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J/ψ gluonic dissociation revisited: I. Fugacity, flux and formation time effects

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Abstract. We revisit the standard treatment due to Xu et al. of J/ψ suppression due to gluonic bombardment in an equilibrating quark–gluon plasma. Effects arising from gluon fugacity, relative g–ψ flux, and ψ meson formation time are explicitly incorporated in the formulation of the gluon number density, velocity-weighted cross section, and the survival probability. Our new formulae are applied to a numerical study of the pattern of J/ψ suppression in the central rapidity region at RHIC/LHC energies. The temperature and transverse momentum dependence of our graphs have noticeable differences from those of Xu et al.

1 Introduction

Relativistic heavy-ion collision experiments at CERN SPS/LHC and BNL/RHIC are believed to have led to a phase transition from the hadronic world into a deconfined and/or chirally symmetric state of free quarks and gluons, the so called quark–gluon plasma (QGP) [1–6]. However, as of now, no conclusive evidence of QGP formation has been discerned. Among the most hotly debated theoretically proposed signatures in this context are the erstwhile J/ψ suppression due to the medium's influence [6] and recent J/ψ enhancement via dynamical regeneration [7]. The well known mechanisms responsible for J/ψ suppression are summarized briefly in the appendix for the sake of completeness.

Attention in the sequel will be focused on the break-up of the J/ψ owing to bombardment with energetic gluons [8]. For this mechanism, Xu et al. [2] employed statistical mechanics coupled with phenomenological QCD to calculate the J/ψ survival probability in a temporally evolving parton gas. It is well recognized that to evaluate the magnitude of quarkonium suppression in nuclear collisions one has to use a number of model assumptions, and their validity must be continually challenged or explored further. It is only in this modest spirit that the present paper aims at extending or modifying the work of Xu et al. in the following various respects.

(i) Gluon fugacity effect: Xu et al. ((7) in [2]) used a *fudge* factor λ_g^{Xu} to account for the deviation of the initial gluon distribution (as computed in the mini-jet model) from the thermal one. It is natural to ask what would happen if a true fugacity λ_g (stemming e.g. from a chemical potential) were used in the *conventional* form of the Bose–Einstein distribution. In Sect. 2 below we derive a new formula for the gluon number density n_g valid for general $\lambda_g \le 1$, appropriate for a system not yet in chemical equilibrium.

(ii) Relative flux effect: Xu et al. ((8) in [2]) were interested in the invariant product $\Gamma = v_{\mathrm{rel}}\sigma$ where the gluon–ψ break-up cross section σ was written in the ψ meson rest frame, but unfortunately their relative flux $v_{\mathrm{rel}}^{\mathrm{Xu}}$ was evaluated in the fireball frame so that their expression $v_{\mathrm{rel}}^{\mathrm{Xu}} = 1 - (\mathbf{k}\cdot\mathbf{p}_{\mathrm{T}})/k^0 M_{\mathrm{T}}$ remains *different* from unity, i.e., c. In Sect. 3 we modify this procedure by treating the product Γ strictly in the ψ rest frame so that the corresponding relative flux $v_{\mathrm{rel}}^{\mathrm{Our}}$ remains $c = 1$ always.

(iii) Formation time effect: For computing their survival probability $S(p_{\mathrm{T}})$ Xu et al. ((14) in [2]) used an integration over τ_ψ having the lower limit 0, where τ_ψ is the proper time measured in the J/ψ rest frame. This is inconvenient because the gluon density $n_g(t)$ and the thermally-averaged cross section $\langle v_{\mathrm{rel}}\sigma\rangle$ are natural functions of the usual time t in the fireball rest frame. In Sect. 4 we write a modified expression for $S(p_{\mathrm{T}})$ using a t integration where formation times of the QGP as well as J/ψ are explicitly included. Of course, in our numerical results in Sects. 2, 3 and 4 the actual velocity profiles of hydrodynamic flow are ignored. Finally, our main conclusions appear in Sect. 5.

2 Number density

2.1 Preliminaries

Assuming thermal equilibrium and working in the fireball rest frame let the symbol T denote the absolute temperature, K the gluon four momentum, 16 the spin-color degeneracy factor, $\lambda_g \le 1$ the gluon fugacity, and

Table 1. Initial values for the temperature, time, fugacities etc. at RHIC(1), LHC(1) only [5]

	T (GeV)	$t_i = \tau_0$ (fm)	λ_g	λ_q	n_g^{Xu} (fm)$^{-3}$	n_g^{Our} (fm)$^{-3}$
RHIC(1)	0.55	0.70	0.05	0.008	2.11	1.76
LHC(1)	0.82	0.5	0.124	0.02	17.34	14.66

$f = f(K^0, T, \lambda_g)$ the one-body gluon distribution function. Then the gluon number density n_g is obtained from

$$n_g = 16 \int \frac{\mathrm{d}^3 K}{(2\pi)^3} f = \frac{8}{\pi^2} \int_0^\infty \mathrm{d}K^0 {K^0}^2 f. \qquad (1)$$

2.2 Xu procedure

It is known that in the early stage of the evolution the quark–gluon plasma is not in chemical equilibrium. Hence Xu et al. ((7) in [2]) employed a factorized gluon distribution function,

$$f^{\mathrm{Xu}} = \frac{\lambda_g^{\mathrm{Xu}}}{\mathrm{e}^{K^0/T} - 1} = \lambda_g^{\mathrm{Xu}} \sum_{n=1}^{\infty} \mathrm{e}^{-nK^0/T}, \qquad (2)$$

where λ_g^{Xu} is a fudge factor accounting for the deviation of the initial distribution (as computed in the mini-jet model) from the thermal one. This led them to a number density linearly depending on λ_g^{Xu} through

$$n_g^{\mathrm{Xu}} = \frac{16}{\pi^2} T^3 \lambda_g^{\mathrm{Xu}} \zeta(3). \qquad (3)$$

2.3 Our proposal

We *do not* contest the validity of the above view. Rather we ask the following independent question: What would happen if one employs the standard Bose–Einstein distribution

$$f^{\mathrm{Our}} = \frac{\lambda_g}{\mathrm{e}^{K^0/T} - \lambda_g} = \sum_{n=1}^{\infty} \lambda_g^n \mathrm{e}^{-nK^0/T}, \qquad (4)$$

where λ_g is the true fugacity stemming e.g. from a chemical potential? The expression (4) is certainly of interest since it is familiar from text books of statistical mechanics. Moreover, (4) has the advantage that the denominator $\mathrm{e}^{K^0/T} - \lambda_g \geq 1 - \lambda_g > 0$ as long as $\lambda_g < 1$, implying that any possibility of Bose condensation is avoided. Insertion of (4) into (1) guides us to a number density containing the fugacity in a power series via

$$n_g^{\mathrm{Our}} = \frac{16}{\pi^2} T^3 \sum_{n=1}^{\infty} \frac{\lambda_g^n}{n^3}. \qquad (5)$$

Remembering that a $1/n^3$ type series converges rapidly with n, numerical comparison of (3) and (5) is easily done via the ratio

$$\frac{n_g^{\mathrm{Xu}}}{n_g^{\mathrm{Our}}} \sim \left(1 + \frac{1}{8}\right) \Big/ \left(1 + \frac{\lambda_g^2}{8}\right), \qquad (6)$$

where we have set $\lambda_g^{\mathrm{Xu}} \approx \lambda_g$ in the leading order. Of course, the algebraic reason for the inequality $n_g^{\mathrm{our}} < n_g^{\mathrm{Xu}}$ is the fact that the distribution function $f^{\mathrm{Our}} < f^{\mathrm{Xu}}$ as long as $\lambda_g < 1$. In other words, the effect of actual fugacity (before chemical equilibration) is to *reduce* the number density of gluons below the value of Xu et al.

2.4 Initial conditions

It is well recognized that the scenario resulting from relativistic heavy-ion collisions is rapidly time-dependent. Quick scattering among the partons drives the QGP to thermal equilibrium in the fireball rest frame within a time $t_i = \tau_0 \sim 1/\Lambda \sim 0.5\,\mathrm{fm}/c$ where the suffix i stands for "initial" and Λ is the QCD energy scale. The initial conditions predicted by the HIJING Monte Carlo simulation are summarized in Table 1. There gluon densities computed via the Xu procedure (3) and our proposal (5) are also listed. Clearly the *relative* difference between n_{gi}^{Xu} and n_{gi}^{Our} is of the order of $1/8 \sim 12\%$, which is significant.

2.5 Temporal evolution

The thermally equilibrated QGP produced at the instant $t_i = \tau_0$ undergoes rapid expansion (accompanied with cooling) while partonic reactions tend to drive the plasma towards chemical equilibrium. In Bjorken's boost-invariant longitudinal expansion scenario the fugacities and temperature are known [16] to evolve through the following master rate equations:

$$\begin{aligned}
&\frac{\dot{\lambda}_g}{\lambda_g} + 3\frac{\dot{T}}{T} + \frac{1}{\tau} = R_3 \left(1 - \lambda_g\right) - 2R_2 \left(1 - \frac{\lambda_g^2}{\lambda_q^2}\right), \\
&\frac{\dot{\lambda}_q}{\lambda_q} + 3\frac{\dot{T}}{T} + \frac{1}{\tau} = R_2 \frac{a_1}{b_1} \left(\frac{\lambda_g}{\lambda_q} - \frac{\lambda_q}{\lambda_g}\right), \\
&\left(\lambda_g + \frac{b_2}{a_2}\lambda_q\right)^{3/4} T^3 \tau = \mathrm{const.}
\end{aligned} \qquad (7)$$

Here τ is the medium proper time, λ_q the quark fugacity, N_f the number of flavours, and the remaining symbols are defined by

$$R_2 = 0.5 n_g \langle v\sigma_{gg\longrightarrow q\bar{q}} \rangle, \quad R_3 = 0.5 n_g \langle v\sigma_{gg\longrightarrow ggg} \rangle,$$

$$a_1 = 16\zeta(3)/\pi^2, \quad a_2 = 8\pi^2/15,$$
$$b_1 = 9\zeta(3)N_f/\pi^2, \quad b_2 = 7\pi^2 N_f/20. \qquad (8)$$

Their solutions on the computer yield the functions $T(t)$, $\lambda_g(t)$, $n_g(t)$ in terms of the fireball time t. The lifetime (or freeze-out time) t_{life} of the plasma is the instant when the temperature drops to $T(t_{\text{life}}) = 200\,\text{MeV}$, say.

3 Flux-weighted rate

3.1 Preliminaries

Next, the question of applying statistical mechanics to gluonic break-up of the J/ψ becomes relevant. In the fireball frame consider a ψ meson of mass m_ψ, four momentum p_ψ, three velocity $\mathbf{v}_\psi = \mathbf{p}_\psi/p^0_\psi$, and dilation factor $\gamma_\psi = p^0_\psi/m_\psi$. If q is the gluon four momentum measured in ψ meson rest frame, then by Lorentz transformations

$$K^0 = \gamma_\psi(q^0 + \mathbf{v}_\psi \cdot \mathbf{q}) = \gamma_\psi q^0(1+ \mid \mathbf{v}_\psi \mid \cos\theta_{q\psi})$$
$$\mathrm{d}^3K = (K^0/q^0)\ \mathrm{d}^3q, \qquad (9)$$

where $\theta_{q\psi}$ is the angle between the $\hat{q}$ and $\hat{v}_\psi$ unit vectors.

The invariant quantum mechanical dissociation rate for the g–ψ collision can be written compactly as

$$\Gamma = v_{\text{rel}}\ \sigma, \qquad (10)$$

where v_{rel} is the relative flux and σ the cross section written in any chosen frame. Its thermal average over gluon momentum in the fireball frame reads

$$\langle\Gamma\rangle = \frac{16}{n_g}\int \frac{\mathrm{d}^3q}{(2\pi)^3}\frac{K^0}{q^0}\ \Gamma\ f, \qquad (11)$$

with f being the distribution function already encountered in (1).

3.2 Xu procedure

Xu et al. ((8) in [2]) worked with their relative flux in the fireball frame, viz.

$$v^{Xu}_{\text{rel}} = q^0/K^0\gamma_\psi = 1 - (\mathbf{K}\cdot\mathbf{p}_\psi)/K^0 p^0_\psi, \qquad (12)$$

which, of course, is different from unity. At the same time their dissociation cross section σ_{Rest} was written in the ψ meson rest frame based on the standard QCD value [20]

$$\sigma_{\text{Rest}} = B(Q^0-1)^{3/2}/{Q^0}^5, \quad q^0 > \epsilon_\psi,$$
$$Q^0 = \frac{q^0}{\epsilon_\psi}, \quad B = \frac{2\pi}{3}\left(\frac{32}{3}\right)^2 \frac{1}{m_c(\epsilon_\psi m_c)^{1/2}}, \qquad (13)$$

where ϵ_ψ is the J/ψ binding energy and m_c the charm quark mass. Insertion into (11) led them to

$$\langle\Gamma^{\text{Xu}}\rangle = \frac{16}{n_g^{\text{Xu}}}\int \frac{\mathrm{d}^3q}{(2\pi)^3}\frac{1}{\gamma_\psi}\sigma_{\text{Rest}}\ \lambda_g \sum_{n=1}^{\infty} \mathrm{e}^{-nK^0/T}, \qquad (14)$$

where the approximate f^{Xu} given by (2) has been recalled. The simple angular integration over $\mathrm{d}\cos\theta_{q\psi}$ can be done by taking the polar axis along $\hat{v}_\psi$ to yield

$$\langle\Gamma^{\text{Xu}}\rangle = \frac{8\epsilon_\psi^3\lambda_g}{\pi^2\gamma_\psi n_g^{\text{Xu}}}\sum_{n=1}^{\infty}\int_1^{\infty}\mathrm{d}Q^0 {Q^0}^2 \times \sigma_{\text{Rest}}\mathrm{e}^{-C_nQ^0}\left(\frac{\sinh D_nQ^0}{D_nQ^0}\right) \qquad (15)$$
$$= \frac{4\epsilon_\psi^2\lambda_g T}{\pi^2\gamma_\psi^2 \mid \mathbf{v}_\psi \mid n_g^{\text{Xu}}}\sum_{n=1}^{\infty}\frac{1}{n}\int_1^{\infty}\mathrm{d}Q^0\ Q^0 \times \sigma_{\text{Rest}}\left(\mathrm{e}^{-A_n{}^-Q^0} - \mathrm{e}^{-A_n{}^+Q^0}\right). \qquad (16)$$

Here the following abbreviations have been introduced:

$$\gamma_\psi = p^0_\psi/m_\psi, \quad \mathbf{v}_\psi = \mathbf{p}_\psi/p^0_\psi,$$
$$C_n = n\epsilon_\psi\gamma_\psi/T, \quad D_n = \mid \mathbf{v}_\psi \mid C_n,$$
$$A_n^\pm = C_n \pm D_n = C_n\,(1\pm \mid \mathbf{v}_\psi \mid). \qquad (17)$$

For J/ψ produced in the central rapidity region, Xu et al. have drawn elaborate curves showing the dependence of $\langle\Gamma^{\text{Xu}}\rangle$ on T and p_{T}.

3.3 Our proposal

We set up our Γ entirely in the ψ meson rest frame where $v^{\text{Our}}_{\text{rel}} = c = 1$. This is because when a *massless* gluon is incident upon a ψ meson at rest, the corresponding relative flux must become unity in sharp contrast to the value due to Xu et al., (12). Thereby (11) becomes

$$\langle\Gamma^{\text{Our}}\rangle = \frac{16}{n_g^{\text{Our}}}\int\frac{\mathrm{d}^3q}{(2\pi)^3}\gamma_\psi\,(1+ \mid \mathbf{v}_\psi \mid \cos\theta_{q\psi}) \times \sigma_{\text{Rest}}\sum_{n=1}^{\infty}\lambda_g^n\mathrm{e}^{-nK^0/T}, \qquad (18)$$

where the exact f^{Our} given by (4) has been recalled. The slightly complicated angular integration over $\mathrm{d}\cos\theta_{q\psi}$ can be performed by choosing the polar axis along $\hat{v}_\psi$, yielding

$$\langle\Gamma^{\text{Our}}\rangle = \frac{8\epsilon_\psi^3\gamma_\psi}{\pi^2 n_g^{\text{Our}}}\sum_{n=1}^{\infty}{\lambda_g}^n\int_1^{\infty}\mathrm{d}Q^0\ {Q^0}^2 \times \sigma_{\text{Rest}}\mathrm{e}^{-C_nQ^0}\left[I_0(\rho_n) - \mid \mathbf{v}_\psi \mid I_1(\rho_n)\right]$$
$$= \frac{4\epsilon_\psi^2 T}{\pi^2 \mid \mathbf{v}_\psi \mid n_g^{\text{Our}}}\sum_{n=1}^{\infty}\frac{\lambda_g^n}{n}\int_1^{\infty}\mathrm{d}Q^0\ Q^0 \times \sigma_{\text{Rest}}\left[\left(1- \mid \mathbf{v}_\psi \mid \left(1-\frac{1}{\rho_n}\right)\right)\mathrm{e}^{-A_n^-Q^0} - \left(1+ \mid \mathbf{v}_\psi \mid \left(1+\frac{1}{\rho_n}\right)\right)\mathrm{e}^{-A_n^+Q^0}\right], \qquad (19)$$

where

$$\rho_n = D_nQ^0 = D_nq^0/\epsilon_\psi, \quad I_0(\rho_n) = (\sinh\rho_n)/\rho_n,$$
$$I_1(\rho_n) = (\cosh\rho_n)/\rho_n - (\sinh\rho_n)/\rho_n^2. \qquad (20)$$

Clearly the dependence of $\langle\Gamma^{\text{Our}}\rangle$ on λ_g and γ_ψ is more involved than that of $\langle\Gamma^{\text{Xu}}\rangle$ given by (15) and (16).

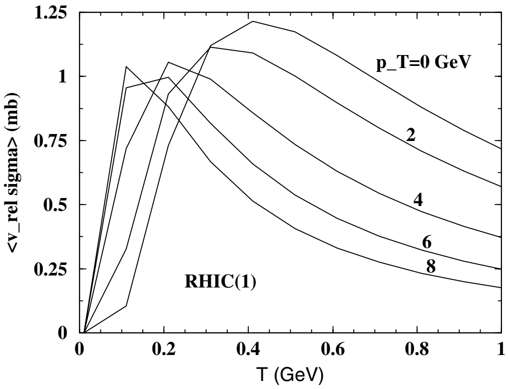

Fig. 1. The thermally-averaged gluon–J/ψ dissociation cross section $\langle v_{\rm rel}\sigma\rangle$ as a function of temperature at different transverse momenta $p_{\rm T}$ as done by Xu et al. ((16) in [2]). The initial gluon fugacity is given in Table 1 at RHIC energy

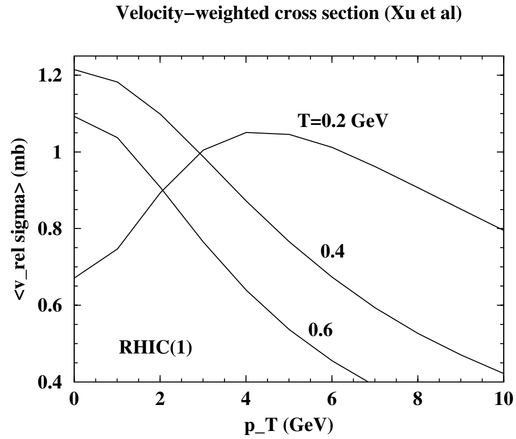

Fig. 3. The thermally-averaged gluon–J/ψ dissociation cross section $\langle v_{\rm rel}\sigma\rangle$ as a function of transverse momentum at different temperatures as done by Xu et al. [2]

3.4 Numerical work

The initial thermally-averaged rates $\langle\Gamma^{\rm Xu}\rangle$ (16) and $\langle\Gamma^{\rm Our}\rangle$ (19) are depicted in Figs. 1, 3 and Figs. 2, 4, respectively. The physics of the dependence of the peak on the temperature and transverse momentum has already been discussed in [2]. Here we wish to focus attention only on the striking *similarity* between Figs. 1 and 2 in spite of the different fugacities and fluxes employed. For this purpose, we first go back to the Lorentz transformation (9) and observe that the relative flux due to Xu et al. receives a dominant contribution from the antiparallel ($\cos\theta_{q\psi}=-1$) configuration. Indeed, then

$$v_{\rm rel}^{\rm Xu}\equiv\frac{q^0}{\gamma_\psi K^0}\sim\frac{1}{\gamma_\psi^2(1-\mid{\bf v}_\psi\mid)}\sim 1+\mid{\bf v}_\psi\mid\,. \tag{21}$$

Now let us consider the ratio

$$\frac{\langle\Gamma^{\rm Xu}\rangle}{\langle\Gamma^{\rm Our}\rangle}=\left[\frac{n_g^{\rm Our}}{n_g^{\rm Xu}}\right]\left[\frac{\text{phase space integral of } {\rm v}_{\rm rel}^{\rm Xu}\ \sigma_{\rm Rest}}{\text{phase space integral of c } \sigma_{\rm Rest}}\right]. \tag{22}$$

Due to the fugacity effect the number density ratio in (22) is somewhat smaller than unity as already mentioned in Sect. 2. On the other hand, in a near antiparallel configuration, the relative flux $v_{\rm rel}^{\rm Xu}>c$, hence the ratio of the phase space integrals, is somewhat larger than unity. These two effects tend to partially compensate each other in (22) so that the relative difference between the curves of Figs. 1 and 2 is not more than about 5–6%. However, the influence of $1+\mid{\bf v}_\psi\mid$ becomes more pronounced at high transverse momentum, causing a noticeable difference between the curves of Figs. 3 and 4.

4 Survival probability

4.1 Preliminaries

Consider a cylindrical coordinate system in the fireball frame where the ψ meson was created at the time-space point $(t_{\rm I},r_\psi^{\rm I},\phi_\psi^{\rm I})$ with transverse velocity ${\bf v}_\psi{}^{\rm T}$. The

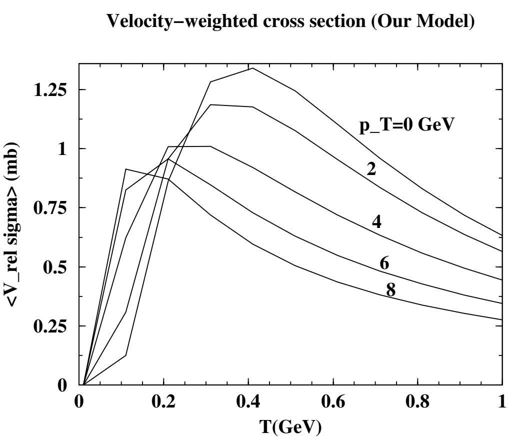

Fig. 2. Same as in Fig. 1, but these curves are obtained by our (19)

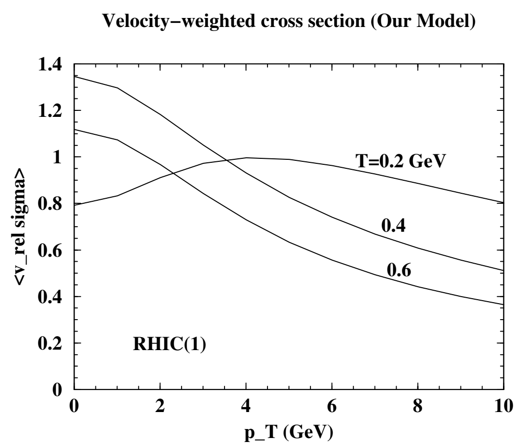

Fig. 4. Same as in Fig. 3 except that the results are obtained by us

plasma is supposed to be contained within a cylinder of radius R, expanding longitudinally till the end of its lifetime $t_{\rm life}$. The ψ meson's trajectory will hit the said cylinder after covering a distance d_{RI} in the time interval t_{RI} such that

$$d_{RI} = -r_\psi^{\rm I} \cos\phi_\psi^{\rm I} + \sqrt{R^2 - {r_\psi^{\rm I}}^2 \sin^2\phi_\psi^{\rm I}},$$
$$t_{RI} = d_{RI} / \mid {\bf v}_\psi{}^{\rm T} \mid \quad , \tag{23}$$

and the full temporal range of interest is obviously

$$t_{\rm I} \le t \le t_{\rm II} \ ; \ t_{\rm II} = \min\left(t_{\rm I} + t_{RI}, \quad t_{\rm life}\right). \tag{24}$$

The corresponding survival probability of J/ψ averaged over its initial production configuration extending over the transverse area A becomes

$$S(p_{\rm T}) = \int_A {\rm d}^2 r_\psi^{\rm I} \left(R^2 - {r_\psi^{\rm I}}^2\right) {\rm e}^{-W} \Bigg/ \int_A {\rm d}^2 r_\psi^{\rm I} \left(R^2 - {r_\psi^{\rm I}}^2\right),$$
$$W = \int_{t_{\rm I}}^{t_{\rm II}} {\rm d}t \ n_g(t) \ \langle \varGamma(t) \rangle. \tag{25}$$

4.2 Xu procedure

Xu et al. ((14) in [2]) did not take into account the formation time of the coulombic bound state, i.e., they chose the instant of production as

$$t_{\rm I}^{\rm Xu} = t_i = \tau_0. \tag{26}$$

Also, they seem to have used as integration variable the proper time $\tau_\psi = (t - t_i)/\gamma_\psi$ measured in the ψ meson rest frame. This procedure is inconvenient since the gluon number density $n_g^{\rm Xu}$ was best known in the fireball frame.

4.3 Our proposal

We do take into account the formation time $\tau_{\rm F}$, say, of the bound state in the $c\bar{c}$ barycentric frame. Remembering the dilation factor γ_ψ we choose

$$t_{\rm I}^{\rm Our} = t_i + \gamma_\psi \tau_{\rm F} \tag{27}$$

and retain the fireball time t for integration in (25). Of course, the value of $\tau_{\rm F}$ is not unique since many possible prescriptions for the same are available in the literature as outlined below.

Firstly, a naive estimate of $\tau_{\rm F}$ is given by $\tau_{\rm F} \approx 1/(m'_\psi - m_\psi) \approx 0.3\,{\rm fm}/c$. Secondly, Blaizot and Ollitrault [25] define $\tau_{\rm F} \approx 0.7\,{\rm fm}/c$ as the time spent by the heavy quark in going between two WKB classical turning points. Thirdly, Karsch and coworkers [27] use $\tau_{\rm F} \approx 0.89\,{\rm fm}/c$ as the time consumed by a quark to traverse a distance equal to the radius of the quarkonium in its rest frame. Finally, the dispersion-theoretic analysis of a correlation function by Kharzeev and Thews [28] yields $\tau_{\rm F} \sim 0.44\,{\rm fm}/c$. In the numerical work reported below we employ $0.89\,{\rm fm}/c$ [27] and $0.44\,{\rm fm}/c$ [28] as two *representative* values of the formation time.

There still remains an important question to be answered, viz. "why cannot the unformed $c\bar{c}$ pair interact before producing a fully-developed resonance?" Actually, such a preformation interaction is possible but its treatment will require tedious time-dependent wave packets in which the $c\bar{c}$ pair may scatter in the color singlet/octet states and one or more additional gluons may also be present. Our modest aim here was just to incorporate the effect of $\tau_{\rm F}$ via the simplest prescription in the original treatment of $S(p_{\rm T})$ by Xu et al.

4.4 Numerical work

In Figs. 5 and 6 the J/ψ survival probability has been plotted as a function of the transverse momentum based on the general formula (25). The solid and dashed-dotted curves denote our result using (27) with two representative values of $\tau_{\rm F}$ while the dashed curve is that of Xu et al. employing (26). Clearly, the J/ψ's survival chance is much higher (i.e.. their suppression is substantially lower) in our case compared to Xu et al.'s one. The reason for this can be understood by examining the integral appearing in (25), viz.

$$W = \int_{t_{\rm I}}^{t_{\rm II}} {\rm d}t \left[\begin{array}{l} \text{phase spaceintegral of} \\ v_{\rm rel}\sigma \text{ over } f \text{ at time } t \end{array} \right]. \tag{28}$$

First, we recall from Sect. 2 that $f^{\rm Our} < f^{\rm Xu}$ due to the fugacity effect. Next, we know from (21) that $v_{\rm rel} \equiv c < v_{\rm rel}^{\rm Xu}$ due to the flux effect. Finally, (26) and (27) tell us that the time interval available for dissociation $t_{\rm II} - t_{\rm I}^{\rm Our} <$

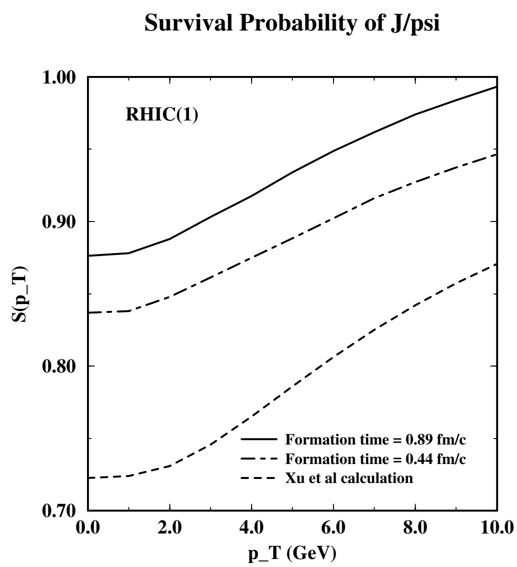

Fig. 5. The survival probability of J/ψ in an equilibrating parton plasma at RHIC energy with initial conditions given in Table 1 [2]. The solid and dashed-dotted curves are our result, while the dashed curve is the result obtained by Xu et al. [2]

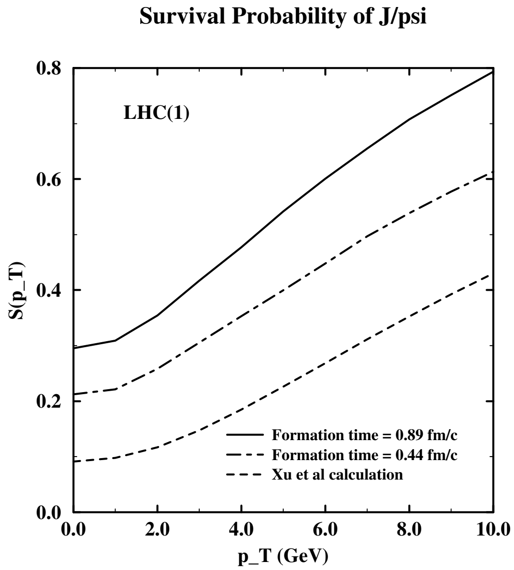

Fig. 6. Same as in Fig. 5 but at LHC energy

$t_{\rm II} - t_{\rm I}^{\rm Xu}$ due to the formation time effect. These three mechanisms operate *cooperatively* to make $W^{\rm Our} < W^{\rm Xu}$ resulting in less suppression.

5 Conclusions

(i) In this paper we have extended the work of Xu et al. [2] concerning the gluonic break-up of the J/ψ's created in an equilibrating QGP. Our theoretical formulae for number density (5), flux-weighted cross section (19), and survival probability (27) are new.

(ii) Our numerical results are also significant as compared to those of Xu et al. Since gluon *fugacity* is less than unity before chemical equilibration, our number density $n_g^{\rm Our}(t)$ of hard gluons (which are primarily instrumental in dissociating the J/ψ's) is lower as shown in Table 1.

(iii) Next, since our *g–ψ relative flux* in the meson rest frame is only $v_{\rm rel}^{\rm Our} = 1$ (and not $1+ \mid \mathbf{v}_\psi \mid$ of the fireball frame), our thermally-averaged rate $\langle \Gamma^{\rm Our}(t) \rangle$ is also smaller as depicted in Fig. 4.

(iv) Since we properly take into account the *production time* of the J/ψ's, the temporal span available for their break-up becomes shorter. These three effects act in a cooperative manner to reduce substantially the amount of J/ψ suppression (i.e. to increase noticeably their survival chance $S^{\rm Our}(p_{\rm T})$ as demonstrated by Figs. 5 and 6.

(v) Apart from possible $c\bar{c}$ recombination [7] another important effect not considered in the present paper is the *transverse* hydrodynamic expansion of the QGP. Mathematically such an expansion demands that the gluon statistical mechanics must be done in a local comoving frame, while physically the temperature will drop more quickly with time. This highly non-trivial problem is under investigation at present and its results will be published in a future communication.

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Appendix

J/ψ suppression mechanisms summarized

In relativistic heavy-ion collision the heavy quark–antiquark pairs (leading potentially to J/ψ mesons) are produced on a very short time scale, $\simeq 1/2m_c \simeq 10^{-24}$ s, with m_c being mass of the charmed quark. The pair develops into the physical resonance over a formation time $\simeq 0.89\,{\rm fm}/c$ in its own rest frame. This J/ψ traverses the deconfined plasma together with the hadronic matter before leaving the interaction region to decay into a dimuon which is finally detected. However, this chain of events can be prevented via any of the following mechanisms.

Even before the $c\bar{c}$ bound state is created it may be absorbed by the nucleons streaming past (Glauber/normal absorption [6]). Or, by the time the resonance is formed the Debye screening of the color forces in the plasma may be sufficient to kill it [9]. Or, an energetic parton could hit and dissociate the J/ψ [8]. Or, the Brownian motion of the J/ψ through the medium could cause its sufficient swelling/ionization [3]. The extent of suppression will be decided by the competition between the J/ψ momentum and the rate of hydrodynamic expansion (with associated cooling) of the plasma [11]. Of course, the entire above picture will be substantially modified if the J/ψ's are regenerated via $c\bar{c}$ recombination [7].

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