J/ψ Gluonic Dissociation Revisited: I. Fugacity, Flux And Formation Time Effects

B. K. Patra¹ and V. J. Menon²

¹ Dept. of Physics, Indian Institute of Technology, Roorkee 247 667, India
² Dept. of Physics, Banaras Hindu University, Varanasi 221 005, India

Abstract

We revisit the standard treatment [Xu, Kharzeev, Satz and Wang, Phys. Rev. C 53, 3051 (1996)] 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 correctly incorporated in the formulation of the gluon number density, velocity-weighted cross section, and the survival probability. Our new formulae are applied to numerically study 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.

PACS numbers: 12.38M
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 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 $J/\psi$ suppression due to medium influence [6] and recent $J/\psi$ enhancement via dynamical regeneration [7]. The well known mechanisms responsible for $J/\psi$ 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 $J/\psi$ owing to bombardment with energetic gluons [8]. For this mechanism, Xu et al. [2] employed statistical mechanics coupled with phenomenological QCD to calculate the $J/\psi$ survival probability in a temporally evolving parton gas. The aim of the present paper is to extend/modify the work of Xu et al in the following multifold respects:

(i) **Gluon fugacity effect:** For large momentum gluons which are responsible for the $\psi$ meson dissociation Xu et al. [2, Eq.(7)] replaced the gluon fugacity $\lambda_g$ by unity in the denominator of the Bose-Einstein distribution function. However, in the early stage of evolution the system may be quite far from chemical equilibrium implying that $\lambda_g$ need not be close to unity. In Sec. 2 below we derive a new formula for the gluon number density $n_g$ valid for general $\lambda_g \leq 1$.

(ii) **Relative flux effect:** Xu et al. [2, Eq.(8)] were interested in the product $\Gamma = v_{\text{rel}}\sigma$ where the gluon-$\psi$ break-up cross section $\sigma$ was written in the $\psi$ meson rest frame, but unfortunately their relative flux $v_{\text{rel}}$ was evaluated in the fireball frame. In Sec. 3 below we modify this procedure by treating the product $\Gamma$ strictly in the $\psi$ rest frame.

(iii) **Formation time effect:** For computing their survival probability $S(p_T)$ Xu et al. [2, Eq.(14)] used an integration over $\tau_\psi$ having lower limit 0, where $\tau_\psi$ is the proper time measured in $J/\psi$ rest frame. This is very inconvenient because the gluon density $n_g(t)$ and the thermal-averaged cross section $\langle v_{\text{rel}}\sigma \rangle$ are natural functions of the usual time $t$ in fireball rest frame. In Sec. 4 below we write a modified expression for $S(p_T)$ using $t$ integration where formation times of the QGP as well as $J/\psi$ are explicitly included. Of course, in our numerical results in Sections 2, 3, 4 the explicit velocity profiles of hydrodynamic flow are ignored. Finally, our main conclusions appear in Sec. 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-colour degeneracy factor, \( \lambda_g \leq 1 \) the gluon fugacity, and \( f = f(K^0, T, \lambda_g) \) the one-body gluon distribution function. Then the gluon number density \( n_g \) is obtained from

\[
n_g = 16 \int \frac{d^3K}{(2\pi)^3} f = \frac{8}{\pi^2} \int_0^\infty dK^0 K^0 f
\]

2.2 Xu Procedure

For near chemical equilibration Xu et al. [2, Eq.(7)] employed an approximate, factorized Bose-Einstein distribution

\[
f^{\text{Xu}} = \frac{\lambda_g}{e^{K^0/T} - 1} = \lambda_g \sum_{n=1}^\infty e^{-nK^0/T}
\]

which led them to a number density depending on the fugacity linearly through

\[
n_g^{\text{Xu}} = \frac{16}{\pi^2} T^3 \lambda_g \zeta(3)
\]

2.3 Our Proposal

In order to tackle the possibility of gluon chemical non-equilibration we use the full Bose-Einstein form

\[
f^{\text{Our}} = \frac{\lambda_g}{e^{K^0/T} - \lambda_g} = \sum_{n=1}^\infty \lambda_g^n e^{-nK^0/T}
\]

which guides us to a number density containing the fugacity in a power series via

\[
n_g^{\text{Our}} = \frac{16}{\pi^2} T^3 \sum_{n=1}^\infty \frac{\lambda_g^n}{n^3}
\]

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

\[
\frac{n_g^{\text{Xu}}}{n_g^{\text{Our}}} \sim (1 + \frac{1}{8})/(1 + \frac{\lambda_g^2}{8})
\]
which, however, requires the knowledge of $\lambda_g$ at various times. Of course, the algebraic reason for the inequality $n_{g}^{\text{Our}} < n_{g}^{\text{Xu}}$ is the fact that the distribution function $f^{\text{Our}} < f^{\text{Xu}}$ as long as $\lambda_g < 1$. In other words, the effect of correct 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 $t_i = \tau_0 \sim 1/\Lambda \sim 0.5\text{fm}/c$ where the suffix $i$ stands for “initial” and $\Lambda$ is the QCD energy scale. The initial conditions predicted by HIJING Monte Carlo simulation are summarized in Table 1. There gluon densities computed via Xu procedure (Eq.3) and our proposal (Eq. 5) are also listed. Clearly the relative difference between $n_{g}^{\text{Xu}}$ and $n_{g}^{\text{Our}}$ is of the order of $1/8 \sim 12\%$ which is significant.

|          | $T$ (GeV) | $t_i = \tau_0$ (fm) | $\lambda_g$ | $\lambda_q$ | $n_{g}^{\text{Xu}}(\text{fm})^{-3}$ | $n_{g}^{\text{Our}}(\text{fm})^{-3}$ |
|----------|-----------|---------------------|-------------|-------------|-------------------------------|-------------------------------|
| 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                         |

### 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:

\[
\frac{\dot{\lambda}_g}{\lambda_g} + 3 \frac{\dot{T}}{T} + \frac{1}{\tau} = R_3 (1 - \lambda_g) - 2 R_2 \left(1 - \frac{\lambda^2}{\lambda_g^2}\right),
\]

\[
\frac{\dot{\lambda}_q}{\lambda_q} + 3 \frac{\dot{T}}{T} + \frac{1}{\tau} = R_2 \frac{a_1}{b_1} \left(\frac{\lambda_g}{\lambda_q} - \frac{\lambda_q}{\lambda_g}\right),
\]
\[ \left( \lambda_g + \frac{b_2}{a_2} \lambda_q \right)^{3/4} T^3 \tau = \text{const} \]  

(7)

Here \( \tau \) is the medium proper time, \( \lambda_q \) the quark fugacity, \( N_f \) the number of flavours, and remaining symbols are defined by

\[
R_2 = 0.5n_g \langle v \sigma_{gg \rightarrow q\bar{q}} \rangle, \quad R_3 = 0.5n_g \langle v \sigma_{gg \rightarrow ggg} \rangle
\]

\[
a_1 = 16\zeta(3)/\pi^2, \quad a_2 = 8\pi^2/15
\]

\[
b_1 = 9\zeta(3)N_f/\pi^2, \quad b_2 = 7\pi^2 N_f/20
\]

(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) \( t_{\text{life}} \) 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 \( J/\psi \) becomes relevant. In the fireball frame consider a \( \psi \) meson of mass \( m_\psi \), four momentum \( p_\psi \), three velocity \( \vec{v}_\psi = \vec{p}_\psi/p_0 \psi \), and dilation factor \( \gamma_\psi = p_0 \psi / m_\psi \). If \( q \) is the gluon four momentum measured in \( \psi \) meson rest frame then by Lorentz transformations

\[
K^0 = \gamma_\psi (q^0 + \vec{v}_\psi \cdot \vec{q}) = \gamma_\psi q^0 (1 + |\vec{v}_\psi| \cos \theta_{q\psi})
\]

\[
d^3K = \left( \frac{K^0}{q^0} \right) d^3q
\]

(9)

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

The invariant quantum mechanical dissociation rate for \( g-\psi \) collision can be written compactly as

\[
\Gamma = v_{\text{rel}} \sigma
\]

(10)

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

\[
\langle \Gamma \rangle = \frac{16}{n_g} \int \frac{d^3q}{(2\pi)^3} \frac{K^0}{q^0} \Gamma f
\]

(11)

with \( f \) being the distribution function already encountered in Eq.(1).
3.2 Xu Procedure

Xu et al. [2, Eq.(8)] worked with their relative flux $v_{\text{rel}}^{Xu} = q^0/(K^0\gamma_\psi)$ in the fireball frame, but unfortunately the cross section $\sigma_{\text{Rest}}$ was in $\psi$ meson rest frame based on the standard QCD value [20]

$$\sigma_{\text{Rest}} = B(Q^0 - 1)^{3/2}/Q^{05} ; \quad q^0 > \epsilon_\psi$$

(12)

$$Q^0 = \frac{q^0}{\epsilon_\psi} , \quad B = \frac{2\pi}{3} \left( \frac{32}{3} \right)^2 \frac{1}{m_c(\epsilon_\psi m_c)^{1/2}}$$

(13)

where $\epsilon_\psi$ is the $J/\psi$ binding energy and $m_c$ the charm quark mass. Insertion into Eq. (11) led them to

$$\langle \Gamma^{Xu} \rangle = 16 n_{g_{Xu}}^2 \int \frac{d^3 q}{(2\pi)^3} \frac{1}{\gamma_\psi} \sigma_{\text{Rest}} \lambda_g \sum_{n=1}^{\infty} e^{-nK_0^0/T}$$

(14)

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

$$\langle \Gamma^{Xu} \rangle = \frac{8\epsilon_\psi^3 \lambda_g}{\pi^2 \gamma_\psi n_{g_{Xu}}} \sum_{n=1}^{\infty} \int_1^\infty dQ^0 Q^{02} \sigma_{\text{Rest}} e^{-C_n Q_0^0} \left( \frac{\sinh D_n Q^0}{D_n Q^0} \right)$$

(15)

$$= \frac{4\epsilon_\psi^2 \lambda_g T}{\pi^2 \gamma_\psi^2} \sum_{n=1}^{\infty} \int_1^\infty dQ^0 Q^0 \sigma_{\text{Rest}} \left( e^{-A_n^+ Q^0} - e^{-A_n^- Q^0} \right)$$

(16)

Here the following abbreviations have been introduced:

$$\gamma_\psi = p^0_\psi/m_\psi , \quad \bar{v}_\psi = \bar{p}_\psi/p^0_\psi , \quad C_n = n\epsilon_\psi \gamma_\psi / T$$

$$D_n = |\bar{v}_\psi| , \quad A^\pm_n = C_n \pm D_n = C_n (1 \pm |\bar{v}_\psi|)$$

(17)

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

3.3 Our Proposal

We set up our $\Gamma$ entirely in the $\psi$ meson rest frame where $v_{\text{rel}}^{\text{Our}} = c = 1$. Thereby Eq. (11) becomes

$$\langle \Gamma^{\text{Our}} \rangle = \frac{16}{n_{g_{\text{Our}}}^2} \int \frac{d^3 q}{(2\pi)^3} \gamma_\psi (1 + |\bar{v}_\psi| \cos \theta_{q\psi}) \sigma_{\text{Rest}} \sum_{n=1}^{\infty} \lambda_g^n e^{-nK_0^0/T}$$

(18)
where the exact $f^{\text{Our}}$ given by Eq.(4) has been recalled. The slightly complicated angular integration over $d \cos \theta_{q}q \psi$ can be performed by choosing the polar axis along $\hat{v}_{\psi}$ yielding

$$
\langle \Gamma^{\text{Our}} \rangle = \frac{8 \epsilon_{q}^{3} \gamma_{q} \sum_{n=1}^{\infty} \lambda_{g}^{n} \int_{1}^{\infty} dQ^{0} Q^{0} \sigma_{\text{Rest}} e^{-Q_{\text{Rest}}^{0}} [I_{0}(\rho_{n}) - |\vec{v}_{\psi}| I_{1}(\rho_{n})]}
= \frac{4 \epsilon_{q}^{2} T}{\pi^{2} |\vec{v}_{\psi}| n_{g}^{\text{Our}} \int_{1}^{\infty} dQ^{0} Q^{0} \sigma_{\text{Rest}} \left[ \left( 1 - |\vec{v}_{\psi}| \left( 1 - \frac{1}{\rho_{n}} \right) \right) e^{-A_{n}^{0} Q_{\text{Rest}}} - \left( 1 + |\vec{v}_{\psi}| \left( 1 + \frac{1}{\rho_{n}} \right) \right) e^{-A_{n}^{0} Q_{\text{Rest}}} \right]}
$$

(19)

where

$$
\rho_{n} = D_{n} Q^{0} = D_{n} q_{0}^{0}/\epsilon_{q} \gamma_{q} K_{0} \sim 1/\gamma_{q}^{2} (1 - |\vec{v}_{\psi}|) \sim 1 + |\vec{v}_{\psi}|
$$

(20)

Clearly the dependence of $\langle \Gamma^{\text{Our}} \rangle$ on $\lambda_{g}$ and $\gamma_{q} \psi$ is more involved than that of $\langle \Gamma^{\text{Xu}} \rangle$ given by Eqs. (15, 16).

### 3.4 Numerical Work

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

$$
v^{\text{Xu}}_{\text{rel}} \equiv \frac{q^{0}}{\gamma_{q} K_{0}} \sim \frac{1}{\gamma_{q}^{2} (1 - |\vec{v}_{\psi}|)} \sim 1 + |\vec{v}_{\psi}|
$$

(21)

Now let us consider the ratio

$$
\frac{\langle \Gamma^{\text{Xu}} \rangle}{\langle \Gamma^{\text{Our}} \rangle} = \left[ \frac{n_{g}^{\text{Our}}}{n_{g}^{\text{Xu}}} \right] \left[ \frac{\text{phase space integral of } v^{\text{Xu}}_{\text{rel}} \sigma_{\text{Rest}}}{\text{phase space integral of } c \sigma_{\text{Rest}}} \right]
$$

(22)

Due to fugacity effect the number density ratio in Eq.(22) is somewhat smaller than unity as already mentioned in Sec.2. On the other hand, since in near antiparallel configuration, the relative flux $v^{\text{Xu}}_{\text{rel}} > c$, hence the ratio of the phase space integrals is somewhat larger than unity. These two effects tend to partially compensate each other.
in Eq. (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 $1+ |\vec{v}_\psi|$ becomes more pronounced at high transverse momentum, causing noticeable difference between the curves of Figs. 3, 4.

4 Survival Probability

4.1 Preliminaries

Consider a cylindrical coordinate system in the fireball frame where the $\psi$ meson was created at the time-space point $(t_I, r_{\psi}^I, \phi_{\psi}^I)$ with transverse velocity $\vec{v}_\psi^T$. The plasma is supposed to be contained within a cylinder of radius $R$, expanding longitudinally till the end of its lifetime $t_{\text{life}}$. The $\psi$ meson’s trajectory will hit the said cylinder after covering a distance $d_{RI}$ in the time interval $t_{RI}$ such that

$$d_{RI} = -r_{\psi}^I \cos \phi_{\psi}^I + \sqrt{R^2 - r_{\psi}^I^2 \sin^2 \phi_{\psi}^I},$$

$$t_{RI} = \frac{d_{RI}}{|\vec{v}_\psi^T|},$$

(23)
Figure 2: Same as in Fig.1, but these curves are obtained by our Eq.(19).

the full temporal range of interest is obviously

\[ t_I \leq t \leq t_{II} ; \quad t_{II} = \min(t_i + t_{RI}, t_{\text{life}}) \] (24)

The corresponding survival probability of \( J/\psi \) averaged over its initial production configuration extending over the transverse area \( A \) becomes

\[
S(p_T) = \int_A d^2r_{\psi}^I \left( R^2 - r_{\psi}^I r_{\psi}^J \right) e^{-W} / \int_A d^2r_{\psi}^I \left( R^2 - r_{\psi}^I r_{\psi}^J \right)
\]

\[
W = \int_{t_I}^{t_{II}} dt \, n_g(t) \langle \Gamma(t) \rangle
\] (25)

4.2 Xu Procedure

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

\[
t^X_I = t_i = \tau_0
\] (26)

Also, they seem to have used as integration variable the proper time \( \tau_\psi = (t - t_i) / \gamma_\psi \) measured in \( \psi \) meson rest frame. This procedure is inconvenient since the gluon number density \( n^X_g \) was best known in the fireball frame.
Figure 3: The thermal-averaged gluon-$J/\psi$ dissociation cross section $\langle v_{\text{rel}} \sigma \rangle$ as a function of transverse momentum at different temperatures as done by Xu et al. [2].

4.3 Our Proposal

We do take into account the formation time $\tau_F \sim 0.89\text{fm}/c$ of the bound state in the $c\bar{c}$ barycentric frame. Remembering the dilation factor $\gamma_\psi$ we choose

$$t_I^{\text{Our}} = t_i + \gamma_\psi \tau_F$$  \hspace{1cm} (27)

and retain the fireball time $t$ for integration in Eq.(25).

4.4 Numerical Work

In Figs.5, 6 the $J/\psi$ survival probability has been plotted as a function of the transverse momentum based on the general formula (25). The solid curve denotes our result using Eq.(27) while the dashed curve is that of Xu et al employing Eq.(26). Clearly, the $J/\psi$'s survival chance is much more (i.e., their suppression is substantially less) in our case compared to Xu et al.'s. Its reason can be understood by examining the integral appearing in Eq.(25) viz.

$$W = \int_{t_I}^{t_{II}} dt \ [\text{phase space integral of } v_{\text{rel}} \sigma \text{ over } f \text{ at time } t]$$  \hspace{1cm} (28)
First, we recall from Sec.2 that $f^{\text{Our}} < f^{\text{Xu}}$ due to the fugacity effect. Next, we know from Eq.(21) that $v_{\text{rel}} \equiv c < v_{\text{rel}}^{\text{Xu}}$ due to the flux effect. Finally, Eqs.(26, 27) tell that the time interval available for dissociation $t_{II} - t_{I}^{\text{Our}} < t_{II} - t_{I}^{\text{Xu}}$ due to the formation time effect. These three mechanisms operate *cooperatively* to make $W^{\text{Our}} < W^{\text{Xu}}$ 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 $J/\psi$’s created in an equilibrating QGP. Our theoretical formulae on number density (Eq. 5), flux-weighted cross section (Eq.19), and survival probability (Eq. 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, hence our number density $n_{g}^{\text{Our}}(t)$ of hard gluons (which are primarily instrumental in dissociating the $J/\psi$’s) is lower as shown in Table 1.

(iii) Next, since our $g - \psi$ *relative flux* in meson rest frame is only $v_{\text{rel}}^{\text{Our}} = 1$ (and not $1 + |\vec{v}_\psi|$ of the fireball frame) hence our thermally-averaged rate $\langle \Gamma^{\text{Our}}(t) \rangle$ is also
Figure 5: The survival probability of $J/\psi$ in an equilibrating parton plasma at RHIC energy with initial conditions given in Table 1 [2]. The solid curve is our result, while the dashed curve is the result obtained by Xu et al. [2].

smaller as depicted in Fig.4.

(iv) Since we properly take into account the production time of the $J/\psi$’s, hence the temporal span available for their break-up becomes shorter. These three effects act in a cooperative manner to reduce substantially the amount of $J/\psi$ suppression (i.e. to increase noticeably their survival chance $S^{\text{Out}}(p_T)$ as demonstrated by Figs. 5, 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 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 nontrivial problem is under investigation at present and its results will be published in a future communication.

**ACKNOWLEDGEMENTS**

VJM thanks the UGC, Government of India, New Delhi for financial support. We thank Dr. Dinesh Kumar Srivastava for useful discussions during this work.
Figure 6: Same as in Fig. 5 but it is at LHC energy.

References

[1] Helmut Satz, Rept.Prog.Phys. 63, 1511 (2000).

[2] Xiao-Ming Xu, D. Kharzeev, H. Satz, and Xin-Nian Wang, Phys. Rev. C 53, 3051 (1996).

[3] B. K. Patra and V. J. Menon, Nucl. Phys. A 708, 353 (2002).

[4] B. K. Patra, D. K. Srivastava, Phys. Lett. B 505, 113 (2001).

[5] X.-N Wang and M. Gyulassy, Phys. Rev. D 44, 3501 (1991).

[6] C. Gerschel and J. Hüfner, Phys. Lett. B 207, 253 (1988).

[7] R. L. Thews, M. Schroedter, and J. Rafelski, , Phys. Rev. C 63, 054905 (2001).

[8] D. Kharzeev and H. Satz, Phys. Lett. B 334, 155 (1994).

[9] T. Matsui and H. Satz, Phys. Lett. B 178, 416 (1986).

[10] H. Satz and D. K. Srivastava, Phys. Lett. B 475, (2000).

[11] D. Pal, B. K. Patra, and D. K. Srivastava, Eur. Phys. Jour. C 17, 179 (2000).
[12] D. Kharzeev and H. Satz, Phys. Lett B 366, 316 (1996).

[13] J. P. Blaizot and J. Y. Ollitrault, Phys. Lett. B 199 (1987) 499; F. Karsch and H. Satz, Z. Phys. C51 (1991) 209.

[14] K. Geiger, Phys. Rep. 258 (1995) 237.

[15] K. J. Eskola, B. Müller, and X.-N. Wang, Phys. Lett. B 374 (1996) 20.

[16] T. S. Biro, E. van Doorn, M. H. Thoma, B. Müller, and X.-N. Wang, Phys. Rev. C 48 (1993) 1275.

[17] D. K. Srivastava, M. G. Mustafa, and B. Müller, Phys. Rev. C 56 (1997) 1064.

[18] K. J. Eskola, K. Kajantie, P. V. Ruuskanen, and K. Tuominen, Nucl. Phys. B 570 (2000) 379.

[19] K. Kajantie et al., Phys. Rev. Lett. 17 (1997) 3130.

[20] M. E. Peskin, Nucl. Phys. B 156 (1979) 365; G. Bhanot and M. E. Peskin, Nucl. Phys. B 156 (1979) 391.

[21] J. P. Blaizot and J. Y. Ollitrault, in Quark Gluon Plasma Ed. R. C. Hwa, World Scientific, Singapore, p. 531.

[22] L. Gerland, L. Frankfurt, and M. I. Strikman, H. Stöcker, and W. Greiner, Nucl. Phys. A 663 (2000) 1019.

[23] J. Guinon and R. Vogt, Nucl. Phys. B 492 (1997) 301.

[24] D. Kharzeev and H. Satz, in Quark Gluon Plasma 2 Ed. R. C. Hwa, World Scientific, Singapore, 1995, p.395.
Appendix

$J/\psi$ Suppression Mechanisms Summarized

In relativistic heavy ion collision the heavy quark-antiquark pairs (leading potentially to $J/\psi$ mesons) are produced on a very short time scale $\simeq 1/2m_c \simeq 10^{-24}$ sec with $m_c$ being mass of the charmed quark. The pair develops into the physical resonance over a formation time $\simeq 0.89\text{fm}/c$ in its own rest frame. This $J/\psi$ 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 colour forces in the plasma may be sufficient to kill it [9]. Or, an energetic parton could hit and dissociate the $J/\psi$ [8]. Or, the Brownian motion of the $J/\psi$ through the medium could cause its sufficient swelling/ionization [3]. The extent of suppression will be decided by a competition between the $J/\psi$ 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 $J/\psi$’s are regenerated via $c\bar{c}$ recombination [7].