J/ψ suppression in an equilibrating parton plasma

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Short-distance QCD is employed to calculate the J/ψ survival probability in an equilibrating parton gas whose evolution is governed by a set of master rate equations. Partons in the early stage of high-energy nuclear collisions may initially not be in equilibrium, but their average transverse momentum is sufficiently high to break up a $Q\bar{Q}$ bound state. Such a breakup during the evolution of the parton gas is shown to cause a substantial J/ψ suppression at both RHIC and LHC energies, using realistic estimates of the initial parton densities. The transverse momentum dependence of the suppression is also shown to be sensitive to the initial conditions and the evolution history of the parton plasma. [S0556-2813(96)06205-X]

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I. INTRODUCTION

It is generally believed, and confirmed by lattice QCD calculations [1], that hadronic matter under extreme conditions will form a plasma in which quarks and gluons are no longer confined to individual hadrons and are in both thermal and chemical equilibrium. To search for such a quark-gluon plasma, it is proposed to study collisions of heavy nuclei at extremely high energies. Recent work using models based on perturbative QCD indeed shows that at high energy a dense partonic system can be produced [2–6]. Though it is not all clear if such partonic systems will reach thermal and chemical equilibrium before hadronization [7–11], the partons, mainly gluons, are certainly in a deconfined state at such high densities [12].

Many signals could arise from this deconfined state, such as charm quark enhancement [13–15] or enhanced photon and dilepton production [16,17]. Charm quarks, for example, cannot be easily produced during the mixed and hadronic phases of dense matter, due to their large masses compared to the temperature. They can be readily produced only during the early stage of the evolution, when partonic degrees of freedom are relevant. In this paper, we will discuss how preequilibrium J/ψ suppression can be used to probe the early deconfined state of the partonic system and the dynamics governing its evolution toward equilibrium.

J/ψ suppression due to color screening has been proposed to probe deconfinement [18]. This requires that the interactions of J/ψ with hadrons and deconfined partons be different [19]. Because of its small size, a heavy quarkonium can probe the short-distance properties of light hadrons. It is thus possible to make a parton-based calculation of the J/ψ-hadron cross section via an operator product expansion method similar to that used in deeply inelastic lepton-hadron scatterings [19–21]. The resulting J/ψ-hadron cross section can be related to the distribution function of gluons inside a hadron. The energy dependence of the cross section near the threshold of the breakup of a J/ψ is determined by the large-x behavior of the gluon distribution function, giving rise to a very small breakup cross section at low energies. Only at very high energies will this cross section reach its asymptotic value of a few mb. In other words, dissociation can only occur if the gluon from the light hadron wave function is hard in the J/ψ rest frame, i.e., its energy is high enough to overcome the binding energy threshold. A hadron gas with temperature below 0.5 GeV certainly cannot provide such energetic gluons to break up the J/ψ. Therefore, a slow J/ψ is very unlikely to be absorbed inside a hadron gas of reasonable temperature [19]. This conclusion does not seem to be affected substantially by nonperturbative effects, analyzed in Ref. [22].

On the other hand, a deconfined partonic system contains much harder gluons which can easily break up a J/ψ [19,23]. A study of the energy dependence of the gluon-J/ψ inelastic cross section [23] shows a strong peak just above the breakup threshold of the gluon energy, $\epsilon_0 = 2M_D - M_{J/\psi}$, where $M_{J/\psi}$ and M_D are the J/ψ and D meson masses, respectively. In the preequilibrium stage, i.e., before the partons have reached equilibrium, the average parton transverse momentum is sufficiently large [24] to break up a J/ψ, provided the partons are deconfined. The dissociation of J/ψ will continue during the whole equilibration process until the effective temperature drops below a certain value or the beginning of hadronization, whichever takes place first. Therefore measurements of J/ψ suppression can probe the deconfinement of the early partonic system and shed light on the subsequent equilibration process, provided that possible nuclear effects on the production of $Q\bar{Q}$ pairs and on preresonance charmonium states are understood and taken into account.

In the following we will first calculate the thermal gluon-J/ψ dissociation cross section at different temperatures and for different J/ψ transverse momenta. We then follow the evolution of an initially produced parton gas toward equilibrium and calculate the resulting total survival probability of a J/ψ and its P_T dependence.

II. J/ψ DISSOCIATION BY GLUONS

The operator product expansion allows one to express the hadron-J/ψ inelastic cross section in terms of the convolution of the gluon-J/ψ dissociation cross section with the gluon distribution inside the hadron [19]. The gluon-J/ψ dissociation cross section is given by [23]

$$\sigma(q^0)=\frac{2\pi}{3}\left(\frac{32}{3}\right)^2\left(\frac{16\pi}{3g_s^2}\right)\frac{1}{m_Q^2}\frac{(q^0/\epsilon_0-1)^{3/2}}{(q^0/\epsilon_0)^5}, \quad (1)$$

where g_s is the coupling constant of a gluon and c quark, m_Q the c quark mass, and q^0 the gluon energy in the J/ψ rest frame; its value must be larger than the J/ψ binding energy ϵ_0. Since for the tightly bound ground state of quarkonium the binding force between the heavy quark and antiquark is well approximated by the one-gluon-exchange Coulomb potential, the $Q\bar{Q}$ bound state is hydrogenlike and the Coulomb relation holds,

$$\epsilon_0=\left(\frac{3g_s^2}{16\pi}\right)^2 m_Q. \quad (2)$$

The cross section thus can be rewritten as

$$\sigma(q^0)=\frac{2\pi}{3}\left(\frac{32}{3}\right)^2\frac{1}{m_Q(\epsilon_0 m_Q)^{1/2}}\frac{(q^0/\epsilon_0-1)^{3/2}}{(q^0/\epsilon_0)^5}. \quad (3)$$

As shown in Monte Carlo simulations [24], the parton density in the early stage of high-energy heavy-ion collisions has an approximate Bjorken-type [25] scaling behavior. We will only consider J/ψ suppression in the central rapidity region ($y_{J/\psi}\simeq 0$). In this case, the J/ψ will move in the transverse direction with a four-velocity

$$u=(M_T,\vec{P}_T,0)/M_{J/\psi}, \quad (4)$$

where $M_T=\sqrt{P_T^2+M_{J/\psi}^2}$ is defined as the J/ψ transverse mass. A gluon with a four-momentum $k=(k^0,\vec{k})$ in the rest frame of the parton gas has an energy $q^0=k\cdot u$ in the rest frame of the J/ψ. The thermal gluon-J/ψ dissociation cross section is then defined as

$$\langle v_{\rm rel}\sigma(k\cdot u)\rangle_k=\frac{\int d^3k v_{\rm rel}\sigma(k\cdot u)f(k^0;T)}{\int d^3k f(k^0;T)}, \quad (5)$$

where the gluon distribution in the rest frame of the parton gas is defined as

$$f(k^0;T)=\frac{\lambda_g}{e^{k^0/T}-\lambda_g}, \quad (6)$$

with $\lambda_g\leqslant 1$ specifying the deviation of the system from chemical equilibrium. For large momentum gluons which are responsible for the J/ψ dissociation, we can approximate the above distribution by a factorized Bose-Einstein distribution

$$f(k^0;T)\approx\frac{\lambda_g}{e^{k^0/T}-1}. \quad (7)$$

The relative velocity $v_{\rm rel}$ between the J/ψ and a gluon is

$$v_{\rm rel}=\frac{P_{J/\psi}\cdot k}{k^0 M_T}=1-\frac{\vec{k}\cdot\vec{P}_T}{k^0 M_T}. \quad (8)$$

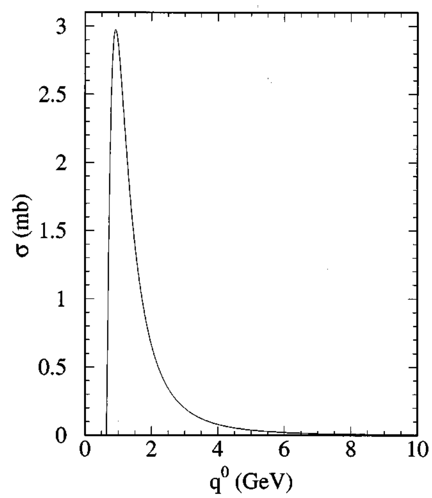

FIG. 1. Gluon-J/ψ dissociation cross section as a function of the gluon energy q^0 in the rest frame of the J/ψ.

Changing the variable to the gluon momentum, $q=(q^0,\vec{q})$, in the rest frame of the J/ψ, the integral in the numerator of Eq. (5) can be rewritten as

$$\int d^3q\frac{M_{J/\psi}}{M_T}\sigma(q^0)f(k^0;T), \quad (9)$$

where

$$k^0=(q^0M_T+\vec{q}\cdot\vec{P}_T)/M_{J/\psi}. \quad (10)$$

One can carry out the integral in the denominator, $\int d^3k f(k^0;T)=8\pi\zeta(3)\lambda_g T^3$, and the angular part in the numerator, to get

$$\langle v_{\rm rel}\sigma(k\cdot u)\rangle_k=\left(\frac{8}{3}\right)^3\frac{\pi}{\zeta(3)}\frac{M_{J/\psi}^2}{P_T M_T T^3}\left(\frac{\epsilon_0}{m_Q}\right)^{3/2}\sum_{n=1}^{\infty}T_n \times\int_1^\infty dx\frac{(x-1)^{3/2}}{x^4}(e^{-a_n^- x}-e^{-a_n^+ x}), \quad (11)$$

with $T_n=T/n$ and

$$a_n^\pm=\frac{\epsilon_0}{T_n}\frac{M_T\pm P_T}{M_{J/\psi}}. \quad (12)$$

In order to understand the temperature and P_T dependence of the thermal gluon-J/ψ dissociation cross section, we first plot in Fig. 1 the cross section $\sigma(q^0)$ of Eq. (3) as a function of the gluon energy in the J/ψ rest frame. It decreases strongly toward the threshold and is broadly peaked around $q^0=10\epsilon_0/7=0.92$ GeV, with a maximum value of about 3 mb. Low-momentum gluons do not have the resolution to distinguish the heavy constituent quarks or the energy to excite them to the continuum. On the other end, high-momentum gluons also have small cross section with a J/ψ since they cannot see the large size.

We can also express the cross section as a function of the center-of-mass energy of gluons and the J/ψ, $\sigma(q^0)=\sigma(s/2M_{J/\psi}-M_{J/\psi}/2)$, where $s=(k+P_{J/\psi})^2$. One

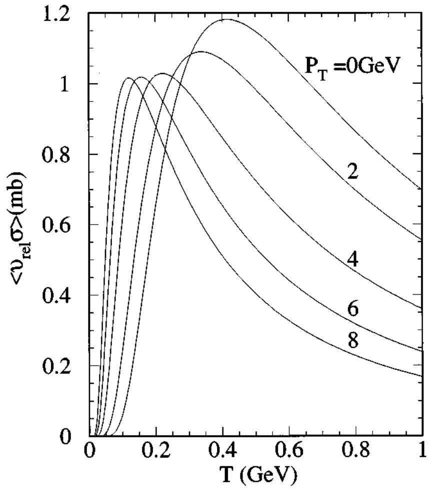

FIG. 2. The thermal-averaged gluon-J/ψ dissociation cross section $\langle v_{\rm rel}\sigma\rangle$ as a function of the temperature at different transverse momenta P_T.

can thus translate the energy dependence in Fig. 1 into temperature and P_T dependences after thermal average, since the thermally averaged $\langle s\rangle$ is proportional to both P_T and temperature T. In Fig. 2 we plot the thermally averaged gluon-J/ψ dissociation cross section as a function of temperature for different values of the J/ψ transverse momentum P_T. We observe the same kind of peak structure, with a decreased maximum value due to the thermal average. The position of the peak also shifts to smaller values of T when P_T is increased, corresponding to a fixed value of the averaged center-of-mass energy $\langle s\rangle$. A similar behavior is expected if one plots the thermal cross section as a function of P_T at different temperatures, as done in Fig. 3. However, in this case, the peak simply disappears at high enough temperatures, because the averaged $\langle s\rangle$ will be above the threshold value even for $P_T=0$. These features will have considerable consequences for the survival probability of a J/ψ in an equilibrating parton gas, especially the P_T dependence. We should also mention that the use of the factorized approximation of the Bose-Einstein distribution function in Eq. (7) has an effect of about 20% on the thermal cross section, compared to that obtained with a Boltzmann distribution [23].

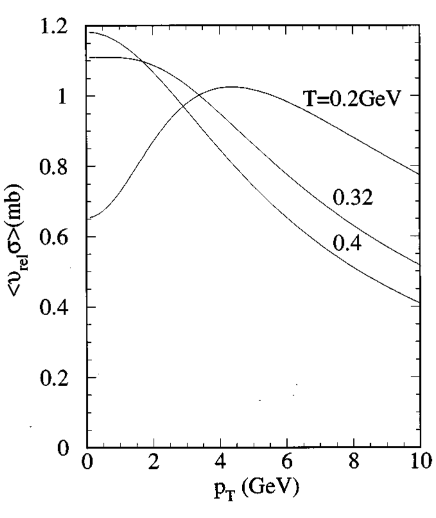

FIG. 3. The thermal-averaged gluon-J/ψ dissociation cross section $\langle v_{\rm rel}\sigma\rangle$ as a function of the transverse momentum P_T at different temperatures.

III. J/ψ SUPPRESSION IN AN EQUILIBRATING PARTON GAS

Using the thermal cross section just obtained, we can now calculate the survival probability of J/ψ in an equilibrating parton plasma. In this paper, we will neglect the transverse expansion and consider only longitudinal expansion. We will also only consider J/ψ suppression in the central rapidity region. A J/ψ produced at point $\vec{r}$ with velocity $\vec{v}$ in the transverse direction will travel a distance

$$d=-r\cos\phi+\sqrt{R_A^2-r^2(1-\cos^2\phi)} \tag{13}$$

in the time interval $t_\psi=M_T d/P_T$ before it escapes from a gluon gas of transverse extension R_A; here, $\cos\phi=\hat{v}\cdot\hat{r}$. Suppose the system evolves in a deconfined state until the temperature drops below a certain value, which we assume to be 200 MeV. The total amount of time the J/ψ remains inside a deconfined parton gas is the smaller one of the two times t_ψ and t_f, the lifetime of the parton gas. Assume that the initial production rate of the J/ψ is proportional to the number of binary nucleon-nucleon interactions at impact parameter r, $N_A(r)=A^2(1-r^2/R_A^2)/2\pi R_A^2$. The survival probability of the J/ψ averaged over its initial position and direction in an equilibrating parton gas is

$$S(P_T)=\frac{\int d^2r(R_A^2-r^2)\exp[-\int_0^{t_{\rm min}}d\tau n_g(\tau)\langle v_{\rm rel}\sigma(k\cdot u)\rangle_k]}{\int d^2r(R_A^2-r^2)}, \tag{14}$$

where

$$t_{\rm min}=\min(t_\psi,t_f), \tag{15}$$

and $n_g(\tau)$ is the gluon number density at a given time τ.

In Eq. (14), both the gluon number density $n_g(\tau)$ and the thermal cross section $\langle v_{\rm rel}\sigma(k\cdot u)\rangle_k$ depend on the temperature, which in turn is a function of time. In addition, n_g is proportional to gluon fugacity which also evolves with time. To evaluate the survival probability we need to know the entire evolution history of the parton system. Roughly speaking, one can divide this history into two stages: (1) First there is kinetic thermalization, mainly through elastic scatterings and expansion. The kinematic separation of free-streaming partons gives us $\tau_0\sim0.5$–0.7 fm/c as an estimate of the time when local isotropy in momentum distribution is reached [9,24]. (2) The parton gas now further evolves toward chemical equilibrium through parton proliferation and gluon fusion; it does so until hadronization or freeze-out, whichever happens first. This evolution can be determined by a set of master rate equations which give us the time dependence of the temperature and fugacities.

In this paper we will not address the question of preresonance J/ψ suppression, which has been discussed and shown to be responsible for the J/ψ suppression observed in p-A and S-U collisions [26–29]. We will here consider the suppression of fully formed physical J/ψ states as it should take place if nuclear collision produces dense partonic system.

Following Ref. [9], we characterize the nonequilibrium of the system by gluon and quark fugacities which are less than unity. The dominant reactions leading to chemical equilibrium are assumed to be the following two processes:

$$gg \leftrightarrow ggg, \quad gg \leftrightarrow q\overline{q}. \tag{16}$$

Assuming that elastic parton scatterings are sufficiently rapid to maintain local thermal equilibrium, the evolution of the parton densities can be given by the master rate equations. Combining these master equations together with one-dimensional hydrodynamic equation, one can get the following set of equations [9]:

$$\frac{\dot{\lambda}_g}{\lambda_g} + 3\frac{\dot{T}}{T} + \frac{1}{\tau} = R_3(1-\lambda_g) - 2R_2\left(1 - \frac{\lambda_q^2}{\lambda_g^2}\right), \tag{17}$$

$$\frac{\dot{\lambda}_q}{\lambda_q} + 3\frac{\dot{T}}{T} + \frac{1}{\tau} = R_2\frac{a_1}{b_1}\left(\frac{\lambda_g}{\lambda_q} - \frac{\lambda_q}{\lambda_g}\right), \tag{18}$$

$$\left(\lambda_g + \frac{b_2}{a_2}\lambda_q\right)^{3/4} T^3\tau = \text{const}, \tag{19}$$

where $a_1 = 16\zeta(3)/\pi^2 \approx 1.95$, $a_2 = 8\pi^2/15 \approx 5.26$, $b_1 = 9\zeta(3)N_f/\pi^2 \approx 2.20$, and $b_2 = 7\pi^2 N_f/20 \approx 6.9$. The density and velocity weighted reaction rates

$$R_3 = \frac{1}{2}\langle\sigma_{gg\to ggg}v\rangle n_g, \quad R_2 = \frac{1}{2}\langle\sigma_{gg\to q\overline{q}}v\rangle n_g \tag{20}$$

can be found in Refs. [9,13]. After taking into account parton screening [12] and the Landau-Pomeranchuck-Migdal effect in induced gluon radiation [30,31], R_3/T and R_2/T are found to be functions of only λ_g. Solving the above rate equations as shown in Ref. [9], one finds that the parton gas cools considerably faster than predicted by Bjorken's scaling solution ($T^3\tau$ = const), because the production of additional partons approaching the chemical equilibrium state consumes an appreciable amount of energy. The accelerated cooling, in turn, slows down the chemical equilibration process. For the initial conditions given by HIJING Monte Carlo simulations [4], at RHIC energy the parton system can hardly reach its equilibrium state, since the effective temperature here drops below $T_c \approx 200$ MeV in a time of some 1–2 fm/c. At LHC energy, however, the parton gas comes very close to equilibrium, since the system may exist in a deconfined state for as long as 4–5 fm/c. In the following we will use the numerical results for the time evolution of the temperature and fugacities obtained from the above master equations to calculate the survival probability of a J/ψ in such an equilibrating parton system.

Shown in Fig. 4(a) are the J/ψ survival probabilities in the deconfined and equilibrating parton plasma at RHIC and LHC energies with initial conditions given by HIJING Monte Carlo simulations [4], denoted as set 1 in Table I. We find that there is stronger J/ψ suppression at LHC than at RHIC energy, due both to the higher initial parton densities and longer lifetime of the parton plasma. The increase of the survival probabilities with J/ψ transverse momentum is a consequence of the decrease of the thermal cross section with increasing P_T at high temperatures, as shown in Fig. 3, and the shorter time spent by a higher-P_T J/ψ inside the parton plasma, an effect first considered in Ref. [32]. J/ψ dissociation during the late stage of the evolution should have a peak and increase a little with P_T when the temperature drops below 0.3 GeV, as illustrated in Fig. 3. This behavior flattens the total survival probability at small values of P_T when we integrate over the entire history of the evolution from high initial temperatures. For a parton system with a low initial temperature (below 300 MeV), the P_T dependence of the survival probability should be even flatter. One can therefore use the P_T dependence to shed light on the initial temperature and the evolution history of the system.

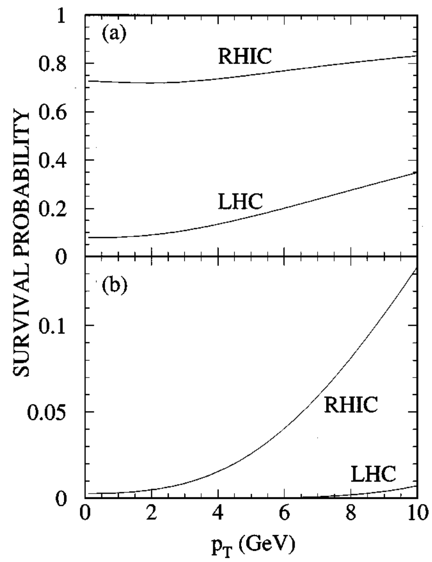

FIG. 4. (a) The survival probability of J/ψ in an equilibrating parton plasma at RHIC and LHC energies with initial conditions given as set 1 in Table I and (b) for an initially equilibrated plasma at the same temperatures.

To demonstrate the effects of the chemical nonequilibrium in the initial system, we also show in Fig. 4(b) the survival probabilities in an ideal parton gas with the same initial temperatures as before, but with full chemical equilibrium (unit initial fugacities). In this case, the temperature simply decreases like $T(\tau) = T(\tau_0)(\tau_0/\tau)^{1/3}$. We see that the J/ψ is now much more suppressed than in the case of an equilibrating parton plasma, because of both the higher parton density and the longer lifetime of the system.

TABLE I. Different sets of initial conditions of the temperature, fugacities, and parton number densities at $\tau_0 = 0.7$ fm/c for RHIC and $\tau_0 = 0.5$ fm/c for LHC.

	RHIC(1)	LHC(1)	RHIC(2)	LHC(2)	RHIC(3)	LHC(3)
T (GeV)	0.55	0.82	0.55	0.82	0.4	0.72
λ_g	0.05	0.124	0.2	0.496	0.53	0.761
λ_q	0.008	0.02	0.032	0.08	0.083	0.118
n_g (fm^{-3})	2.15	18	8.6	72	8.6	72
n_q (fm^{-3})	0.19	1.573	0.76	6.29	0.76	6.29

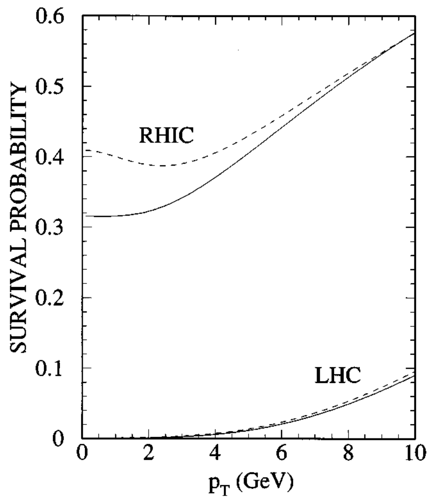

FIG. 5. The survival probability of J/ψ in an equilibrating parton plasma with initial conditions given as set 2 (solid) and set 3 (dashed) in Table I.

Since there is considerable uncertainty in the estimate of the initial parton production by HIJING Monte Carlo simulations, as discussed in Ref. [9], we would like to test the sensitivity of J/ψ suppression to this. We therefore multiply initial parton number densities by a factor of 4, thus increasing initial parton fugacities. We denote such initial conditions as set 2 in Table I. If the uncertainties in initial conditions are caused by soft parton production from the color mean fields, the initial effective temperature will decrease. Therefore, we can alternatively increase the initial parton density by a factor of 4 and at the same time decrease T_0 to 0.4 and 0.72 GeV at RHIC and LHC energies, respectively. This leads to higher initial fugacities, listed as set 3 in Table I. The corresponding survival probabilities calculated with these two sets of initial conditions are shown in Fig. 5. We can see that the J/ψ suppression is much stronger if the initial parton densities are higher. Comparing the solid and dashed lines, however, shows that the J/ψ suppression is less sensitive to the variation of the initial temperature and fugacities as far as the parton densities are fixed.

IV. CONCLUSIONS

To summarize, we have used the cross section of J/ψ dissociation by gluons to calculate the J/ψ suppression in an equilibrating parton gas produced in high-energy nuclear collisions. The large average momentum in the hot gluon gas enables gluons to break up the J/ψ, while hadron matter at reasonable temperature does not provide sufficiently hard gluons. We find a substantial J/ψ suppression in such a nonequilibrium partonic medium; however, it is smaller than that in a fully equilibrated parton plasma. In particular, in an equilibrating plasma the behavior of the J/ψ-gluon cross section at high gluon momenta reduces the J/ψ suppression at large P_T.

In addition to the J/ψ dissociation during the equilibration of the parton plasma, there are other possible sources of suppression for the actually observed J/ψ's. As already noted, nuclear modifications of the $Q\bar{Q}$ production process, e.g., through modified gluon distributions in a nucleus [33,34], multiple scattering accompanied by energy loss [35], or a suppression of the nascent J/ψ before it forms an actual physical resonance [29] must be taken into account. Such effects would cause J/ψ suppression in addition to what we have obtained from the equilibrating parton plasma and modify the transverse momentum dependence of J/ψ suppression [36]. Moreover, interactions between gluons and $c\bar{c}$ bound states before the kinetic thermalization of the partons could also lead to a substantial J/ψ suppression; this would depend on the time needed to achieve local momentum isotropy. On the other hand, gluon fusion could also result in J/ψ production during the evolution of the parton system, similar to the enhancement of open charm. Although studies of preequilibrium open charm production indicate [13,15] that J/ψ production during the parton evolution is not significant compared to primary production, a consistent study of J/ψ suppression should include the preequilibrium production in a form of a master rate equation.

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