Quarkonium Dissociation at Finite Chemical Potential

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Abstract

We have studied the dissociation of quarkonia states in a deconfined medium of quarks and gluons at large baryon chemical potential and small temperature region. The aim of this study is to probe the dense baryonic medium expected to be produced at FAIR facility, GSI Darmstadt. This is done by correcting both the short and long-distance terms of the Cornell potential by a dielectric function, embodying the effects of deconfined quarks and gluons, at finite baryon chemical potential and temperature. It is found that \( J/\psi \) is dissociated approximately at 2 \( \mu_c \) in the temperature range 20-50 MeV, which can indirectly help to locate the point on QCD phase diagram at large chemical potential and low temperature zone.

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1 Introduction

The relativistic heavy ion collisions endows with the unique opportunity of creating a hot and dense nuclear matter in the laboratory. According to statistical quantum chromodynamics (QCD), a nuclear matter may undergo a color deconfined partonic phase, the quark gluon plasma (QGP) at sufficiently high temperature and/or density. Over the last few decades, strenuous efforts have been invested to devise clean and experimentally viable signals that can unambiguously identify

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the existence of QCD phase transition and trace out its signatures. Charmonium (a bound state of charm and anti-charm quark) suppression had been predicted as a signature for the deconfinement transition since long back[1], where the Debye screening of the colored partons was the dominant mechanism of suppression. Experimental investigations revealed a noticeable reduction of the charmonium production in proton-nucleus (p-A) collisions compared to the hadronic collisions (scaled with binary nucleon-nucleon (p-p) collisions)[2, 3, 4, 5]. Thorough understanding of this normal nuclear suppression is essential in order to establish a robust baseline reference, with respect to the anomalous suppression [6, 7], pertaining to the formation of a deconfined partonic medium observed at RHIC and LHC energies that mainly scan the high temperature and almost vanishing baryon density region of QCD phase diagram.

Depending upon the magnitude of energy deposited in center of mass frame of colliding nuclei there are in general two scenario in heavy ion collision. The first one is possibility of complete stopping of colliding nuclei if energy density, $\epsilon \approx 5 \text{ GeV/fm}^3$ creating sufficiently high baryon density (parton density) to produce QGP. Full transparency of colling nuclei is other scenario takes place if $\epsilon \sim 100 \text{ GeV/fm}^3$ creating favorable environment to produce QGP by vacuum polarization, leading to vanishing chemical potential. Exploration of high baryon densities and moderate temperature region of QCD phase diagram is possible with the upcoming compressed baryonic matter (CBM) experiment at the Facility for Anti-proton and Ion Research (FAIR). At FAIR, the light and heavy-ions can be collided in the beam energy range $E_{\text{Lab}} = 10-40 \text{ A GeV}$ to create an environment favorable to investigate extremely dense nuclear matter in the laboratory through the measurements of bulk as well as rare probes. Model calculations based on transport as well as hydrodynamical equations [8] predict that the highest net baryon densities ($\rho_B$) produced in the center of collision is $\sim 6$ to 12 times the density of normal nuclear matter for the most central collision (b=0).

The FAIR will open avenues to have deeper insight on some of the fundamental but yet enigmatic issues of strong interaction thermodynamics at large baryon density. Some of the key issues in this region are the study of hadronic properties in dense nuclear matter, deconfinement phase transition from hadronic to quark-gluon matter driven by high baryon densities, the nuclear equation of state (EOS) at high baryon densities etc. Thus the behavior of QCD at high baryon density and moderate/low temperature is interesting and has potential applications to cosmology, to the astrophysics of neutron stars, and to heavy ion collisions. Although the motive of FAIR is to produce the baryon rich QGP but it is uncertain that at what point on the phase diagram on the
Since the relative velocity of heavy quark bound state (\(v_{Q\bar{Q}} \ll 1\)) is very small, so in-medium dynamics of \(Q\bar{Q}\) bound states has been extensively studied with phenomenological potential models [9], where the temperature-dependent potential carries all medium effects with nonperturbative terms borrowed from lattice simulations. Although the derivation of such models from QCD is not established, free energies and other quantities [10, 11] from lattice calculations obtained using the correlation functions of Polyakov loop are often taken as input for the potential. These quantities have been thought to be related to the color-singlet and color-octet heavy quark potentials at finite temperature [11, 12] however, a precise answer is still missing [13]. It was until recently that the smallness of the velocity (\(v_{Q\bar{Q}} \sim \alpha_s < 1\)) has opened up the possibility of studying heavy quark bound state at finite temperature, where both the non-relativistic scales, \(v_{Q\bar{Q}}\), the heavy quark mass (\(m_Q\)), the momentum exchange (\(m_Qv_{Q}\)), the binding energy (\(m_Qv_{Q}^2\)) etc. and the thermal scales, \(T\), \(gT\), \(g^2T\) are hierarchically ordered: By exploiting the aforesaid separation of scales, a sequence of low-energy effective field theories (EFTs) [14, 15] viz. non-relativistic QCD (NRQCD), potential non-relativistic QCD (pNRQCD) etc. have been synthesized by integrating out the successive energy scales. In this context, pNRQCD and its thermal version are of particular interest because it describes the quarkonium dynamics through potentials and low-energy interactions [16, 17]. The color-screening phenomenon [16] and thermal width are thought to be an outcome of thermal corrections to the real and imaginary parts of the color-singlet potential, respectively. Out of various mechanisms that might be contributing to the imaginary part of the potential, Landau-damping [18] is the dominant one, which is responsible for quarkonium dissociation at weak coupling. It makes the quarkonium to be dissociated even at temperatures, where the possibility of color screening is negligible. Moreover, lattice studies have also shown that even at strong coupling [19, 20, 21] the potential may have a sizable imaginary part and such contributions are related to quarkonium decay processes in the plasma.

The works referred above are restricted to finite temperature only. However, the color screening effect at finite temperature and finite density has been studied in thermo-field dynamics approach to calculate the Debye screening mass, \(m_D(T, \mu)\), where authors in Refs. [22] and [23] have used the phenomenological potential model [9] and error function-type confined force with color screened Coulomb-type potential, respectively. Applying the effective perturbation theory for gauge theories at finite temperature [24], authors in Ref.[25] extended it to finite chemical potential for studying...
the collisional energy-loss of heavy quark in QGP. Photon production at finite chemical potential in QCD plasma at leading-order in strong coupling has been computed in [26]. Furthermore the dissipative hydrodynamic effects on QGP at finite density has been studied by developing a causal dissipative hydrodynamic model at finite baryon density for RHIC and LHC energies to study the net-baryon rapidity distribution [27]. Although, finite density lattice QCD calculations are seriously affected by the sign problem, recently color screening in heavy quark potential at finite density with Wilson fermions has been studied in lattice QCD in both real and imaginary chemical potential regions [28].

In the present work we have studied quarkonium dissociation in the quark matter of high baryon density and low temperature. For that we have first obtained the heavy quark potential by correcting both perturbative and non-perturbative terms in the Cornell potential through the dielectric function at finite temperature and chemical potential and then calculate the binding energy of the quarkonium states. Rest of the work is organized as: We first discuss the heavy quark potential at finite chemical potential in section 2. Then focus on the estimation of binding energy of quarkonium at finite chemical potential and thereby obtained the dissociation chemical potential for quarkonium ground states in section 3. Finally we conclude in section 4.

2 Heavy quark potential at finite chemical potential and running coupling constant

The interaction potential between $Q$ and $\bar{Q}$ is subjected to modify when the pair $(Q\bar{Q})$ is placed in a hot and dense QCD medium and it plays an important role in understanding the status of $Q\bar{Q}$ bound states in such medium. This issue has been well taken up by various authors and are reported in several reviews [29, 30], by dealing on both fronts, the phenomenology [9] as well as the lattice QCD [19, 20, 21, 31]. All these studies assume that the phase transition from a hadronic matter to a QGP phase melts the string thereby the string tension vanishes at the transition point. However lattice results [32] indicate that there is no genuine phase transition at vanishing baryon density, it is rather a crossover, and there may not be immediate melting of string at the deconfinement temperature. Thus one should consider its effect on the potential even above the transition temperature, in addition to the coulomb term.

The large mass of heavy quark meets the requirements: (i) $m_Q \gg \Lambda_{QCD}$ and (ii) $T \ll m_Q$
for the description of the interactions between $Q\bar{Q}$ at finite temperature and density. It is thus possible to obtain the quantum mechanical potential, $V(r,T)$ by correcting short as well as long-distance parts of the $Q\bar{Q}$ potential through dielectric permittivity, $\epsilon(k_0)$, embodying the effects of deconfined medium as:

$$V(r,T,\mu) = \int \frac{d^3k}{(2\pi)^{3/2}} \left( e^{ikr} - 1 \right) \frac{V(k)}{\epsilon(k)}.$$  \hspace{1cm} (1)

The term $V(k)$ in above equation is the Fourier Transform (FT) of the potential. However, obtaining the FT of the linear term is a tricky part and needs regulation, which is carried out by multiplying by an exponential damping factor and is switched off after the evaluation of FT. We have taken the same screening scale for both the terms $^3$ and the FT of the Cornell potential is obtained as [33]

$$V(k) = -\sqrt{\frac{2}{\pi}} \frac{\alpha + 2\sigma}{k^2},$$  \hspace{1cm} (2)

where $\alpha$ and $\sigma$ are the coupling and string tension, respectively. We have taken $\alpha$ as function of chemical potential ($\mu_q$) and temperature up to two loop [34] as

$$\alpha(T,\mu_q) = \frac{6\pi}{(33 - 2N_f) \ln \sqrt{\frac{T^2}{\Lambda_T^2} + \frac{\mu^2}{\pi^2T^2}}} \left[ 1 - \frac{3(153 - 19N_f) \ln \left( 2 \ln \sqrt{\frac{T^2}{\Lambda_T^2} + \frac{\mu^2}{\pi^2T^2}} \right)}{(33 - 2N_f)^2 \ln \sqrt{\frac{T^2}{\Lambda_T^2} + \frac{\mu^2}{\pi^2T^2}}} \right],$$  \hspace{1cm} (3)

where $\Lambda_T$ is the QCD scale-fixing parameter which characterizes the strength of the interaction.

To obtain the dielectric permittivity, $\epsilon(k)$, one has to evaluate the self-energies and the corresponding static propagators in weak coupling Hard Thermal Loop (HTL) approximation. In real-time formalism using Keldysh representation, the retarded (R), advanced (A) and symmetric (F) propagators are written as the linear combination of the components of the $(2 \times 2)$ matrix propagator:

$$D^0_R = D^0_{11} - D^0_{12}, \hspace{0.5cm} D^0_A = D^0_{11} - D^0_{21}, \hspace{0.5cm} D^0_F = D^0_{11} + D^0_{22},$$  \hspace{1cm} (4)

where only the symmetric component involves the distribution functions and is of particular advantage for the HTL diagrams where the terms containing distribution functions dominate. Similar relations hold good for the retarded ($\Pi^0_R$), advanced ($\Pi^0_A$) and symmetric ($\Pi^0_F$) self energies. The

$^3$However, in the framework of classical Debye-H"uckel theory, different screening functions were employed, viz. $f_c$ and $f_s$ for the Coulomb and string terms, respectively to obtain the free energy [35] whereas in Ref. [36], different scales for the Coulomb and linear pieces have been employed through a dimension-two gluon condensate.
The resummation of the propagators can be obtained using the Dyson-Schwinger equation

\[
D_{R, A} = D^0_{R, A} + D^0_{R, A} \Pi_{R, A} D_{R, A},
\]
\[
D_{F} = D^0_{F} + D^0_{R} \Pi_{R} D_{F} + D^0_{F} \Pi_{A} D_{A} + D^0_{R} \Pi_{F} D_{A}.
\]

For the static potential, we need only the temporal component (“00” ≡ L) of the propagator, whose evaluation is easier in the Coulomb gauge. Thus the above resummation (5) can be recast through its temporal component as

\[
D^L_{R, A} = D^{L(0)}_{R, A} + D^{L(0)}_{R, A} \Pi^{L}_{R, A} D^L_{R, A}.
\]

The leading contribution to temporal component of the retarded/advanced and symmetric gluon self-energy in the HTL-approximation can be written as

\[
\Pi^L_{R, A}(k) = m^2_D \left( \frac{k_0}{2p} \ln \frac{k_0 + k \pm i\epsilon}{k_0 - k \pm i\epsilon} - 1 \right)
\]

and

\[
\Pi^L_{F}(k) = -2\pi \imath m^2_D \frac{T}{k} \Theta(k^2 - k_0^2),
\]

respectively. Thus the gluon self-energy is composed of real and imaginary parts which are responsible for the Debye screening and the Landau damping, respectively. The real part of potential can be obtained from the retarded (or advanced) self energy with the prescriptions +i\epsilon (−i\epsilon), respectively.

The term \( m_D \) in above equation is known as electric screening mass (Debye screening mass) which encodes the medium effects in terms of temperature and baryon chemical potential. The properties of QCD at finite temperature and density are studied in [37] and the electric screening mass can be obtained in the static infrared limit of “00” component of gluon self energy \( \Pi^{00}(q_0 = 0, \vec{q} \to 0) \). Taking the dynamical quark mass \( m_Q = 0 \), the electric screening mass at finite temperature and vanishing baryon chemical potential \( (\mu = \mu_b = 0) \) based on perturbation theory in high temperature limit is \[38\],

\[
\Pi^{00}(q_0 = 0, \vec{q} \to 0) = \left( \frac{N_c}{3} + \frac{N_f}{6} \right) gT,
\]

where \( N_c \) and \( N_f \) are number of colors and flavors of color group \( \text{SU}(N_c) \) respectively. For a hot and dense plasma, the resummation technique to the perturbation theory proposed by [24] works
well in momentum scale $gT$. The leading-order Debye screening mass from HTL \cite{25, 39, 40, 41} at finite temperature and chemical potential is given by

$$m^2_D(T, \mu_q) = g^2(T) T^2 \left( \frac{N_C}{3} + \frac{N_f}{6} + \frac{1}{2\pi^2} \sum_f \mu^2_q \frac{T}{T^2} \right),$$

(11)

where the quark chemical potential is related to the baryon chemical potential by $\mu_q = \mu_b/3$. For zero chemical potential ($\mu_q=0$), the Debye mass is reduced to

$$m_D(T) = g(T) T \sqrt{\frac{N_c}{3} + \frac{N_f}{6}}$$

(12)

The next-to-leading contribution to Debye mass ($\delta m_D(T, \mu_q)$) comes from the resumed one-loop gluon diagrams \cite{24}. However, we restrict ourselves only up to the leading-order term in the Debye mass. The variations of the Debye mass in leading-order with the chemical potential and temperature are shown in Fig. 1. The left panel of Fig.1 shows that the Debye mass increases with the chemical potential and is more for higher temperature. This implies that the contribution to the Debye mass at lower temperature and high baryon chemical potential mainly come from chemical potential, which can be seen in the right panel of Fig.1. For low temperature region (20-50 MeV) the increase in Debye mass is negligible but increases with $\mu_b$. The temporal component of the retarded (or advanced) propagator in the static limit \cite{42} is

$$D_{R,A}^{00}(0,k) = \frac{1}{(k^2 + m^2_D(T, \mu))},$$

(13)

Hence we can now obtain the dielectric permittivity from the static limit of the “00”-component of gluon propagator

$$\epsilon^{-1}(k) = -\lim_{\omega \to 0} k^2 D_{11}^{00}(\omega, k) ; \quad D_{11}^{00}(\omega, k) = \frac{1}{2} (D_{R}^{00} + D_{A}^{00})$$

(14)

Thus the medium modified potential is obtained by substituting Eqs. (2 and 14) in Eq.(1)

$$V(r, T, \mu) = \int \frac{d^3k}{(2\pi)^{3/2}} (e^{ik\cdot r} - 1) \left( -\frac{\sqrt{2/\pi}}{k^2} \left( \alpha + \frac{2\sigma}{k^2} \right) \right) \left[ \frac{k^2}{(k^2 + m^2_D)} \right]$$

(15)

Solving the integral in above equation, one can to obtain the medium modified potential at finite
density and temperature (with $\hat{r} = rm_D(T, \mu)$) as:

$$V(r; T, \mu) = \frac{2\sigma}{m_D(T, \mu)} \left( \frac{e^{-\hat{r}} - 1}{\hat{r}} + 1 \right) - \alpha m_D(T, \mu) \left( \frac{e^{-\hat{r}}}{\hat{r}} + 1 \right),$$  \hspace{1cm} (16)

where the chemical potential and temperature dependencies are introduced through the Debye mass, $m_D(T, \mu)$. The potential has a long range Coulombic tail in addition to the standard Yukawa term and the constant terms are introduced to yield the correct limit of $V(r, T, \mu)$ as $T \to 0, \mu \to 0$ (it reduces to the Cornell potential).

It is worth to note that the potential in a hot QCD medium is not the same as the lattice parametrized heavy quark free-energy in the deconfined phase which is basically a screened Coulomb [43, 44] because one-dimensional Fourier transform of the Cornell potential in the medium yields the similar form as used in the lattice QCD to study the quarkonium properties which assumes the one-dimensional color flux tube structure [45]. However, at finite temperature and density that may not be the case since the flux tube structure may expand in more dimensions[43], hence the three-dimensional form of the medium modified Cornell potential may be the better option.

The potential (16) leads to an analytically solvable Coulomb potential if one neglect Yukawa term in the limit $r >> 1/m_D$ and the product $\alpha m_D$ will be much greater than $2\sigma/m_D$ for large values of chemical potentials.

$$V(r, T, \mu) \sim -\frac{2\sigma}{m_D^2(T, \mu)r} - \alpha m_D(T, \mu)$$  \hspace{1cm} (17)
To see the effects of a hot and dense medium on $Q\bar{Q}$ potential, we have evaluated the potential at large chemical potential (in the low temperature range 20-60 MeV) *viz.* at 600 MeV, 800 MeV and 1000 MeV *etc.* as shown in Fig.2. We have found that the potential rises with increasing $r$, but slower than in the linear way (long-distance QCD force becomes short-distance) and at some large enough $r$ it simply flattens out (due to the breaking of QCD string, *i.e.* a heavy meson splits into two heavy-light mesons). This observation agrees with the lattice calculations [46, 47]. Thus the deconfinement is reflected in the large-distance behavior of heavy quark potential at large chemical potential. That is why the in-medium behavior of heavy quark bound states is used to probe the state of matter in QCD thermodynamics at finite density and/or temperature.

### 3 Binding energy and dissociation chemical potential

The medium-modified effective potential obtained in previous section has been employed to study the in-medium properties of the heavy quark bound states such as, binding energies and dissociation chemical potentials. We studied the binding energy and dissociation chemical potential for the ground states of $c\bar{c}$. In order to understand bound state properties of quarkonium states in hot and dense QCD medium one need to solve the Schrödinger equation using the $Q\bar{Q}$ potential. It can be seen from Eq.(16) that in the short-distance limit, the vacuum contribution of $Q\bar{Q}$ potential supersedes over the in-medium contribution whereas the in-medium contribution affects the
potential in the long-distance limit. As we have already seen the potential (17) in long-distance \((rm_D \gg 1)\) limit and for large values of chemical potential (or temperature) is Coulomb-like potential after identifying \(2\sigma/m_D^2\) with the square of color charge, \(g_s^2\), so the energy eigenvalues are read as in hydrogen atom problem:

\[
E_{\text{bin}} \equiv \frac{r \gg 1}{m_D^2(T,\mu) n^2 + \alpha m_D(T,\mu)} ; \quad n = 1, 2 \cdots.
\]

(18)

However, in the intermediate-distance \((rm_D(T,\mu) \simeq 1)\) scale, the interaction becomes nontrivial and the potential does not look simpler in contrast to the asymptotic limits, thus the potential in general needs to be dealt numerically to obtain the binding energies. The matrix method serves the purpose and the stationary Schrödinger equation can be solved in a matrix form through a discrete basis, instead of the continuous real-space position basis spanned by the states \(|\vec{x}\rangle\). Subdividing the confining potential \(V\) into \(N\) discrete wells having potentials \(V_1\) through \(V_{N+2}\), such that \(V = V_i\) for \(x_{i-1} < x < x_i; i = 2, 3, \ldots, (N + 1)\) is the \(i\)th boundary potential. To exists a bound state there must be exponentially decaying wave function in the region \(x > x_{N+1}\) as \(x \to \infty\) and the wave function in this region reads like:

\[
\Psi_{N+2}(x) = P_E \exp[-\gamma_{N+2}(x-x_{N+1})] + Q_E \exp[\gamma_{N+2}(x-x_{N+1})],
\]

(19)

where, \(P_E = \frac{1}{2}(A_{N+2} - B_{N+2})\), \(Q_E = \frac{1}{2}(A_{N+2} + B_{N+2})\) and, \(\gamma_{N+2} = \sqrt{2\mu(V_{N+2} - E)}\). The energy eigenvalues can be identified by scanning the zeros of \(Q_E\). (Therefore, using the matrix method, we have obtained the binding energies as a function of chemical potential for small temperature range \((20 - 40 \text{ MeV})\) in Fig. 3. It is found that the binding energy decreases with the increase of chemical potential, which can be understood by the fact that the screening gets stronger with the increase of the baryon chemical potential, hence the inter-quark potential becomes weaker compared to \(\mu = 0\). Thus the study of the binding energies at finite temperature and chemical potential provides a information about the dissociation pattern of quarkonium states in the region of high baryon chemical potential and low temperature.

Thus one can obtain the dissociation chemical potential of quarkonia when their binding energies in the dense baryonic medium are of the order of the baryon chemical potential. It is noticed that with the rise in temperature the dissociation chemical potential deceases (Table 1) i.e. for example at \(T = 20\ \text{MeV}\), the \(J/\psi\) will be dissolve at \(\mu_D/\mu_C = 2.2\) whereas for slightly higher temperature, \(T = 30\ \text{MeV}\), \(\mu_D/\mu_c\) becomes slightly smaller (2.1). This feature can be used to locate
the pin-point on QCD phase diagram in the high-baryon density limit, \textit{i.e.} where the confined nuclear matter cross to the phase boundary and becomes baryon-rich deconfined quark matter.

### 4 Conclusion

We have studied the quarkonium dissociation in hot and dense QCD medium by correcting both the perturbative and nonperturbative terms of the Cornell potential through the dielectric permittivity. Thereafter the chemical potential- and temperature-dependent potential is plugged into Schrödinger equation to study the properties of quarkonia states in high baryon density and low temperature region.

It is noticed that the screening mass increases dominantly with the chemical potential rather than with the temperature in the high baryon density and low temperature region. As a consequence the binding energies decrease with the chemical potential. The dissociation chemical potentials of $J/\psi$ at $T=40$ and $50$ MeV is found to be $1412$ MeV and $1255$ MeV, respectively. Finally we conclude that the dissociation of quarkonia states in large baryon density may help to
trace out the point on QCD phase diagram at which the baryon rich QGP (which is expected at FAIR energies) will be created.

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