Electromagnetic response of quark-gluon plasma in heavy-ion collisions

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We study the electromagnetic response of the quark-gluon plasma in AA-collisions at RHIC and LHC energies for a realistic space-time evolution of the plasma fireball. We demonstrate that for a realistic electric conductivity the electromagnetic response of the plasma is in a quantum regime when the induced electric current does not generate a classical electromagnetic field, and can only lead to a rare emission of single photons.

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I. INTRODUCTION

Prediction of the chiral magnetic effect \[1\] in AA-collisions stimulated studies of magnetic field generated in heavy-ion collisions. In the noncentral AA-collisions the magnetic field perpendicular to the reaction plane can reach the values \(eB \sim 3m^2_\pi\) for RHIC (\(\sqrt{s} = 200\) GeV) and a factor of 15 bigger for LHC (\(\sqrt{s} = 2.76\) TeV) conditions \[1,3\]. In the initial stage the magnitude of the magnetic field falls rapidly with time (\(|B_0| \propto t^{-3}\), \(y\)-axis being perpendicular to the reaction plane). It was suggested \[1,3\] that the presence of the hot quark-gluon plasma (QGP) may increase the lifetime of the strong magnetic field. This may be important for a variety of new phenomena, such as the anomalous transport effects (for recent reviews, see \[6,7\]), the magnetohydrodynamics effects \[8,9\], the magnetic field induced photon production \[10,11\].

The effect of the QGP on the evolution of the electromagnetic field in AA-collisions has been estimated under the approximation of a uniform static matter in \[4,5,8,12\]. The difference between the calculations of \[4,5,8\] and that of \[12\] is that in \[4,5,8\] the nuclei all the time move in the matter, and in \[12\] it was assumed that the matter exists only after the AA-collision at \(t > 0\). In \[4,5,8\] a strong increase of the lifetime of the magnetic field in the presence of the QGP was found. But one can expect that in the model of \[4,5,8\] the matter effects should be underestimated, since there is an infinite time for the formation of the electromagnetic field around the colliding nuclei. In \[12\] it was obtained that for reasonable values of the conductivity the matter does not increase the lifetime of the strong \((eB/m^2_\pi \sim 1)\) magnetic field, and a significant effect was found only for the long-time evolution where \(eB/m^2_\pi \ll 1\). The model of \[12\] seems to be more realistic, but nevertheless it also may be too crude, since in reality the matter does not occupy the whole space at \(t > 0\). The plasma fireball is formed only in the region inside the light-cone \(t > |z|\) between the flying apart remnants of the colliding nuclei, and in a restricted transverse region of the overlap of the colliding nuclei. Evidently, it is highly desirable to evaluate the electromagnetic response for a realistic space-time evolution of the matter.

II. THEORETICAL FRAMEWORK

The electromagnetic field tensor satisfies the Maxwell equations

\[
\frac{\partial F_{\mu\nu}}{\partial x^\lambda} + \frac{\partial F_{\nu\lambda}}{\partial x^\mu} + \frac{\partial F_{\lambda\mu}}{\partial x^\nu} = 0,
\]

(1)
\[
\frac{\partial F^{\mu\nu}}{\partial x^\mu} = -j^{\nu}.
\]

For AA-collisions the current \( J^\mu \) may be decomposed into two physically different pieces:

\[
J^\mu = J^\mu_{\text{ext}} + J^\mu_{\text{in}}.
\]

Here the term \( J^\mu_{\text{ext}} \), which we call the external current, is the contribution of the fast right and left moving charged particles, which are mostly protons of the colliding nuclei. And the term \( J^\mu_{\text{in}} \) is the induced current generated in the created hot QCD matter. We decompose the field tensor also into the external and the induced pieces:

\[
F^{\mu\nu} = F^{\mu\nu}_{\text{ext}} + F^{\mu\nu}_{\text{in}}.
\]

Both \( F^{\mu\nu}_{\text{ext}} \) and \( F^{\mu\nu}_{\text{in}} \) separately satisfy the first Maxwell equation (1) and the following Maxwell equations with sources:

\[
\frac{\partial F^{\mu\nu}_{\text{ext}}}{\partial x^\mu} = -j^\nu_{\text{ext}},
\]

\[
\frac{\partial F^{\mu\nu}_{\text{in}}}{\partial x^\mu} = -j^\nu_{\text{in}}.
\]

We assume that Ohm’s law is valid in the fireball. Then the induced current reads

\[
j^\mu_{\text{in}} = \rho u^\mu + \sigma (F^{\mu\nu}_{\text{ext}} + F^{\mu\nu}_{\text{in}}) u^\nu,
\]

where \( \sigma \) is the electric conductivity of the QCD matter, \( \rho \) is its charge density, and \( u^\mu \) is the four-velocity of the matter. For the Bjorken 1+1D expansion of the fireball \( u^\mu = (t/\tau, 0, 0, z/\tau) \), where \( \tau = \sqrt{t^2 - z^2} \) is the proper time.

The induced current (7) couples \( F^{\mu\nu}_{\text{ext}} \) to \( F^{\mu\nu}_{\text{in}} \). And \( F^{\mu\nu}_{\text{ext}} \) does not depend on the fireball evolution at all. We approximate \( J^\mu_{\text{in}} \) simply by the currents of the two colliding nuclei with the velocities \( V_R = (0, 0, V) \) and \( V_L = (0, -b/2) \) and the impact parameters \( b_R = (0, -b/2) \) and \( b_L = (0, b/2) \) as shown in Fig. 1. We assume that in the center of mass frame of the AA-collision the trajectories of the centers of mass of the colliding nuclei in the longitudinal direction \( z \) are \( z_{R,L} = \pm Vt \). The contribution of each nucleus to \( F^{\mu\nu}_{\text{ext}} \) is given by the Lorentz transformation of its Coulomb field. We write the electric and magnetic fields of a nucleus with the velocity \( V = (0, 0, V) \) and the impact vector \( b \) as

\[
E_T(t, \rho, \zeta) = \frac{E_A(r') (\rho - b)}{r'},
\]

\[
E_z(t, \rho, \zeta) = \frac{E_A(r') \zeta'}{r'},
\]

\[
B(t, \rho, \zeta) = |V \times E|.
\]

Here \( \gamma = 1/\sqrt{1-V^2} \) is the Lorentz factor, \( r'^2 = (\rho - b)^2 + z'^2, z' = \gamma(z-Vt), \) and

\[
E_A(r) = \frac{1}{r} \int_0^r d\xi \xi^2 \rho_A(\xi)
\]

is the electric field of the nucleus in its rest frame, \( \rho_A \) is the nucleus charge density. In our calculations we used for \( \rho_A \) the Woods-Saxon parametrization. From (8)–(11) one can obtain that at \( t'^2 \geq (R_A^2 - b^2/4)/\gamma^2 \) (here \( R_A \) is the nucleus radius, and \( b \) is assumed to be \( < 2R_A \)) and \( r = 0 \) the only nonzero \( y \)-component of the magnetic field for the two colliding nuclei is approximately

\[
B_y(t, r = 0) \approx \frac{\gamma Ze_b}{4\pi(b^2/4 + \gamma^2 V^2 t^2)^{3/2}}.
\]

From (8)–(11) one can obtain that at \( t \gg R_A/\gamma \)

\[
B_y(t, \rho, z = 0) \approx Ze_b/4\pi R_A^3.
\]

The quantity \( R_A/\gamma \) is very small: \( \sim 0.06 \) for Au+Au collisions at RHIC energy \( \sqrt{s} = 200 \text{ GeV} \) and \( \sim 0.004 \) fm for Pb+Pb collisions at LHC energy \( \sqrt{s} = 2.76 \text{ TeV} \).

At \( t'^2 \geq (R_A^2 - b^2/4)/\gamma^2 \) the \( t \)-dependence of \( B_y(r = 0) \) flattens and at \( t = 0 \) one can obtain

\[
B_y(t = 0, r = 0) \approx \gamma Ze_b/4\pi R_A^3.
\]

Here the right-hand side corresponds to the spherical nuclei. For the realistic Woods-Saxon distribution of the protons the result is just a bit (\( \sim 5\% \)) smaller.

III. MODEL OF THE FIREBALL

The interaction of the Lorentz-contracted nuclei lasts for a short time from \( t \sim -R_A/\gamma \) to \( R_A/\gamma \). As the nuclei fly apart after the collision a hot fireball is created. It is widely accepted that the creation of the plasma fireball goes through the thermalization of the plasma longitudinal color fields created after multiple color exchanges between the colliding nuclei. We performed the calculations for the Bjorken longitudinal expansion of the fireball that gives the \( T \)-dependence of the entropy density \( s \propto 1/T \). We also performed the calculations accounting for the corrections to the Bjorken picture from the transverse and the additional longitudinal expansions of the fireball treating them perturbatively as described in \[13\]. We assume that these corrections come into play at \( \tau_0 = 0.5 \) fm. Roughly such \( \tau_0 \) is often used in the hydrodynamical simulations of AA-collisions (for a recent review, see [16]). But we observed that these corrections give a negligible effect.

For simplicity as in [14] we parametrize the initial entropy density profile at the proper time \( \tau_0 \) in a Gaussian form

\[
s(x, y, \eta_0) \propto \exp \left(-\frac{x^2}{2\sigma_x^2} - \frac{y^2}{2\sigma_y^2} - \frac{\eta_0^2}{2\sigma_\eta^2}\right).
\]
Here $\sigma_x$ and $\sigma_y$ are the root mean square widths of the fireball in the transverse directions, and $\sigma_\eta$ is the root mean square width in the space-time rapidity. $\eta_s = \frac{1}{2} \ln \left( \frac{1 + \tau}{1 - \tau} \right)$. We adjusted the parameters $\sigma_{x,y}(\tau_0)$ using the entropy distribution in the transverse coordinates at $\eta_s = 0$ given by

$$\frac{dS(\eta_s = 0)}{dN_{ch}/d\eta} = \frac{dS(\eta_s = 0)}{d\eta_s} - \frac{\alpha N_{part}}{\alpha N_{part} + (1 - \alpha) N_{coll}} \frac{dN_{coll}}{d\eta_s},$$

(16)

where $dN_{part}/d\eta$ and $dN_{coll}/d\eta$ are the well known Glauber distributions of the participant nucleons and of the binary collisions (see, for instance, [16]). We used in (16) $\alpha = 0.95$. It allows to reproduce well the centrality dependence of the data on the pseudorapidity density $dN_{ch}/d\eta$ from STAR [17] for Au+Au collisions at $\sqrt{s} = 200$ GeV and from ALICE [18] and CMS [19] for Pb+Pb collisions at $\sqrt{s} = 2.76$ TeV. To fix the normalization of the entropy density we used the entropy/multiplicity ratio $\sigma_b$ obtained in [20]. Making use of the data on $dN_{ch}/d\eta$ [17, 19] we obtained for the impact parameter $b = 6$ fm $\sigma_x(\tau_0) \approx 2.3$ fm and $\sigma_y(\tau_0) \approx 3.92$ fm for RHIC and $\sigma_x(\tau_0) \approx 2.42$ fm and $\sigma_y(\tau_0) \approx 3.13$ fm for LHC. We take $\sigma_\eta(\tau_0) = 2.63$ and 4.03 for RHIC and LHC, respectively, that allow to reproduce qualitatively the experimental $\eta$-dependence of $dN_{ch}/d\eta$. In evaluating the temperature through the entropy we used the ideal gas formula for the number of flavors $N_f = 2.5$. It gives the temperature at the center of the fireball at $\tau = 0.5$ fm: $T \approx 400$ MeV for RHIC and $T \approx 520$ MeV for LHC.

It seems likely that the model of the QGP as a conducting matter makes sense at $\tau \gtrsim \tau_0$ when the hydrodynamics is assumed to be applicable. At present, the details of the thermal and chemical equilibration of the matter at early times $\tau \lesssim \tau_0$ are unclear. Often it is assumed that the plasma thermalization starts with the gluon dominated stage and the production of quarks is somewhat delayed (see, for instance, [21, 22]). But it is possible that immediately after the AA-collision the amount of quarks are close to that for the chemically equilibrated QGP [24]. However, even in this case it is hardly possible to describe the matter in terms of the equilibrium conductivity because anyway the thermalization requires some time. Nevertheless, since we would like to demonstrate that the electromagnetic response of the QGP is too small for applicability of the classical treatment, we will consider a maximally optimistic scenario. We assume that already at $\tau \gtrsim R_A/\gamma$ the conductivity makes sense and equals to that for the equilibrium QGP with the entropy density $\propto 1/\tau$, as in the Bjorken model [13]. We solve the Maxwell equations (11) (for $F_{\mu\nu}^{em}$) and (14) with the initial condition $F_{\mu\nu}^{em} = 0$ at $\tau = R_A/\gamma$.

In our analysis we use the conductivity obtained in the most recent lattice calculations for $N_f = 3$ [25] for $T \sim 140 - 350$ MeV. This analysis gives $\sigma/C_{em} T$ which rises smoothly from $\sim 0.07$ at $T = 150$ MeV to $\sim 0.32$ at

$T = 350$ MeV. We parametrize the results of [25] in the form

$$\sigma = C_{em} T f(T/1 \text{ GeV}),$$

(17)

$$f(x) = \begin{cases} f_1 \frac{x_1(x_2-x_1)+x_2(x_1-x)}{x_2-x_1}, & \text{if } x \leq x_1, \\ f_2 \frac{x_1(x_2-x_1)+x_2(x_1-x)}{x_2-x_1}, & \text{if } x_1 < x < x_2, \\ f_3 \frac{x_1(x_2-x_1)+x_2(x_1-x)}{x_2-x_1}, & \text{if } x \geq x_2, \end{cases}$$

(18)

$f_{1-3} = 0.0662, 0.2153, 0.3185$ and $x_{1-3} = 0.1747, 0.234, 0.5516$. The results of [25] agree qualitatively with that obtained within the Dyson-Schwinger equation approach [26].

IV. RESULTS AND DISCUSSION

To numerically solve the Maxwell equations (11), (14) we rewrote them in the Milne coordinates $\{\tau, x, y, \eta_s\}$ which are convenient for imposing the initial condition at a given $\tau$ that we need. We used the Yee algorithm [27]. We found that it works well in the Milne coordinates as well. We performed the calculations for the impact parameter $b = 6$ fm for Au+Au collisions at RHIC energy $\sqrt{s} = 200$ GeV and for Pb+Pb collisions at LHC energy $\sqrt{s} = 2.76$ TeV.

In Fig. 2 we show the $t$-dependence of $B_y$ at $r = 0$. And in Fig. 3 we present the $x$-profile of $B_y$, $E_x$, and $E_z$ at $y = z = 0$ for $t = 1, 2$, and 4 fm. We show separately the results for the total (external plus induced) and for the external fields. We present the curves for the Bjorken model. We observed that the corrections to the Bjorken model due to the transverse and the longitudinal
field is approximately proportional to $h$, and rescaled the results of $[5,8]$ by the factor $6/7$. Unfortunately, there is some problem with comparing of our results with that of $[12]$ (for Au+Au collisions at $\sqrt{s} = 200$ and $b = 6$ fm). By examining the external magnetic field shown in Fig. 1 of $[12]$ we have found that prediction for the external field given there is clearly wrong. Indeed, from $[13]$ one obtains $eB_y(t = 0, r = 0)/m_e^2 \approx 2.73$, and accurate calculations with the Woods-Saxon density give for this quantity a bit smaller value 2.6, while Fig. 1 of $[12]$ gives $eB_y(t = 0, r = 0)/m_e^2 \approx 7$. Thus at $t = 0$ $[12]$ overestimates the field by a factor of $\sim 2.7$. The $t$-dependence of the external field in $[12]$ is also wrong, say Fig. 1 gives $eB_y(t = R_A, r = 0)/m_e^2 \approx 1.8 \cdot 10^{-4}$, while formula (13) gives for this quantity a value smaller by a factor of $\sim 80$. It is possible that the above strange behavior of the external field in Fig. 1 of $[12]$ is just a consequence of some errors in axis variables. We will assume that it is really the case. Then using the correct value $eB_y(t = 0, r = 0)/m_e^2 \approx 2.6$ for normalization of the total (external plus induced) field in the region where in Fig. 1 of $[12]$ it flattens, we obtain $eB_y(t \sim 1\text{fm}, r = 0)/m_e^2 \sim 0.017$. It is about a factor of 3.7 smaller than prediction of $[5]$ (rescaled by a factor 6/7 accounting for difference in $b$), and by a factor of 13 bigger than our prediction $eB_y(t \sim 1\text{fm}, r = 0)/m_e^2 \approx 0.0013$.

The magnitude of the difference in predictions of $[5]$ and $[12]$ seems to be quite reasonable since in $[5]$ there is a contribution from the unreacted region at $t < 0$. But the difference between $[12]$ and our results seems to be too big to be explained by the difference in the space-time distribution of the conductivity. One can expect that the latter could only give a factor of $\sim 2 - 4$. So there must be a mechanism which enhances the medium effect in the model of $[12]$ as compared to our one. It seems likely that it is the difference in the induced current for the static matter and the matter with the Bjorken longitudinal expansion. It can be seen by comparing the form of the induced current in these two models. Our calculations show that the effect of the induced electromagnetic field in the induced current is practically negligible. So we can consider in the induced current only the term with $F_{\mu \nu \tau \eta}^{\mu \nu}$. For the static matter the dominating transverse component of the current for each of the colliding nuclei reads $J_T = \sigma E_T$. So each nucleus produces a running pancake-like distribution of the induced current. It acts as an antenna radiating the induced electromagnetic field. One can easily show that for the Bjorken expansion of the matter the electromagnetic field in the current is replaced by the transverse electric field in the comoving frame. The latter is suppressed by a factor $\exp(-|\eta_s|)$ as compared to $E_T$ in the center mass frame. In the vicinity of the nuclei this factor may be $\sim 1/\gamma$. Of course, the induced field acquires the contributions radiated from the points with different $\eta_s$, and the resulting suppression factor should be bigger than $1/\gamma$. Nevertheless, it is clear that the suppression effect may be quite strong. The finite size of the fireball in the rapidity also can give an ad-

![Diagram](image-url)
dential suppression of the induced field. Unfortunately, the induced field for the conditions of Ref. [12] cannot be computed directly with our code because it is written in the Milne coordinates and works only inside the region \( t > |z| \) (for finite values of the proper time \( \tau = \sqrt{t^2 - z^2} \)). However, we tested that in our formulation for the finite fireball with zeroth longitudinal velocity and a flat distribution of the conductivity in the rapidity the induced field is really enhanced by a factor of \( \sim 8 \) and the results turns out to be qualitatively similar to that of [12]. In the formulation of [3, 8], where an unrestricted region of time and transverse space is involved into the formation of the induced field at a given space-time point, the enhancement may be considerably bigger. Therefore the observed disagreement of our results with that of [3, 8] do not seem to be unrealistic.

From the results shown in Figs. 2, 3 one can show that the electromagnetic response of the plasma fireball is in reality in a quantum regime. Indeed, it is known [28] that the classical treatment of an electromagnetic field is valid when

\[
|E|, |B| \gg 1/\Delta t^2, \quad (19)
\]

where \( \Delta t \) is the typical time of the observation. For \( \Delta t \) one can simply take the typical time of the variation of the fields [28]. The inequality (19) follows from the condition that for the classical description the occupation numbers should be large. It is especially transparent for a free field occupying (at a given instant) a restricted region of space. If the size of the region is \( L \), then \( \Delta t \sim L \), and the dominating Fourier component should have a frequency \( \omega \gtrsim 1/\Delta t \). Then by requiring that the energy of the field, which is \( \sim L^3(E^2 + B^2)/2 \), is much bigger than the typical one photon energy \( \omega \) one obtains (19). In our case one can take \( \Delta t \sim t \). From the curves shown in Figs. 2, 3 one can easily see that the induced fields are much smaller than \( 1/t^2 \). It means that the typical photon occupations numbers are much smaller than unity. In this situation one cannot talk about classical fields, and the electromagnetic response of the QGP should be described as radiation of single photons. This fact is quite evident from the x-profile of the magnetic and electric fields shown in Fig. 3. On the one hand, one can see that the typical wave vector is of the order of \( 1 \text{ fm}^{-1} \). On the other hand, if we estimate the total energy of the field \( U \), say, at \( t \sim 4 \text{ fm} \) taking (with a large excess) for the volume \( V \sim 1000 \text{ fm}^3 \) we obtain in fm units \( U \lesssim 0.01 \text{ fm}^{-1} \). It is much smaller than the expected typical photon energy \( \sim 1 \text{ fm}^{-1} \). This means that we have a situation of the deep quantum regime when the electromagnetic response of the QGP is a very rare emission of the single photons which freely leave the fireball. The fact that the photons are not absorbed in the fireball is evident from calculation of the photon attenuation length \( l_\alpha \approx 2/\sigma \), which, in our case, turns out to be very large \( l_\alpha \sim 100 \text{ fm} \) for \( T \sim 250 \text{ MeV} \). This is why we have found the negligible effect of switching off the conductivity at \( \tau > 1 \text{ fm} \). An experimental observation of the photons radiated by the induced current is practically impossible due to a huge background from other mechanisms of the photon production. Note that since the electromagnetic response of the QGP consists of emission of the single photons, it cannot contribute to the magnetohydrodynamics effects [3, 9] and to the magnetic field induced photon production [10, 11].

One remark is appropriate at this point. From Figs. 2, 3 (or from Eq. (13)) one can see that the external fields also do not satisfy the criterion (19). It seems to be in contradiction with the estimate of the photon occupation numbers for the electromagnetic field of a fast nucleus. Indeed, one can easily show that for \( \gamma \gg 1 \) the typical occupation numbers are

\[
N \sim Z^2 \alpha_{em}/4\pi. \quad (20)
\]

From (20) one sees that for \( Z \sim 100 \) the classical approximation should work well, at least except for the tail regions (in the longitudinal direction) where the field becomes very small and the situation may be a quantum one. This puzzling situation with the contradiction of the criterion (19) and the estimate (20) is related to the fact that in [28] in deriving (19) it was implicitly assumed that the distribution of the photon modes is more or less isotropic. This is clearly not true for the field of a fast nuclei, when the field has a pancake-like form and the modes are strongly collimated in the direction of the nucleus velocity. One can easily show that in this situation the criterion (19) should be replaced by

\[
|E|, |B| \gg 1/\Delta t \Delta \rho, \quad (21)
\]

where \( \Delta \rho \) is the typical scale of variation of the fields in the transverse directions. One can see that, taking \( \Delta \rho \sim \rho \), where \( \rho \) is the transverse distance from the nucleus, and \( \Delta t \sim \rho/\gamma \), the criterion (21) and the estimate (20) of the occupation numbers give the same condition for the validity of the classical description. For the induced field which occupies the whole fireball region and is not collimated in one direction the criterion (19) should work. Our estimate of the occupation numbers based on the comparison of the energy with the typical wave vector confirms this. We would like to emphasize that, in any case, our conclusion about the quantum character of the electromagnetic response is completely independent of the situation (classical or quantum) with the external field at later times where we apply our energy argument because in this region the external field becomes very small as compared to the induced field and can simply be ignored.

Note that our main result, that the electromagnetic response of the fireball is in a quantum regime, persists for a wide range of the electric conductivity. We checked that for \( \sigma = 7C_{em}T \) from the early analysis [29], which is by a factor of \( \sim 20 \) (at \( T \sim 300 \text{ MeV} \)) larger than that of [27], the induced magnetic and electric fields also violate the inequality (19). It is hardly possible that \( \sigma \) can be bigger than that of [29], since even it seems to
be too large. Indeed, using the Drude formula one can show that, in terms of the quark collisional time $\tau_c$, $\sigma$ from $[29]$ corresponds to $\tau_c/\tau \sim 10 - 20$ for $\tau \sim 0.5 - 1$ fm and $\tau_c/\tau \sim 4 - 8$ for $\tau \sim 2 - 4$ fm. Such large ratios say that the quarks are in a ballistic regime. This cannot be reconciled with the successful hydrodynamical description of the flow effects in $AA$-collisions $[15]$. The results of $[25]$ seem to be more realistic. In this case $\tau_c/\tau \sim 0.2 - 0.3$ for $\tau \sim 0.5 - 1$ fm and $\tau_c/\tau \sim 0.05 - 0.1$ for $\tau \sim 2 - 4$ fm which look quite reasonable (from the viewpoint of the applicability of the hydrodynamics and of the model of a conducting matter).

In summary, we have studied the electromagnetic response of the QGP in the noncentral $AA$-collisions at RHIC and LHC energies by solving the Maxwell equations for a realistic space-time evolution of the plasma fireball. We demonstrate that the resulting induced electromagnetic field turns out to be too small for applicability of the classical treatment, and in reality the electromagnetic response is in a quantum regime, when the induced electric current in the plasma fireball cannot generate a classical electromagnetic field at all. In this regime the electromagnetic response consists only of a rare emission of the single photons. Thus, the emerging physical picture of the electromagnetic response of the QGP differs qualitatively from that assumed previously.

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[1] D.E. Kharzeev, L.D. McLerran, and H.J. Warringa, Nucl. Phys. A803, 227 (2008).
[2] V. Skokov, A.Yu. Illarionov, and V. Toneev, Int. J. Mod. Phys. A24, 5925 (2009).
[3] A. Bzdak and V. Skokov, Phys. Lett. B710, 171 (2012).
[4] K. Tuchin, Adv. High Energy Phys. 2013, 490495 (2013).
[5] K. Tuchin, Phys. Rev. C88, 024911 (2013).
[6] D.E. Kharzeev, Prog. Part. Nucl. Phys. 75, 133 (2014).
[7] J. Liao, arXiv:1401.2500.
[8] U. Gursoy, D. Kharzeev, and K. Rajagopal, arXiv:1401.3805.
[9] M. Hongo, Y. Hirono, and T. Hirano, arXiv:1309.2823.
[10] G. Basar, D. Kharzeev, and V. Skokov, Phys. Rev. Lett. 109, 202303 (2012).
[11] K. Tuchin, Phys. Rev. C87, 024912 (2013).
[12] L. McLerran and V. Skokov, arXiv:1305.0774.
[13] J.D. Bjorken, Phys. Rev. D27, 140 (1983).
[14] J.-Y. Ollitrault, Eur. J. Phys. 29, 275 (2008).
[15] U. Heinz and R. Snellings, Ann. Rev. Nucl. Part. Sci. 63, 123 (2013).
[16] D. Kharzeev, C. Lourenco, M. Nardi, and H. Satz, Z. Phys. C74, 307 (1997).
[17] B.I. Abelev et al. [STAR Collaboration], Phys. Rev. C79, 034909 (2009).
[18] K. Aamodt et al. [ALICE Collaboration], Phys. Rev. Lett. 106, 032303 (2011).
[19] S. Chatrchyan et al. [CMS Collaboration] JHEP 1108, 141 (2011).
[20] B. Müller and K. Rajagopal, Eur. Phys. J. C43, 15 (2005).
[21] M. Chiu, T.K. Hemmick, V. Khachatryan, A. Leonidov, J. Liao, and L. McLerran, Nucl. Phys. A900, 16 (2013).
[22] A. Monnai and B. Müller, arXiv:1403.7310.
[23] L. McLerran and B. Schenke, arXiv:1403.7462.
[24] F. Gelis, K. Kajantie, and T. Lappi, Phys. Rev. Lett. 96, 032304 (2006).
[25] A. Amato, G. Aarts, P. Giudice, S. Hands, and A. Lindert, arXiv:1510.7466.
[26] Si-xue Qin, arXiv:1307.4587.
[27] K. Yee, IEEE Transactions on Antennas and Propagation, 14, 302 (1966).
[28] V.B. Berestetski, E.M. Lifshits and L.P. Pitaevski, Quantum Electrodynamics (Landau Course of Theoretical Physics Vol. 4), Oxford, Pergamon Press, 1979.
[29] S. Gupta, Phys. Lett. B597, 57 (2004).