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Nonlinear excitation of longitudinal photons and plasmons by high-power, short pulse lasers

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Modulational excitation of longitudinal photons and electron Langmuir waves, as well as ion sound waves by an incoherent strong and superstrong radiation (high-power short pulse lasers, nonthermal equilibrium cosmic field radiation, etc.), in plasmas, is investigated. A simultaneous generation of longitudinal photons and plasmons is demonstrated. Furthermore, the kinetic instability is considered when low-frequency longitudinal photons are generated alone. The growth rates of these modes are obtained.

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Electromagnetic (EM) radiation in plasmas is a fundamental physical system which has played a crucial role in opening up frontiers in physics, such as the fast ignition in laser fusion [1], plasma-based high-energy particle acceleration [2,3], electron-positron and neutrino-antineutrino pair production [4–7], nonlinear optics [8], optically induced nuclear fission [9], etc. One of the most salient phenomena in the above-mentioned problems of laboratory plasmas, as well as astrophysical and space plasmas, are relativistic parametric and modulational instabilities, which gained prominence largely due to their importance in connection with the strongly nonlinear structures in plasmas, the deposition of EM energy through different modes into the plasma, and heating and/or acceleration of plasma particles to very high energies. Relativistic parametric instabilities have a long history starting with the work of Tsintsadze [10]. Since then, a great number of theoretical and simulation results have been reported [11]. Recent progress in the development of high-power, short pulse lasers has renewed the interest in this phenomenon [12–15]. The above treatments were restricted to the case of monochromatic EM waves. For ultrashort pulses the bandwidth of coherent waves is increasingly broad. Even if the bandwidth may be initially narrow, its spectrum may eventually broaden, either as a result of several kinds of instability processes or as the result of other nonlinear wave-wave interaction processes.

In order to study the interaction of spectrally broad relativistically intense EM waves with a plasma, in previous papers [16,17] starting from the fully relativistic equations, we have derived a general kinetic equation for the photon gas incorporating two forces of distinct nature. One force appears due to the redistribution of electrons in space, ∇n_e, and time, $\frac{\partial n_e}{\partial t}$. The other force arises by the variation of the shape of wave packets. In other words, this force originates from alteration of the average kinetic energy of the electron oscillating in a rapidly varying field of EM waves and is proportional to $\nabla\frac{1}{\gamma}$ and $\frac{\partial}{\partial t}\frac{1}{\gamma}$, where γ is the relativistic Lorentz factor.

In the field of superstrong femtosecond pulses, it is expected that the character of the nonlinear response of a medium would radically change. At high intensities the motion of free electrons near the focal volume would be extremely relativistic. Thus, the relativistic nonlinear effect, which is basically associated with an increase in the electron mass, will tend to determine the dynamics of EM pulses. Currently, lasers produce pulses whose intensity approaches 10^{22} W/cm^2 [18]. At these intensities, the highly nonlinear and complicated relativistic dynamics of the laser-plasma system gives rise to a number of interesting phenomena—for example, Bose-Einstein condensation (BEC) and a intermediate state of the photon gas [19]. It was shown in Ref. [20] that the behavior of photons in a plasma is radically different from the one in a vacuum. Namely, plasma particles perform oscillatory motion in the field of EM waves affecting the radiation field. Photons acquire the rest mass m_γ and become one of the bosons in plasmas and possess all characteristics of nonzero rest mass, i.e., we may say that the photon is an elementary particle of the optical field. This difference leads to certain phenomena, such as longitudinal photons (photonikos) [21–23], which originate from the decay process. Namely, the photon passing through the photon bunch absorbs and emits longitudinal photons, with frequencies $\Omega = \mp(\omega-\omega')$ and wave vectors $\vec{q}=\mp(\vec{k}-\vec{k}')$. In Refs. [19,21] we have shown that under certain conditions photon-photon interactions dominate photon-plasma particle interactions. So that under such conditions the variation of the plasma density can be neglected in comparison with the variation of the photon density. In recent two-dimensional fully relativistic particle-in-cell simulations (Fourier space code EPIC3D-AP) [24], which is based on a potential form, simultaneous emission of longitudinal photons and plasmons by the laser pulse was observed, as well as low-frequency photons that can be associated with BEC photons were seen. Thus, it is of interest to examine these modes.

In this Rapid Communication, we study analytically the nonlinear excitation of longitudinal photons and plasmons by high-power, short pulse lasers. As we will see in the following, there are cases when both longitudinal photons and plasmons are simultaneously generated that confirm the results of recent simulations [24]. For our purpose, we employ the general dispersion relation obtained in Refs. [16,17], which reads

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$$\varepsilon\left(1+\frac{\omega_{Le}^2}{2\gamma_0^2}\int\frac{d\mathbf{k}}{(2\pi)^3}\frac{\mathbf{P}_0^+-\mathbf{P}_0^-}{\mathbf{qk}c^2-\Omega\omega(k)}\right)+(1+\delta\varepsilon_i)\delta\varepsilon_e\frac{q^2c^2}{2\gamma_0^2}\int\frac{d\mathbf{k}}{(2\pi)^3}\frac{\mathbf{P}_0^+-\mathbf{P}_0^-}{\mathbf{qk}c^2-\Omega\omega(k)}=0, \quad (1)$$

where

$$\varepsilon=1+\delta\varepsilon_e+\delta\varepsilon_i, \quad \delta\varepsilon_\alpha=\frac{4\pi e^2}{q^2}\int\frac{(\mathbf{q}\partial f_{0\alpha}/\partial\mathbf{p})}{\Omega-\mathbf{qv}_\alpha}d\mathbf{p},$$

$$\omega_{L\alpha}^2=\frac{4\pi e^2 n_{0e}}{m_{0\alpha}\gamma_0},$$

$$\mathbf{P}_0^\pm=\mathbf{P}_0(\mathbf{k}\pm\mathbf{q}/2), \quad (2)$$

$\gamma_0=\sqrt{1+Q_0}=\sqrt{1+2\int\frac{dk}{(2\pi)^3}P_0}$ and $P_0=\frac{e^2A(k)A^*(k)}{(m_{0e}c^2)^2}$ is the spectral function. Equation (1) is rewritten as

$$\varepsilon\left(1+\frac{1}{A}\right)+(1+\delta\varepsilon_i)\delta\varepsilon_e\frac{q^2c^2}{\omega_{Le}^2}=0, \quad (3)$$

where

$$A=\frac{\omega_{Le}^2}{2\gamma_0^2}\int\frac{d\mathbf{k}}{(2\pi)^3}\frac{\mathbf{P}_0^+-\mathbf{P}_0^-}{\mathbf{qk}c^2-\Omega\omega(k)}=\frac{\omega_{Le}^2}{2\gamma_0^2}\int\frac{d\mathbf{k}}{(2\pi)^3}P_0\left\{\frac{1}{\Omega-\mathbf{qu}+q^2c^2/2\omega}-\frac{1}{\Omega-\mathbf{qu}-q^2c^2/2\omega}-\imath\pi\left[\delta\left(\Omega-\mathbf{qu}+\frac{q^2c^2}{2\omega}\right)-\delta\left(\Omega-\mathbf{qu}-\frac{q^2c^2}{2\omega}\right)\right]\right\}. \quad (4)$$

If $A\gg 1$, then Eq. (3) reduces to Eq. (20) of Ref. [16].

For the photon gas we use the spectral Gaussian distribution function

$$P_0=Q_0(2\pi\sigma_k^2)^{-3/2}\exp\left(-\frac{(\mathbf{k}-\mathbf{k}_0)^2}{2\sigma_k^2}\right), \quad (5)$$

where

$$Q_0=\frac{e^2|A(k_0)|^2}{(m_{0e}c^2)^2}=\gamma_0^2-1.$$

Note that if there is no variation of the plasma density, $\delta n_\alpha=0$, $\delta\varepsilon_\alpha=0$, then we have the equation that was studied in Ref. [21].

We now rewrite Eq. (3) taking into account poles in the integrals:

$$\Omega-\mathbf{qu}=0, \quad (6)$$

$$\Omega-\mathbf{qv}_\alpha=0, \quad (7)$$

where

$$\mathbf{u}=\frac{\mathbf{k}c^2}{\omega(k)}.$$

Using the well-known relation

$$\lim_{\varepsilon\to 0}\frac{1}{x+\imath\varepsilon}=\mathrm{p}\frac{1}{x}-\imath\pi\delta(x) \quad (8)$$

and recalling [21]

$$A=\mathrm{Re}\,A+i\,\mathrm{Im}\,A=A_0+\imath A_1,$$

$$\delta\varepsilon_\alpha=\delta\varepsilon_\alpha'+\imath\delta\varepsilon_\alpha'', \quad \Omega=\Omega'+\imath\Omega'', \quad (9)$$

and assuming that $A_0\gg|A_1|$ and $|\delta\varepsilon_\alpha'|\gg|\delta\varepsilon_\alpha''|$, we rewrite Eq. (3) as

$$\varepsilon\left(1+\frac{1}{A_0}-\imath\frac{A_1}{A_0^2}\right)+(1+\delta\varepsilon_i)\delta\varepsilon_e\frac{q^2c^2}{\omega_{Le}^2}=0. \quad (10)$$

We first consider the excitation of the longitudinal photons and electron Langmuir waves ($\Omega'\sim\omega_{Le}$ and $\Omega\gg\omega_{Li}$ or $\delta\varepsilon_i=0$). In this case,

$$\varepsilon=1+\delta\varepsilon_e'+\imath\delta\varepsilon_e''=1-\frac{\omega_{Le}^2(1+3q^2r_{De}^2)}{\Omega^2}-\imath\frac{4\pi^2e^2m_{0e}}{q^2}\left(\frac{\partial f_{0e}}{\partial p_x}\right)_{v_x=\Omega/q}, \quad (11)$$

$$1+\frac{1}{A_0}=\frac{1}{q^2V_E^2}\{(\Omega-\mathbf{qu})^2-q^2(V_s^2-V_E^2)-\alpha^2q^4\}, \quad (12)$$

$$A_1=-\sqrt{\frac{\pi}{2}}\frac{\omega^3}{(\sigma_kc)^3}\frac{V_E^2}{c^2}\frac{(\Omega-\mathbf{qu})}{qc}\exp\left(-\frac{3}{2}\frac{(\Omega-\mathbf{qu})^2}{q^2V_s^2}\right), \quad (13)$$

where

$$V_E^2=\frac{c^2}{2}\left(\frac{\omega_{Le}}{\omega}\right)^2\frac{\gamma_0^2-1}{\gamma_0^2}, \quad V_s=c\sqrt{3}\frac{\sigma_kc}{\omega(k_0)}, \quad \alpha=\frac{c^2}{2\omega(k_0)}, \quad (14)$$

and r_{De} is the Debye length for electrons.

We now neglect the small imaginary term in Eq. (10) and examine the dispersion relation

$$\varepsilon_e'\left(1+\frac{1}{A_0}\right)+\frac{q^2c^2}{\omega_{Le}^2}\delta\varepsilon_e'=0, \quad (15)$$

or more explicitly

$$\{\Omega^2-\omega_{Le}^2(1+3q^2r_{De}^2)\}\{(\Omega-\mathbf{qu})^2-q^2U^2-\alpha^2q^4\}=q^4V_E^2c^2(1+3q^2r_{De}^2), \quad (16)$$

where

$$U^2=V_s^2-V_E^2, \quad \mathbf{u}(\mathbf{k}_0)=\frac{\mathbf{k}_0c^2}{\omega(k_0)}, \quad \mathbf{qu}=qu\cos\Theta.$$

The maximum growth rate is obtained when

$$V_s^2+\alpha^2q^2=V_E^2, \quad (17)$$

in which case from Eq. (16) for $\Omega=\omega_{Le}(1+\frac{3}{2}q^2r_{De}^2)+\delta\approx\mathbf{qu}+\delta$, we get

$$\delta^3 = \frac{1}{2\omega_{Le}} q^4 V_E^2 c^2 (1 + 3q^2 r_{De}^2)^{1/2} \tag{18}$$

or

$$\mathrm{Im}\ \delta = \frac{\sqrt{3}}{2} \left(\frac{qc}{2\omega_{Le}} \frac{V_E^2}{c^2} (1 + 3q^2 r_{De}^2)^{1/2} \right)^{1/3} qc. \tag{19}$$

If relation (17) does not hold, then Eq. (16) has an unstable solution such as

$$\Omega = \omega_{Le}\sqrt{1 + 3q^2 r_{De}^2} + \delta, \quad \Omega = \mathbf{qu} - q\sqrt{U^2 + \alpha^2 q^2} + \delta. \tag{20}$$

From Eq. (16) follows, for the imaginary part of δ,

$$\mathrm{Im}\ \delta = \frac{qV_E}{2\omega_{Le}} qc \left(\frac{\mathbf{qu}}{\omega_{Le}\sqrt{1 + 3q^2 r_{De}^2}} - 1 \right)^{-1/2}. \tag{21}$$

Note that $\mathbf{qu} > \omega_{Le}\sqrt{1+3q^2 r_{De}^2}$ always, as follows from relation (20). Equations (19) and (21) demonstrate the modulational excitation of longitudinal photons and plasmons simultaneously.

We now consider the kinetic instability for the case when $q^2 v_{tre}^2 < \Omega'^2 < \omega_{Le}^2$, which means that in this case the low-frequency longitudinal photons are generated alone. From Eq. (10) a simple calculation gives for the imaginary part of $\Omega(q)$ the expression

$$\Omega'' = -\sqrt{\frac{\pi}{8}} \left\{ (\Omega' - \mathbf{qu}) \frac{\omega_{Le}}{\Omega'} \left(\frac{\omega_{Le}}{qc} \right)^3 \left(\frac{V_E}{c} \right)^2 \left(\frac{\omega}{\sigma_k c} \right)^3 \times \exp\left(-\frac{3}{2} \frac{(\Omega' - \mathbf{qu})^2}{2q^2 V_s^2} \right) + \Omega' \left(\frac{\Omega'}{q v_{tre}} \right)^3 \exp\left(-\frac{\Omega'^2}{2q^2 v_{tre}^2} \right) \right\}. \tag{22}$$

From Eq. (22) it follows that Ω'' changes sign and becomes positive if the Cherenkov condition is satisfied:

$$u > \frac{\Omega'}{q\cos\Theta} \left\{ 1 + \frac{\Omega'}{\omega_{Le}} \left(\frac{qc}{\omega_{Le}} \right)^3 \left(\frac{c}{V_E} \right)^2 \left(\frac{\sigma_k c}{\omega} \right)^3 \times \exp\left(-\frac{\Omega'^2}{2q^2 v_{tre}^2} \right) \right\}, \tag{23}$$

so that this instability leads to the excitation of longitudinal photons. Here we have supposed that $|\Omega' - \mathbf{qu}| < qV_s$.

We next study the range of frequencies

$$q^2 v_{tri} < \Omega' < q^2 v_{tre} \tag{24}$$

and the case when $\delta n_i \neq 0$. In this case,

$$\varepsilon = \varepsilon' + \imath\delta\varepsilon'' = 1 + \frac{1}{q^2 r_{De}^2} + \imath\sqrt{\frac{\pi}{2}} \frac{\omega_{Le}^2 \Omega}{(q v_{tre})^3} - \frac{\omega_{pi}^2}{\Omega^2}. \tag{25}$$

First, we neglect the small imaginary terms in Eqs. (25) and (4), and use Eq. (10) to obtain the dispersion relation

$$(\Omega^2 - \Omega_s^2)\{(\Omega - \mathbf{qu})^2 - q^2(U^2 + \alpha^2 q^2)\} + \frac{q^4 c^2 V_E^2 (\Omega^2 - \omega_{pi}^2)}{\omega_{Le}^2 (1 + q^2 r_{De}^2)} = 0, \tag{26}$$

where $\Omega_s = \frac{q v_s}{\sqrt{1+q^2 r_{De}^2}}$ is the ion sound frequency.

If the length of waves is shorter than the Debye length—i.e., $\lambda \ll r_{De}$—then $\Omega_s \to \omega_{pi}$ and Eq. (26) reduces to

$$(\Omega - \mathbf{qu})^2 - q^2(V_s^2 + \alpha^2 q^2) + q^2 V_E^2 \left(1 + \frac{c^2}{v_{tre}^2} \right) = 0. \tag{27}$$

This relation demonstrates that ions play no role in the instabilities; whereas, in the presence of the hot electrons, the growth rate becomes large, because of the coupling term [the last term in Eq. (26)], as c^2/v_{tre}^2, and the imaginary part of the frequency is

$$\mathrm{Im}\ \Omega = \Omega'' \simeq q V_E \frac{c}{v_{tre}}.$$

In the opposite case—i.e., $q^2 r_{De}^2 \ll 1$—we have the modulational excitation of longitudinal photons and ion sound waves simultaneously. That is, the photon flux triggers both waves. To show this, we first consider the case when the condition (17) is satisfied. In this case for the coinciding roots $\Omega = \Omega_s + \delta = \mathbf{qu} + \delta$, where $\Omega_s \gg |\delta|$, we obtain from Eq. (26)

$$\mathrm{Im}\ \delta = \frac{\sqrt{3}}{2} qc \left(\frac{m_e}{2m_i} \frac{V_E^2}{c v_s} \right)^{1/3}. \tag{28}$$

Next, if relation (17) does not hold, then Eq. (26) has an unstable solution such as

$$\Omega = \Omega_s + \delta, \quad \Omega = \mathbf{qu} - q\sqrt{U^2 + \alpha^2 q^2} + \delta, \tag{29}$$

and we get the growth rate from Eq. (26) as

$$\mathrm{Im}\ \delta = \frac{qc}{2} \left(\frac{V_E}{v_{tre}} \right) \frac{1}{(1 + q^2 r_{De}^2)^{1/2}} \frac{1}{(\mathbf{qu}/\Omega_s - 1)^{1/2}}, \tag{30}$$

which is true for the Cherenkov condition $u > \frac{\Omega_s}{q\cos\Theta} = \frac{v_s}{\cos\Theta}$. Furthermore, Eq. (30) is valid when $V_E \ll v_{tre}$, since $\Omega_s = q v_s > \mathrm{Im}\ \delta$.

To summarize, we have shown a simultaneous nonlinear excitation of longitudinal photons and plasmons in the interaction of relativistically intense nonmonochromatic radiation bunches with a nonmagnetized plasma. Cases without and with ion dynamics are discussed. The generation of only low-frequency longitudinal photons is also confirmed. The growth rates of these new modes, which have no counterpart in the case of monochromatic EM waves, are obtained. These investigations may play an essential role in advanced fusion concepts and advanced accelerators, as well as for the description of extremely complex phenomena that appear only in energetic astrophysical systems and in experiments modeling high-energy density astrophysics in the laboratory. It was argued in Ref. [25] that m_γ would lead to a catastrophic emission of longitudinal photons. This gas may play a decisive role in the expansion of the Universe. In addition, in the fireball model of γ-ray bursts, the afterglow may be due to the decay process discussed in this and previous papers [19,21]. A cursory examination of burst profiles indicates that some are chaotic and spiky with large fluctuations on all time scales, while others show rather simple structures with few peaks. However, some bursts are seen with both characteristics present within the same burst [26].

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