Smith-Purcell radiation on a surface wave

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Abstract

We consider the radiation from an electron in flight over a surface wave of an arbitrary profile excited in a plane interface. For an electron bunch the conditions are specified under which the overall radiation essentially exceeds the incoherent part. It is shown that the radiation from the bunch with asymmetric density distribution of electrons in the longitudinal direction is partially coherent for waves with wavelengths much shorter than the characteristic longitudinal size of the bunch.

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1 Introduction

Surface waves have wide applications in various fields of science and technology. In the present talk, based on [1]-[4], we discuss the radiation from a charged particle flying over surface acoustic wave. The physics of this phenomenon is similar to that for the radiation of particle flying over a diffraction grating (Smith-Purcell radiation). The latter is used for the generation of the radiation in the range of millimeter and submillimeter waves.

2 Radiation from a single electron

Let a charge q move with constant velocity v parallel to a surface wave excited in the plane interface between homogeneous media with permittivities ε0 and ε1 (see figure [1]). If the axis z is aligned with the particle trajectory, then the equation of the interface has the form

\[ x = x_0(z, t) = -d + f(k_0 z + \omega_0 t), \] (1)

where \(k_0\) and \(\omega_0\) are the wave number and cyclic frequency, \(d\) is the distance from the non-excited interface, and \(f(u)\) is the function describing the wave profile, \(f(u + 2\pi) = f(u)\). We assume that the particle moves in the medium with permittivity \(\varepsilon_0\) and will consider the radiation in this medium.

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Figure 1: Electron flying over surface wave excited on the plane interface between homogeneous media.

From the symmetry properties of the problem it follows that the emission angle \( \theta \) (with respect to the particle velocity \( \mathbf{v} \)) and the frequency \( \omega \) of the emitted photon are related by the formula

\[
\omega = \frac{m(k_0 v \mp \omega_0)}{1 - \beta \sqrt{\varepsilon_0} \cos \theta}, \quad \beta = \frac{v}{c},
\]

where \( m \) is an integer. The dependence of the radiation intensity on the distance of the electron trajectory from the non-excited interface, \( d \), is determined by the factor \( \exp(-2\omega d \text{Re} \sigma/v) \), where

\[
\sigma = \left[ (1 - \beta^2 \varepsilon_0) \frac{\omega_1^2}{\omega^2} + \beta^2 \varepsilon_0 \sin^2 \theta \sin^2 \varphi \right]^{1/2}, \quad \omega_1 = \omega \pm m \omega_0,
\]

with \( \varphi \) being the polar angle with respect to the \( x \) axis in the plane perpendicular to the particle trajectory. This factor does not depend on the specific form of the wave profile in (1) and is determined by the dependence of the field spectral components for a uniformly moving particle on distance \( d \). For relativistic particles with \( \beta \varepsilon_0 \ll 1 \), one has \( \omega_1 \approx \omega \) and the radiation at large azimuthal angles is exponentially suppressed and the radiation distribution is strongly anisotropic. For the radiation in the vacuum at \( \varphi \ll 1 \) the factor is equal to \( \exp(-2\omega_1 dv/\gamma) \), with \( \gamma = 1/\sqrt{1 - \beta^2} \) been the Lorentz factor. In the case \( \beta \sqrt{\varepsilon_0} < 1 \), the quantity \( \sigma \) is real and the intensity exponentially decreases with increasing distance. The same is the case for the directions of radiation satisfying the condition \( \sin^2 \theta \sin^2 \varphi > 1 - 1/(\beta^2 \varepsilon_0) \), when \( \beta \sqrt{\varepsilon_0} > 1 \). For \( \beta \sqrt{\varepsilon_0} > 1 \) and \( \sin^2 \theta \sin^2 \varphi < 1 - 1/(\beta^2 \varepsilon_0) \), the radiation intensity does not depend on particle distance from a surface of the periodic structure in the absence of absorption. This corresponds to the reflection of Cherenkov radiation emitted in the first medium.

In order to determine the radiation intensity we have used two independent approximate methods. In the first one it is assumed that \( |\varepsilon_1 - \varepsilon_0| \ll \varepsilon_0 \). The second method, which is more appropriate for the problem under consideration, assumes that the amplitude of the surface wave is small. Assuming that the charge moves in the vacuum \( \varepsilon_0 = 1 \), and under the condition \( \beta \sqrt{\varepsilon_1 - 1} > 1 \), the spectral-angular distribution of the radiation intensity (per unit path length)
in the region $x > 0$ is given by the expression

$$
\frac{dW}{d\omega d\Omega} = \frac{2q^2(\varepsilon_1 - 1)}{\pi \varepsilon^2 v^2} \sum_{m \neq 0} \omega^3 \sin^2 \theta \cos^2 \varphi |f_m|^2 e^{-2\omega \sigma d/v}
$$

$$
\times \frac{A_1 \sigma^2 + A_2 \sin^2 \theta \sin^2 \varphi + A_3 \beta^2 \sin^4 \theta \sin^4 \varphi \delta\left(\cos \theta - \frac{1}{\beta} + \frac{mk_0 c}{\omega}\right)}{(1 + \beta^2 \sin^2 \theta \sin^2 \varphi)(\sigma_2 + \varepsilon_1 \sin \theta \cos \varphi)^2}, \quad (4)
$$

where $d\Omega = \sin \theta d\theta d\varphi$, and we have assumed that $v \gg \omega_0/k_0$. In (4), $f_m = \frac{1}{2\pi} \int_{-\pi}^{+\pi} du f(u)e^{-imu}$ is the Fourier transform of the profile function,

$$
\sigma_1 = \sqrt{\beta^2(\varepsilon_1 - \sin^2 \theta \sin^2 \varphi) - 1}, \quad \sigma_2 = \sqrt{\varepsilon_1 - 1 + \sin^2 \theta \sin^2 \varphi}, \quad (5)
$$

and

$$
A_1 = \left[\sigma_1 \sigma_2 \cos \theta - \varepsilon_1(1 - \sin^2 \theta \cos^2 \varphi)\right]^2
$$

$$
+ \sin^2 \theta[\sigma_1(\sigma_2 \cos \varphi + \sin \theta \sin^2 \varphi) - \varepsilon_1 \cos \theta \cos \varphi]^2,
$$

$$
A_2 = \sigma^2(\sigma_1 \cos \theta + \varepsilon_1 \sin \theta \cos \varphi)^2
$$

$$
+ \beta^2(\sigma_2 \sin \theta \cos \varphi + \beta \cos \theta \sin^2 \theta \sin^2 \varphi + \cos^2 \theta)^2,
$$

$$
A_3 = \sigma_2^2(1 - \beta \cos \theta)^2 + \left[\cos \theta + \beta \sin \theta(\sigma_2 \cos \varphi + \sin \theta \sin^2 \varphi)\right]^2. \quad (6)
$$

For a sinusoidal surface wave, $f(u) = a \sin u$, one has $f_m = \pm(a/2i)\delta_{m, \pm 1}$.

![Figure 2: Spectral density of radiation intensity as a function of the frequency of the radiated photon. The numbers near the curves correspond to the values of the ratio $d/\lambda_0$.](image)

We have numerically evaluated the spectral-angular distribution of the radiation intensity for various values of the parameters in the case of sinusoidal profile. In this case the only contribution to the radiation intensity comes from the harmonic $m = 1$. The results of these calculations show that the parameters of the radiation may be effectively controlled by tuning the characteristics of the surface wave. For an illustration we plot in figure 2 the spectral distribution of the radiation intensity

$$
\frac{dW}{d\omega} = \int_0^\pi d\theta \sin \theta \int_{-\pi/2}^{\pi/2} d\varphi \frac{dW}{d\omega d\Omega}. \quad (7)
$$
in the case of a sinusoidal surface wave as a function of the frequency of the radiated photon for various values of the ratio $d/\lambda_0$ (numbers near the curves), where $\lambda_0 = 2\pi/k_0$ is the wavelength of the surface wave. The full (dashed) curves correspond to the electron energy $E_e = 100$ MeV ($E_e = 500$ MeV). Note that $\omega/k_0c = \lambda_0/\lambda$ with $\lambda$ being the wavelength of the radiated photon. The corresponding radiation angle is related to the frequency by the relation $\cos \theta = 1/\beta - k_0c/\omega$.

In accordance with this relation, for relativistic electrons, large values of $\omega/(k_0c)$ correspond to small angles $\theta$. For $\theta \gg \gamma^{-1}$ the intensity is relatively insensitive to the particle energy, whereas for $\theta \lesssim \gamma^{-1}$ the intensity strongly increases with increasing energy. From figure 2 the suppression of the radiation intensity with increasing $d$ is well seen. In figure 3 we present the radiation intensity

$$\frac{dW}{d\omega d\varphi} = \int_0^\pi d\theta \sin \theta \frac{dW}{d\omega d\Omega},$$

as a function of the azimuthal angle $\varphi$ for various values of $\theta$ (numbers near the curves) and for $d/\lambda_0 = 0.15$. As in the case of figure 1, the full (dashed) curves correspond to the electron energy $E_e = 100$ MeV ($E_e = 500$ MeV). As we have mentioned, for $\theta \gg \gamma^{-1}$ the dependence of the radiation intensity on the particle energy is weak and for $\theta = 30^\circ, 90^\circ$ the full and dashed curves coincide. For small angles $\theta$ the radiation is emitted mainly along $\varphi \lesssim \gamma^{-1}$.

Figure 3: Spectral-angular density of radiation intensity as a function of the azimuthal angle of the radiated photon for various values of $\theta$.

3 Radiation from an electron bunch

In this section, as a source of the radiation we shall consider a cold bunch consisting of $N$ electrons and moving with constant velocity $v$ along the $z$ axis (for coherence effects in the Smith-Purcell radiation see also [5]). The density of the current in the bunch can be written in the form

$$j = qv \sum_{j=1}^N \delta (r - R_j - vt),$$

with $R_j = (X_j, Y_j, Z_j)$ being the position of the $j$-th particle at the initial moment $t = 0$. The spectral density of the radiation energy flux for a given $m$, defined by the relation

$$\frac{c}{4\pi} \int_{-\infty}^{+\infty} dt [\mathbf{E} \cdot \mathbf{H}] = \int_0^\infty d\omega \sum_m P_m^{(N)} (\omega),$$

for }
with \( \mathbf{E} \) and \( \mathbf{H} \) being the electric and magnetic fields, can be written in the form

\[
\mathbf{P}_m^{(N)}(\omega) = \mathbf{P}_m^{(1)}(\omega) S_N.
\]

In (11), \( \mathbf{P}_m^{(1)}(\omega) \) is the corresponding function for the radiation from a single electron with coordinates \( x = y = z = 0 \) at the initial moment \( t = 0 \), and

\[
S_N = \left| \sum_{j=1}^{N} \exp \left[ -\omega_1 \sigma X_j/v - ik_y Y_j - i\omega_1 Z_j/v \right] \right|^2.
\]

Assuming that the coordinates of the \( j \)-th particle are independent random variables and averaging the quantity (12) over the positions of a particle in the bunch, we obtain

\[
\langle \mathbf{P}_m^{(N)}(\omega) \rangle = \mathbf{P}_m^{(1)}(\omega) \langle S_N \rangle,
\]

where we have introduced the notations

\[
h = \langle \exp(2\omega X \text{Re } \sigma/v) \rangle, \quad h_l = \langle \exp(iK_l) \rangle, \quad l = x, y, z,
\]

\[
K_x = -i\omega \sigma/v, \quad K_y = k_y = (\omega \sqrt{\varepsilon_0/c}) \sin \theta \sin \varphi, \quad K_z = \omega_1/v.
\]

The functions \( |h_l|^2 \) determine the bunch form factors in the corresponding directions. In (13), the term proportional to \( N^2 \) determines the contribution of coherent effects. Conventionally it is assumed that the coherent radiation is produced at wavelengths equal and longer than the electron bunch length. However, as we shall see below, this conclusion depends on the distribution of electrons in the bunch.

As it is seen from (13), the form factors in \( y \) and \( z \) directions are determined by the Fourier transforms of the corresponding bunch distributions. First let us consider a Gaussian distribution:

\[
f_l = \frac{1}{\sqrt{2\pi \delta_l^2}} \exp \left( -\frac{l^2}{2\delta_l^2} \right), \quad l = x, y, z,
\]

with \( b_x, b_y \) and \( b_z \) being the corresponding characteristic sizes of the bunch. Assuming that all particles of the bunch are in the medium with permittivity \( \varepsilon_0 \) and, therefore, \( b_x < d - a \), with \( a \) being the surface wave amplitude, one finds

\[
\langle S_N \rangle = N \exp \left[ \frac{2\omega^2}{v^2} (\text{Re } \sigma)^2 b_x^2 \right] \left[ 1 + (N - 1) \exp \left( -\frac{\omega^2}{v^2} |\sigma|^2 b_x^2 - k_y b_y^2 - \frac{\omega^2 b_z^2}{v^2} \right) \right],
\]

where the second term in the square brackets determines the relative contribution of the coherent effects. For a non-relativistic bunch with \( b_y > b_{x,z} \), the corresponding exponent is equal to

\[
\exp[-(2\pi m/\lambda_0)^2(b_x^2 + b_y^2) - (2\pi/\lambda)^2 b_y^2 \sin^2 \theta \sin^2 \varphi]
\]

(here we consider the case \( \varepsilon_0 = 1 \)). Note that in this case for the wavelength of the radiation one has \( \lambda \sim \lambda_0/\beta m \) and the coherent effects are exponentially suppressed when \( \lambda \lesssim b_l/\beta \). For a relativistic bunch the relative contribution of coherent effects is given by

\[
N \exp\{-(2\pi/\lambda)^2[(b_x^2 + b_y^2) \sin^2 \theta \sin^2 \varphi + b_z^2])\},
\]

for \( \sin \theta \sin \varphi > \gamma^{-1} \), and by

\[
N \exp[-(2\pi b_x/\lambda \gamma)^2 - (2\pi/\lambda)^2 (b_z^2 + b_y^2 \sin^2 \theta \sin^2 \varphi)],
\]
for the radiation with \( \sin \theta \sin \varphi \lesssim \gamma^{-1} \). As we see, in this case the transverse form factor is strongly anisotropic. It follows from (16) that for a real \( \sigma \), fixed electron number and fixed distance of the bunch center from the surface wave the radiation intensity exponentially increases with increasing \( b_\perp \). This is because the number of electrons passing close to the surface wave increases. The number of electrons with long distances will also increase. But the contribution of close electrons surpasses the decrease of the intensity due to distant ones.

As we have seen, for a Gaussian distribution the relative contribution of coherent effects is exponentially suppressed in the case \( b_i > \lambda \). This result is a consequence of the mathematical fact that for a function \( f(x) \in C^\infty(R) \) one has the estimate

\[
F(u) \equiv \int_{-\infty}^{+\infty} f(l)e^{ilu}dl = O(u^{-\infty}), \quad u \to +\infty,
\]

where \( u = 2\pi b_\perp/\lambda \) for the longitudinal form factor of the relativistic bunch.

We have considered the case of Gaussian distribution in the bunch. However, it should be noted that due to various beam manipulations the bunch shape can be highly non-Gaussian. In the estimate (19) the continuity condition for the function \( f(l) \) and for infinite number of its derivatives is essential. It can be seen that when \( f(l) \in C^{n-1}(R) \), and the derivative \( f^{(n)}(l) \) is discontinuous at point \( l_1 \), we have the asymptotic estimate

\[
F(u) = (-iu)^{-n-1} \left[ f^{(n)}(l_1+) - f^{(n)}(l_1-) \right], \quad u \to +\infty.
\]

Unlike the case of (19), now the form factor for the short wavelengths decreases more slowly, as power-law, \( (2\pi b_\perp/\lambda)^{-n-1} \). In the coherent part of the radiation this form factor is multiplied by a large number, \( N \), of particles per bunch and the coherent effects dominate under the condition

\[ 2\pi b_\perp/\lambda < N^{1/2(n+1)}. \]

The radiation intensity is enhanced by the factor \( N(\lambda/2\pi b_\perp)^{2(n+1)} \). Since conventionally there are \( 10^8 - 10^{10} \) electrons per bunch, the condition (21) can be easily met even in the case \( 2\pi b_\perp/\lambda > 1 \). For instance, in the case of \( n = 1 \), \( N \sim 10^{10} \), the coherent radiation dominates for \( 100\lambda > b_i \).

As an example, let us discuss the case when the electrons are normally distributed in the \( x \) and \( y \) directions and the distribution function in the \( z \) direction has an asymmetric Gaussian form

\[
f(z) = \frac{2}{\sqrt{2\pi}(1+p)b_\perp} \left[ \exp \left( -\frac{l_0^2}{2p^2b_\perp^2} \right) \theta(-l) + \exp \left( -\frac{l_0^2}{2b_\perp^2} \right) \theta(l) \right],
\]

where \( \theta(l) \) is the unit step function, \( l_0 = (1+p)b_\perp \) is the characteristic bunch length, parameter \( p \) determines the degree of bunch asymmetry. Now the expression for \( \langle S_N \rangle \) can be written as

\[
\langle S_N \rangle = N \exp \left( \frac{2\omega^2}{v^2} \sigma^2 b_\perp^2 \right) \left[ 1 + (N-1) \exp \left( -\frac{\omega^2}{v^2} \sigma^2 b_\perp^2 - k_y^2 b_y^2 \right) \right] |F(\omega_1/v)|^2,
\]

where

\[
F(u) = \frac{1}{p+1} \left\{ e^{-t^2} + pe^{-p^2t^2} - \frac{2i}{\sqrt{\pi}} [W(t) - pW(pt)] \right\},
\]

with the notation

\[
W(t) = \int_0^t exp(l^2 - t^2)dl, \quad t = \frac{ub_\perp}{\sqrt{2}}
\]

The expression \(|F(u)|^2\) is invariant with respect to the replacement \( p \to 1/p, b_i \to b_i p \) that corresponds to the mirror reversal of the bunch. When the electron distribution is symmetric.
(p = 1), the second summand in the square brackets of (24) vanishes and, as was mentioned earlier, the form factor exponentially decreases for short wavelengths $\lambda < 2\pi b_1/\beta$. For $pt \gg 1$ the asymptotic behaviour of the function $F(u)$ is found from (24) and has the form

$$F(u) \sim i \sqrt{\frac{2}{\pi}} \frac{1 - p}{u^{3/2}p}. \quad (26)$$

In the case of a relativistic bunch and for $\varphi > \gamma^{-1}$, the exponent of the second summand in the square brackets of (23) is of the order $(2\pi b_1/\lambda)^2$, $i = x, y$. When the transverse size of the bunch is shorter than the radiation wavelength, then the relative contribution of the coherent effects is $\sim N |F(\omega_1/v)|^2$. For an asymmetrical bunch this contribution can be dominant even in the case when the bunch length is greater than the radiation wavelength. Indeed, according to (26) even for weakly asymmetrical bunch we have $N |F|^2 \sim N(v/\omega_1b_2)^6$, and the radiation is coherent for $b_x \lesssim \lambda N^{1/6}/(2\pi)$, where we have taken into account that for a relativistic bunch $\omega \gg m\omega_0$ and therefore $\omega_1 \approx \omega$, as was mentioned above. In the case of the bunch with $N \sim 10^{10}$ and for $2\pi b_2/\lambda \lesssim 10$ the radiation is coherent. For the radiation $\varphi \lesssim \gamma^{-1}$, the exponent of the second summand in square brackets of Eq. (23) for $\varphi_0 = 1$ is of the order $(2\pi b_2/\lambda)^2$. Hence, for a relativistic bunch the radiation in directions $\varphi \lesssim \gamma^{-1}$ can be coherent even in the case when the transverse size of the bunch is greater than the wavelength. For this it is sufficient to have the conditions

$$b_x \lesssim \frac{\gamma \lambda}{2\pi}, \quad b_z \lesssim \frac{\lambda N^{1/6}}{2\pi}. \quad (27)$$

The second of these conditions is written for weakly asymmetrical bunches. In the case of strong asymmetry the corresponding conditions are less restrictive: $b_z \lesssim \lambda N^{1/2}/2\pi$ for $pb_2 \ll \lambda$.

Let $f(l, a)$ be a continuous distribution function depending on the parameter $a$, and $\lim_{a \to 0} f(l, a) = f(l)$. The integral $F(u, a)$ for $f(l, a)$ uniformly converges and hence $\lim_{a \to 0} F(u, a) = F(u)$. It follows from here that the estimate presented above is valid for continuous functions as well if they are sufficiently close to the corresponding discontinuous function (the corresponding derivative is sufficiently large, see below). Aiming to illustrate this, we consider the asymmetric distribution

$$f(z, a_l, b_l, l_0) = \frac{1}{4l_0} \left[ \text{th} \left( \frac{l + l_0}{a_l} \right) - \text{th} \left( \frac{l - l_0}{b_l} \right) \right]. \quad (28)$$

For $l_0 > a_l, b_l$ this function describes a rectangular bunch having exponentially decreasing asymmetric tails with characteristic sizes $a_l$ and $b_l$. In the limit $a_l, b_l \to 0$ one obtains the rectangular distribution with the bunch length $2l_0$. The explicit evaluation of the expression (19) with the function (28) leads to

$$F(u, a_l, b_l, l_0) = \frac{i}{2ul_0} \left( \frac{\overline{a} e^{-iul_0}}{\sinh \overline{a}} - \frac{\overline{b} e^{iul_0}}{\sinh \overline{b}} \right), \quad (29)$$

with the notations $\overline{a_l} \equiv \pi u a_l/2$ and $\overline{b_l} \equiv \pi u b_l/2$. From here it follows that if $\overline{a_l} \sim 1$ and $\overline{b_l} \sim 1$ then $F \sim (ul_0)^{-1}$ for $ul_0 \gg 1$. In the limit $a_l, b_l \to 0$, from (29) one obtains the well known form factor for the rectangular distribution: $F_{\text{rect}}(u, l_0) = \sin(ul_0)/(ul_0)$. As it is seen, the rectangular distribution is a good approximation for (28) if $a_l, b_l \ll \lambda$. The main contribution to (29) comes from the bunch tails, i.e. from the parts of bunch with large derivatives of the distribution function. This is the case for the general case of distribution function as well: if $ul_0 \gg 1$ the main contribution comes from the parts of the bunch where $df/dl(l/l_0) \gtrsim u$ and in this case $F(u) \sim 1/u$, $u \to \infty$. This can be generalized for higher derivatives as well: if $d^i f/dl(l/l_0)^i \ll u$, $i = 1, \ldots, n - 1$, and $d^n f/dl(l/l_0)^n \gtrsim u$ then $F(u) \sim (ul_0)^{-n}$.
For a relativistic bunch one has $u \sim \frac{2\pi}{\lambda}$ for the form factors in $y$ and $z$ directions and the conclusion can be formulated as follows. If for the distribution function one has

$$\lambda \frac{d^i f}{d(l/l_0)^i} \ll 2\pi, \quad i = 1, \ldots, n - 1; \quad \lambda \frac{d^n f}{d(l/l_0)^n} \gtrsim 2\pi, \quad (30)$$

then the relative contribution of coherent effects into the radiation intensity is proportional to $N(\lambda/2\pi l_0)^{2n}$ and the radiation is partially coherent in the case $\lambda < l_0$ but $\lambda > 2\pi l_0 N^{-1/2n}$, with $l_0$ being the characteristic bunch size in the corresponding direction. In this case the main contribution into the radiation intensity comes from the parts of the bunch with large derivatives of the distribution function in the sense of the second condition in (30). For example, in the case of asymmetric distribution (22), when $ub_l \gg 1$ and $p ub_l < 1$, the main contribution comes from the left Gaussian tail with $l < 0$. For this tail $df/d(l/b_l) \sim u/(p ub_l) \gtrsim u$, at $l \sim pb_l$ and therefore $F(u) \sim 1/(ub_l)$. This can be seen directly from the exact relation (24) as well by using the asymptotic formula for the function $W(t)$.

4 Conclusion

In the present talk we have considered the radiation from a single electron and from an electron bunch of arbitrary structure flying over the surface wave excited in a plane interface. For small amplitudes of the surface wave the spectral-angular distribution of the radiation intensity from a single electron is given by expression (4). We have shown that at large angles with respect to the electron trajectory the dependence of the radiation intensity on the electron energy is relatively weak, whereas at small angles the intensity strongly increases with increasing energy. It is demonstrated that the radiation from a bunch can be partially coherent in the range of wavelengths much shorter than the characteristic longitudinal size of the bunch and the main contribution to the radiation intensity comes from the parts of the bunch with large derivatives of the distribution function. In this case the short wavelengths the relative contribution of coherent effects decreases as a power-law instead of exponentially decreasing. The corresponding conditions for the distribution function are specified. The coherent effects lead to an essential increase in the intensity of the emitted radiation.

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