Charmonium Productions: Nucleon-Nucleon versus Nucleus-Nucleus Collisions

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Abstract

Charmonium productions in $p-p$ and $A-A$ collisions have been estimated within the ambit of colour evaporation model (CEM). The model parameters have been fixed by fitting the theoretical results with CDF data. The method is then applied to RHIC and LHC energies to obtain the transverse momentum distributions of $J/\psi$, $\psi'$ and $\chi_c$. Suppression due to Debye screening in a quark gluon plasma (QGP) is estimated at various centrality cuts. The final $p_T$ distributions of various resonances are then predicted convoluting with the survival probability.

I. Introduction

Ever since the possibility of creating quark gluon plasma (QGP) in relativistic heavy ion collision was envisaged, numerous signals were proposed to probe the properties of such an exotic state of matter. In this context Satz and Matsui [1] had suggested that the production of heavy quark resonances ($J/\psi$) will be suppressed as a result of colour Debye screening in a hot and dense system of quarks, anti-quarks and gluons. This suppression could be detected experimentally through the dileptonic decay mode of these resonances. ALICE dimuon spectrometer [2] is dedicated to look for this type of signal. However, it is a daunting task to disentangle the contributions of the heavy quarkonium states to muon spectrum due to the background from several other sources, \textit{e.g.} Drell-Yan, semileptonic decay of open heavy flavoured mesons ($D\bar{D}, B\bar{B}$) etc. Low energy muons from kaons and pions also constitute a large background.

In this work we shall estimate the hard charmonium productions (\textit{i.e.} $J/\psi$ produced from initial hard process, will be called hard $J/\psi$ hereafter) both in $p-p$ and $A-A$ collisions. The production of heavy resonances proceeds via the following
two steps: (i) the production of heavy quark-antiquark pairs (perturbative), (ii) their resonance interactions to form the bound state (non-perturbative).

The initial state in relativistic heavy ion collisions consists of either hadronic matter or QGP depending on the incident energies of the colliding nuclei. At LHC energies the formation of QGP is unavoidable. Even if the system is formed in QGP phase it will revert back to hadronic phase due to the cooling of the expanding QGP system and hence the interaction of the hard $J/\psi$ with the hadronic matter is inevitable. Therefore, in addition to the suppression due to Debye screening, one needs to consider the survival probability of those $J/\psi$ due to its interactions with the hadronic medium. In this work we have neglected the suppressions due to co-movers as this effect is found to be negligibly small.

The paper is organized as follows. In section II we shall describe the formalism for charmonium productions in the CEM along with the survival probability as a function of centrality. In section III results of our calculations will be presented followed by summary and discussion in section IV.

II. Charmonium Productions

a. Nucleon-Nucleon Collisions

The CEM (also known as local duality approach) \cite{3, 4, 5, 6, 7, 8} states that a heavy quark pair with mass $M_{Q\bar{Q}} < 2M_q$ transforms ($q$ designates light quark, while $Q$ stands for $c$ or $b$ quarks), independent to its colour and spin, to a $Q\bar{Q}$ bound state. The production of a particular bound state depends on the dynamical details of hadronization procedure. The bound state formation probability ($F[nJ^{PC}]$; the so called normalization factor) cannot be calculated from first principle. This is treated as a parameter and can be extracted by fitting the model with the experimental data. $F[nJ^{PC}]$ depends on the particular resonance state under consideration, mass of the heavy flavour, the order (LO or NLO) of the $Q\bar{Q}$ production and the type of parton distribution function (PDF) used. Another important point is that the produced $Q\bar{Q}$ is not constrained to have the proper spin-parity and colour neutrality. The pair sheds its colour non-perturbatively to evolve into asymptotic state without affecting the cross-section. This model cannot predict the absolute cross-section of
heavy resonance production, it however, predicts, their \( p_T \) and \( \sqrt{s} \) dependence. As mentioned, CEM does not bother about the spin-parity and colour neutrality of the bound state. This constraint can be removed by considering an improved version of the duality approach, where one assumes that only the colour singlet part of the cross-section contributes to the bound state production. It might be recalled here that the CEM contains some features of non-relativistic QCD, due to the inclusion of colour octet processes.

In spite of these limitations of CEM, it is capable of explaining \( p_T \) distribution of charmonium in hadron-hadron collisions reasonably well. As mentioned before, the production of charmonium consists of two stages; production of a \( c \bar{c} \) pair (perturbative process) and subsequent non-perturbative evolution into asymptotic states. We have considered those hard processes which can contribute to \( c \bar{c} \) productions irrespective of their colour and spin-parity. The colour neutralization occurs by the interactions (one or more soft gluon emission) with the surrounding colour fields and this step is considered to be non-perturbative. In CEM quarkonium production is treated identically to open heavy flavour production with the exception that in the case of quarkonium, the invariant mass of the heavy quark pair is restricted below the open charm/bottom mesons threshold (see eq. (2) below), which is twice the mass of the lowest meson mass that can be formed with the heavy quark. Depending on the quantum numbers of the \( Q \bar{Q} \) pair different matrix elements are needed for various resonances. The effects of these non-perturbative matrix elements are combined into the factor \( F[nJ^{PC}] \) which is universal [9] i.e. process and kinematics independent. It describes the probability that the \( Q \bar{Q} \) pair forms a quarkonium of given spin (\( J \)), parity (\( P \)) and charge conjugation (\( C \)). The production cross section for a \( J/\psi \) or any charmonium state is therefore given by [9]

\[
\sigma[R(nJ^{PC})] = F[nJ^{PC}] \tilde{\sigma}[Q\bar{Q}],
\]

where the non-perturbative (long distance) factor can be written in terms of the probability to have colour singlet state \((1/9)\) and the fraction \( \rho_R \) of each specific charmonium state. The perturbative contribution (short distance) is given by

\[
\tilde{\sigma}[Q\bar{Q}] = \int_{2m_Q}^{2m_{D/B}} dM_{Q\bar{Q}}^2 \frac{d\sigma[Q\bar{Q}]}{dM_{Q\bar{Q}}^2}.
\]
The contributions to heavy quark production in leading order come from \( q\bar{q} \rightarrow Q\bar{Q} \) and \( gg \rightarrow Q\bar{Q} \). The differential cross-section for heavy quark pair production in hadron-hadron collision is given by [10]

\[
\frac{d\sigma}{dM_{Q\bar{Q}}^2 dy}[h_A h_B \rightarrow Q\bar{Q}X] = \frac{H(x_a, x_b, Q^2)}{s}, \tag{3}
\]

where

\[
H(x_a, x_b, Q^2) = \sum_f \left[ q_f^A(x_a, Q^2)\bar{q}_f^B(x_b, Q^2) + \bar{q}_f^A(x_a, Q^2)q_f^B(x_b, Q^2) \right] \hat{\sigma}_{q\bar{q} \rightarrow Q\bar{Q}} + g^A(x_a, Q^2)g^B(x_b, Q^2)\hat{\sigma}_{gg \rightarrow Q\bar{Q}}. \tag{4}
\]

\( x_{a,b} = M_{Q\bar{Q}}^2 e^{\pm y}/\sqrt{s} \), \( \sqrt{s} \) being the centre of mass energy of the hadronic system, \( y \) stands for rapidity and \( M_{Q\bar{Q}} \) is the invariant mass of the pair. \( q_f \)'s and \( g \)'s are the PDFs for quarks and gluons respectively, these are to be taken from either CTEQ or MRST or GRV [11]. Results presented in this work have been obtained with CTEQ(LO) distribution function. Combining eqs.(1),(2) and (3) we obtain the cross-section for resonance production per unit rapidity as,

\[
\frac{d\sigma}{dy}[h_A h_B \rightarrow RX] = F[nJ^{PC}] \int_{2m_{D/B}}^{2m_{D/B}} dM_{Q\bar{Q}}^2 \frac{H(x_a, x_b, M_{Q\bar{Q}}^2)}{s}. \tag{5}
\]

The above equation can be used to calculate the longitudinal momentum dependence of the quarkonium production cross section as shown in Ref. [10].

The leading order (LO) calculation does not have any \( p_T \) dependence as the heavy quark pairs are produced with \( p_T = 0 \) (assuming there is no \( p_T \) broadening of the initial state partons). In order to obtain the \( p_T \) distribution one has to go beyond LO. The dominant production mechanism of heavy quark pairs with large \( p_T \) and invariant mass near the threshold is the large \( p_T \) gluon splitting with the probability given by [5]

\[
\frac{d\text{Prob}}{dM_{Q\bar{Q}}^2} = \frac{\alpha_s}{6\pi} \frac{1}{M_{Q\bar{Q}}^2}, \tag{6}
\]

and

\[
\frac{d\sigma_{NN}^{NN}}{d^2p_T dy} = F[nJ^{PC}] \int_{2m_{Q\bar{Q}}}^{2m_{D/B}} dM_{Q\bar{Q}}^2 \frac{d\sigma}{d^2p_T dy} \frac{\alpha_s}{6\pi} \frac{1}{M_{Q\bar{Q}}^2}. \tag{7}
\]
where $N$ stands for nucleon, $d\sigma^g/d^2p_Tdy$ is the inclusive gluon $p_T$ and $y$ distribution in $NN$ collisions calculated in LO. It is given by $(a b \rightarrow g c)$

$$
\frac{d\sigma^g}{d^2p_Tdy} = \frac{1}{16\pi^2s^2} \sum_{a b} \int G_{a/N}(x_a, Q^2) G_{b/N}(x_b, Q^2) dy_4 \frac{\langle M \rangle^2}{x_a x_b}, \quad (8)
$$

where $x_a = p_T(e^y + e^{y_4})/\sqrt{s}$, $x_b = p_T(e^{-y} + e^{-y_4})/\sqrt{s}$, $\langle M \rangle^2$ is the matrix element for the process and $G_s$ are the parton distribution function. $p_T$ and $y$ are the transverse momentum and rapidity of the gluon which splits into $Q\bar{Q}$ pair and $y_4$ is the rapidity of the particle $c$. The processes that contribute to the heavy quark pair productions are $q\bar{q} \rightarrow gg$, $gg \rightarrow gg$ and $gq(\bar{q}) \rightarrow gq(\bar{q})$. Here the final state gluon splits into a heavy pair. Our analyses show that $F[J/\psi] = 0.045$ reproduces the CDF [12] data quite well (see later).

b. Nucleus-Nucleus collisions

The charmonium production, is observed to be suppressed both in $p-A$ and $A-B$ collisions (compared to the scaled $p-p$ scattering). This could be either due to nuclear absorption or due to the reduction of $Q-\bar{Q}$ interaction range in a QCD plasma caused by Debye screening. While the former is known as the normal nuclear suppression, the latter driven by the plasma effect is dubbed as anomalous suppression.

To understand the mechanism of anomalous suppression one introduces the concept of quarkonium formation time ($\tau_{0f}$) and the dissociation temperature $T_d$ determined from the condition at which the $Q-\bar{Q}$ interaction range becomes equal to the size of the quarkonium. The corresponding time when plasma attains a temperature $T = T_d$, is denoted as $\tau_d$. $\tau_{0f}$ on the other hand is the time required for the $Q-\bar{Q}$ pair to evolve into a physical charmonium state. This in the plasma rest frame would be Lorentz dilated and reads as $\tau_{\text{form}} \equiv \gamma \tau_{0f} = \tau_{0f} \sqrt{1 + p_T^2/m_R^2}$. High $p_T$ quarkonium states can evade suppression under two circumstances: (i) if $Q-\bar{Q}$ pair materializes into a bound state when the plasma has cooled down below $T_d$ or (ii) when the bound state is formed outside the plasma. The first condition, i.e. $\tau_{\text{form}} > \tau_d$ implies no suppression for $p_T > p_{1T}^{\text{crit}}$ while $p_T > p_{2T}^{\text{crit}}$ for no suppression stems from the condition, $|\vec{r} + \tau_{0f} \vec{p}_T/m_R| > R_s$. Here $R_s$ is the radius of the screening zone (see later) and $\vec{r}$ is the position where the heavy quark
pair is produced. Therefore anomalous suppression will be realized for quarkonium momenta less than \( \min \{ P_T^{\text{crit}}, P_T^{\text{crit}} \} \).

Next we consider \( J/\psi(\psi', \chi_c) \) production in \( p-A \) and \( A-B \) collisions. To this end, we first briefly mention the necessary formulae in Glauber model [13, 14]. The total inelastic cross-section in \( A-B \) collisions at an impact parameter \( b \) is given by

\[
\frac{d\sigma^{AB}_{\text{in}}}{db} = 1 - \left[ 1 - T_{AB}(b) \sigma^{NN}_{\text{in}} \right]^{AB} \equiv 1 - P_0(b),
\]

where \( T_{AB}(b) \) is the nuclear overlap function given by

\[
T_{AB}(b) = \int ds T_A(b) T_B(b - s)
\]

The nuclear thickness functions are normalized to unity, i.e. \( \int db T_A(b) = \int db T_B(b) = 1 \), where \( \rho_A \) is the nuclear density distribution.

Generally we are interested in the cross-sections for a set of events in a given centrality range defined by the trigger settings. Centrality selection corresponds to a cut on the impact parameter, \( b \), of the collisions. The sample of events in a given centrality range \( 0 \leq b \leq b_m \), contains a fraction of the total inelastic cross-section. This fraction is defined by [15]

\[
f(b_m) = \frac{\int_0^{b_m} db \int \sigma^{AB}_{\text{in}} \rho_A(b, z) \exp\left[-(A-1) \int_0^\infty \sigma_{\text{abs}} \rho_A(b, z') dz' \right] \sigma^{NN}_{J/\psi}}{\int_0^\infty db \int \sigma^{AB}_{\text{in}} \rho_A(b, z) \exp\left[-(A-1) \int_0^\infty \sigma_{\text{abs}} \rho_A(b, z') dz' \right] \sigma^{NN}_{J/\psi}}.
\]

Now we discuss the \( J/\psi \) survival probability when, after production, it propagates through the target/projectile nucleus. After creation, the \( J/\psi \) meson can interact with other nucleons in the target and the projectile and may get destroyed mainly due to \( J/\psi - N \) interactions. The cross-section for \( J/\psi \) production in \( p-A \) collisions can be written as [16]

\[
\sigma_{J/\psi}^{pA}(b_m) = A \int_0^{b_m} db d\rho_A(b, z) \exp\left[-(A-1) \int_0^\infty \sigma_{\text{abs}} \rho_A(b, z') dz' \right] \sigma^{NN}_{J/\psi},
\]

where \( \sigma^{NN}_{J/\psi} \) is obtained from eq.(5). The interpretation of the above equation is as follows. The resonance is formed at \( \vec{r} = (b, z) \) where the density of the target nucleus is \( \rho_A(\vec{r}) \). It can travel in forward direction (z) at constant impact parameter and its intensity is attenuated due to \( J/\psi - N \) inelastic collisions. The exponential factor accounts for this attenuation.
The generalization of eq.(12) in nucleus-nucleus collisions is straightforward. The \( J/\psi \) production cross-section in \( A - B \) collisions at an impact parameter \( b \) can be written as

\[
\frac{d\sigma_{J/\psi}^{AB}}{d^2b dp_T}(b) = \frac{d\sigma_{J/\psi}^{NN}}{dp_T} AB \int ds dz_1 dz_2 \rho_A(s, z_1) \rho_B(b - s, z_2) \\
\times \exp \left( -(A - 1) \int_{z_1}^{\infty} \sigma_{\text{abs}} \rho_A(s, z') dz' \right) \\
\times \exp \left( -(B - 1) \int_{z_2}^{\infty} \sigma_{\text{abs}} \rho_B(b - s, z') dz' \right)
\]

(13)

The observed charmonium suppression data at SPS energies [17] could be explained either by the Debye screening in a QGP or by co-mover scattering due to hadronic matter in the initial state [16, 18]. But the co-mover interpretation of heavy quarkonium suppression at LHC energies does not seem to be plausible, as the initial temperatures achieved at these energies could be considerably larger than the transition temperature, \( T_c \). Thus, the heavy quarkonium will encounter the hadronic comover at a much later time. It should be noted here that the formation of QGP does not necessarily guarantee the suppression of the heavy resonances. The temperature of the system must be greater than the dissociation temperature of the particular resonance in order to get suppressed. The dissociation temperature has recently been calculated using lattice QCD [19, 20]. The most recent values of \( T_d \) for different heavy flavour resonances are shown in Table I. Thus as long as the plasma temperature remains greater than \( T_d \) bound state cannot be formed. The time \( \tau_d \) is obtained by Bjorken scaling law [21]:

\[
\tau_d(b) = \tau_i \left( \frac{T_i(b)}{T_d} \right)^3,
\]

(14)

where \( \tau_i \) and \( T_i(b) \) (defined later) are the plasma formation time and initial temperature of the plasma respectively. In order to obtain the centrality dependence of the suppression we estimate the initial temperature by assuming the isentropic expansion of the system, namely,
Table I: Dissociation temperatures of charmonium and bottomonium system in pure gluonic plasma.

| Resonance | $T_d/T_c$ [19] | $T_d/T_c$ [19] | $\langle T_d \rangle/T_c$ |
|-----------|----------------|----------------|-------------------------|
| $J/\psi$  | 1.1            | 2.0            | 1.5                     |
| $\chi_c$  | 0.74           | 1.1            | 0.9                     |
| $\psi'$   | 0.1 − 0.2      | 1.1            | 0.625                   |
| $\Upsilon$| 2.31           | 4.5            | 3.4                     |
| $\chi_b$  | 1.13           | 2.0            | 1.55                    |
| $\Upsilon'$| 1.1            | 2.0            | 1.55                    |
| $\chi_b'$ | 0.83           | −              | −                       |
| $\Upsilon''$| 0.74           | −              | −                       |

$T_i^3(b) = \frac{2\pi^4}{45\zeta(3)} \frac{1}{4a_i\pi R_t^2(b)\tau_i} \frac{dN_{AB}}{dy}(b).$  \hspace{1cm} (15)

$a_i = \pi^2/90(21N_F/2 + 16),$ is the degeneracy in the initial state, $dN_{AB}/dy$ is the total multiplicity of produced hadrons measured experimentally and $R_t$ is the transverse dimension of the plasma.. One would expect that the total multiplicity in nucleus-nucleus collisions at RHIC and LHC energies will come from both soft as well as hard collisions. It has been shown recently [22] that about 10% of the total multiplicity at RHIC energies comes from hard collisions. Therefore, the total hadron multiplicity in nucleus-nucleus collisions at a given impact parameter can be written as (the so-called two-component model)

$$\frac{dN_{AB}}{dy}(b) = \left[ (1 - x) \frac{N_{\text{part}}(b)}{2} + x N_{\text{coll}}(b) \right] \frac{dN_{NN}}{dy},$$ \hspace{1cm} (16)

In the above $x$ is the fraction of hard collisions. For hadron rapidity density $(dN_{NN}/dy)$ in nucleon-nucleon collision we use the following form [23]: $dN_{NN}/dy = 1.5(2.5 - 0.25 \ln(s) + 0.023 \ln^2(s)).$ $N_{\text{part}}(b)$ and $N_{\text{coll}}(b)$ are the numbers of participants and number of collisions respectively at an impact parameter $b$. In non-central collisions the plasma transverse radius can be approximated as $R_t \sim 1.2[0.5N_{\text{part}}(b)]^{1/3}$[24].

The average initial temperatures and the dissociation times for various centrality classes have been shown in Table II.

The critical radius $R_s$, beyond which there will be no suppression, is given by [25]

$$R_s = R_t \left[ 1 - \left( \frac{\tau_0}{\tau_d} \right)^4 \right]^{1/2},$$ \hspace{1cm} (17)
Table II: Initial temperatures, transverse radii, and dissociation times at different centralities for $J/\psi$ with $T_d \sim 1.1T_c$ and number of flavour, $N_F = 0$. All the averages here are over the impact parameter.

The survival probability, which is the ratio of bound states produced by the escaped pairs relative to the number of bound states that would be formed in the absence of QGP, is given by [25]

$$S_Q(p_T, b) = \frac{3}{\pi R_t^2} \phi_{\text{max}} \left[ 1 - \left( \frac{r}{R_t} \right)^2 \right]^{1/2} r \, dr,$$

where

$$\phi_{\text{max}} = \begin{cases} \pi & z \leq -1, \\ \cos^{-1}|z| & -1 \leq z \leq 1, \\ 0 & z \geq 1 \end{cases}$$

with $z = [(R_s^2 - r^2)m_R - \tau_{0f}^2/m_R]/(2r\tau_{0f}p_T)$. We assume a pure gluonic plasma with $T_c \sim 270$ MeV [26].

Suppression due to Debye screening continues till the temperature of the QGP drops below $T_d$. The QGP starts hadronizing at $T_c$ (equivalently, at $\tau_Q = T_c^3/\tau_i$, the time when phase transition starts). During the time interval $\tau_Q - \tau_d$, the remaining resonances will not be dissociated as long as $T_c$ is lower than $T_d$. During the mixed phase interval ($\tau_H - \tau_Q$) QGP part does not contribute to the suppression for $T_c < T_d$. Here $\tau_H = g_Q\tau_Q/g_H$, where $g_Q$ and $g_H$ are the statistical degeneracy for QGP and hadronic phases respectively. However, the co-moving absorption starts at $\tau_Q$ and the density of the co-moving hadrons is determined by the temperature $T_c$. Following Ref. [27] we have checked that this effect is very small in the present scenario. Hence we do not consider suppression due to co-mover absorption here.

Finally, the $p_T$-distribution of heavy resonances is given by

$$\frac{dN_{J/\psi}}{dp_T} = \frac{d\sigma_{J/\psi}^{NN}}{dp_T} \int_0^{R_t} d^2b S_Q(p_T, b) ABT_{AB}(b) \int_0^{\tau_{0f}} d^2b [1 - P_0(b)].$$

(20)
However, with nuclear absorption the above equation will be modified to:

\[
\frac{dN_{J/\psi}}{dp_T} = \int_0^{b_m} \frac{d\sigma_{J/\psi}}{db d^2p_T} d^2b S_Q(p_T, b) \frac{1 - P_0(b)}{\int_0^{b_m} d^2b [1 - P_0(b)]},
\]

(21)

Similar expression can be obtained for \(\psi'\) and \(\chi_c\).

### III. Results

In Fig. (1) we compare our results with CDF data where \(m_c = 1.3\) GeV, \(Q^2 = m_c^2 + p_T^2\) have been used. It is clear that the data is reproduced reasonably well with the values of \(F\) given in Table III. Any effect due to \(k_T\)-broadening is ignored here. However, it might be mentioned that the low \(p_T\) data on \(\Upsilon\) production could be explained in color evaporation model by assuming substantial \(k_T\) smearing [7].

| \(F[J/\psi]\) | \(F[\psi']\) | \(\sum J B(\chi_cJ \rightarrow J/\psi X) F[\chi_cJ]\) |
|---|---|---|
| 0.045 | 0.0126 | 0.005 |

Table III: Values of the universal factor used in the calculation (no feed down).

Next we consider the \(J/\psi\) production at LHC energies. First of all, in Fig. (2), we show \(d\sigma/dp_T\) of \(J/\psi\) and \(\psi'\) from \(p - p\) collisions at \(\sqrt{s} = 5500\) GeV. Here the
Figure 2: Differential cross-sections of $J/\psi$ and $\psi'$ in $p - p$ collisions at $\sqrt{s} = 5500$ GeV. Only direct contributions are shown.

yields correspond to the direct productions i.e. feed down from higher resonance states have been ignored with $m_c = 1.3$ GeV. The feed down contribution can be estimated by multiplying the yield with appropriate branching ratio.

The differential $p_T$ distributions of $J/\psi$ and $\psi'$ from $p - p$ scattering (as shown in Fig. 2) provide the baseline for $Pb - Pb$ collisions at LHC energies. To obtain charmonium $p_T$ distributions at LHC energies we use eq.(20) which gives $p_T$ spectra without nuclear and comover absorptions. The amount of nuclear absorption can be calculated using $\sigma_{abs} = \sigma_{abs}(\sqrt{s_0})(s/s_0)^{\Delta/2}$, where $s_0 = 17.3$ GeV, $\sigma_{abs}(s_0) = 5 \pm 0.5$ mb and $\Delta = 0.125$ [28]. The nuclear absorption of heavy resonances at SPS energies is well studied and at LHC energies it can be eliminated using $p - A$ collisions with the $\sqrt{s}$-dependent cross-section mentioned above. In the present work, our motivation is to see the anomalous suppression due to Debye screening, we do not include nuclear suppression in our calculation. We also note here that the co-moving absorption will start when the plasma begins to hadronize. We have seen that the co-moving absorption is minimal with the parameters used in Ref. [27]. The extrapolation of these parameters to high energies is not straightforward. Therefore, in our calculation we do not include the hadronic absorptions. The expected $p_T$
distributions of $J/\psi$ at LHC energies are shown in Fig. (3) for $0 \leq b \leq 3$ fm and $6 \leq b \leq 9$ fm centralities with and without Debye screening. For $T_d \sim 1.1 T_c$ we see that the $J/\psi$s with $p_T < 15(10)$ GeV are suppressed for $0 \leq b \leq 3(6 \leq b \leq 9)$ fm centrality. However, the other value of $T_d$ (see Table I) the $J/\psi$s will not be suppressed for the same value of the hard fraction $x$. We take $T_d = 1.5 T_c$ in Fig. (4). It is seen that suppression occurs below $p_T \sim 6$ GeV. The value of the hard fraction $x$ is not known at LHC energies at this stage. The survival probability decreases as $x$ increases, since in that case, the initial temperature will be high (see eqs.(15) and (16)) as demonstrated in fig. (5).

For $\psi'$ the $T_d$ values are small compared to that of $J/\psi$ and hence $\psi'$ will be suppressed substantially. This is shown in Fig. (6) with $T_d \sim 0.15 T_c$ at two different centralities. It shows that even a very high $p_T (\sim 30$ GeV) $\psi'$ is suppressed. We also plot the $p_T$ spectra of $\psi'$ for average values of $T_d$ at three centralities in Fig. (7) to show the centrality dependence of heavy resonance suppression. We observe that the suppression is small in going from most central to peripheral collisions as expected.
Figure 4: Same as Fig. (3) for $T_d = 1.5T_c$.

Figure 5: Same as Fig. (3) for $T_d = 1.1T_c$ for different values of the hard fraction with impact parameter in the range $0 \leq b \leq 3$ fm.
Figure 6: Same as Fig. (3) for $\psi'$ with $T_d = 0.15T_c$.

Figure 7: Same as Fig. (6) with $T_d = 0.63T_c$. 
Figure 8: Same as Fig. (3) at $\sqrt{s} = 200$ GeV, $\tau_i = 0.6$ fm/c, $x = 0.1$ and $0 \leq b \leq 3$ fm.

The prediction at RHIC energies is shown in Fig. (8). It is seen that $J/\psi$ is not suppressed even for most central collisions as the initial temperature is very close to the dissociation temperature. However, $\psi'$ is suppressed and the critical values of $p_T$ depend on the the dissociation temperature used.

IV. Summary and discussions

In this work we have calculated the $p_T$ distribution of charmonium by using colour evaporation model (LO with gluon splitting) for RHIC and LHC energies. To validate our formalism we first reproduce the CDF data for $p_T$ distribution of $J/\psi, \psi'$, and $\chi_c$ using CEM. It is assumed that at LHC energies QGP is formed and the heavy resonances are suppressed in a QGP due to Debye screening. We have calculated the survival probability using the prescription of Ref. [25] and convoluted it with the $p_T$ distribution of various resonances. First, the $p_T$ spectra of various charmonium states are obtained in nucleon-nucleon collision and then these are convoluted with Glauber model to get the same in nucleus-nucleus collisions. The centrality dependence of transverse momentum distribution is also estimated. We have seen that $J/\psi$ with $p_T > 15(10)$ GeV will not be suppressed in nuclear collisions for the
impact parameter in the range $0 \leq b \leq 3$ fm ($6 \leq b \leq 9$) fm. We do not include the co-moving suppression due to hadrons, as this is found to be small within the existing model parameters [27].

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