad \gamma(\psi_s) = \frac{\mathcal{E}(\psi_s)}{mc^2}, \quad T_n = \frac{k_B T_e}{mc^2},
\]
and the dimensionless relativistic intensity parameters of EM wave,
\[
\xi(\psi_s) = \xi_0(\hat{e}_1 \cos \psi_s + \hat{e}_2 \zeta \sin \psi_s).
\]
The scaled relativistic distribution function is
\[
f(p_0) = \frac{1}{4\pi T_n K_2(T_n^{-1})} \exp \left( -\frac{\gamma_0}{T_n} \right),
\]
and Coulomb logarithm
\[
L_{cb} = L_0 + \ln \left( \frac{\bar{p}^3(\psi_s)}{\gamma^2(\psi_s)} \right)^2,
\]
where \(L_0 = \ln \left( \lambda/2\pi r_e Z \right)^2\). In the near-infrared and optical domain of frequencies, at the \(Z = 1 - 10\), \(L_0 \approx 30\). Taking into account that normalized absorption coefficient (31) depends on \(Z\) and \(\omega\) through logarithm \(L_0\), for the numerical simulations we will not concretize \(Z\) and \(\omega\), assuming \(L_0 \approx 30\).

It is well known that the kinematics of an electron in the field of a strong EM wave essentially depends on the polarization of the wave [10]. In particular, for the particle initially at rest, in the CPW, energy \(\gamma(\psi_s)\) and momentum \(\bar{p}(\psi_s)\) are constants, since \(\xi^2(\psi_s) = \text{const}\). Meanwhile in the LPW \(p(\psi_s)\) oscillates and as a consequence small values of \(\bar{p}(\psi_s)\) give the main contribution in Eq. (31). The latter leads to more complicated behavior of the dynamics of SB at the linear polarization of a stimulating strong wave. Besides in the case of CPW, thanks to azimuthal symmetry one can make a step forward in analytical calculation and obtain explicit formula for the absorption coefficient at superstrong laser fields. Thus, taking into account azimuthal symmetry in the case of CPW, one can make integration over phase \(\psi_s\), which results in
\[
\frac{\alpha}{\alpha_0} = \frac{1}{2} \int d\bar{P}_0 \bar{f}(\bar{P}_0) \left( \gamma(\bar{P}_0) \bar{\Lambda} + \frac{p_0}{\xi_0} \sin \vartheta_0 \cos \varphi_0 \right)
\]
\[
\times \frac{\gamma(\bar{P}_0)}{\bar{\Lambda} \bar{p}^3(\bar{P}_0)} L_{cb},
\]
where
\[
\gamma(\bar{P}_0) = \gamma_0 + \frac{1}{2\bar{\Lambda}} \left[ \xi_0^2 + 2p_0 \xi_0 \sin \vartheta_0 \cos \varphi_0 \right],
\]
\[
\bar{p}(\bar{P}_0) = \sqrt{\bar{\gamma}^2(\bar{P}_0) - 1}.
\]
At the large $\xi_0^2 >> 1$, taking into account Eqs. (35) and (36), from Eq. (34) one can obtain

$$\alpha = \frac{\alpha_0}{\xi_0^2} \int d\mathbf{p}_0 \mathcal{J}(\mathbf{p}_0) \frac{1}{\Lambda} L_{cb}^{(c)},$$

(37)

where

$$L_{cb}^{(c)} = L_0 + \ln \left( \frac{\xi_0^2}{2\Lambda} \right)^2.$$

The formula (37) shows the suppression of the SB rate with increase of the wave intensity. Ignoring weak logarithmic dependence, we see that absorption coefficient inversely proportional to laser intensity: $\alpha \sim 1/\xi_0^2$. For the large $\xi_0^2$ the dependence of the absorption coefficient on temperature comes from $\Lambda$ in Eq. (37). In particular, for initially nonrelativistic plasma $T_n << 1$ in Eq. (37) one can put $\Lambda \simeq 1$, which gives

$$\alpha \equiv \alpha_C = \frac{\alpha_0}{\xi_0^2} \ln \left( \frac{c}{Zr_e\omega} \right)^2.$$

(38)

The relation for the absorption coefficient in the case of LPW is complicated and even for large $\xi_0$ one can not integrate it analytically. Therefore, for the analysis we have performed numerical investigations, making also analytic interpolation.

The results of numerical investigations of Eq. (31) are illustrated in Figures 2-7 both for the CPW and LPW.

To show the dependence of the inverse-bremsstrahlung absorption rate on the laser radiation intensity in Fig. 2 it is shown scaled rate $\alpha/\alpha_0$ versus relativistic invariant parameter of the wave intensity for various plasma temperatures. The wave is assumed to be circularly polarized. As is seen from this figure the SB rate is suppressed with increase of the wave intensity and for the large values of $\xi_0$ it exhibits a tenuous dependence on the plasma temperature. The behavior is also seen from Fig. 3, where total scaled rate of inverse-bremsstrahlung absorption of CPW in plasma, as a function of the plasma temperature $T_n$ is shown for various wave intensities. Here for large
Figure 3. (Color online) Total scaled rate of inverse-bremsstrahlung absorption (in arbitrary units) of circularly polarized laser radiation in plasma, as a function of the plasma temperature (in units of an electron rest energy $mc^2$) is shown for various wave intensities.

Figure 4. (Color online) Density plot of the total rate of inverse-bremsstrahlung absorption scaled to asymptotic rate $\alpha_C$ (in arbitrary units), as a function of the plasma temperature (in units of an electron rest energy $mc^2$) and the dimensionless relativistic invariant parameter of the circularly polarized laser beam.

Figure 5. (Color online) Total scaled rate of inverse-bremsstrahlung absorption (in arbitrary units) of linearly polarized laser radiation in plasma versus the dimensionless relativistic invariant parameter of wave intensity for various plasma temperatures.
values of $\xi_0$ we have a weak dependence on temperature, which is a result of laser modified relativistic scattering of electrons irrespective of the initial state of electrons. The absorption coefficient $\alpha$ decreases as $1/\xi_0^2$ in accordance with analytical result (37). To clarify the range of applicability of the asymptotic formula (38) in Fig. 4 density plot of the total rate of inverse-bremsstrahlung absorption scaled to asymptotic rate $\alpha_L$, as a function of the plasma temperature and the relativistic invariant parameter $\xi_0$ is shown for CPW. As is seen in the wide range of $T_n$ and $\xi_0$ one can apply asymptotic formula (38).

In Fig. 5 and 6 it is shown the total scaled rate of inverse-bremsstrahlung absorption of LPW in plasma versus the dimensionless relativistic invariant parameter of wave intensity and temperature, respectively. With the interpolation we have seen that $\alpha$ decreases as $1/\xi_0^{5/4}$ and exhibits a tenuous dependence on the plasma temperature.

**Figure 6.** (Color online) Total scaled rate of inverse-bremsstrahlung absorption (in arbitrary units), as a function of the plasma temperature (in units of an electron rest energy $mc^2$) is shown for various wave intensities. The wave is assumed to be linearly polarized.

**Figure 7.** (Color online) Density plot of the total rate of inverse-bremsstrahlung absorption scaled to asymptotic rate $\alpha_L$ (in arbitrary units), as a function of the plasma temperature (in units of an electron rest energy $mc^2$) and the dimensionless relativistic invariant parameter of the linearly polarized laser beam.
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Making analogy with the case of CPW, for the large \( \xi_0 \) we interpolate \( \alpha \) by the following formula:

\[
\alpha \simeq \alpha_L = \frac{\alpha_0}{2\xi_0^{5/4}} \ln \left( \frac{c}{Zr_{\phi}\omega} \frac{\xi_0^2}{4} \right)^2.
\]  \hspace{1cm} (39)

As is seen from Fig. 7, in the case of LPW and for the moderate temperatures, with the well enough accuracy one can apply asymptotic rate (39).

4. Conclusion

We have presented a theory of inverse-bremsstrahlung absorption of an intense laser radiation in relativistic Maxwellian plasma in the relativistic low-frequency approximation. The coefficient of nonlinear inverse-bremsstrahlung absorption has been calculated for relativistic Maxwellian plasma. The simple analytical formulae have been obtained for absorption coefficient at asymptotically large values of laser fields both for circularly and linearly polarized radiations. The obtained results demonstrate that the SB rate is suppressed with the increase of the wave intensity and for large values of \( \xi_0 \) absorption coefficient \( \alpha \) decreases as \( 1/\xi_0^2 \) for circularly and as \( 1/\xi_0^{5/4} \) for the linearly polarized one in contrast to nonrelativistic case where one has a dependence \( 1/\xi_0^3 \) [17].

The SB rate is suppressed with increase of the plasma temperature but for the relativistic laser intensities it exhibits a tenuous dependence on the plasma temperature.

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