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

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ABSTRACT

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.

1. 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 $eB \sim 3m_\pi^2$ 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 ($|B_y| \propto t^{-3}$, y-axis being perpendicular to the reaction plane). It was suggested [4,5] 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 overestimated, 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 ($eB/m_\pi^2 \sim 1$) magnetic field, and a significant effect was found only for the long-time evolution where $eB/m_\pi^2 \ll 1$. 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.

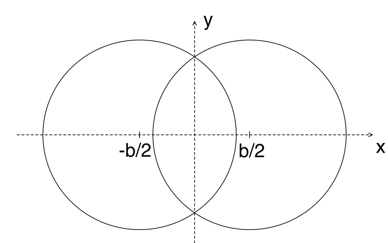

Fig. 1. The transverse plane of a noncentral AA-collision with the impact parameter b.

In this Letter we study the electromagnetic response of the QGP in the noncentral AA-collisions for a realistic expanding plasma fireball which is created inside the light-cone $t > |z|$ in the almond-shaped transverse overlap of the colliding nuclei as shown in Fig. 1. We demonstrate that the physical picture of the electromagnetic response is qualitatively different from the one assumed in previous studies. Our numerical results show that for a realistic electric conductivity the induced electromagnetic field generated in the fireball turns out to be too small for applicability of the classical treatment. We show that for both RHIC and LHC energies the electromagnetic response is essentially in the deep quantum regime when one cannot talk about a classical electromagnetic field at all. In this regime the induced current in the QGP can just produce single photons which freely leave the fireball without generation of an additional induced current in the QGP. The

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probability of the photon emission from this mechanism is very small, and, due to a huge background from other mechanisms of the photon production, an experimental observation of the photons from this mechanism is practically impossible.

2. Theoretical framework

The electromagnetic field tensor satisfies the Maxwell equations

$$\frac{\partial F_{\mu\nu}}{\partial x^\lambda}+\frac{\partial F_{\nu\lambda}}{\partial x^\mu}+\frac{\partial F_{\lambda\mu}}{\partial x^\nu}=0, \tag{1}$$

$$\frac{\partial F^{\mu\nu}}{\partial x^\nu}=-J^\mu. \tag{2}$$

For AA-collisions the current J^μ may be decomposed into two physically different pieces:

$$J^\mu=J^\mu_{ext}+J^\mu_{in}. \tag{3}$$

Here the term J^μ_{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^μ_{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:

$$F^{\mu\nu}=F^{\mu\nu}_{ext}+F^{\mu\nu}_{in}. \tag{4}$$

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

$$\frac{\partial F^{\mu\nu}_{ext}}{\partial x^\nu}=-J^\mu_{ext}, \tag{5}$$

$$\frac{\partial F^{\mu\nu}_{in}}{\partial x^\nu}=-J^\mu_{in}. \tag{6}$$

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

$$J^\mu_{in}=\rho u^\mu+\sigma\left(F^{\mu\nu}_{ext}+F^{\mu\nu}_{in}\right)u_\nu, \tag{7}$$

where σ is the electric conductivity of the QCD matter, ρ is its charge density, and u^μ is the four-velocity of the matter. For the Bjorken $1+1$D expansion [13] 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 $F^{\mu\nu}_{in}$ to $F^{\mu\nu}_{ext}$. And $F^{\mu\nu}_{ext}$ does not depend on the fireball evolution at all. We approximate J^μ_{ext} simply by the currents of the two colliding nuclei with the velocities $\mathbf{V}_R=(0,0,V)$ and $\mathbf{V}_L=(0,0,-V)$ and with the impact parameters $\mathbf{b}_R=(0,-b/2)$ and $\mathbf{b}_L=(0,b/2)$ 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 $F^{\mu\nu}_{ext}$ is given by the Lorentz transformation of its Coulomb field. We write the electric and magnetic fields of a nucleus with the velocity $\mathbf{V}=(0,0,V)$ and the impact vector $\mathbf{b}$ as

$$\mathbf{E}_T(t,\boldsymbol{\rho},z)=\gamma\frac{E_A(r')(\boldsymbol{\rho}-\mathbf{b})}{r'}, \tag{8}$$

$$E_z(t,\boldsymbol{\rho},z)=\frac{E_A(r')z'}{r'}, \tag{9}$$

$$\mathbf{B}(t,\boldsymbol{\rho},z)=[\mathbf{V}\times\mathbf{E}]. \tag{10}$$

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

$$E_A(r)=\frac{1}{r^2}\int_0^r d\xi\,\xi^2\rho_A(\xi) \tag{11}$$

is the electric field of the nucleus in its rest frame, ρ_A is the nucleus charge density. In our calculations we used for ρ_A the Woods–Saxon parametrization. From (8)–(11) one can obtain that at $t^2\gtrsim(R_A^2-b^2/4)/\gamma^2$ (here R_A is the nucleus radius, and b is assumed to be $<2R_A$) and $\mathbf{r}=0$ the only nonzero y-component of the magnetic field for the two colliding nuclei is approximately

$$B_y(t,\mathbf{r}=0)\approx\frac{\gamma Zeb}{4\pi(b^2/4+\gamma^2V^2t^2)^{3/2}}. \tag{12}$$

From (8)–(11) one can obtain that at $t\gg R_A/\gamma$ $B_y(t,\boldsymbol{\rho},z=0)$ in the region $\rho\ll t\gamma$ takes a simple ρ-independent form

$$B_y(t,\boldsymbol{\rho},z=0)\approx Zeb/4\pi\gamma^2t^3. \tag{13}$$

The quantity R_A/γ is very small: ~ 0.06 for Au + Au collisions at RHIC energy $\sqrt{s}=200$ GeV, and ~ 0.004 fm for Pb + Pb collisions at LHC energy $\sqrt{s}=2.76$ TeV. At $t^2\lesssim(R_A^2-b^2/4)/\gamma^2$ the t-dependence of $B_y(\mathbf{r}=0)$ flattens and at $t=0$ one can obtain

$$B_y(t=0,\mathbf{r}=0)\approx\gamma Zeb/4\pi R_A^3. \tag{14}$$

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.

3. 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/γ. 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 glasma longitudinal color fields created after multiple color exchanges between the colliding nuclei. We performed the calculations for the Bjorken longitudinal expansion [13] of the fireball that gives the τ-dependence of the entropy density $s\propto 1/\tau$. 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 [14]. We assume that these corrections come into play at $\tau_0=0.5$ fm. Roughly such τ_0 is often used in the hydrodynamical simulations of AA-collisions (for a recent review, see [15]). 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 τ_0 in a Gaussian form

$$s(x,y,\eta_s)\propto\exp\left(-\frac{x^2}{2\sigma_x^2}-\frac{y^2}{2\sigma_y^2}-\frac{\eta_s^2}{2\sigma_\eta^2}\right). \tag{15}$$

Here σ_x and σ_y are the root mean square widths of the fireball in the transverse directions, and σ_η is the root mean square width in the space–time rapidity $\eta_s=\frac{1}{2}\ln(\frac{t+z}{t-z})$. We adjusted the parameters $\sigma_{x,y}(\tau_0)$ using the entropy distribution in the transverse coordinates at $\eta_s=0$ given by

$$\frac{dS(\eta_s=0)}{d\eta_s d\boldsymbol{\rho}}=\frac{dS(\eta_s=0)}{d\eta_s}\cdot\frac{\alpha\frac{dN_{part}}{d\boldsymbol{\rho}}+(1-\alpha)\frac{dN_{coll}}{d\boldsymbol{\rho}}}{\alpha N_{part}+(1-\alpha)N_{coll}}, \tag{16}$$

where $dN_{part}/d\boldsymbol{\rho}$ and $dN_{coll}/d\boldsymbol{\rho}$ 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