Dynamic Debye screening for a heavy-quark-antiquark pair traversing a quark-gluon plasma

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We investigate dynamical plasma screening effects on a heavy-quark–antiquark pair traversing a quark-gluon plasma. The screened potential created by a test charge (heavy quark) moving in the plasma medium is calculated in the test charge frame by solving the transport equation for a collisionless ultrarelativistic plasma interacting via Abelian gauge field. It is shown that the screened potential becomes strongly anisotropic as the velocity of the plasma medium increases. Possible implications of this effect for charmonium production in relativistic heavy-ion collisions are discussed.

I. INTRODUCTION

In a high-temperature quark-gluon plasma, deconfinement and the plasma screening effect are expected to modify the heavy-quark–antiquark ($Q\bar{Q}$) potential and above a certain threshold temperature this may lead to a dissolution of $Q\bar{Q}$ bound states (charmonium). It has been suggested[1] that the observation of this effect through a strong and systematic suppression of the J/ψ peak in the dilepton mass spectrum can be used to test the formation of a quark-gluon plasma in ultrarelativistic nucleus-nucleus collisions.

The recent experimental results by the NA38 Collaboration at the CERN SPS have shown[2] an apparent but systematic suppression of J/ψ production in the collisions of oxygen and sulfur beams at 200 GeV/nucleon on heavy nuclear targets. The measured peak-to-continuum ratio in the dilepton mass spectrum exhibits a systematic dependence on the transverse energy deposited by the collision and on the transverse momentum of dileptons: this ratio is more suppressed for events with large transverse energy and for dileptons with small transverse momenta. It has been shown that the observed magnitude and pattern of J/ψ suppression in this experimental condition can be well explained in terms of a simple model[3–5] which incorporates the finite J/ψ formation time and the finite plasma lifetime due to the longitudinal scaling expansion.[6]

In order to confirm the observed J/ψ suppression as a signature of quark-gluon-plasma formation, we must however, rule out other possible nonplasma effects which may mimic the plasma suppression: such effects include nuclear absorptions.[7,8] The collisional loss of J/ψ in dense hadron or parton gas,[9–12] and the distortion of quark/gluon distribution by precollision[13] or by initial-state interactions.[14–16] These different suppression mechanisms lead to different results in the magnitude and the pattern of J/ψ suppression (transverse-energy dependence, J/ψ transverse-momentum dependence, nuclear mass dependence, incident-energy dependence, etc.), and hence it could in principle be possible to disentangle these different effects from genuine plasma effect by more detailed systematic study both theoretically and experimentally. For such a purpose, it is desirable to seek other dynamical effects characteristic of the plasma formation.

In this paper we study the dynamical plasma screening effect which may arise when the $Q\bar{Q}$ pair is moving in the plasma medium. Such a situation would be relevant for the formation of J/ψ with large transverse momentum since a $c\bar{c}$ pair with large transverse momentum can escape the plasma region forming J/ψ (Refs. 3–5). Let us recall briefly the standard derivation of the Debye screening for a static test charge Q placed in a plasma. In this case, the electrostatic potential $\varphi(x)$ is determined by Poisson's equation $\Delta\varphi(x)=Q\delta(x)+\rho_{\rm ind}(x)$, where $\rho_{\rm ind}$ is the induced polarization charge density. By using the linear response relation between the induced charge density and the potential $\rho_{\rm ind}(x)=\Pi_{00}\varphi(x)$, where Π_{00} is the time component of the photon polarization tensor in plasma, one finds the well-known screened potential

$$\varphi(r)=\frac{Q}{4\pi r}e^{-m_e r}\ , \tag{1}$$

where $m_e=\sqrt{\Pi_{00}}$ is the inverse screening length or the electric screening mass. This derivation of the Debye screening is adequate only for the static test charge in plasma at rest. It is well known that this result is modified when the test charge is moving in the plasma.[17] One can find good examples of the manifestation of the dynamic screening effect, for instance, in the Coulomb explosion of molecular ions transmitted through metallic foil[18] and in the calculation of thermonuclear fusion cross section in stellar plasma.[19]

For simplicity, we consider here a model relativistic plasma whose massless constituents are interacting via an Abelian gauge field A^μ, and we study the problem within the framework of the kinetic theory. Since we are interested in the screening of the potential acting between a heavy quark and its antiquark *moving together* in the plasma, our objective is to find how the plasma screening is seen *by an observer moving together with the test charge.* It will be shown that the screened potential becomes anisotropic and acquires nonzero vector components $\mathbf{A}$ which generate (static) magnetic field around the test charge. We make some speculations at the end for the possible consequences of our results for the charmonium production in nucleus-nucleus collisions.

II. SCREENED POTENTIALS IN MOVING PLASMA

Let us consider a very heavy test charge Q traversing a plasma with velocity v on the z axis in the $-z$ direction. We assume for simplicity that the test charge is very massive so that we can ignore the energy transfer between the test charge and the plasma. This implies that in the rest frame of the test charge the field is static and can be obtained by solving Poisson's equation for the static potential $A^\mu(\mathbf{r})$:

$$\Delta A^\mu(\mathbf{r}) = j^\mu_{\text{ext}}(\mathbf{r}) + j^\mu_{\text{ind}}(\mathbf{r}) , \tag{2}$$

where $j^\mu_{\text{ext}} = -Q\delta(\mathbf{r})(1,0,0,0)$ is the external current associated with the test charge placed at $\mathbf{r}=0$ and j^μ_{ind} is the induced (static) current in the plasma which is now in a steady motion with velocity $\mathbf{v}=(0,0,v)$. Note that in contrast with the usual Debye screening where only the scalar potential $\phi = A^0$ appears, here one must include the vector potential as well since, as we shall see in a moment, the test charge will induce not only a polarization charge density but also a (steady) polarization charge current in the moving plasma medium. The latter generates nonzero vector potentials and hence nonzero magnetic fields around the test charge.

In order to solve Eq. (2) we have to first obtain the induced current j^μ_{ind} as a function of the potential A^μ. In the kinetic theory the charge current in the plasma is given in terms of the one-body distribution function $f_i(p,x)$ as $j^\mu = \sum_i q_i \int d\Gamma\, p^\mu f_i(p,x)$, where we have used a compact notation $d\Gamma \equiv d^3p/(2\pi)^3E$ for the phase-space integral, and i denotes the particle species that carries the charge q_i. We assume that the plasma is in equilibrium and locally neutral before the test charge is inserted. The induced polarization current is thus given by

$$j^\mu_{\text{ind}} = \sum_i q_i \int d\Gamma\, p^\mu \delta f_i(p,x) , \tag{3}$$

where $\delta f_i = f_i - f_i^0$ is the change in the distribution function from the equilibrium distribution caused by the inserted test charge.

To determine the change of the distribution function caused by the test charge we use the collisionless Boltzmann-Vlasov equation

$$p^\mu \partial_\mu f_i - q_i F_{\mu\nu} p^\mu \partial f_i / \partial p_\nu = 0 , \tag{4}$$

where the second term represents the influence of the test charge, which generates nonzero self-consistent fields $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ in the plasma. Equation (4) can be solved by perturbation expansion in the potentials in the case of weak fields $q_i \beta A^\mu \ll 1$, where $\beta = 1/T$ is the inverse temperature of the plasma. To obtain the lowest-order solution we replace the distribution function on the right-hand side by the equilibrium distribution $f_i^0(\mathbf{p}) = (e^{\beta p^\mu u_\mu} \pm 1)^{-1}$ with $u^\mu = \gamma(1,\mathbf{v})$ being the plasma's four flow velocity, and use the following identities: $p^\mu \partial_\mu f_i^0 = 0$ and $\partial f_i^0/\partial p_\mu = -u^\mu \beta f_i^0(1 \pm f_i^0)$. Then the Boltzmann equation (3) gives, for the Fourier components of δf_i,

$$\delta \tilde{f}_i(p,k) = q_i \left[\tilde{A}\cdot u - \frac{(k\cdot u)(\tilde{A}\cdot p)}{k\cdot p} \right] \times \beta f_i^0 (1 \pm f_i^0) + O(\tilde{A}^2) . \tag{5}$$

The current induced due to this change in the distribution function can be obtained by inserting (5) into (3), and to lowest order in $\tilde{A}$ the result may be expressed as

$$\tilde{j}^\mu_{\text{ind}}(k) = \Pi^{\mu\nu}(k;T,u)\tilde{A}_\nu(k) , \tag{6}$$

where $\Pi^{\mu\nu}(k;T,u)$ is the polarization tensor for the plasma moving with arbitrary four-velocity u^μ. According to Lorentz covariance and current conservation, $k_\mu \Pi^{\mu\nu}(k) = k_\nu \Pi^{\mu\nu}(k) = 0$, the polarization tensor can be decomposed into the following manifestly covariant form:

$$\Pi^{\mu\nu}(k;T,u) = \Pi_1(k^2, k\cdot u)\left[g^{\mu\nu} - \frac{k^\mu k^\nu}{k^2} \right] + \Pi_2(k^2, k\cdot u)\left[u^\mu - \frac{(k\cdot u)k^\mu}{k^2} \right] \times \left[u^\nu - \frac{(k\cdot u)k^\nu}{k^2} \right] , \tag{7}$$

where the metric tensor is chosen here as $g^{\mu\nu} = \text{diag}(1,-1,-1,-1)$: Π_1 and Π_2 are scalar functions of two scalar variables k^2 and $k\cdot u \equiv k^\mu u_\mu$.

The calculation of the polarization tensor can be most easily performed in the plasma rest frame where $k^2 = \omega^2 - \mathbf{k}^2$ and $k\cdot u = \omega$, and one finds

$$\begin{aligned} \Pi_1(\omega^2 - \mathbf{k}^2, \omega) &= -m_e^2 \Phi_T(\omega/|\mathbf{k}|) , \\ \Pi_2(\omega^2 - \mathbf{k}^2, \omega) &= m_e^2 \left[\frac{\omega^2}{\mathbf{k}^2} - 1 \right] \times [\Phi_T(\omega/|\mathbf{k}|) - \Phi_L(\omega/|\mathbf{k}|)] , \end{aligned} \tag{8}$$

where

$$\begin{aligned} \Phi_T(x) &= \tfrac{1}{2}x^2 + \tfrac{1}{4}(1-x^2)\left[x \ln\left| \frac{1+x}{1-x} \right| - i\pi x \theta(1-x) \right] , \\ \Phi_L(x) &= (1-x^2)\left[1 - \frac{1}{2}\left[x \ln\left| \frac{1+x}{1-x} \right| - i\pi x \theta(1-x) \right] \right] . \end{aligned} \tag{9}$$

Here the electric screening mass m_e is determined by $m_e^2 = \sum_i q_i^2 c_i T^2$ with numerical constant $c_i = \frac{1}{12}$ $(\frac{1}{6})$ for fermions (bosons). For example, in the case of an ultrarelativistic electron-positron plasma which consists of electrons and positrons of electric charge $\pm e$ with spin up and down, $m_e^2 = \frac{1}{3}e^2 T^2$. This result was first derived by Silin.[20] (For the quark-gluon plasma, we can reproduce the electric screening mass obtained by the lowest-order perturbation calculation[21] $[m_e^2 = g^2 T^2(1+N_F/6)]$ by the following replacements: $\sum_{\text{quark}} q_i^2 \to 4N_F(g/2)^2 \text{Tr}(\lambda^a \lambda^b) = 2N_F g^2 \delta_{ab}$ for quarks (and antiquarks) and $\sum_{\text{gluon}} q_i^2 \to 2g^2 f_{abc} f_{a'bc} = 6g^2 \delta_{aa'}$ for gluons, where λ^a and f_{abc} are the generators and the structure constants of the color-SU(3) group, respectively, and N_F is the number of light flavors.)

The functional form of the covariant expression of $\Pi_1(k^2,k\cdot u)$ and $\Pi_2(k^2,k\cdot u)$ in (7) can be readily obtained from (8) by the following replacements: $\omega \to k\cdot u$ and $|k| \to [(k\cdot u)^2-k^2]^{1/2}$. This results in

$$\Pi_1(k^2,k\cdot u)=-m_e^2\Phi_T(k\cdot u/\sqrt{(k\cdot u)^2-k^2}) ,$$
$$\Pi_2(k^2,k\cdot u)=\frac{m_e^2k^2}{(k\cdot u)^2-k^2}[\Phi_T(k\cdot u/\sqrt{(k\cdot u)^2-k^2}) -\Phi_L(k\cdot u/\sqrt{(k\cdot u)^2-k^2})] . \tag{8'}$$

With the expression (6) for the induced current, the Fourier transform of Poisson's equation (2) reads

$$[\mathbf{k}^2g^{\mu\nu}-\Pi^{\mu\nu}(k)]\tilde{A}_\nu(k)=-Q\delta(\omega)(1,0,0,0) . \tag{10}$$

The solution of this equation is given by

$$\tilde{A}_\mu(k)=-\left[\frac{Q}{\mathbf{k}^2\epsilon(k)}\right]_{\mu 0}\delta(\omega) , \tag{11}$$

where we have introduced the dielectric tensor $\epsilon(k;u)$ in the plasma moving frame defined by $\epsilon_{\mu\nu}(k;u)=g_{\mu\nu}-(1/\mathbf{k}^2)\Pi_{\mu\nu}(k;u)$. The explicit form of the potentials can be calculated by setting $u^\mu=\gamma(1,0,0,v)$ and $k^\mu=(\omega,|\mathbf{k}|\cos\phi\sin\theta,|\mathbf{k}|\sin\phi\sin\theta, |\mathbf{k}|\cos\theta)$. We list below our final results for the Fourier components of the potential:

$$\tilde{A}^0(k)=2\pi Q\delta(\omega)\left[\frac{1-\gamma^2(1-z^2)}{\mathbf{k}^2\epsilon_T(\mathbf{k})}+\frac{\gamma^2(1-z^2)}{\mathbf{k}^2\epsilon_L(\mathbf{k})}\right] , \tag{12a}$$

$$\tilde{A}^1(k)=2\pi Q\delta(\omega)z\gamma\left[1-\frac{z^2\gamma^2}{\gamma^2-1}\right]^{1/2}\times\cos\phi\frac{1}{\mathbf{k}^2}\left[\frac{1}{\epsilon_T(\mathbf{k})}-\frac{1}{\epsilon_L(\mathbf{k})}\right] , \tag{12b}$$

$$\tilde{A}^2(k)=2\pi Q\delta(\omega)z\gamma\left[1-\frac{z^2\gamma^2}{\gamma^2-1}\right]^{1/2}\times\sin\phi\frac{1}{\mathbf{k}^2}\left[\frac{1}{\epsilon_T(\mathbf{k})}-\frac{1}{\epsilon_L(\mathbf{k})}\right] , \tag{12c}$$

$$\tilde{A}^3(k)=2\pi Q\delta(\omega)\gamma\frac{(1-z^2)\gamma^2-1}{\sqrt{\gamma^2-1}}\frac{1}{\mathbf{k}^2}\left[\frac{1}{\epsilon_T(\mathbf{k})}-\frac{1}{\epsilon_L(\mathbf{k})}\right], \tag{12d}$$

where

$$\epsilon_T(\mathbf{k})=1+\frac{m_e^2}{\mathbf{k}^2}\Phi_T(z) ,\qquad \epsilon_L(\mathbf{k})=1+\frac{m_e^2}{\mathbf{k}^2}\Phi_L(z) \tag{13}$$

are the transverse and longitudinal dielectric "constants," respectively, which satisfy $\det(\epsilon)=-\epsilon_T^3\epsilon_L$. The quantity z which has appeared in (12) and (13) is defined by

$$z=\left[\frac{k\cdot u}{\sqrt{(k\cdot u)^2-k^2}}\right]_{\omega=0}=\frac{v\cos\theta}{\sqrt{1-v^2\sin^2\theta}} .$$

It is easy to check that in the limit $v \to 0$, the potential (11) reduces to the Fourier transform of the ordinary isotropic screened potential (1):

$$\lim_{v\to 0}\tilde{A}^\mu(k)=2\pi Q\delta(\omega)\delta_{\mu 0}\frac{1}{\mathbf{k}^2+m_e^2} .$$

At finite v, the potential (12) not only becomes anisotropic but also acquires nonzero magnetic components.

III. DISCUSSION

We now examine our result and discuss its possible implications on the fate of a $c\bar{c}$ pair in ultrarelativistic nucleus-nucleus collisions.

We present in Fig. 1 the modified Coulomb potential in the configuration space calculated from (12a). In a cylindrical coordinate system the potential takes a form of

$$A^0(r)=\frac{Qm_e}{4\pi}F(m_e\rho,m_ez) , \tag{14}$$

where $F(x,y)$ is a dimensionless function. In the Fourier transformation of (12a) both the real and imaginary parts of the momentum-space potential $\tilde{A}^0(k)$ contribute to the real part of $A^0(\rho,z)$ but no imaginary part appears in the configuration-space potential. The equipotential surfaces are plotted in Fig. 1 at several different values of v. Note that Fig. 1(a) for $v=0$ corresponds to the usual static screened potential. At nonzero v, the potential becomes anisotropic and loses forward-backward symmetry with respect to the direction of the plasma motion (z direction). The potential slope is steeper downstream than upstream. This implies that a particle with opposite charge to the test charge would feel stronger attraction downstream.

The origin of the forward-backward asymmetry of the potential can be traced back to the imaginary part of the dielectric constant, which physically arises due to the dissipation of the field energy into the excitation of the plasma medium. It is noteworthy here that no singular behavior appears in the potential at any value of v. This is ensured by the absence of the collective branch in the dispersion relation $\det\epsilon(k)=0$ in the spacelike momentum region ($k^2<0$). This is the property characteristic of massless relativistic plasmas that the dispersion relation of all plasmon modes is always timelike, and hence the plasmon cannot be excited by the test charge passing through the plasma medium as in a nonrelativistic plasma.

We must note that (14) is *not* the two-body potential between $Q\bar{Q}$ pair moving in the plasma. The effective $Q\bar{Q}$ potential due to the modified Coulomb interaction (14) can be obtained by calculating the static Coulomb energy[19] $\frac{1}{2}\int d^3r' j_\mu^{\text{ext}}(\mathbf{r}')A_{\text{dip}}^\mu(\mathbf{r}')$ associated with a dipole external current $j_\mu^{\rm ext}$ which is the sum of two currents created by the Q and $\bar{Q}$ separated by the distance $\mathbf{r}$. Here $A_{\rm dip}^{\mu}(\mathbf{r}')$ is the dipole field created by $j_\mu^{\rm ext}$ which in the linear approximation as we took to derive the screened potential (14) is just the sum of two screened potentials created by Q and $\bar{Q}$ independently in the absence of other charge. Upon the subtraction of the divergent self-energy terms, this results in

$$V_e(\mathbf{r})=\tfrac{1}{2}[QA_{\bar{Q}}^{0}(-\mathbf{r})+\bar{Q}A_{Q}^{0}(\mathbf{r})]$$
$$=\frac{4\alpha_s}{3}\frac{m_e}{2}[F(\rho m_e, zm_e)+F(\rho m_e,-zm_e)]\ . \qquad (15)$$

In the last line we have replaced $Q\bar{Q}/(4\pi)$ by the color-singlet strength of the one-gluon-exchange force $4\alpha_s/3$. This potential is invariant by $\mathbf{r}=\mathbf{r}_Q-\mathbf{r}_{\bar{Q}}\to-\mathbf{r}$, as it should be.

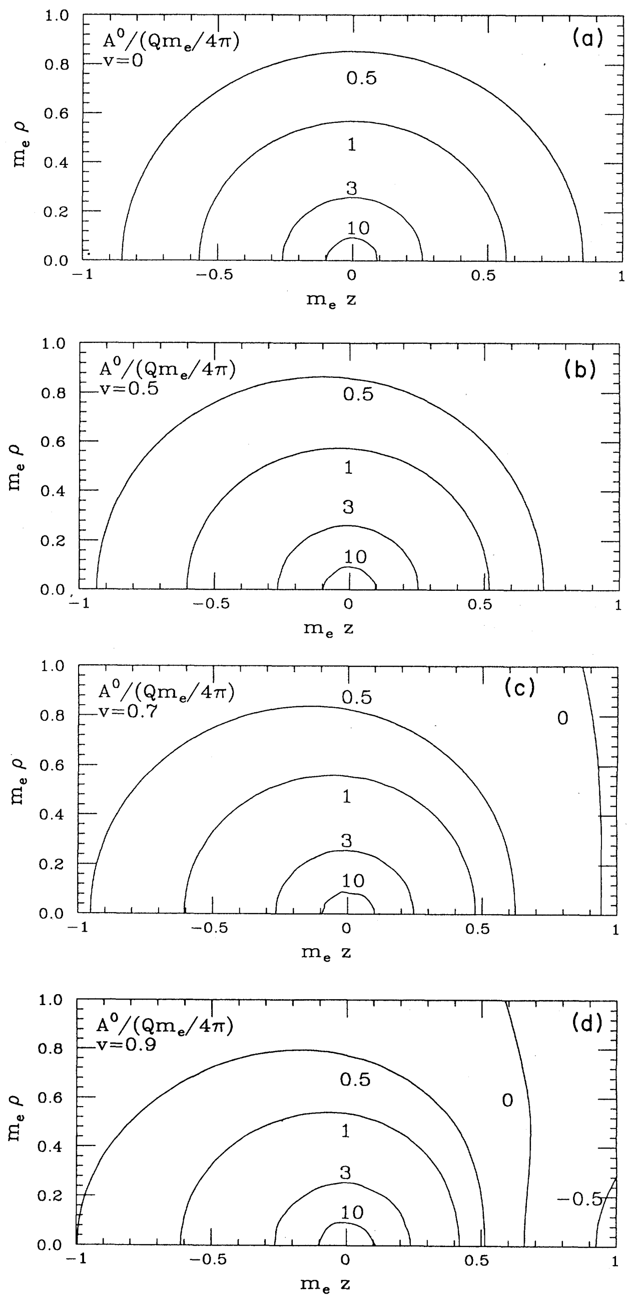

FIG. 1. Electric potential around a test charge (Q) at the origin, with the plasma flowing in the $+z$ direction with velocity v. Four different cases are shown here: $v=0,0.5,0.7,0.9$. The z and ρ axes are scaled by the screening mass m_e.

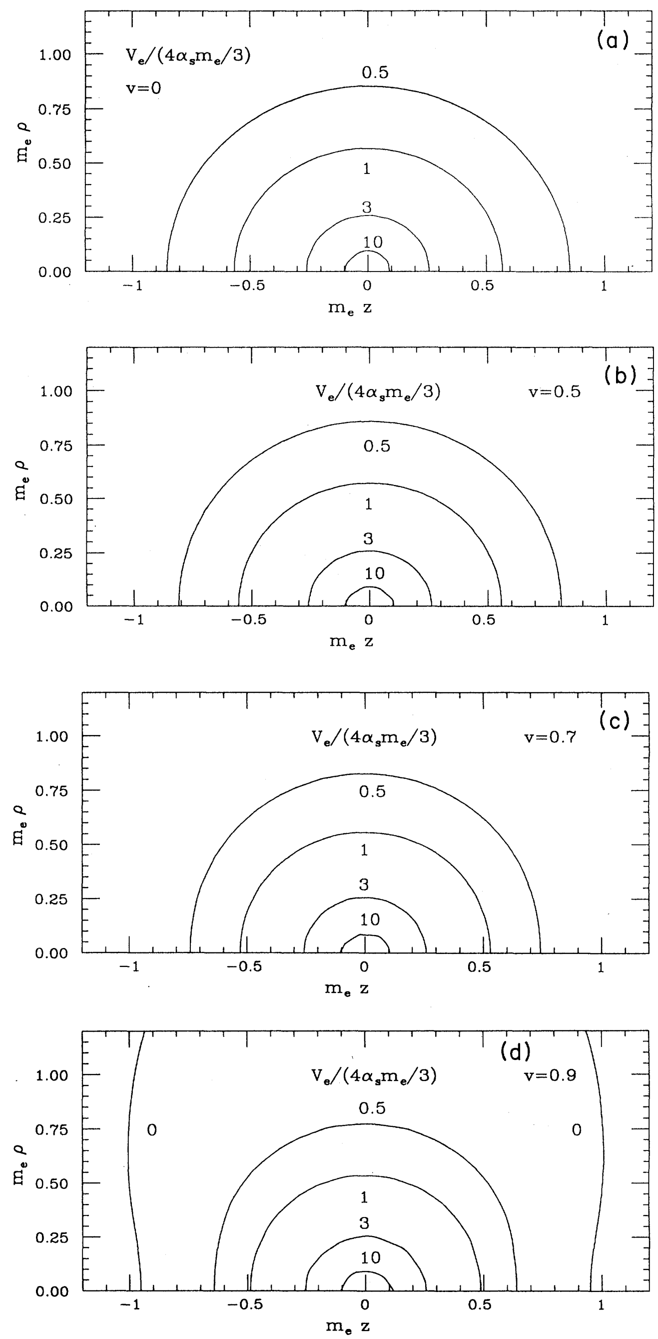

FIG. 2. Two-body potential between a pair of test charges calculated by estimating static modified Coulomb energy of the pair in the plasma which is flowing in the $+z$ direction with four different velocities: $v=0,0.5,0.7,0.9$.

We plot in Fig. 2 the equipotential surfaces of the effective two-body potential given by (15) at several different values of v. The shape of the two-body potential is identical to that of the screened Coulomb potential in the case of $v=0$ since $V_e(\mathbf{r})=QA^0(\mathbf{r})$, but they differ at nonzero v due to the averaging procedure in (15). It is seen that the strong anisotropy observed in the screened Coulomb potential $A^0(r)$ is significantly reduced in the two-body potential. We note however that the potential strength is further weakened when the plasma is moving and the equipotential surfaces are squashed in the longitudinal direction.

One immediate consequence of this modification of two-body potential would be the modification of the p_T dependence of the J/ψ suppression. Since the $Q\bar{Q}$ potential is more effectively screened when a pair is moving in the plasma, it requires less energy density to suppress the J/ψ formation at large p_T. *We hence expect that the transverse-momentum dependence of the J/ψ suppression would be slightly "flattened" compared with the previous results obtained without taking into account this effect.*[3–5]

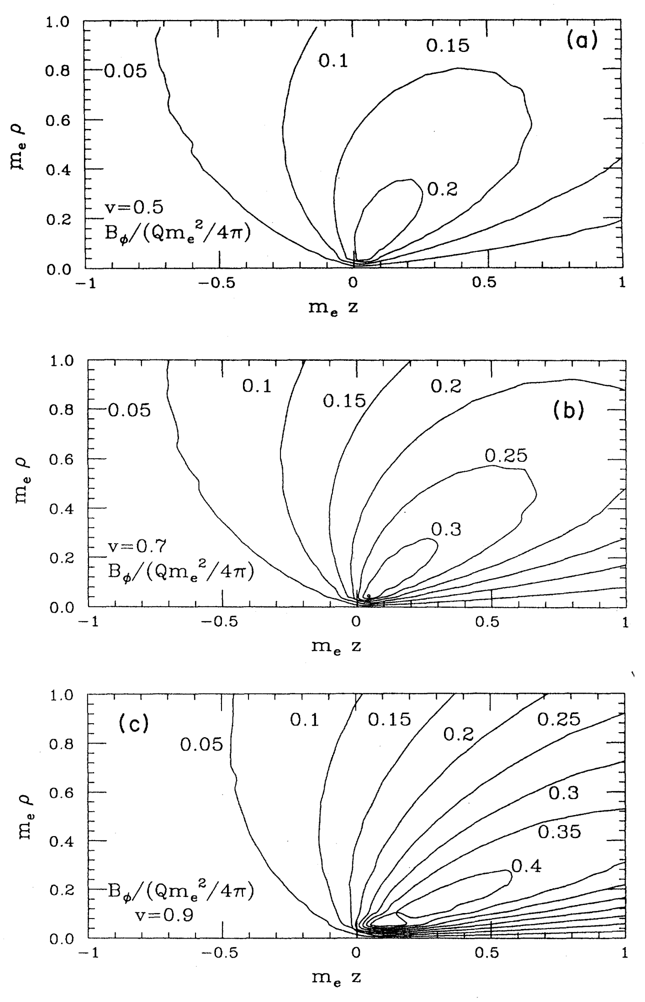

FIG. 3. Contour plots of the strength of the azimuthal magnetic fields B_ϕ around a test charge at the origin. A plasma is flowing in the $+z$ direction with velocities $v=0.5, 0.7, 0.9$.

At present experimental conditions, the relevant value of the transverse velocity v of a $c\bar{c}$ pair which can barely escape from the high-density plasma region is $v=p_T/E_T \simeq 0.7$. Here we have used $p_T=3$ GeV for the transverse momentum of the pair which is observed not to suffer significant suppression.[2] At such values of v, we see only little modification in the two-body potential from the static Debye screened potential. The effect of the dynamic screening is enhanced at higher values of v. Since the escape velocity of the charmonium becomes larger as the lifetime of the plasma increases,[3–5] we expect that the effect of the dynamic screening becomes more important when a long-lived plasma is formed in the future experimental conditions.

We finally examine the effect of the induced color-magnetic field created by nonzero vector potential (12b)–(12d):

$$\mathbf{B}(\mathbf{r})=\frac{Q}{4\pi}m_e^2 G(m_e\rho, m_e z; v)\mathbf{e}_\varphi , \tag{16}$$

where $\mathbf{e}_\varphi=(\sin\varphi, \cos\varphi)$ is a unit vector perpendicular both to the direction of the plasma motion $\mathbf{v}$ and to the transverse position vector $\boldsymbol{\rho}$. The equistrength surfaces of the magnetic field are shown in Fig. 3. The magnetic field vanishes on the z axis and its strength becomes maximum at the distance of $\rho \simeq 1/m_e$. Again we see that the magnetic field is not symmetric by the space inversion $z \to -z$ reflecting the plasma flow pattern.

This magnetic field will couple to the intrinsic magnetic moment of heavy quark and antiquark and generate a spin-dependent effective two-body potential given by

$$\begin{aligned} V_m(\mathbf{r}) &= \mu_Q \mathbf{B}_{\bar{Q}}(\mathbf{r}) - \mu_{\bar{Q}} \mathbf{B}_Q(-\mathbf{r}) \\ &= -\frac{4\alpha_s}{3}\frac{m_e^2}{2m_Q}[\sigma_Q \mathbf{e}_\varphi G(\rho m_e, z m_e) \\ &\qquad - \sigma_{\bar{Q}} \mathbf{e}_\varphi G(\rho m_e, -z m_e)] , \end{aligned} \tag{17}$$

where $\mu_Q = Q\sigma_Q/(2m_Q)$ are the color-magnetic moments of the heavy quark. The potential (17) is not the usual type of the two-body interaction one finds in atomic physics. At the first glance, it may look like a spin-orbit interaction, which arises in the case of atomic physics due to the coupling of the electron magnetic moment to the apparent magnetic field seen by the electron moving in the electrostatic Coulomb field of the nucleus. In the present case, the interaction is caused by the relative motion of $Q\bar{Q}$ pair and the plasma medium instead of the relative motion of $Q\bar{Q}$. The magnitude of the potential (17) is small compared with the modified Coulomb interaction by the factor m_e/m_Q but appears to be greater than the usual spin-orbit interaction which is proportional to $1/m_Q^2$. However, the expectation value of (17) vanishes for all wave functions with definite orbital angular momentum due to the vanishing azimuthal angle integral, and hence the effect would arise only in higher-order corrections due to the coupling of different orbital angular momentum states. We expect hence that the effect of the magnetic interaction is very small.

In conclusion, we have studied the dynamic Debye screening for a test charge traversing a collisionless relativistic plasma of massless particles which are interacting via an Abelian gauge field, and we found that the usual Debye screened Coulomb potential is significantly modified when the plasma is moving with respect to the test charge. We have indicated a few possible interesting consequences of these effects on the fate of a heavy-quark–antiquark pair created in ultrarelativistic nucleus-nucleus collision. Further work is needed to draw more quantitative conclusions on the significance of these effects. Especially, we note that the perturbative calculation of the gluon self-energy at finite temperature shows nonanalytic peculiar behavior in the long-wavelength limit[21,22] which did not show up in the present treatment based on the semiclassical transport theory. It would be very interesting to examine the consequence of such an effect in the dynamical screening considered here. It is also interesting to see how our results are modified by inclusion of the dynamical evolution of the plasma medium,[23] especially the effect of the collisions in plasma on the screening.

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