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Velocity dependence of charmonium dissociation temperature in high-energy nuclear collisions

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ABSTRACT

In high-energy nuclear collisions, heavy quark potential at finite temperature controls the quarkonium suppression. Including the relaxation of the medium induced by the relative velocity between quarkonia and the deconfined expanding matter, the Debye screening is reduced and the quarkonium dissociation takes place at a higher temperature. As a consequence of the velocity-dependent dissociation temperature, the quarkonium suppression at high transverse momentum is significantly weakened in high-energy nuclear collisions at RHIC and LHC.

Heavy quarkonia J/ψ and Υ are tightly bound hadronic states. Their dissociation temperature T_d is, in general, higher than the critical temperature T_c for the deconfinement phase transition [1] in high-energy nuclear collisions [2–4]. Therefore, the measured cross sections of quarkonia carry the information of the early stage hot and dense medium. They have long been considered as a signature of the formation of the new state of matter, the so-called quark–gluon plasma [5,6].

The quarkonium dissociation in a static deconfined quark matter is generally described in terms of the screening effect. The heavy quark potential, which is normally taken as the Cornell form [7] and can be calculated through a non-relativistic quantum chromodynamic potential [8] and lattice simulations [9], is reduced to a Yukawa-like potential due to the Debye screening. When the screening radius becomes smaller than the quarkonium size, the bound state dissociates. Substituting the screened potential, extracted from lattice simulations [10,11], into the Schrödinger equation for the wave function of the quarkonium state, one obtains the dissociation temperature that corresponds to the zero binding energy and infinite size of the di-quark system [12,13]. For charmonia, while the excited states χ_c and ψ' start to dissociate already around T_c, the calculated dissociation temperature for the ground state J/ψ is much higher than the critical temperature [12].

The quarkonia produced in relativistic heavy ion collisions are, however, not at rest in the medium. There exists a relative velocity between the quarkonia and the expanding medium. The question is what is the velocity dependence of the heavy quark potential at finite temperature [14–16]. The screening effect is due to the rearrangement of the charged particles when a pair of heavy quarks (source) is present in the medium. For a moving source, it will take a longer time for the source to interact with the medium, comparing with that of a stationary source. This 'delay' of the response reduces the screening charges around the source and thus weakens the screening effect. In relativistic heavy ion collisions, the average transverse momentum of the initially produced J/ψs is about 2 GeV at RHIC energy [17] and 3 GeV at LHC energy. [18], corresponding to an averaged relative velocity above 0.5c. A significant modification of the Debye screening is expected for such fast moving J/ψs, especially for those produced in higher transverse momentum region [19,20]. In this Letter, we study the velocity dependence of the heavy quark potential and the quarkonium dissociation temperature in a transport approach. The velocity induced effects on charmonium suppression at both Relativistic Heavy Ion Collider (RHIC) and Large Hadron Collider (LHC) will be discussed. In the following calculations we take the speed of light $c=1$.

For a static source located at $\mathbf{r}=0$, the ambient charge density $\rho_0(\mathbf{r})$ is modified by the screening potential $V_0(\mathbf{r})$ at finite temperature T [21],

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$$\rho_0(\mathbf{r}) = \sum_i q_i f_i e^{-q_i V_0(\mathbf{r})/T} \approx -\sum_i \left(q_i^2 f_i/T\right) V_0(\mathbf{r}), \tag{1}$$

where q_i is the charge of the particles of species i, and f_i is the initial particle density without the source. The neutrality condition for the total charge has been considered here. The solution of (1) with the assumption of small V gives the Debye screening of the potential. At large distance, the potential is weak, so that the approximation in (1) is appropriate, while at small distance, the solution of (1) means a small correction to the original potential, as the lattice simulations indicated [10,11].

For a source moving with velocity $\boldsymbol{\upsilon}$ with respect to the medium, the non-equilibrium charge density $\rho(\mathbf{r}, t)$ in the source-rest frame satisfies the transport equation in the relaxation time approximation,

$$\partial_t \rho - \boldsymbol{\upsilon} \cdot \nabla \rho = -(\rho - \rho_0)/\tau, \tag{2}$$

where τ is the relaxation time of the medium. Taking the limit $t \to \infty$, the final distribution $\rho_f(\mathbf{r}, \mathbf{L}) \equiv \lim_{t\to\infty} \rho(\mathbf{r}, t)$ becomes stable and is characterized by the equation

$$\mathbf{L} \cdot \nabla \rho_f = \rho_f - \rho_0, \tag{3}$$

where we have introduced the relaxation length defined as $\mathbf{L} \equiv \boldsymbol{\upsilon}\tau$ which controls the velocity dependence of the Debye screening. Since the screening charge distribution is proportional to the screening potential, see Eq. (1), the potentials V_0 and V corresponding to a stationary and moving source, respectively, satisfy the same equation (3) which can be solved analytically,

$$V(\mathbf{r}, \mathbf{L}) = \int_0^\infty V_0(\mathbf{r} + \lambda \mathbf{L}) e^{-\lambda} d\lambda. \tag{4}$$

It is obvious that for a static source with $L = 0$ we have $V(\mathbf{r}, \mathbf{0}) = V_0(\mathbf{r})$.

With the potentials V_0 and V, the screening radius r_d can be expressed as

$$r_d(\mathbf{L}) = \frac{1}{2} \frac{\int d^3\mathbf{r}\, r \rho_f(\mathbf{r}, \mathbf{L})}{\int d^3\mathbf{r}\, \rho_f(\mathbf{r}, \mathbf{L})} = \frac{1}{2} \frac{\int d^3\mathbf{r}\, r V(\mathbf{r}, \mathbf{L})}{\int d^3\mathbf{r}\, V_0(\mathbf{r})}. \tag{5}$$

For the second equality, we have used the total charge conservation $\int d^3\mathbf{r}\, \rho_f(\mathbf{r}, \mathbf{L}) = \int d^3\mathbf{r}\, \rho_0(\mathbf{r})$ for any $\mathbf{L}$ which is guaranteed by integrating Eq. (3) over the whole coordinate space. The screening radius $r_d(\mathbf{L})$ is in general an angle-dependent function. However, for a spherically symmetric potential $V_0(r)$, the integration over the angles in the numerator of Eq. (5) can be analytically done, and the averaged screening radius can be effectively expressed as $r_d(L) = \int dr\, r^3 V(r, L)/(2 \int dr r^2 V_0(r))$ with the factorized averaged potential

$$V(r, L) = V_0(r) W(r/L), \tag{6}$$

where the modification factor W is defined as

$$W(y) = 1 + \frac{2 + y^3 Z(y) - (y^2 - y + 2)e^{-y}}{3y^2},$$
$$Z(y) = \int_y^\infty \frac{dt}{t} e^{-t}. \tag{7}$$

We emphasize that the general potential $V(\mathbf{r}, \mathbf{L})$ can be simplified as $V(r, L)$ only in the sense of the screening radius (5). For a more detailed calculation about a general potential, one may refer to the Refs. [20,22] based on the linear response theory.

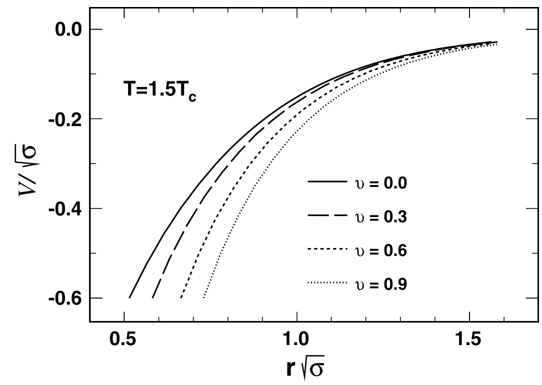

Fig. 1. The velocity dependence of the screening potential at a fixed temperature $T = 1.5T_c$. σ is the string tension and the stationary potential is taken as the free energy [10–12].

Now we apply the above transport solutions to the quarkonium dissociation in hot and dense matter created in high-energy nuclear collisions. The interaction between two quarks in vacuum can be well characterized by the Cornell potential [12] $V_0(r) = -\alpha/r + \sigma r$ with coupling constant $\alpha = \pi/12$ and string tension $\sigma = 0.2\ \text{GeV}^2$. At finite temperature, the screening potential for a stationary pair of heavy quarks can be written as [23,12]

$$V_0(r) = -\frac{\alpha}{r} e^{-\mu r} - \frac{\sigma}{2^{\frac{3}{4}} \Gamma(\frac{3}{4})} \left(\frac{r}{\mu}\right)^{1/2} K_{\frac{1}{4}}\left((\mu r)^2\right), \tag{8}$$

where Γ and K are the Gamma and modified Bessel functions. The temperature of the medium is hidden in the screening mass $\mu(T)$ which can be extracted [12] from lattice QCD calculated free energy [10,11].

To establish a unique mapping between the relaxation length and the velocity, we estimate the relaxation time of the hot and dense matter by considering its electric analogue. When an electric charge is put into a conducting medium, the medium is neutralized in a time scale of $\tau = 1/(4\pi\sigma_e\alpha_e)$, where σ_e is the electric conductivity of the medium, and α_e is the fine-structure constant. We replace σ_e by the conductivity $\sigma_s \approx 0.4T$ for a strong field, estimated from hot quenched lattice QCD [24], and α_e by α, the relaxation length becomes $L = 15\upsilon/(2\pi^2 T)$.

In Fig. 1 one sees the velocity induced change in the heavy quark potential at a fixed temperature $T = 1.5T_c$. The stationary potential is taken from the lattice simulation Eq. (8). Since the screening length is proportional to the velocity and inversely proportional to the temperature of the medium, the potential well becomes deeper and screening becomes less effective, when the quarkonium velocity relative to the medium increases.

With the known potentials $V_0(T, r)$ and $V(T, r, L)$, the screen radius $r_d(T, L)$ at finite temperature T can be calculated through (5), where the temperature T-dependent inherited from potential V is written explicitly. In the rest frame of the di-quark system, the condition for dissociating a quarkonium should not depend on its relative velocity, namely the critical screening radius is a constant,

$$r_d\big(T_d, L(\upsilon, T_d)\big) = C, \tag{9}$$

where the constant C can be calculated directly at $\upsilon = 0$,

$$C = r_d\big(T_d(\upsilon = 0), L = 0\big) = \frac{1}{\mu} \frac{1 + \frac{\pi}{16\Gamma^2(3/4)} \frac{\sigma}{\alpha\mu^2}}{1 + \frac{1}{4} \frac{\sigma}{\alpha\mu^2}} \tag{10}$$

with μ the screening mass at $T_d(\upsilon = 0)$. Thus when the dissociation temperature of a J/ψ at rest is given, the constant C