Kinetics of Photon Radiation off an $e^-e^+$- Plasma created from the Vacuum in a Strong Laser Field

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We consider the one-photon annihilation mechanism in a electron - positron quasiparticle plasma (EPP) created from the vacuum in a strong subcritical laser field due to the dynamical Schwinger mechanism. On the basis of a kinetic theory approach we show that the secondary photons have a radiation spectrum proportional to $1/k$ (flicker noise). This effect is very small for EPP excitations in the optical spectrum but can reach quite observable values in the $\gamma$ - ray region.

1 Introduction

The effect of vacuum particle creation in a strong electric field was predicted long ago [1] but is not confirmed experimentally up to now. Theoretical research on this effect in the framework of kinetic theory [2] has shown the existence of a quasiparticle plasma of high density already in subcritical electric fields $E \ll E_c = m^2/e$ [3]. Such a quasiparticle plasma does not survive after switching off the external field in contrast to the real (observed) plasma that remains after applying the pulse of a strong field $E \geq E_c$ [4], see also [5] for the case $E \sim E_c$. Different proposals have been made for observing the quasiparticle EPP generated in the focal spot of counter propagating laser beams (see, e.g., Refs. [6, 7] and literature cited therein).

One of these effects is based on the possibility of one-photon and two-photon annihilation of the quasiparticle pairs. The first reassuring estimations were performed on the basis of the S-matrix formalism for modern high-power lasers with an intensity of $I \sim 10^{20}$ W/cm$^2$ [8]. However, there are some doubts in the validity of the S-matrix methods for the short-lived quasiparticles in the presence of a strong time dependent electric field. This stimulates the development of more adequate methods such as a kinetic description of the EPP, taking into account also its photon component. The first steps in the direction of including the photon sector of the EPP into the kinetic description were done in the works [9][10].

In present work we develop the kinetic theory of the EPP photon sector on the basis of the one-photon annihilation process which is no longer forbidden in the presence of strong fields [11]. This requires a derivation of the second level equations of the BBGKY chain. In the case of an oscillating electric field we show that the photon production rate is determined by joint action of ponderomotive forces and multi-photon processes. The spectrum of the annihilation photons has the $1/k$ behaviour of flicker noise. This effect is very small for optical excitations of the EPP but can reach quite observable values in the $\gamma$ - ray region. We calculate the intensity of such processes on the basis of the "photon count method" [6]. We have the methodical basis for more advanced calculations of the photon production rate for the two-photon annihilation mechanism requiring the fourth level equations of the BBGKY chain.

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Let us stress that the present estimates are actual especially in view of the planned experiment for testing the subcritical EPP production in vacuum with the Astra-Gemini high power laser [12].

2 One-photon annihilation process

An external electric field generates a vacuum instability due to EPP creation accompanied by the appearance of internal currents and electromagnetic fields (backreaction problem). Quantum fluctuations of this internal field are interpreted as photon excitations, which can leave the active zone (focal spot) and can be detected. Below the KE of the photon component will be examined on the basis of the one-photon annihilation mechanism, thus going beyond the investigation of the KE in the fermionic quasiparticle sector in previous works [13] [14] [15].

The single-time, two-point photon correlation function in momentum space is

\[ F_{rr'}(k, k', t) = \langle A_r^{(+)}(k, t) A_{r'}^{(-)}(k', t) \rangle, \]

with the photon vacuum \( \langle A_r^{(+)}(k, t) \rangle = 0 \). Let us write the first equation of the BBGKY hierarchy

\[ \dot{F}_{rr'}(k, k', t) = i \int \frac{d^3p_1 d^3p_2}{(2\pi)^6} \left\{ \delta(p_1 - p_2 - k) \langle \bar{u}v \rangle_a \langle \bar{u}p_1, p_2, k; t \rangle \langle b_{\alpha}(-p_2, t) a_{\alpha}(p_1, t) A_r^{(-)}(k', t) \rangle \right. \]

\[ + \left. \delta(p_1 - p_2 + k') \langle \bar{u}p_1, p_2, k; t \rangle \langle b_{\beta}(-p_2, t) a_{\beta}(p_1, t) A_r^{(+)}(k, t) \rangle \right\}, \] (2)

where \( a_{\alpha}, b_{\beta}, A_r^{(-)}(a_{\alpha}^{+}, b_{\beta}^{+}, A_r^{(+)}) \) denote the annihilation (creation) operators for electrons, positrons and photons, respectively [10]. For obtaining a closed photon KE a truncation procedure for the correlators entering this equation is required. The equations of the second order for the correlators in Eq. (2) can be obtained using the Heisenberg-like equations of motion, e.g.,

\[ \left\{ \frac{\partial}{\partial t} + i[\omega(p_1, t) + \omega(p_2, t) - k] \right\} \langle b_{\alpha}(-p_2, t) a_{\beta}(p_2, t) A_r^{(+)}(k, t) \rangle = -ie \int \frac{d^3p' d^3k'}{(2\pi)^6} \right\} \delta(p' - p_1 + k') \times \left[ \tilde{u}u_{\alpha\beta}(p', p_1, k'; t) \langle a_{\beta}^{+}(p', t) a_{\beta}(p_2, t) A_r(k', t) A_r^{(+)}(k, t) \rangle \right. \]

\[ + \left. \langle \tilde{u}u_{\alpha\beta}(p', p_1, k'; t) \rangle \langle b_{\beta}(-p_2, t) a_{\beta}(p_1, t) A_r(k', t) A_r^{(+)}(k, t) \rangle \right\} \]

\[ - \delta(p_2 - p' + k') \cdot [\tilde{u}\tilde{u}_{\alpha\beta}(p_2, p', k'; t) \langle b_{\alpha}(-p_1, t) a_{\beta}(p', t) A_r(k', t) A_r^{(+)}(k, t) \rangle] \]

\[ + \left. \langle \tilde{u}\tilde{u}_{\alpha\beta}(p_2, p', k'; t) \rangle \langle b_{\beta}(-p_1, t) a_{\beta}(p_1, t) A_r(k', t) A_r^{(+)}(k, t) \rangle \right\} \]

\[ + S_{\alpha\beta}(p_1, p_2, k; t) + U_{\alpha\beta}(p_1, p_2, k; t). \] (3)

On the r.h.s. of this equation there is a set of terms defined by the vacuum polarization effects in the presence of the quantized electromagnetic field (summed up in \( U^{(r)} \)). The other set of terms (summed up in \( S^{(r)} \)) is connected with different transformations in the \( e^- e^+ \gamma \) - plasma without participation of the quasiparticle photon excitations. These processes have an indirect influence on the photon distribution due to the secondary momentum redistribution of quasiparticles by means of collisions. These \( U^{(r)} \) and \( S^{(r)} \) terms will be omitted below (the approximation A1).

The remaining processes are written out explicitly in Eq. (3). The first and the fourth terms describe the mutual influence of the electron and positron quasiparticle excitations and photons. Only these last terms will be kept subsequently in the first step (the approximation A2). Some basis for such selection is the following reasoning: in RPA the truncation of correlators of the type

\[ \langle a_{\beta}^{+}(p', t) a_{\beta}(p_2, t) A_r(k', t) A_r^{(+)}(k, t) \rangle \sim \langle a_{\beta}^{+}(p', t) a_{\beta}(p_2, t) \rangle \langle A_r(k', t) A_r^{(+)}(k, t) \rangle \] (4)
leads to a nonvanishing contribution already in the lowest order of perturbation theory. In this approximation Eq. (3) gets the form

\[
\frac{\partial}{\partial t} + i[\omega(p_1, t) + \omega(p_2, t) - k] \langle b_\alpha(-p_1, t)a_\beta(p_2, t)A^{(+)\dagger}(k,t) \rangle \approx -ie\int \frac{d^3p'}{(2\pi)^3} \frac{d^3k'}{2k} \left\{ \delta(p' - p_1 + k')[\bar{u}u]_{\alpha\beta}'(p', p_1, k'; t) \right. \\
\times \langle a_{\beta'}^+(p', t)a_\beta(p_2, t)A^{(-)}(k', t)A^{(+)\dagger}(k,t) \rangle - \delta(p_2 - p' + k') \right. \\
\left. \times [\bar{u}u]_{\beta'\beta}(p_2, p', k', t) \rangle b_\alpha(-p_1, t)b_{\beta'}^+(-p', t)A^{(-)}(k', t)A^{(+)\dagger}(k,t) \rangle \right\} .
\] (5)

The processes of the instantaneous radiation of two photons was omitted here, i.e., in Eq. (3) the substitution

\[ A_{r}(k', t) \rightarrow A^{(-)}_{r}(k', t) \]

has been made (the approximation A3). Their inclusion requires the consideration of the next (the third) level equation of the BBGKY chain under the correlators of the type

\[ \langle a_{\beta'}^+(p', t)a_\beta(p_2, t)A^{(+)}_{r}(k', t)A^{(+)}_{r}(k,t) \rangle . \]

Let us remark that the two-photon annihilation process with the correlator of the type (the second and the third terms on the r.h.s. of Eq. (3) \langle b_{\beta'}(-p', t)a_\beta(p_2, t)A^{(+)}_{r}(k', t)A^{(+)\dagger}(k,t) \rangle \) is also found among the processes omitted in this section. Thus, the selected simplest approximation here corresponds to taking into account the one-photon annihilation process only. Let us rewrite Eq. (5) in integral form

\[
\langle b_\alpha(-p_1, t)a_\beta(p_2, t)A^{(+)\dagger}(k,t) \rangle = -\frac{ie}{\sqrt{2k}} \int_{t_0}^{t} \frac{dt'}{(2\pi)^3} \exp \left\{ -i \int_{t'_0}^{t} d\tau [\omega(p_1, \tau) + \omega(p_2, \tau) - k] \right. \\
\times \left. \left\{ [\bar{u}u]_{\alpha\beta}'(p_2, p_1, k; t')f_{\beta\beta'}(p_2, t') - [\bar{u}u]_{\beta'\beta}(p_2, p_1, k; t')f_{\alpha\alpha'}(p_1, t') \right\} \right. \\
\left. \times (1 + F_r(k, t')) \delta(p_2 - p_1 + k) \right\} .
\] (6)

In order to obtain Eq. (6), the following truncation procedure (the approximation A4) has been applied

\[ \langle a_{\beta'}^+(p', t)a_\beta(p_2, t')A^{(-)}_{r}(k', t')A^{(+)\dagger}(k,t') \rangle = \delta_{r'r}\delta(p' - p_2)\delta(k - k')f_{\beta\beta'}(p_2, t')\{1 + F_r(k, t')\} , \]

and use was made of the relation \[ \langle a_{\beta'}^+(p', t)a_\beta(p_2, t) \rangle = \delta(p - p')f_{\alpha\beta}(p, t) . \] The relation (6) combined with Eq. (2) and the system of equations of the fermion sector leads to a formally closed KE system. However, keeping in mind the estimation of the radiated photon spectrum, we will introduce a set of the additional simplifying assumptions. An essential simplification is reached if we neglect the spin effect in Eq. (6), i.e., \[ f_{\alpha\beta} \rightarrow f_{\delta\alpha\delta} \] (the approximation A5). This leads to the following KE of non-Markovian type

\[
\hat{F}(k, t) = \frac{e^2}{2(2\pi)^3} \int d^3p \int_{t_0}^{t} d\tau \cos \left\{ \int_{t'_0}^{t} d\tau' [\omega(p, \tau) + \omega(p + k, \tau) - k] \right. \\
\times K(p, p + k, \tau,t') \{f(p + k, t') + f(p, t') - 1\}\{1 + F_r(k, t')\} ,
\] (7)

where the kernel \( K \) is defined in the nonstationary spinor basis \[ 13 \] \[ \] \[ 14 \] \[ \] \[ \] \[ \] \[ \] as the spinor construction

\[ K(p_1, p_2, k|t, t') = [\bar{u}u]_{\alpha\beta}'(p_1, p_2, k; t)\bar{u}u_{\alpha\beta}'(p_2, p_1, k; t') . \] (8)

We assume that the photon distribution has not a fixed direction of polarization, i.e., \( F_1 = F_2 = F \) (the approximation A6). It has also been taken into account that \( f^c = 1 - f \) due to electroneutrality of the vacuum.

### 3 The case of a weak external field

If \( E(t) \ll E_c \), it can be expected that the density of the observable photons is small, \( F(k, t) \ll 1 \), and thus their distribution function can be neglected on the r.h.s. of Eq. (7), the photon production rate. Thus, the photon generation is defined, in the first place, by the statistical factor \( 1 - f(p, t') - f(p + k, t') \), which describes the quasiparticle EPP as the material medium generating real photons and the fast oscillating factor with the phase

\[
\Phi(t', t') = \int_{t'_0}^{t} d\tau [\omega(p, \tau) + \omega(p + k, \tau) - k] ,
\] (9)
The basic idea is the compensation of the mismatch (11) with help of the multiphoton process in order to eliminate the external field. For subcritical fields absorption coefficient of the EPP created from the vacuum in the infrared region \([6]\).

The other photon source is the fastly oscillating factor with the phase (9) which depends on the parameters of the external field. That is the reason why the radiation of a real photon can be interpreted as the multiphoton process. The photon number \(n\nu\) of an external field. That is the reason why the radiation of a real photon can be interpreted as the multiphoton process. The photon number \(n\nu\) is very large and the energy conservation law is violated for the one-photon annihilation process in the absence of an external field. That is the reason why the radiation of a real photon can be interpreted as the multiphoton process.

In order to generalize this result to the case of arbitrary energies for the real photons, we perform the Fourier decomposition of the functions on the r.h.s of Eq. (10) (method of photon count \([6]\)). This is a rather complicated problem because the time scales are defined by two parameters with dimension of energy (mass \(\nu\) and some weak oscillations \(\nu\)). Apparently, this leads to a divergence of the momentum space integral on the r.h.s of Eq. (11) from the “1” in the statistical factor. Subtracting the corresponding contribution we obtain the renormalized production rate of real photons

\[
\hat{F}(k, t) = \frac{e^2 K_0}{2(2\pi)^3k} \int d^3p \int_0^t dt' \{ f(p, t') + f(p + k, t') \} \cos \Phi(t, t').
\]  

(10)

The fastly oscillating function in Eq. (10) is the characteristic behaviour of flicker noise (e.g., \([17]\)), which arises in the present case from the fluctuations of vacuum oscillations \([16]\). Technically, this result is a direct consequence of the decomposition for the electromagnetic field operators in the region of small \(k\). The phase \(\Phi\) is always very large since the mismatch

\[
\omega(p, \tau) + \omega(p + k, \tau) - k \sim 2m
\]

(11)
is very large and the energy conservation law is violated for the one-photon annihilation process in the absence of an external field. That is the reason why the radiation of a real photon can be interpreted as the multiphoton process. The photon number \(N_\nu\) from the photon condensate of the external quasiclassical photons with the frequency \(\nu\) can be estimated if it has to compensate the mismatch (11) in the energy of \(N_\nu\) quasiclassical photons. This leads to \(N_\nu \sim 2m/\nu\). For optical lasers it is a huge number and therefore such a fluctuation event is very rare. This conclusion about the role of multiphoton processes is correlated with the analysis of the absorption coefficient of the EPP created from the vacuum in the infrared region \([6]\).

In order to generalize this result to the case of arbitrary energies for the real photons, we perform the Fourier decomposition of the functions on the r.h.s. of Eq. (10) (method of photon count \([6]\)). This is a rather complicated problem because the time scales are defined by two parameters with dimension of energy (mass \(m\) and the frequency \(\nu\) of the external periodical field) under the arbitrary third parameter \(k\), the energy of the real photon. The basic idea is the compensation of the mismatch (11) with help of the multiphoton process in order to eliminate the beating of the phase (9) (the resonance approximation). In Eq. (10) the source of the photon excitations with the energy \(k\) is the fermion distribution function which is an implicit function defined by the KE in the fermion sector.

The other photon source is the fastly oscillating factor with the phase (9) which depends on the parameters of the external field. For subcritical fields \(E \ll E_c\) we can write approximately \((\omega_0(p) = \omega(p, t)|_{\lambda=0}) \omega(p, t) \simeq \omega_0(p) - \frac{\epsilon E_0}{\nu \omega_0(p)} \cos \nu t, \frac{\epsilon E_0}{\nu \omega_0(p)} \ll 1\), if \(A(t) = (E_0/\nu) \cos \nu t\). Then the phase (9) is given by

\[
\Phi(t, t') \simeq \Omega_0(p, k)(t - t') + a(p, k)[\sin \nu t - \sin \nu t'],
\]

(12)

where the mismatch is

\[
\Omega_0(p, k) = \omega_0(p) + \omega_0(p + k) - k
\]

(13)

and

\[
a(p, k) = -\frac{\epsilon}{\nu^2} \{ |E_0(p)|/\omega_0(p) + |E_0(p + k)|/\omega_0(p + k) \}.
\]

The fastly oscillating function in Eq. (10) allows then the following representation

\[
\cos \Phi(t, t') \simeq e^{i\Omega_0(t-t')} \sum_{n, n'} J_n(a) J_n'(a) e^{in\nu(nt - nt')} + c.c.,
\]

(14)

where \(J_n(a)\) is the Bessel function of order \(n\). Thus, the representation (14) contains two high frequency harmonics \(\omega_0(p)\) and \(\omega_0(p + k)\) (see Eq. (13)), the set of the low frequency harmonics \(n\nu\) and the \(k\)-harmonic corresponding to the radiating photon. This structure of the Eq. (14) and the discussion of the mismatch (11) suggest the following twofold decomposition of the fermion distribution function \((f_{-n-l} = f_{n'}^{k'})\):

\[
f(p, t) = \sum_{n,l} f_{n,l}(p)e^{in\nu t + imlt},
\]

(15)
which has two modes: the soft ("breathing") one with the frequencies $n\nu$ and the hard ("trembling") one with the harmonics $lm$. Substituting Eqs. (14) and (15) into Eq. (10) leads to the relation

$$\hat{F}(k, t) = \frac{i e^2 K_0}{4(2\pi)^3 k} \sum_{n, n'} \int d^3 p J_n(a) J_{n'}(a) \{f_{n+1}(p) + f_{n+1}(p+k)\} \exp \left\{i \left[ (m + \nu(n + n' - n')) t + \frac{2\pi a_k}{\Omega_0 - plm - \nu(n - n') + i\varepsilon} \right] \right\}. \quad (16)$$

The shift $i\varepsilon$ in the complex plane is a consequence of the adiabatic hypothesis guaranteeing the convergence of the integral for $t_0 \to -\infty$.

Let us use now the resonance approximation which allows to select from the sum on r.h.s. of Eq. (16) the constant component $lm + (n + n' - n') \nu = 0$. At $l = 0$ the decomposition (16) contains the basic low frequency breathing modes only. Below we will consider this case only (the case $l = 1$ must be investigated separately also). The photon production rate (16) has then the form

$$\hat{F}(k) = \frac{i e^2 K_0}{4(2\pi)^3 k} \sum_{n, n'} \int d^3 p J_n(a) J_{n'}(a) \frac{f_{n+1}(p) + f_{n+1}(p+k)}{\Omega_0 + n\nu + i\varepsilon} \frac{\exp \left\{i \left[ (n \nu(n + n' - n')) + \frac{2\pi a_k}{\Omega_0} \right] t \right\}}{\Omega_0 + n\nu + i\varepsilon}.$$ 

Using the textbook formula $\delta[\phi(x)] = \sum_i \{\phi'(x_i)\}^{-1}\delta(x - x_i), \phi(x_i) = 0$, we obtain

$$\hat{F}(k) = \frac{\alpha K_0}{2k} \sum_{n, n'} \xi(k, n) J_n(a_0) J_{n'}(a_0) \{f_{n+1}(p_0) + f_{n+1}(p_0 + k)\} \quad (18)$$

where $a_0 = a(p_0), \alpha = e^2/4\pi$ and according to Eq. (13)

$$\xi(k, n) = \frac{p_0}{\Omega_0(p_0)} = \frac{\omega_0(p_0) \sqrt{\omega_0^2(p_0) + k^2}}{\Omega_0(p_0) + \omega_0^2(p_0) + k^2} \quad (19)$$

and $p_0$ is the positive root of the equation $\Omega_0(p_0) - n\nu = 0$,

$$p_0 = \left\{ \frac{(n\nu)^2(n\nu - 2k)^2}{4(n\nu - k)^2} - m^2 \right\}^{1/2} \geq 0 \quad (20)$$

and

$$n = \frac{1}{\nu} \{\omega_0(p) + \omega_0(p + k) - k\}.$$ 

One can see that $n \in [\omega_0(p), 2\omega_0(p)], \text{i.e.} \ n$ is very large in the optical region and can be $\sim 1$ for $\nu \sim m$.

Now we take into account the well known property of the fermion distribution function: the basic breathing mode corresponds to the second harmonic ($n = \pm 2$) in Eq. (15). Keeping only these modes ($n' - n = \pm 2$) in Eq. (18), we obtain

$$\hat{F}(k) = \frac{\alpha K_0}{2k} \sum_{n \geq n_0} \xi(k, n) J_n(a_0) [J_{n+2}(a_0) + J_{n-2}(a_0)] \{f_{2,0}(p_0) + f_{2,0}(p_0 + k)\} \quad (22)$$

where according to Eq. (21)

$$n_0 = n(p = 0) = \frac{1}{\nu} \{m + \omega(k) - k\}.$$ 

Leaving in Eq. (22) only the leading term ($n = n_0 + 1$), we obtain

$$\hat{F}(k) = \frac{\alpha K_0}{2k} \xi(k, n_0 + 1) J_{n_0+1}(a_0) [J_{n_0+3}(a_0) + J_{n_0-1}(a_0)] \{f_{2,0}(p_0) + f_{2,0}(p_0 + k)\} \quad (24)$$

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where
\[ \xi(k, n_0 + 1) = p_0(n_0 + 1) \frac{m_\omega_0(k)}{m + \omega_0(k)} \]  
(25)

and \( p_0(n_0 + 1) \) can be found from Eqs. (20) and (23). Using the decompositions of the Bessel functions \( J_0(a_0) = 1, \ J_{n_0}(a_0) = (a_0/2)^{n_0}/n_0 \sim \alpha_0^{n_0}/2 \), valid for \( a_0 \ll 1 \), we obtain the final estimates. The photon production rate \( \dot{F} \) in the case of the one-photon annihilation mechanism is very small \( \dot{F} \sim \alpha_0^{n_0}/2 \) for the optical laser radiation with \( n_0 \gg 1 \), see Eq. (21). This means that the excitation of a single observable photon requires a huge number of optical "laser" photons so that the probability of such process is negligible. However, in the \( \gamma \)-ray region \( \nu \sim m \) and then \( n_0 \sim 1 \). In this case the photon production rate can be quite observable. For example, for \( n_0 = 1 \) (it corresponds to \( \nu = 2m \)) and for soft photon radiation \( k \ll m \) we obtain \( \dot{F} \sim \alpha^2 \).

4 Summary

We have shown that the one-photon annihilation mechanism acting in the quasiparticle EPP can lead to the excitation of observable photons by multi-photon processes. In the infrared region \( k \ll m \) the radiation spectrum behaves as \( 1/k \) (the flicker noise). The increase \( k \) results in a more complicated spectrum, Eq. (24). In the case of subcritical laser fields \( E \ll E_c \) and optical excitations of the EPP this effect is very small. However, in the \( \gamma \)-ray region the photon production rate can increase up to quite observable values. In this connection one can expect that the absence of the multi-photon excitations mechanism in the two-photon annihilation channel can lead to the generation of an experimentally significant photon production intensity in the optical laser domain.

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