Diffusive radiation in one-dimensional Langmuir turbulence

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We calculate spectra of radiation produced by a relativistic particle in the presence of one-dimensional Langmuir turbulence which might be generated by a streaming instability in the plasma, in particular, in the shock front or at the shock-shock interactions. The shape of the radiation spectra is shown to depend sensitively on the angle between the particle velocity and electric field direction. The radiation spectrum in the case of exactly transverse particle motion is degenerate and similar to that of spatially uniform Langmuir oscillations. In the case of oblique propagation, the spectrum is more complex, it consists of a number of power-law regions and may contain a distinct high-frequency spectral peak. The emission process considered is relevant to various laboratory plasma settings and for astrophysical objects as gamma-ray bursts and collimated jets.

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Various kinds of two-stream instability including longitudinal, transverse, and oblique modes [1–4] are common for many astrophysical objects [5–7] and laboratory plasma settings [8,9]. These instabilities are efficient to produce a high level of magnetic and/or electric turbulence in the plasma. We note that the electric turbulence is of primary importance for many applications, in particular, it is capable of charged particle acceleration [7,10,11] in contrast to purely magnetic turbulence, which is only efficient for angular scattering of the charged particle while it cannot change the particle energy directly. Remarkably, recent theoretical and numerical studies [12–14] demonstrate that the energy density going to the electric (Langmuir) turbulence can be very large, exceeding, for example, the energy density of the initial regular magnetic field in the plasma.

Relativistic electrons propagating in the plasma with developed Langmuir turbulence will random walk due to random Lorenz forces produced by these turbulent electric fields. Apparently, during this random diffusive motion the electron will generate electromagnetic emission. We will refer to this radiative process as *diffusive* radiation in Langmuir waves (DRL) to emphasize the key role of the diffusive random walk of the particle in the stochastic electric fields.

Curiously, the theory of DRL has not been developed in sufficient detail yet, although a few particular issues related to this radiative process (called also *electrostatic Bremsstrahlung*) have been considered [15–20]. However, the detailed treatment of the DRL spectral shape in various regimes is currently unavailable. This Brief Report attempts to remedy the situation by calculating the DRL spectrum within the perturbation theory and determines the region of applicability of the perturbative treatment.

Electromagnetic emission produced by a charged particle can be calculated within the perturbation theory when the particle moves almost rectilinearly with almost constant velocity. Apparently, the nonzero acceleration of the particle in the external field should be taken into account to obtain a nonzero radiation intensity. This perturbative treatment is widely used because of its simplicity. Frequently, one calculates first the particle acceleration $\mathbf{w}(t)$ due to a given field along the rectilinear trajectory and then uses this expression obtained for $\mathbf{w}(t)$ to find the radiation spectrum. In the case of a random external field, however, when $\mathbf{w}(t)$ is also a random function of time t it is more convenient to express the radiation intensity via the spatial and temporal spectrum of the external electric and/or magnetic field directly [21]. Accordingly, the spectral and angular distribution of the emission produced by a single particle in a plasma with random field has the form

$$W^{\perp}_{\Omega,\omega}=\frac{(2\pi)^3Q^2}{M^2c^3\gamma^2V}\left(\frac{\omega}{\omega'}\right)^2\left[1-\frac{\omega}{\omega'\gamma_*^2}+\frac{\omega^2}{2\omega'^2\gamma_*^4}\right]\times\int dq_0d\mathbf{q}\,\delta(\omega'-q_0+\mathbf{q}\cdot\mathbf{v})|\mathbf{F}_{q_0,\mathbf{q}\perp}|^2, \tag{1}$$

where $\gamma_*=\left(\gamma^{-2}+\frac{\omega_{pe}^2}{\omega^2}\right)^{-1/2}$, Q, M, and γ are the charge, mass, and Lorenz-factor of the emitting particle, c is the speed of light, ω_{pe} is the plasma frequency, V is the volume of the emission source, $\mathbf{F}_{q_0,\mathbf{q}\perp}$ is the temporal and spatial Fourier component of the Lorenz force transverse to the emitting particle velocity,

$$\omega'=\frac{\omega}{2}\left(\gamma^{-2}+\theta^2+\frac{\omega_{pe}^2}{\omega^2}\right), \tag{2}$$

θ is the emission angle relative to the particle velocity vector $\mathbf{v}$, and ω is the frequency of the emitted wave. Contribution $W^{\perp}_{\Omega,\omega}$ (marked with the superscript $\perp$) is provided by a component of the particle acceleration transverse to the particle velocity. In the case of the electric $\mathbf{E}$ (in contrast to magnetic) field, there is also a component of the acceleration along the particle velocity. The corresponding contribution has the form

$$W^{\parallel}_{\Omega,\omega}=\frac{2(2\pi)^3Q^4}{M^2c^3\gamma^6V}\left(\frac{\omega}{\omega'}\right)^3\left[1-\frac{\omega}{2\omega'\gamma_*^2}\right]\times\int dq_0d\mathbf{q}\,\delta(\omega'-q_0+\mathbf{q}\cdot\mathbf{v})|\mathbf{E}_{q_0,\mathbf{q}\parallel}|^2, \tag{3}$$

which is typically small by a factor γ^{-2} compared with the transverse contribution. Nevertheless, there exist special cases when the transverse contribution is zero or very small and the parallel contribution comes to play. In particular, for the considered here one-dimensional turbulent electric field the parallel contribution will dominate for a particle moving along the field direction.

Let us start with the case when the particle moves at a large angle to the Langmuir turbulence direction $\vartheta \gg \gamma^{-1}$, so the standard transverse contribution dominates. Following the derivation given in [21], but with the Lorenz force $\mathbf{F} = Q\mathbf{E}$ specified by electric $\mathbf{E}$ in the place of magnetic $\mathbf{B}$ field, it is easy to find

$$|\mathbf{F}_{q_0,\mathbf{q}\perp}|^2 = Q^2|\mathbf{E}_{q_0,\mathbf{q}\perp}|^2 = \frac{TV}{(2\pi)^4}Q^2\left(\delta_{\alpha\beta} - \frac{v_\alpha v_\beta}{v^2}\right)K_{\alpha\beta}(q_0,\mathbf{q}), \tag{4}$$

where $K_{\alpha\beta}(q_0,\mathbf{q}) = C_{\alpha\beta}K(q_0,\mathbf{q})$, T is the total time of emission; $C_{\alpha\beta}$ describes the longitudinal nature of the Langmuir waves, i.e., $C_{\alpha\beta} = q_\alpha q_\beta / q^2$, while $K(q_0,\mathbf{q})$ is the temporal and spatial spectrum of the Langmuir turbulence.

The developed approach allows for arbitrary anisotropy of the turbulence, although we have to specify the shape of the turbulence spectrum to promote further the calculation of the DRL spectrum. Analytical and numerical studies of the two-stream instabilities suggest that the Langmuir turbulence produced is frequently highly anisotropic [1,3,4,12–14], which is confirmed also by available *in situ* observations, e.g., performed in the Earth's magnetosphere [22]. Although this anisotropy can be reduced at later stages of the nonlinear turbulence evolution due to randomization of the wave vector directions, here we assume that the Langmuir turbulence is highly anisotropic, namely, one-dimensional. As an example, we can suppose that all the wave vectors are directed along the shock normal $\mathbf{n}$, therefore $C_{\alpha\beta} = n_\alpha n_\beta$, while the spectrum $K(q_0,\mathbf{q})$ can be approximated by a power law over $q_\parallel$ above a certain critical value k_0:

$$K(q_0,\mathbf{q}) = \frac{a_\nu k_0^{\nu-1}\langle E_L^2\rangle}{(k_0^2 + q_\parallel^2)^{\nu/2}}\delta(\mathbf{q}_\perp)\,\delta(q_0 - \omega_{\rm pe}). \tag{5}$$

Here, the presence of the second δ-function is related to the assumption that the electric turbulence is composed of Langmuir waves all of which oscillate in time with the same frequency $\omega_{\rm pe}$; the normalization constant a_ν is set up by the condition $\int K(q_0,\mathbf{q})dq_0 d\mathbf{q} = \langle E_L^2\rangle$, where $\langle E_L^2\rangle$ is the mean square of the electric field in the Langmuir turbulence.

Now, substituting Eq. (4) with Eq. (5) into general expression (1), taking the integrals over dq_0, $dq_\parallel$, and $dq_\perp$ with the use of three available δ-functions, and dividing by the total (infinite) time of emission T we find

$$I^{\perp}_{\Omega,\omega} = \frac{a_\nu k_0^{\nu-1}\langle E_L^2\rangle Q^4 \sin^2\vartheta}{2\pi M^2 c^4 \gamma^2 |\cos\vartheta|}\left(\frac{\omega}{\omega'}\right)^2 \times\left[1 - \frac{\omega}{\omega'\gamma_*^2} + \frac{\omega^2}{2\omega'^2\gamma_*^4}\right]\left\{\left(\frac{\omega' - \omega_{\rm pe}}{\mathbf{n}\cdot\mathbf{v}}\right)^2 + k_0^2\right\}^{-\nu/2}, \tag{6}$$

where ϑ is the angle between the particle velocity $\mathbf{v}$ and vector $\mathbf{n}$. This expression diverges formally when $\cos\vartheta \to 0$. Accordingly, for a particle moving transversely to the Langmuir turbulence direction we need to recalculate the integrals in Eq. (1) taking into account that $\delta(\omega' - q_0 + \mathbf{q}\cdot\mathbf{v}) \to \delta(\omega' - q_0)$ for $\mathbf{q}\cdot\mathbf{v} = 0$, which yields

$$I^{\perp}_{\Omega,\omega} = \frac{Q^4\langle E_L^2\rangle}{2\pi M^2 c^3 \gamma^2}\left(\frac{\omega}{\omega'}\right)^2\left[1 - \frac{\omega}{\omega'\gamma_*^2} + \frac{\omega^2}{2\omega'^2\gamma_*^4}\right]\delta(\omega' - \omega_{\rm pe}), \tag{7}$$

in full agreement with the results of [17] obtained for spatially uniform Langmuir oscillations. This is not a random coincidence. Indeed, since the one-dimensional Langmuir waves experience spatial variations along only one direction (of vector $\mathbf{n}$), the particle moving transversely to this direction "feels" a spatially uniform field pattern such as that considered in [17].

There is another special geometry when intensity (6) is insufficient to describe the DRL spectrum: it is the case of particle motion along vector $\mathbf{n}$. Indeed, the "parallel" contribution becomes important for $\vartheta \lesssim \gamma^{-1}$. Substituting Eq. (5) into Eq. (3) and taking the integrals similarly to the derivation of Eq. (6) we find

$$I^{\parallel}_{\Omega,\omega} = \frac{a_\nu k_0^{\nu-1}\langle E_L^2\rangle Q^4|\cos\vartheta|}{\pi M^2 c^4\gamma^6}\left(\frac{\omega}{\omega'}\right)^3 \times\left[1 - \frac{\omega}{2\omega'\gamma_*^2}\right]\left\{\left(\frac{\omega' - \omega_{\rm pe}}{\mathbf{n}\cdot\mathbf{v}}\right)^2 + k_0^2\right\}^{-\nu/2}. \tag{8}$$

Apparently, spectra (6) and (8) look rather differently compared with the spectrum (7) produced by a particle moving transversely to vector $\mathbf{n}$. Given, in particular, that the frequency ω' is directly linked with the emission angle θ, the δ-function $\delta(\omega' - \omega_{\rm pe})$ permits emitting a single frequency only in each direction. By comparison, no δ-function enters Eqs. (6) and (8), thus a continuum spectrum rather than distinct frequencies is emitted along any direction. Clearly, there remains a distinct contribution to the emission intensity when $\omega' \approx \omega_{\rm pe}$. However, the range of the parameter space where this resonant condition holds is relatively narrow, so the "nonresonant" contribution from the remaining part of the parameter space where $\omega' \neq \omega_{\rm pe}$ can easily dominate the resonant contribution. To see this explicitly, consider the radiation intensity into the full solid angle by integration of Eq. (6) over $d\Omega = \sin\theta d\theta d\varphi \approx 2\pi d(\omega'/\omega)$ that yields

$$I^{\perp}_{\omega} = \frac{a_\nu k_0^{\nu-1}\langle E_L^2\rangle Q^4\sin^2\vartheta}{M^2c^4\gamma^2|\cos\vartheta|}\int_{1/2\gamma_*^2}^{\infty} d\left(\frac{\omega'}{\omega}\right)\left(\frac{\omega}{\omega'}\right)^2 \times\left[1 - \frac{\omega}{\omega'\gamma_*^2} + \frac{\omega^2}{2\omega'^2\gamma_*^4}\right]\left\{\left(\frac{\omega' - \omega_{\rm pe}}{\mathbf{n}\cdot\mathbf{v}}\right)^2 + k_0^2\right\}^{-\nu/2}. \tag{9}$$

Let us analyze essential properties of the DRL spectrum on the basis of the asymptotic evaluation. At low frequencies $\omega \ll \omega_{\rm pe}\gamma^2$, we can discard ω' in Eq. (9) everywhere in the braces except a narrow region of parameters when $\omega' \approx \omega_{\rm pe}$. This means that for $\omega \ll \omega_{\rm pe}\gamma^2$ the integral is composed of two contributions. The first of them, a nonresonant one,

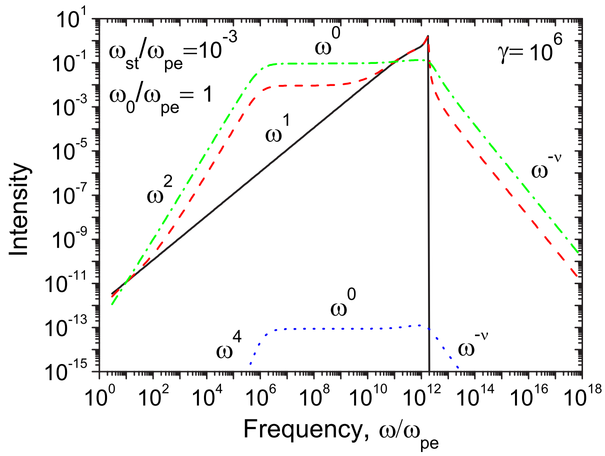

FIG. 1. (Color online) DRL spectra produced by a particle with $\gamma=10^6$ in a plasma with developed one-dimensional Langmuir turbulence for various particle propagation directions: $\cos\vartheta=0$, solid (black) curve; $\cos\vartheta=10^{-3}$, dashed (red) curve; $\cos\vartheta=0.5$, dash-dotted (green) curve; and $\cos\vartheta=1$, dotted (blue) curve. Parameters are given in the figure. The "parallel" contribution [blue (dotted) curve] is very small (10^{-12}) for the highly relativistic particle, although it becomes competing for moderately relativistic particles.

arises from integration over the region, where $\omega' \ll \omega_{pe}$. Here, the emission is beamed within the characteristic emission angle of $\vartheta \sim \gamma^{-1}$ along the particle velocity. The integral converges rapidly, and so it may be taken along the infinite region, which produces a flat radiation spectrum, $I_\omega \propto \omega^0$ or $I_\omega \propto \omega^2$ at lower frequencies, $\omega < \omega_{pe}\gamma$. However, as far as ω' approaches ω_{pe}, a resonant contribution comes into play. Now, in a narrow vicinity of ω_{pe}, we can adopt

$$\left\{\left(\frac{\omega' - \omega_{pe}}{\mathbf{n}\cdot\mathbf{v}}\right)^2 + k_0^2\right\}^{-\nu/2} \propto \delta(\omega' - \omega_{pe}), \qquad (10)$$

which results in a single-wave-like contribution, $I_\omega \propto \omega^1$. The full spectrum at $\omega < \omega_{pe}\gamma^2$, therefore, is just a sum of these two contributions. At high frequencies, $\omega \gg \omega_{pe}\gamma^2$, the term ω' dominates in the braces, so other terms can be discarded. Thus a power-law tail $I_\omega \propto \omega^{-\nu}$ arises in this spectral range.

Although this integral cannot be taken analytically in a general case, it is easy to obtain corresponding spectra numerically. Figure 1 displays a numerically integrated example of the DRL spectra for various angles between $\mathbf{v}$ and $\mathbf{n}$ for a highly relativistic particle with $\gamma=10^6$. The black (solid) curve shows the DRL spectrum arising as the particle moves strictly perpendicular to vector $\mathbf{n}$, which coincides with the spectrum arising in the case of spatially uniform Langmuir oscillations [17]. The spectrum consists of a rising region $I_\omega \propto \omega^1$ at $\omega < 2\omega_{pe}\gamma^2$ and drops abruptly to zero at $\omega > 2\omega_{pe}\gamma^2$. Remarkably, there are prominent differences between this spectrum and those generated for oblique particle propagation even though the spectra can be similar to each other in the immediate vicinity of the spectral peak.

To find the applicability region of the perturbation theory applied above, we should estimate the characteristic deflection angle of the emitting electron on the emission coherence length $l_c=2c\gamma_*^2/\omega$, where the elementary emission pattern is formed [21]. Consider a simple source model consisting of uncorrelated cells with the size $l_0=2\pi c/\omega_0$, each of which contains coherent Langmuir oscillations with the plasma frequency ω_{pe}. Inside each cell the electron velocity can change by the angle $\theta_0 \sim \omega_{st}/(\omega_{pe}\gamma)$ if $\omega_0 \lesssim \omega_{pe}$, where $\omega_{st} = Q\langle E_L^2\rangle^{1/2}/Mc$ (in the other case, $\omega_0 > \omega_{pe}$, the results of [21] apply). Then, after traversing $N=l_c/l_0$ cells, the mean square of the deflection angle is $\theta_c^2=\theta_0^2 N \sim \omega_{st}^2\omega_0/(\omega\omega_{pe}^2)$. The perturbation theory is only applicable if this diffusive deflection angle is smaller than the relativistic beaming angle, γ^{-1}, i.e., it is always valid at sufficiently high frequencies $\omega > \omega_* \equiv \omega_{st}^2\omega_0\gamma^2/\omega_{pe}^2$. Note that the bounding frequency ω_* increases with ω_0, while diffusive synchrotron radiation (DSR) in the random magnetic field displays the opposite trend. The perturbation theory will be applicable to the entire DRL spectrum if the condition $\theta_c^2 < \gamma^{-2}$ holds for the frequency $\omega_{pe}\gamma$ [21], where the coherence length of the emission has a maximum. This happens for the particles whose Lorenz factors obey the inequality

$$\gamma < \omega_{pe}^3/(\omega_{st}^2\omega_0). \qquad (11)$$

Let us compare the DRL spectra calculated numerically (Fig. 1) for the case of a broad turbulence distribution over spatial scales with the DSR spectra in the case of stochastic magnetic fields [21]. Apparently, there is a remarkable difference between these two emission mechanisms, especially in the case of the long-wave turbulence, $\omega_0|\cos\vartheta| \equiv k_0 c|\cos\vartheta| \ll \omega_{pe}$. First, we note that the perturbation theory of DRL has a broader applicability region, in particular, it applies for higher energy electrons than the perturbative version of the DSR theory [21] since the criterion (11) is softer than the corresponding criterion in [21]. This happens because in the presence of the Langmuir turbulence the rapid temporal oscillations of the electric field direction substantially compensate angular deflections of the particle, so the average trajectory is much more similar to a straight line than for the case of the random magnetic field with the same ω_0 and ω_{st}. Then, a distinct spectral peak at $\omega=2\omega_{pe}\gamma^2$ is formed with the linear decrease of the spectrum with frequency, which is not present in case of the DSR. At lower frequencies, however, this falling part of the spectrum gives way to a flat spectrum, which is entirely missing within the one-wave approach [17] and absent for exactly transverse particle motion. Position of the corresponding turning point depends on the $\omega_0\cos\vartheta/\omega_{pe}$ ratio in such a way that for $\omega_0|\cos\vartheta| \gtrsim \omega_{pe}$ the flat spectral region entirely dominates the range from $\omega_{pe}\gamma$ to $\omega_{pe}\gamma^2$. It is worth emphasizing that the deviations of the DRL spectrum from the single-wave spectrum (read, from the transverse case, $\cos\vartheta=0$) are prominent even for oblique propagation angles only slightly different from $\pi/2$, e.g., $\cos\vartheta=10^{-3}$ as in Fig. 1. Therefore the presence of a broad turbulence spectrum considered here in detail results in important qualitative change of the emission mechanism, which cannot generally be reduced to a simplified treatment relying on the single-wave approximation with some rms value of the Langmuir electric field.

Modern computer simulations of shock wave interactions, especially in the relativistic case, suggest that the energy density of the excited Langmuir turbulence can be very large [13], e.g., far in excess of the energy of the initial regular magnetic field. In particular, at the shock wave front the electric field can be as strong as the corresponding wave-breaking limit, i.e., $\omega_{st} \sim \omega_{pe}$ [14]. In this case the random walk of relativistic electrons in the stochastic electric field can give rise to a powerful contribution in the nonthermal emission of an astrophysical object, entirely dominating full radiation spectrum or some broad part of it.

Although any detailed application of the considered emission process is beyond the scope of this Brief Report, we mention that the DRL is a promising mechanism for the gamma-ray bursts and extragalactic jets. In particular, some of the prompt gamma-ray burst emission displays rather hard low-energy spectra with the *photon* spectral index α up to 0. The DRL spectral asymptote $I_\omega \propto \omega^1$, which appears just below the spectral peak at $2\omega_{pe}\gamma^2$, fits well to those spectra. Remarkably, the flat lower-frequency asymptote, $I_\omega \propto \omega^0$, can account for the phenomenon of the x-ray excess [23,24] and prompt optical flashes accompanying some gamma-ray bursts.

In addition, this mechanism along with the DSR in random magnetic fields [25] can be relevant to the UV-x-ray flattenings observed in some extragalactic jets. For example [26], x-ray observations of the jet in 3C 273 look inconsistent with the standard synchrotron of DSR models. Remarkably, the entire UV-to-x-ray spectrum of 3C 273 might be produced by DRL, which can be much flatter than usual DSR in the range $\omega_{Be}\gamma^2 \ll \omega \ll \omega_{pe}\gamma^2$.

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