**J/ψ** suppression in nucleus-nucleus collisions

*B* Binoy Krishna Patra¹, Vinod Chandra², and Vineet Agotiya¹

¹ Department of Physics, Indian Institute of Technology Roorkee, India, 247 667 and ² Department of Physics, Indian Institute of Technology Kanpur, India, 208 016

At high temperatures, strongly interacting matter becomes a plasma of deconfined quarks and gluons. In statistical QCD, deconfinement and the properties of the resulting quark-gluon plasma can be investigated by studying the in-medium behaviour of heavy quark bound states. In high energy nuclear interactions, quarkonia probe different aspects of the medium formed in the collision. So, we first reviewed the fate of quarkonia in the different stages of the (dynamical) system produced at the collision. We have then presented our present work on the dissociation of the heavy quarkonium states in a hot QCD medium by investigating the medium modifications to heavy quark potential. In contrast to the usual screening picture, interestingly our theory gives rise the screening of the charge, not the range of the potential.

**Introduction:** The study of the fundamental forces between quarks and gluons is an essential key to the understanding of QCD and the occurrence of different phases which are expected to show up when going from low to high temperatures \(T\) and/or baryon number densities. For instance, at small or vanishing temperatures quarks and gluons get confined by the strong force while at high temperatures asymptotic freedom suggests a quite different QCD medium consisting of rather weakly coupled deconfined quarks and gluons, the so-called quark gluon plasma (QGP). The anomalous suppression of the J/ψ production in heavy ion collisions which has been experimentally observed in the depletion of the dilepton multiplicity in the region of invariant mass corresponding to the J/ψ meson was proposed long time ago as a possibly unambiguous signal of the onset of deconfinement. Matsui and Satz argued that charmonium states produced before the formation of a thermalized QGP would tend to melt in their path through the deconfined medium, since the binding (colour) Coulomb potential is screened by the large number of colour charges. This, in turn, would produce an anomalous (with respect to normal nuclear absorption) drop in the J/ψ yields.

In this picture it is implicitly assumed that, once the charmonium dissociates, the heavy quarks hadronize by combining with light quarks only (recombination leading to a secondary J/ψ production is neglected). This assumption is certainly justified at the SPS conditions, due to the very small number of \(c\bar{c}\) pairs produced per collision \(N_{c\bar{c}} \sim 0.2\) in a central collision), but at RHIC \(N_{c\bar{c}} \sim 10\) and LHC \(N_{c\bar{c}} \sim 200\) energies it is no longer warranted.

Moreover in a hadronic collisions only about 60% of the observed J/ψ’s are directly produced, the remaining stemming from the decays of excited charmonium states (notably the \(\chi_c\) and the \(\psi’\)). Since each \(c\bar{c}\) bound state dissociates at a different temperature, a model of sequential suppression was developed, with the aim of reproducing the J/ψ suppression pattern as a function of the energy density reached in the heavy ion collision. SPS experimental data for Pb-Pb collisions at different centralities seem indeed to support the dissociation pattern predicted by this model.

The heavy quark pair leading to the J/ψ mesons are produced in nucleus-nucleus collisions on a very short time-scale \(\sim 1/2m_c\), where \(m_c\) is the mass of the charm quark. The pair develops into the physical resonance over a formation time \(\tau_{\psi}\) and traverses the plasma and (later) the hadronic matter before leaving the interacting system to decay (into a dilepton) to be detected. This long ‘trek’ inside
the interacting system is fairly ‘hazardous’ for the $J/\psi$. Even before the resonance is formed it may be absorbed by the nucleons streaming past it. By the time the resonance is formed, the screening of the colour forces in the plasma may be sufficient to inhibit a binding of the $c\bar{c}$ or an energetic gluon or a comoving hadron could dissociate the resonance(s).

Quarkonia at finite temperature are an important tool for the study of quark-gluon plasma formation in heavy ion collisions. Many efforts have been devoted to determine the dissociation temperatures of $Q\bar{Q}$ states in the deconfined medium, using either lattice calculations of quarkonium spectral functions or non-relativistic calculations based upon some effective (screened) potential.

Lattice studies are directly based on quantum chromodynamics and should provide, in principle, a definite answer to the problem. However, in lattice studies the spectral functions have to be extracted — using rather limited sets of data — from the Euclidean (imaginary time) correlators, which are directly measured on the lattice. This fact, together with the intrinsic technical difficulties of lattice calculations, somehow limits the reliability of the results obtained so far, and also their scope, which in fact is essentially limited to the mass of the ground state in each $Q\bar{Q}$ channel. Potential models, on the other hand, provide a simple and intuitive framework for the study of quarkonium properties at finite temperature, allowing one to calculate quantities that are beyond the present possibilities for lattice studies. The main problem of the latter approach is the determination of the effective potential: although at zero temperature the use of effective potentials and their connection to the underlying field theory is well established, at finite $T$ the issue is still open.

Calculations of the $c\bar{c}$ and $b\bar{b}$ dissociation temperatures, using different potential models based upon the lattice free and internal energies, have found on the whole a reasonable agreement with the results from the lattice studies. On the other hand, calculations of Euclidean correlators using a variety of potential models were not able to reproduce the temperature dependence of the lattice correlators.

A precise quantitative agreement with the lattice correlators should not be expected, because of uncertainties coming from a variety of sources. Not only the determination of the effective potential is still an open question but there are also issues tied, e.g., to relativistic effects, to the thermal width of the states or to the contribution of radiative corrections. On the other hand, lattice correlators are also affected by their own uncertainties. These may be due to the use of different lattices (isotropic or anisotropic); to the finite size of the box, which might significantly alter the continuum part of the spectrum, although calculations with boxes of different sizes show discrepancies below 1% [5,6]; or to artifacts in the continuum region of the spectral functions due to the finite lattice spacing [5].

Recently, Umeda and Alberico have shown that the lattice calculations of meson correlators at finite temperature contain a constant contribution, due to the presence of zero modes in the spectral functions. These contributions cure most of the previously observed discrepancies with lattice calculations, supporting the use of potential models at finite temperature as an important tool to complement lattice studies.

Actually, even if the potential supports the existence of bound states, other physical processes may lead to the dissociation of the quarkonium. First, if the $Q\bar{Q}$ binding energy is lower than the temperature — and assuming that the quarkonia have reached the thermal equilibrium with the plasma — a certain fraction of their total number will be thermally excited to resonant states according to a Bose-Einstein distribution: such a process is referred to as thermal dissociation. Furthermore, the collisions with the gluons and the light quarks of the plasma may lead to the collisional dissociation of the quarkonium.

Binoy and Menon revisited the $J/\psi$ suppression due to gluonic bombardment in an expanding quark-gluon plasma in a series of works. First they neatly incorporated the
crucial effects arising from gluon fugacity, relative $g-\psi$ flux, and $J/\psi$ meson formation time and then used these effects in the formulation of the gluon number density, velocity-weighted cross section, and the survival probability in an equilibrated static QGP. This formulation have been used to study the pattern of $J/\psi$ suppression in the central rapidity region at RHIC/LHC energies. Later they explicitly take into account the effect of hydrodynamic (longitudinal) expansion profile on the gluonic breakup of $J/\psi$'s in an (chemically) equilibrating expanding QGP. A novel type of partial-wave interference mechanism is found to operate in the modified dissociation rate. Finally, this formulation has been applied to the case when the medium is undergoing cylindrically symmetric transverse expansion. Compared to the case of longitudinal expansion the new graph of survival probability develops a rich structure at RHIC, due to a competition between the transverse catch-up time and plasma lifetime.

Of course, if the process $g + J/\psi \rightarrow c + \bar{c}$, discussed above, can lead to the dissociation of the charmonium, the same reaction can also occur in the opposite direction. Hence a consistent calculation of $J/\psi$ multiplicity implies the solution of a kinetic rate equation integrated over the lifetime of the QGP phase in which both processes (dissociation and recombination) enter.\footnote{This is of relevance because, as mentioned above, the usual assumption in considering the $J/\psi$ suppression as a signature of deconfinement is that its production can occur only in the very initial stage of the collision. Really, if at SPS the role played by recombination is numerically negligible, this is no longer true at RHIC as pointed out in Ref.\footnote{}}\footnote{\cite{3}}

However, there is a generic question, quite often asked, is that whether we can distinguish between the two mechanisms of dissociation (colour screening and collision with hard gluons) mainly operating in the deconfined phase. Binoy and Srivastava\footnote{\cite{13}} have shown that while the gluonic dissociation of the $J/\psi$ is always possible, the Debye screening is not effective in the case of small systems at RHIC energies. For the larger systems, the Debye screening is more effective for lower transverse momenta, while the gluonic dissociation dominates for larger transverse momenta. At LHC energies the Debye screening is the dominant mechanism of $J/\psi$ suppression for all the cases and momenta studied. As an interesting result, they found the gluonic dissociation to be substantial but the Debye screening to be ineffective for $\Upsilon$ suppression at the LHC energy.

In the context of $J/\psi$ suppression, Langevin dynamics seems to be almost as important as other mechanisms invoked to explain the RHIC and LHC data. Binoy and Menon\footnote{\cite{14}} considered the Brownian motion of a $c\bar{c}$ pair produced in the very early stage of a QGP. They found that, in the weak coupling regime, both the time scales associated with the positional swelling ($\tau_x$) and approach to ionization ($\tau_E$) are positive and less than the frictional relaxation time ($\gamma^{-1}$). Hence Brownian movement can cause a genuine break-up of the $c\bar{c}$ bound by swelling it substantially or Langevin dynamics can cause $c\bar{c}$ to ionize after a time span $\tau_E$. On the other hand, in the strong coupling case, $\tau_x$ is imaginary and $\tau_E$ are negative, i.e., unphysical. Hence random force plus diffusion cannot cause the $c\bar{c}$ bound state to dissociate.

While the short and intermediate distance ($rT \leq 1$) properties of the heavy quark interaction is important for the understanding of in-medium modifications of heavy quark bound states, the large distance property of the heavy quark interaction which is important for our understanding of the bulk properties of the QCD plasma phase viz. the equation of state. In all of these studies deviations from perturbative calculations and the ideal gas behaviour are expected and were indeed found at temperatures which are only moderately larger than the deconfinement temperature. This calls for quantitative non-perturbative calculations. The phase transition in full QCD will appear as an crossover rather than a 'true' phase transition with related singularities in thermodynamic observables (in the high-temperature and low density regime) a cross-over, it can be...
reasonable to assume that the string-tension does not vanish abruptly above $T_c$. So we decide to investigate in our present work what happens to the different quarkonium states if one corrects with a dielectric function encoding the effects of the deconfined medium, the full Cornell potential and not only its Coulomb part as usually done in the literature. We found that with this choice medium effects give rise to a long-range Coulomb potential with a reduced effective charge (inversely proportional to the square of the Debye mass) of the heavy quark, at variance with its usual Debye-screened form employed in most of the literature. With such an effective potential we investigate the effects of different possible choices of the Debye mass on the dissociation temperature of the different quarkonium states.

The Debye masses in hot QCD:

Adequate knowledge of Debye mass is indeed needed to study the medium modifications to heavy quark potential. The Debye mass in QCD unlike QED is generically non-perturbative and gauge invariant. The Debye mass at high temperature in the leading-order in QCD coupling is known from long time and is perturbative. The Debye mass in leading-order from the polarization tensor of a gauge boson derived from the HTL approach can also be obtained from the transport theory. One cannot naively generalize the definition of Debye mass in QED to QCD due to the non-abelian nature of the theory. Rebhan has defined Debye mass through the relevant pole of the static quark propagator instead of the zero momentum limit of the time-time component of the gluon self-energy. The Debye mass thus defined comes out to be gauge independent follows from the fact that the pole of the self-energy is independent of choice of gauge. Braaten and Nieto computed the Debye screening mass for QGP at high temperature to the next-to-leading-order in QCD coupling from the correlator of two Polyakov loops which agreed to the HTL result.

Arnold and Yaffe pointed out that the contribution of order $(g^2T)$ to the Debye mass in QCD needs the knowledge of the non-perturbative physics of confinement of magnetic charges and a perturbative definition of the Debye mass as a pole of gluon propagator no longer holds. They showed how one can define Debye mass in QCD in a manifestly gauge invariant manner (in vector-like gauge theories with zero chemical potential). In the work of Kajantie et. al, the non-perturbative contributions of $O(g^2T)$ and $O(g^3T)$ have been determined from 3-D effective field theory which we consider in the present work. At high temperatures and zero chemical potential Debye mass can be expanded in a power series in QCD coupling:

$$m_D = m_D^{LO} + \frac{N g^2 T}{4 \pi} \ln \frac{m_D^{LO}}{g^2 T} + c_N g^2 T + d_{N,N_f} g^3 T + \mathcal{O}(g^4 T),$$

where $m_D^{LO}$ is the leading-order result. The coefficient $d_{N,N_f}$ have the following dependence on the number of colors $N$ and flavors $N_f$ as:

$$d_{N,N_f} = \frac{b_N}{\sqrt{N/3 + N_f/6}},$$

where the values of $c_N$, $b_N$ have been obtained by fitting the results with the physical 4D finite temperature QCD:

$$\text{SU(3)} : c_N = 2.46 \pm 0.15 \quad b_N = -0.49 \pm 0.15$$

The number $c_N$ captures the non-perturbative 3-D effects, while the $d_{N,N_f}$ is related to the choice of scale in $m_D^{LO}$. We employ the two-loop expression for the QCD coupling constant at finite temperature. We use the following
notations henceforth,

\[ m_{D}^{\text{LO}} = g(T)T \sqrt{\frac{N}{3} + \frac{N_f}{6}} \]
\[ m_{D}^{\text{NLO}} = m_{D}^{\text{LO}} + \frac{N g^2 T}{4\pi} \ln \frac{m_{D}^{\text{LO}}}{g^2 T} \]
\[ m_{D}^{\text{NP}} = m_{D}^{\text{NLO}} + c_N g^2 T + d_{N,N_f} g^3 T \]
\[ m_{D}^2 = 1.4 m_{D}^{\text{LO}} \],

(4)

where \( m_{D}^{\text{LO}} \) is the Debye mass obtained by fitting the (colour-singlet) free energy in lattice QCD [22].

In the weak coupling \((g \ll 1)\) regime, the soft scale \((m_D \simeq gT)\) at the leading-order related to the screening of electrostatic fields is well separated from the ultra-soft scale \((\simeq g^2 T)\) related to the screening of magnetostatic fields. In such regime it appears meaningful to see the contribution of each terms in the the Debye mass (Eq. 1) separately. But when the coupling becomes large enough (which is indeed the case), the two scales are no longer well separated. So while looking for the next-to-leading corrections to the leading-order result from the ultra-soft scale, it is not a wise idea to stop at the logarithmic term (as mentioned in the notation \(m_{D}^{\text{NLO}}\)), since it becomes crucial the number multiplying the factor \(1/g\) to establish the correction to the LO result. In fact we found that the Debye mass in the NLO term \((m_{D}^{\text{NLO}})\) is always smaller than the LO term \((m_{D}^{\text{LO}})\) because of the negative (logarithmic) contribution \((\log(1/g))\) to the leading-order term, while the full correction (all \(g^2 T\) terms) to the Debye mass results positive. So, we will employ only three forms of the Debye masses viz. leading-order term/perturbative result \((m_{D}^{\text{LO}})\), full (non-perturbative) corrections to the leading-order term \((m_{D}^{\text{NP}})\), and lattice parameterized form \((m_{D}^2)\) to study the dissociation phenomena of quarkonium in a hot QCD medium in this work.

We now proceed to investigate in-medium modifications to heavy-quark potential and its application to determine the binding energy and dissociation temperature of the heavy-quark bound states.

### The in-medium heavy-quark potential:

Let us now turn our attention to study the medium modifications to heavy quark potential at \(T = 0\) which is considered as the Cornell potential,

\[ V(r) = -\frac{\alpha}{r} + \sigma r \quad (5) \]

where \(\alpha\) and \(\sigma\) are the phenomenological parameters. The former accounts for the effective coupling between the heavy quark pairs and the latter gives the string coupling. The medium modifications enters in the Fourier transform of the heavy quark potential as follows:

\[ \hat{V}(k) = \frac{V(k)}{\epsilon(k)} \quad (6) \]

where \(\epsilon(k)\) is the dielectric permittivity given in terms of the static limit of the longitudinal part of gluon self-energy [23]

\[ \epsilon(k) = \left( 1 + \frac{\Pi_L(0,k,T)}{k^2} \right) = \left( 1 + \frac{M_D^2}{k^2} \right) \quad (7) \]

The result for the static limit of the dielectric permittivity is the perturbative one. If one assumes that huge non-perturbative effects (like the string tension) survive above \(T_c\) the same could be true also for such a dielectric function. So, there is a caveat that this (linear) relation of dielectric function \(\epsilon\) on \(M_D^2\) may also pick up modifications due to the presence of non-perturbative effects above the deconfinement point. To get rid of the complexity of the problem, we put all the non-perturbative effects (including the non-zero string tension) together in the effective charge \((2\sigma/m_D^2)\) of the medium modified potential which further depends of the Debye mass. The quantity \(V(k)\), the Fourier transform (FT) of the Cornell potential reads [24]:

\[ V(k) = -\sqrt{\frac{2}{\pi}} \frac{\alpha}{k^2} - \frac{4\sigma}{\sqrt{2\pi} k^4} \quad (8) \]

Substituting Eq. (7) and Eq. 8 into Eq. 6 and evaluation of the inverse Fourier-Transform of the RHS of Eq. (6) one obtains
the $r$-dependence of the medium modified potential. The expression thus reads
\[
V(r, T) = \left( \frac{2\sigma}{m_D^2} - \alpha \right) \exp\left( -\frac{m_D r}{\sigma} \right)
- \frac{2\sigma}{m_D^2 r} + \frac{2\sigma}{m_D} - \alpha m_D
\]
(9)

This potential has a long range Coulombic tail in addition to the standard Yukawa term. After taking the high temperature limit, the above potential takes the form:
\[
V(r) \sim -\frac{2\sigma}{m_D^2 r} - \alpha m_D
\]
(10)

The above form (apart from a constant term) is a Coulombic type as in the hydrogen atom problem by identifying the fine structure constant $\alpha$ with the effective charge $2\sigma/m_D^2$. Since the Debye mass $m_D$ always increases with the temperature, the effective charge $2\sigma/m_D^2$ gets waned as the temperature is increased. This makes the potential too shallow to bind $Q, \bar{Q}$. This results the melting of the bound states. It is important to note here the difference between the screening of the charge and the screening of the range of the potential. In the usual picture adopted to study the dissociation of quarkonia through the Cornell potential, linear term vanishes above the critical temperature because string tension vanishes. So, above the critical temperature, the only nonvanishing term in the potential is the attractive coulombic term which gets screened in a Yukawa form making the potential short-ranged. If the range of the potential becomes too short compared to the Bohr radius it will be dissolved into its constituents. However, in our case, this dissociation happens due to screening of the charge, not due to screening of the range of potential. Note that the constant terms in the potential (Eq.9) are needed in computing the masses of the quarkonium states. It is equally important while comparing our effective potential (Eq.9) with the free energy in lattice studies (discussed below). However, the constant terms are not needed while comparing the values of the dissociation temperatures obtained in our model with the values in the lattice spectral studies. This is due to the different criterion is imposed to evaluate the dissociation temperatures (discussed in the next section).

We need to mention that our in-medium effective potential $V(r, T)$ in Eq.9 agrees qualitatively (and also quantitatively) with the singlet part of the free energy in the lattice QCD [23]. This have been checked by plotting $V(r, T)$ with $rT$ for a fixed value of $T/T_c=3.32$ [13].

We shall now systematically study the effects of perturbative and non-perturbative interactions on the binding energies and dissociation temperatures of quarkonium states in a hot QCD medium by employing the three form of the Debye mass (Eq.4). In addition, we take advantage of all the available lattice data, obtained not only in quenched QCD ($N_f = 0$), but also including two and, more recently, three light flavors. We are then in a position to study also the flavor dependence of the dissociation process, a perspective not yet achieved by the parallel studies of the spectral functions, which are only available in quenched QCD.

### Binding energy and dissociation temperatures

Binding energy of a quarkonium state at zero temperature is defined by the energy difference between the mass of the quarkonium and the open charm/bottom threshold. At finite temperature, the binding energy is defined as the distance between the peak position and the continuum threshold, $E_{bin} = 2m_{c,b} + V_{\infty}(T) - M$ with $M$ being the resonance mass [23]. However, our definition is the conventional one viz. the ‘ionization potential’ in the atomic physics.

Finally, Schrödinger equation for the potential (Eq.11) gives the energy eigen values for the ground state and excited states, viz. $J/\psi$, $\psi'$, $\Upsilon$, $\Upsilon'$ etc. for the charmonium and bottomonium states, by the Bohr’s formula:
\[
E_n = -\frac{E_I}{n^2} \quad E_I = \frac{m_q\sigma^2}{m_D}
\]
(11)

Thus $E_n = -E_I, -E_I/4, -E_I/9, \ldots$ are the
allowed energy levels of $Q\bar{Q}$ bound states. These energies are known as ionization potentials/binding energies for the $n$-th bound state. Thus, the binding energy becomes a temperature-dependent quantity through the Debye mass and it decreases with the temperature.

However, there are other states in the charmonium and bottomonium spectroscopy viz $\chi_c$’s and $\chi_b$’s and the binding energies for them are obtained from a variational treatment of the relativistic two-fermion bound state system in quantum electrodynamics \[27\].

$$E(\chi_{c,b}) = \frac{m_{c,b} \sigma^2}{4 m_D^2} \left(1 + \frac{2 \sigma^2}{3 m_D^2}\right). \quad (12)$$

Figures 1-2 show the variation of binding energy with temperature (in units of $T/T_c$) for the $J/\psi$ and $\Upsilon$ states, respectively. Different curves in the figure denote the choice of Debye mass in Eq.\[11\] used to calculate the binding energy from Eq.\[12\] or Eq.\[13\]. We consider three cases for our analysis: pure gluonic medium, 2-flavor and 3-flavor to see the flavor dependence of dissociation pattern in QCD medium.

The binding energy of the $J/\psi(\approx 640$ MeV) is considerably larger than the typical non-perturbative hadronic scale $\Lambda_{QCD}$. As a consequence, perturbative term (leading-order term) in the Debye mass takes care variation of binding energy with temperature. The same argument holds good for $\Upsilon$ also. Once we switch-on the non-perturbative contributions in the Debye mass through the coefficients $c_N$ and $d_{N,N}$, the Debye mass becomes so large that the binding energies of all the quarkonia becomes too small (even at the temperature 100 MeV) compared to their ground state binding energies. Indeed, the value of the Debye mass after inclusion of non-perturbative terms is approximately three times larger than LO and NLO results and twice as large as lattice parametrized Debye mass near $2T_c$. The temperature dependence of the binding energy for other quarkonium states is studied in length in Ref. \[13\]. However, this is not the complete story, the situation may change once the $O(g^4 T)$ non-perturbative contributions to Debye mass are evaluated and utilize to estimate the binding energy for quarkonia states.

Thus the study of temperature dependence of binding energy will help us to determine the dissociation temperatures of the quarkonium states in thermal medium. Mocsy and Petreczky \[26\] have defined the dissociation temperature as the temperature above which the quarkonium spectral function shows no resonance-like structures, meaning that particular state is dissolved.

Physically, (thermal) dissociation of a bound state in a thermal medium can be explained as follows: when the binding energy of a resonance (viz. $J/\psi$) state drops below the mean thermal energy of parton the state have become feebly bound and thermal fluctuations can destroy it by transferring energy and exciting the quark anti-quark pair into its continuum. So, if the binding energy of a $c\bar{c}$ or $b\bar{b}$ state at some temperature becomes equal or smaller than the mean thermal energy then the state is said to be dissociated. Since the (relativistic) thermal energy of the partons is $3T$ hence the dissociation temperature $T_D$ of the $n$-th $Q\bar{Q}$ bound state will be determined by the condition:

$$\frac{1}{n^2 m_D^2(T_D)} = 3T_D \quad (13)$$

The above condition gives the dissociation temperatures after inserting expression for the Debye mass displayed in Eq.\[4\]. However, the choice $3T$ is not rigid because even at low temperatures $T < T_c$ (say) the Bose/Fermi distributions of partons will have a high energy tail with partons of mechanical energy $\epsilon > |E_n|$. While determining the temperature dependence of the binding energy and dissociation temperatures the string tension is chosen to be $\sigma = 0.184 \text{GeV}^2$. The dissociation temperatures for the ground (1S), first excited states (2S), $\chi_c$, and $\chi_b$ (1P) of $c\bar{c}$ and $b\bar{b}$ are listed in the Tables 1 and 2 for the Debye masses in the leading-order and the lattice parametrized form, respectively. We do not put up the list for the dissociation temperatures with the non-perturbative form of the Debye mass be-
FIG. 1: Dependence of $J/\psi$ binding energy on temperature

FIG. 2: Dependence of $\Upsilon$ binding energy on temperature

TABLE I: Dissociation temperatures for various quarkonia (in unit of $T_c$) for $m_{D}^{LO}$.

| Quarkonium state | Pure QCD | $N_f = 2$ | $N_f = 3$ |
|------------------|----------|----------|----------|
| $J/\psi$         | 1.1      | 1.3      | 1.2      |
| $\psi'$          | 0.8      | 0.9      | 0.9      |
| $\chi_c$         | 0.9      | 1.1      | 1.0      |
| $\Upsilon$       | 1.4      | 1.7      | 1.6      |
| $\Upsilon'$      | 1.0      | 1.2      | 1.2      |
| $\chi_b$         | 1.1      | 1.5      | 1.2      |

The dissociation temperatures obtained with the Debye mass in the leading-order (Table 1) is always larger than the dissociation temperatures with the Debye mass parametrized in lattice QCD ($m_{D}^{L}$) (Table 2).

TABLE II: Dissociation temperatures for various quarkonia (in unit of $T_c$) for $m_{D}^{L}$.

| Quarkonium state | Pure QCD | $N_f = 2$ | $N_f = 3$ |
|------------------|----------|----------|----------|
| $J/\psi$         | 0.8      | 0.9      | 0.9      |
| $\psi'$          | 0.5      | 0.7      | 0.6      |
| $\chi_c$         | 0.6      | 0.7      | 0.7      |
| $\Upsilon$       | 1.0      | 1.2      | 1.2      |
| $\Upsilon'$      | 0.7      | 0.9      | 0.8      |
| $\chi_b$         | 0.7      | 0.9      | 0.9      |

cause the values obtained are too small to explain physically. We have taken the values of critical temperatures ($T_c$) 270 MeV, 203 MeV and 197 MeV for pure gluonic, 2-flavor and 3-flavor QCD medium, respectively [28].
for both charmonium and bottomonium states because of the larger value of the Debye mass in lattice compared to LO values. The results shown in the Tables 1-2 lend support with the recent lattice predictions [26, 29]. The upper bound of the dissociation temperatures could be obtained if average thermal energy is replaced by $\sim T$.

Absorption by nucleons and co-movers

So far we have discussed the fate of quarkonia only when the presence of quark gluon plasma is considered. It is very well established that there are several aspects like initial state scattering of the partons, shadowing of partons, absorption of the pre-resonances ($Q\bar{Q}g >$ states) by the nucleons before they evolve into physical quarkonia, and also dissociation of the resonances by the comoving hadrons. It has been argued that the absorption by co-moving hadrons will be important for $\psi'$, due to its very small binding energy, while for more tightly bound resonances it may be weak.

Let us briefly comment on them one-by-one. Shadowing of partons should play an important role in the reduced production of quarkonia, especially at the LHC energies. It is clear that if shadowing is important, we shall witness a larger effect on $J/\psi$ than on $\Upsilon$, because of the smaller values of the $x$ for gluons. At the same time, the effect of shadowing should be similar for different resonances of the charmonium (or bottomonium), as similar $x$ values would be involved for them.

The absorption of the pre-resonances by the nucleons is another source of $p_T$ dependence. It is important, to recall once again that as the absorption is operating on the pre-resonance, the effect should be identical for all the states of the quarkonium which are formed.

This is a very important consideration as it is clear that if we look at the ratio of rates for different states of $J/\psi$ or the $\Upsilon$ family as a function of $p_T$, then in the absence of QGP-effects they would be identical to what one would have expected in absence of nuclear absorption and shadowing, providing a clear pedestal for the observation of QGP.

Thus another aspect of $p_T$ dependence which needs to be commented upon. The (initial state) scattering of partons, before the gluons of the projectile and the target nucleons fuse to produce the $Q\bar{Q}$-pair, leads to an increase of the $< p_T^2 >$ of the resonance which emerges from the collision [17]. The increase in the $< p_T^2 >$, compared to that for $pp$ collisions is directly related to number of collisions the nucleons are likely to undergo, before the gluonic fusion takes place. This leads to a rich possibility of relating the average transverse momentum of the quarkonium to the transverse energy deposited in the collision (which decides the number of participants and hence the number of collisions). Considering that collisions with large $E_T$ may have formation of QGP in the dense part of the overlapping region, the quarkonia, which are produced in the densest part (and hence contributing the largest increase in the transverse momentum) are also most likely to melt and disappear. This may lead to a characteristic saturation and even turn-over of the $< p_T^2 >$ when plotted against $E_T$ when the QGP formation takes place. In absence of QGP, this curve would continue to rise with $E_T$.

Conclusions and Outlook

We have reviewed different theories/mechanisms of dissociation phenomena of heavy quarkonia in different stages of system produced in the relativistic nuclear collisions. Because the system formed just after the collision is not static and (thermally and chemically) equilibrated, it is rapidly expanding, hadronize, and finally produces many pions, photons, leptons etc. detected at the detectors. We discussed the dissociations mainly in three stages: initial state scattering/absorption, plasma interactions, hadronic/comover absorptions. However, we gave much emphasis on the second stage of the dissociation: dissociation in the deconfined medium/plasma. Apart from a comprehensive survey of different approaches in lattice QCD and potential
based studies, we present our very recent work on the dissociation of quarkonia in a hot QCD medium by investigating the in-medium modifications to heavy quark potential. This is something new because, in our formalism, medium modification results the (dynamical) screening of the color charge in contrast to the screening of the range of the potential in the usual screening picture. The screening of the effective charge, in turn, causes the energy of the quarkonium state depends on temperature. We have then systematically studied the temperature dependence of the binding energy of the ground \((1S)\) state, first excited \((2S)\) state, and \(1P\) states of charmonium and bottomonium in the pure and realistic QCD medium. We then determined the dissociation temperatures with the Debye mass in leading-order and in the lattice parametrized form. The results are reasonably close to the finding of other theoretical works based on potential models [26]. On the other hand these values are significantly smaller than the predictions made by others [9, 29].

In the end, we conclude that all the well explored and yet non-QGP effects need to be accounted for, before we can begin to see the suppression of the quarkonium due to the formation of QGP. It seems that this has been achieved at least at the SPS energies.

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