EFFECT OF ELECTRIC FIELD OF THE ELECTROSPHERE ON PHOTON EMISSION FROM QUARK STARS

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Abstract

We investigate the photon emission from the electrosphere of a quark star. It is shown that at temperatures $T \sim 0.1 \div 1$ MeV the dominating mechanism is the bremsstrahlung due to bending of electron trajectories in the mean Coulomb field of the electrosphere. The radiated energy flux from this mechanism exceeds considerably both the contribution from the bremsstrahlung due to electron-electron interaction and the tunnel $e^+e^-$ pair creation.
It is possible that quark stars made of a stable strange quark matter (SQM) [1, 2, 3] (if it exists) may exist without a crust of normal matter [4]. The quark density for bare quark stars should drop abruptly at the scale \( \sim 1 \) fm. The SQM in normal phase and in the two-flavor superconducting (2SC) phase should also contain electrons (for normal phase the electron chemical potential, \( \mu \), is about 20 MeV [2, 5]). Contrary to the quark density the electron density drops smoothly above the star surface at the scale \( \sim 10^3 \) fm [2, 5]. For the star surface temperature \( T \ll \mu \), say \( T \ll 10^{10} \) K \( \sim 1 \) MeV, this “electron atmosphere” (usually called the electrosphere) may be viewed as a strongly degenerate relativistic electron gas [2, 5]. The photon emission from the normal SQM is negligibly small as compared to the black body one at \( T \ll \omega_p \) [6, 7] (here \( \omega_p \sim 20 \) MeV is the plasma frequency of the SQM [6]). However, for the electrosphere the plasma frequency is much smaller than that for the SQM. For this reason the photon emission from the electrosphere may potentially dominate the luminosity of a quark star. Contrary to neutron stars (or quark stars with a crust of normal matter) the photon emission from the electrosphere of bare quark stars may exceeds the Eddington limit, and may be used for distinguishing a bare quark star from a neutron star (or a quark star with a crust of normal matter). For this reason it is of great importance to have quantitative predictions for the photon emission from the electrosphere. This is also of interest in the context of the scenario of the gamma-ray repeaters due to reheating of a quark star by impact of a massive comet-like object [8].

The bremsstrahlung from the electrosphere due to the electron-electron interaction has been addressed in [9, 10]. The authors of [9] used the soft photon approximation and factorized the \( e + e \rightarrow e + e \) cross section in the spirit of Low’s theorem. In [10] it was pointed out that this approximation is inadequate since it neglects the effect of the photon energy on the electron Pauli-blocking which should lead to a strong overestimate of the radiation rate. The authors of [10] have not given a consistent treatment of this problem either. To take into account the effect of the minimal photon energy they suggested some restrictions on the initial electron momenta introduced by hand. In this way they obtained the radiated energy flux from the \( e^- e^- \rightarrow e^- e^- \gamma \) process which is much smaller than that in [9], and than the energy flux from annihilation of positrons produced in the tunnel \( e^+ e^- \) creation in the electric field of the electrosphere [4, 11]. In [12] there was an attempt to include the effect of the mean Coulomb field of the electrosphere on the photon emission. The authors obtained a considerable enhancement of the radiation rate. However, similarly to [9] the analysis [12] treats incorrectly the Pauli-blocking effect.

Thus the theoretical situation with the photon bremsstrahlung from the electrosphere is still controversial and uncertain. The main problem here is an accurate accounting for the photon energy in the Pauli-blocking. In the present paper we address the bremsstrahlung from the electrosphere in a way similar to the Arnold-Moore-Yaffe (AMY) [13] approach to the collinear photon emission from a hot quark-gluon plasma within the thermal field theory. We use a reformulation of the AMY formalism given in [14] which is based on the light-cone path integral (LCPI) approach [15, 16, 17] (for reviews, see [18, 19]) to the radiation processes. For an infinite homogeneous plasma (with zero mean field) the formalism [14] reproduces the AMY results [13]. The LCPI formulation [14] has the advantage that it also works for plasmas with nonzero mean field. It allows to evaluate the photon emission accounting for bending of the electron trajectories in the
mean Coulomb potential of the electrosphere. Contrary to very crude and qualitative methods of [9, 10, 12] the treatment of the Pauli-blocking effects in [13, 14] has robust quantum field theoretical grounds. Of course, our approach is only valid in the regime of collinear photon emission when the dominating photon energies exceed several units of the photon quasiparticle mass. Numerical calculations show that even at $T \sim 0.1$ MeV the effect of the noncollinear configurations is relatively small.

We demonstrate that for the temperatures $T \sim 0.1 \div 1$ MeV the radiated energy flux from the $e^- \rightarrow e^- \gamma$ transition in the mean electric field turns out to be much bigger than contributions from the $e^- e^- \rightarrow e^- e^- \gamma$ process and the tunnel $e^+ e^-$ creation. Our results show that the photon emission from the electrosphere may be of the same order as the black body radiation. For this reason the situation with distinguishing a bare quark star made of the SQM in normal (or 2SC) phase from a neutron star using the luminosity [4, 20] may be more optimistic than in the scenario with the tunnel $e^+ e^-$ creation [4].

2. As in [4, 9, 10] we use for the electrosphere the model of a relativistic strongly degenerate electron gas in the Thomas-Fermi approximation. Then the electron chemical potential (related to the electrostatic potential, $V$, as $\mu = eV$) may be written as [2, 5]

$$\mu(h) = \frac{\mu(0)}{1 + h/H}, \quad \text{(1)}$$

where $h$ is the distance from the quark surface, and $H = \sqrt{3\pi/2\alpha/\mu(0)}$, $\alpha = e^2/4\pi$ (we use units $c = \hbar = k_B = 1$).

We assume that the electrosphere is optically thin. Then the luminosity may be expressed in terms of the energy radiated spontaneously per unit time and volume, $Q_\gamma$, usually called the emissitivity. In the formalism [14] the emissitivity per unit photon energy $\omega$ at a given $h$ can be written as

$$\frac{dQ_\gamma(h, \omega)}{d\omega} = \frac{\omega(k) dk}{4\pi^3} \int d\mathbf{p} \frac{n_F(E)}{p} \frac{dP(p, x)}{dx dL}, \quad \text{(2)}$$

where $k$ is the photon momentum, $E$ and $E'$ are the electron energies before and after photon emission, $n_F(E) = (\exp((E - \mu)/T) + 1)^{-1}$ is the local electron Fermi distribution (we omit the argument $h$ in the functions on the right-hand side of (2)), $x = k/p$ is the photon longitudinal (along the initial electron momentum $\mathbf{p}$) fractional momentum. The function $dP/dxdL$ in (2) is the probability of the photon emission per unit $x$ and length from an electron in the potential generated by other electrons which includes both the smooth collective Coulomb field and the usual fluctuating part. Note that (2) assumes that the photon emission is a local process, i.e. the photon formation length $l_f$ is small compared to the thickness of the electrosphere.

In the LCPI formalism [15, 18] the photon spectrum $dP/dxdL$ can be written as

$$\frac{dP}{dx dL} = 2Re \int_0^{\infty} d\xi \hat{g}(x) \left[ K(\rho', x|\rho, 0) - K_v(\rho', x|\rho, 0) \right]_{\rho' = \rho = 0}. \quad \text{(3)}$$

Here $\hat{g}$ is the spin vertex operator (it can be found in [18]), $K$ is the Green’s function for
the two-dimensional Hamiltonian
\[
\hat{H} = -\frac{1}{2M(x)} \left( \frac{\partial}{\partial \rho} \right)^2 + v(\rho) + \frac{1}{L_0},
\]
where \( M(x) = px(1 - x) \), \( L_0 = 2M(x)/\epsilon^2 \), \( \epsilon^2 = m_\gamma^2 x^2 + (1 - x)m_\gamma^2 \) (\( m_\gamma \) is the photon quasiparticle mass), the form of the potential \( v \) will be given below. In (3), (4) \( \rho \) is the coordinate transverse to the electron momentum \( p \), the longitudinal (along \( p \)) coordinate \( \xi \) plays the role of time. The \( \mathcal{K}_v \) in (3) is the free Green’s function for \( v = 0 \). Note that at low density and vanishing mean field the quantity \( L_0 \) coincides with the real photon formation length \( l_f \) [15].

The potential in the Hamiltonian (4) can be written as \( v = v_m + v_f \). The terms \( v_m \) and \( v_f \) correspond to the mean and fluctuating components of the vector potential of the electron gas. Note that when \( l_f \) is small compared to the scale of variation of \( \mu \) (along the electron momentum) one can neglect the \( \xi \)-dependence of the potential \( v \) in evaluating \( dP/dxdL \). The mean field component is purely real \( v_m = -x\mathbf{f} \cdot \mathbf{\rho} \) with \( \mathbf{f} = e\partial V/\partial \mathbf{\rho} \) [18, 21]. It is related to the transverse force from the mean field. Note that, similarly to the classical radiation [22], the effect of the longitudinal force along the electron momentum is suppressed by a factor \( \sim (m_\epsilon/E)^2 \), and can be safely neglected. The term \( v_f \) can be evaluated similarly to the case of the quark-gluon plasma discussed in [14]. This part is purely imaginary \( v_f(\rho) = -iP(x\rho) \), where
\[
P(\rho) = e^2 \int_{-\infty}^{\infty} d\xi [G(\xi, 0_\perp, \xi) - G(\xi, \rho, \xi)],
\]
\( G(x - y) = u_\mu u_\nu D^{\mu\nu} \), \( D^{\mu\nu} = \langle A^\mu(x)A^\nu(y) \rangle \) is the correlation function of the electromagnetic potential (the mean field is assumed to be subtracted) in the electron plasma, \( u_\mu = (1, 0, 0, -1) \) is the light-cone 4-vector (along the electron momentum). The correlator \( D^{\mu\nu} \) may be expressed in terms of the longitudinal and transverse photon self-energies, \( \Pi_{L,T} \) [13]. In numerical calculations we use for the \( \Pi_{L,T} \) the well known hard dense loop expressions [23, 24].

Treating \( v_f \) as a perturbation one can write
\[
\mathcal{K}(\xi_2, \rho_2|\xi_1, \rho_1) = \mathcal{K}_m(\xi_2, \rho_2|\xi_1, \rho_1) - i \int d\xi d\rho \mathcal{K}_m(\xi_2, \rho_2|\xi, \rho)v_f(\rho)\mathcal{K}_m(\xi, \rho|\xi_1, \rho_1) + \ldots,
\]
where \( \mathcal{K}_m \) is the Green’s function for \( v_f = 0 \). Then (3) can be written as
\[
\frac{dP}{dxdL} = \frac{dP_m}{dxdL} + \frac{dP_f}{dxdL}.
\]
Here the first term on the right-hand side comes from the \( \mathcal{K}_m - \mathcal{K}_v \) in (3) after representing \( \mathcal{K} \) in the form (6). It corresponds to the photon emission in a smooth mean field. The second term comes from the series in \( v_f \) in (6). This term can be viewed as the radiation rate due to electron multiple scattering in the fluctuating field in the presence of a smooth external field. The analytical expression for the Green’s function \( \mathcal{K}_m \) is known (see, for example [25]). The corresponding spectrum is similar to the well known synchrotron
spectrum, and can be written in terms of the Airy function \( \text{Ai}(z) = \frac{1}{\pi} \sqrt{\frac{2}{3}} K_{1/3}(2z^{3/2}/3) \) (here \( K_{1/3} \) is the Bessel function) [21, 26]. In the case of interest, for a nonzero photon quasiparticle mass it reads [21]

\[
\frac{dP_m}{dx dL} = \frac{a}{\kappa} \text{Ai}'(\kappa) + b \int_{\kappa}^{\infty} dy \text{Ai}(y),
\]

where

\[
a = -2e^2 g_1/M, \quad b = M g_2 - e^2 q_1/M, \quad \kappa = e^2/(M^2 x^2 f^2)^{1/3}, \quad g_1 = \alpha(1 - x + x^2/2)/x \quad \text{and} \quad g_2 = \alpha m_e^2 x^3/2M^2.\]

Note that the effective photon formation length for the mean field mechanism is given by \( L_m \sim \min(L_0, L_m) \), where \( L_m = (24M/x^2 f^2)^{1/3} \) [21].

Evaluation of the \( dP_f/dxdL \) for realistic function \( P(\rho) \) and nonzero mean field is a complicated computational problem. In the present work we have performed a qualitative calculation of this term. We evaluated \( dP_f/dxdL \) for zero mean field within the LCPI formalism [15] using the method of [16, 17]. This calculations show that for zero mean field the spectrum is dominated by the leading order term in \( v_f \) on the right-hand side of (6), and the effect of the higher order terms that describe the Landau-Pomeranchuk-Migdal (LPM) suppression is negligible \(^1\). The mean field should suppress the radiation rate. Qualitatively the corresponding suppression factor can be written as the ratio of the formation lengths with and without the mean field, \( i.e. \ S_m \approx L_m/L_0 \). Note that due to reduction of the effective formation length the LPM effect should become even smaller for a nonzero mean field. As will be seen from our numerical results the fluctuation term in (7) is much smaller than the mean field one. For this reason getting of an accurate prediction for \( dP_f/dxdL \) is not important in a pragmational sense. Note that, since the mean field mechanism dominates, the \( l_f \) is simply given by \( L_m \).

\(^3\) In numerical calculations we define the \( k \)-dependent photon quasiparticle mass from the relation \( m_\gamma^2 = \Pi_T(\sqrt{k^2 + m_\gamma^2}, k) \). This gives \( m_\gamma \) rising from \( m_D/\sqrt{2} \) at \( k \ll m_D \) to \( m_D/\sqrt{3} \) at \( k \gg m_D \) with the Debye mass \( m_D^2 = \frac{4\pi}{\alpha}(\mu^2 + \pi^2 T^2/3) \). We ignore the influence of the medium effects on \( m_e \) [27] since the results are not very sensitive to the electron quasiparticle mass.

As we mentioned earlier, the collinear approximation we use becomes invalid for very soft photons with \( k \lesssim m_\gamma \). In this region the formalisms [13, 14, 15] do not apply. In particular, the LCPI approach [15], which assumes that the transverse momentum integration comes up to infinity, should overestimate the photon spectrum at \( k \lesssim m_\gamma \). To take into account (at least, qualitatively) this effect we multiplied \( dP/dxdL \) by the kinematical suppression factor \( S_{\text{kin}}(k) = 1 - \exp(-k^2/m_\gamma^2) \). This factor suppresses the luminosity by \( \sim 10 - 15\% \) at \( T \sim 0.1 \div 0.2 \text{ MeV} \) and \( \sim 1 - 2\% \) at \( T \sim 1 \text{ MeV} \). This says that the errors from the noncollinear configurations are small.

We evaluated the differential, \( dF/d\omega \), and the total energy flux, \( F \). In our approach (the approximation of optically thin electrosphere) the \( dF/d\omega \) reads

\[
\frac{dF}{d\omega} = \int_{0}^{h_{\max}} dh \frac{dQ_\gamma(h, \omega)}{d\omega} \approx \sqrt{\frac{3\pi}{2\alpha}} \int_{\mu_{\text{min}}}^{\mu(0)} d\mu \frac{dQ_\gamma(h(\mu), \omega)}{d\omega} \quad (9)
\]

\(^1\)One can show that a very strong LPM suppression obtained in [9] is due to use of Migdal’s formulas for ordinary materials which become inadequate for the electrosphere.
with \( \mu_{\text{min}} = \mu(h_{\text{max}}) \). We take \( \mu_{\text{min}} = 2m_e \). Of course, the relativistic approximation we made is not good at \( \mu \sim m_e \). However, the contribution of this region is small, and the errors should not be big. We have performed computations for \( \mu(0) = 10 \) and \( \mu(0) = 20 \) MeV. In Fig. 1 we plot the radiation rate \( dF/d\omega \) for the mean field and the fluctuation mechanisms for \( T = 0.2 \) and \( T = 1 \) MeV. For comparison the black body spectrum is also shown. For the fluctuation contribution we show the results with and without the mean field suppression factor \( S_m \). One can see that the Coulomb potential of the electrosphere reduces the fluctuation term by a factor \( \sim 3 - 4 \). From Fig. 1 one can see that the relative contribution of the fluctuation mechanism is very small. Thus, in some sense we have a situation similar to that for an atom with large \( Z \). Note that the form of the spectrum for the mean field mechanism is qualitatively similar to that for the black body radiation.

In Fig. 2 we show the total energy flux \( F = \int_0^\infty d\omega dF/d\omega \) scaled to the black body limit as a function of temperature. For comparison, in Fig. 2 we also plot the energy flux from the \( e^+e^- \) pair production \([4, 11]\). We define it as

\[
F_{e^+e^-} = \int_0^{h_{\text{max}}} dh Q_{e^+e^-}(h) \approx \sqrt{\frac{3\pi}{2\alpha}} \int_{\mu_{\text{min}}}^{\mu(0)} \frac{d\mu}{\mu^2} Q_{e^+e^-}(h(\mu)). \tag{10}
\]

Here \( Q_{e^+e^-} \) is the energy flux from \( e^+e^- \) pairs per unit time and volume. We write it in the form given in \([11]\) \( Q_{e^+e^-} = 2E_{e^+e^-}dN_{e^+e^-}/dtdV \), where \( E_{e^+e^-} \approx 2(m_e + T) \) is the typical energy of \( e^+e^- \) pairs, and \( dN_{e^+e^-}/dtdV \) is the rate of \( e^+e^- \) pair production per unit time and volume defined by the formulas given in \([11]\). From Fig. 2 one sees that in the region \( T \sim 0.1 \div 1 \) MeV the mean field photon emission exceeds considerably both the fluctuation bremsstrahlung and the energy flux from \( e^+e^- \) pair production.

As we mentioned earlier, our assumption that the photon emission is a local process is valid if \( l_f \sim L_m \ll L_{el} \), where \( L_{el} \) is the typical scale of variation of the potential \( v_m \) along the electron trajectory. For the chemical potential \((1)\) it can evidently be defined as \( L_{el} \sim H\mu(0)/\mu(h)\cos\theta \), where \( \theta \) is the angle between the electron momentum and the star surface normal. Evidently the contribution of the configurations with \( L_m \gg L_{el} \) into the photon spectrum will be suppressed by the finite-size suppression factor \( S_{fs} \sim \min(L_{el}, L_m)/L_m \). We have checked numerically that this suppression factor gives a negligible effect. This justifies the local approximation.

Figs. 1, 2 demonstrate that the energy flux from the mean field photon emission may be of the same order of magnitude as the black body radiation. It says that the approximation of optically thin electrosphere is not very good, and the photon absorption and stimulated emission may be important. However, since the radiation rate we obtained does not exceed the black body limit, they cannot modify strongly our results \(^2\).

According to simulation of the thermal evolution of young quark stars performed in \([20]\) the temperature at the star’s surface becomes \( \sim 0.2 \) MeV at \( t \sim 1 \) s. However, in the analysis \([20]\) the mean field bremsstrahlung was not taken into account. In the light

\(^2\)The authors of \([12]\) obtained for \( \mu(0) \sim 10 - 20 \) MeV and \( T \lesssim 1 \) MeV the energy flux considerably exceeding the black body limit. They claim that it is possible for the electrosphere. This statement is obviously incorrect. The violation of the black body limit in \([12]\) is just a signal that the thin medium approximation becomes inadequate at high emissivity. As far as a very large emissivity obtained in \([12]\) is concerned, as we already mentioned, it may be due to incorrect description of the Pauli-blocking and neglect of the photon mass.
of our results one can expect that the cooling of the bare quark star’s surface should go somewhat faster than predicted in [20]. Higher luminosity due to the mean field bremsstrahlung increases the possibility for detecting bare quark stars. From the point of view of the light curves at $t \gtrsim 1$ s it would be interesting to investigate the mean field bremsstrahlung for $T \lesssim 0.1$ MeV as well. However, at such temperatures the photon emission from the nonrelativistic region of the electrosphere may be important, where our formulas become inapplicable. As far as the contribution of the relativistic region $\mu \gg m_e$ is concerned. Extrapolation of the curves shown in Fig. 2 to $T \lesssim 0.1$ MeV allows one to expect that the mean field emission will dominate the energy flux at lower temperatures as well. However, a robust conclusion on the relative contributions of the photon emission and $e^+e^-$ pair production can only be made after calculating the photon bremsstrahlung beyond the collinear approximation (in the relativistic and nonrelativistic regions of the electrosphere).

It is worth noting that for $T \sim 0.1 \div 1$ the form of the differential radiated energy flux and the relative fractions of photons and $e^\pm$ pairs are not important from the point of view of the photon spectrum observed at large distances from the star. One can show that in this temperature region for the energy flux of the order of the black body limit the outflowing wind of photons, electrons and positrons is thermalized at distances much smaller than the star radius. For the thermalized $e^\pm\gamma$ wind the photon distribution seen by a distant observer is close to the black body one, and the fraction of electrons and positrons is negligible [28]. For this reason the specific form of the photon spectrum from $e^+e^-$ annihilation for the tunnel $e^+e^-$ creation mechanism [4] is not important in the investigated temperature window. It may be important only at much smaller temperatures in the regime of a free streaming $e^+e^-\gamma$ wind.

The calculations of the photon emission from bare quark stars in the color flavor locked (CFL) superconducting phase (when electrons are probably absent even near the star surface [29]) performed in [30] give the radiation rate comparable to the black body limit. Since we also obtain the radiation rate comparable to the black body radiation it may be difficult to distinguish a bare quark star in the CFL phase from that in normal (or 2SC) phase.

4. In summary, using the LCPI reformulation [14] of the AMY approach [13] to the photon emission from relativistic plasmas we have calculated the photon emission from the electrosphere of a bare quark star (in normal or 2SC phase). Contrary to the previous qualitative studies [9, 10, 12], it allows, for the first time, to give a robust treatment of the Pauli-blocking effects in the photon bremsstrahlung. We demonstrate that for the temperatures $T \sim 0.1 \div 1$ MeV the dominating contribution to the photon emission is due to bending of electron trajectories in the mean electric field of the electrosphere. The energy flux from the mean field photon emission is of order of the black body limit. Our results show that the contribution of the bremsstrahlung due to electron-electron interaction is negligible as compared to the mean field photon emission.

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3It is worth noting, however, that in the initial hot stage the mean field bremsstrahlung will change only the temperature of the quark star near its surface. While the evolution of the star core temperature is driven by the neutrino emission [20] since for an extended period of time the neutrino luminosity is much larger than the photon (and $e^+e^-$) luminosity [20].
The energy flux related to the mean field bremsstrahlung turns out to be larger than that from the tunnel $e^+e^-$ creation [4, 11] as well. In the light of these results the situation with distinguishing bare quark stars made of the SQM in normal (or 2SC) phase from neutron stars may be more optimistic than in the scenario with the tunnel $e^+e^-$ creation discussed in [20].

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Figure 1: The differential radiated energy fluxes from the electrosphere for the mean field bremsstrahlung (solid line) and for the bremsstrahlung due to electron-electron interaction with (short dashes) and without (long dashes) the mean field suppression. The dotted curves show the black body spectrum.
Figure 2: The total radiated energy fluxes (scaled to the black body radiation) from the electrosphere for the mean field bremsstrahlung (solid line) and for the bremsstrahlung due to electron-electron interaction with (short dashes) and without (long dashes) the mean field suppression. The contribution from the tunnel $e^+e^-$ creation [4, 11] evaluated using (10) is also shown (dash-dotted line).