J/ψ Enhancement by Quark-Gluon-Plasma Formation

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A new mechanism for producing J/ψ's in a quark-gluon plasma is explored. A $c\bar{c}$ pair created in the hard part of a hadronic collision need not be initially formed as a color singlet to result in a charmonium state if a plasma is also present. Instead, a color-octet pair can subsequently scatter off of the plasma to lose its color. It is shown that if a hot plasma is formed in a relativistic heavy-ion collider, then a substantial enhancement of the J/ψ production rate might be expected from this mechanism. Thus this production mechanism acts to counteract the decrease in cross section arising from Debye screening of the color-binding force in the plasma.

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J/ψ production in heavy-ion collisions has attracted much attention in the last year, in large part due to its recently observed suppression at CERN.[1] Matsui and Satz[2] originally predicted such a suppression as a signal for quark-gluon-plasma formation, their idea (suggested from lattice studies of the static Coulomb potential) being that the attractive force of the $c\bar{c}$ pair will be dissolved in a plasma due to Debye screening. However, subsequent analyses have suggested that the same suppression can occur without the existence of a plasma. One mechanism[3] involves charmonium states scattering from a dense hadronic gas, another[4] involves coalescence of the charm quarks with comoving light quarks.

There are many reasons that a plasma need not lead to suppression. First, the lattice studies might be misleading.[5,6] They do not apply to quarks in motion relative to the plasma. Clearly the binding energy of a fast $c\bar{c}$ pair cannot be seriously affected. Whereas one expects Debye screening to be an important factor for slow-moving J/ψ's, the exact demarcation between fast and slow relative velocities is not truly known. In addition, weak attractive forces that do not show up in lattice studies may still be present, holding the heavy quarks together until the plasma itself dissolves and allows for normal binding to occur. Second, a delicate balance concerning the lifetime of the plasma and the formation time of the J/ψ is required for the suppression to work. The $c\bar{c}$ pair need to have drifted far enough apart and then have enough time to interact with the plasma fairly strongly before their attractive forces are screened. It is by no means obvious that such conditions presently prevail at CERN or will be fulfilled at the BNL Relativistic Heavy Ion Collider (RHIC).

Finally, we come to the issue to be addressed here. In all these studies it has been assumed that the initial production of a color-singlet $c\bar{c}$ pair is not affected by plasma formation. That the heavy quarks are formed via some hard interaction in the early stages of the collision, before the plasma has formed, seems uncontestable. However, that they already need to be a color singlet in order to form a charmonium state is not. A plasma provides a large soup in which a passing $c\bar{c}$ color-octet state can rescatter and become a color-singlet state. A plasma thus potentially provides a mechanism by which J/ψ production can be enhanced.

In order to estimate this enhancement, consider the exchange of a single gluon with the plasma, in particular, the three-gluon-fusion graph for direct J/ψ production, Fig. 1 (and all the possible permutations thereof).[7,8] We imagine gluons g_1 and g_2 as part of the colliding hadrons. The third gluon, g_3, is either absorbed from the plasma or emitted (with an increased rate in the presence of the plasma due to stimulated emission). As there are other graphs of the same order by which a J/ψ can be produced, namely those in which a different charmonium state is first formed and which then decays electromagnetically into a J/ψ, the restriction to direct J/ψ production should be understood as truly an estimate and only a lower bound for the possible effect. We will return to this point later. It should be noted that the ideas here for an enhanced J/ψ signal have much in common with, and are in many ways a hybridization of, the works of Appel[9] on the acoplanarity of jets due to a plasma, and Clavelli *et al.*[10] on J/ψ production in nuclear collisions.

The differential cross section for two colliding hadrons

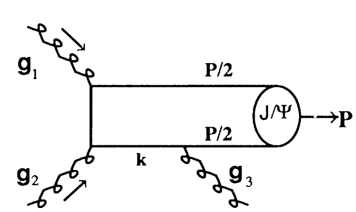

FIG. 1. The three-gluon graph for direct J/ψ production, with all possible permutations of the three gluons understood. Gluons g_1 and g_2 are each part of one of the colliding nuclei, g_3 is either emitted into or absorbed from the plasma.

to produce a J/ψ plus anything $[h_1(P_1)+h_2(P_2)\to J/\psi(P_J)+X]$ by the process in Fig. 1 is

$$\frac{d\sigma^{\rm tot}}{d\hat{s}}=\int dx_1\,dx_2\,\delta(x_1x_2s_{\rm tot}-\hat{s})G(x_1)G(x_2)\int dt\,\frac{d\tilde{\sigma}}{dt}(g_1g_2\to J/\psi(P_J)\pm g_3)\rho^{(a,e)}(E_3)\,, \tag{1}$$

in which a sum over both absorption and emission (signified by a $-g_3$ and a $+g_3$, respectively, in $d\tilde{\sigma}/dt$) is to be understood. x_1 and x_2 are the normal Bjorken variables ($g_1=x_1P_1$, $g_2=x_2P_2$), $G(x_1)$ and $G(x_2)$ are the corresponding gluon structure functions, $s_{\rm tot}$ is the total invariant energy of a hadronic collision $[s_{\rm tot}=(P_1+P_2)^2]$, $\hat{s}$ is the invariant energy of the parton subprocess $[\hat{s}=(g_1+g_2)^2]$, t is the Mandelstam variable $[t=(g_1-P_J)^2]$, and $\rho^{(a,e)}(E_3)$ is the appropriate thermal distribution function for g_3 for the process considered (absorption or emission). The differential cross section of the hard process is given by[7,8]

$$\frac{d\tilde{\sigma}}{dt}(g_1g_2\to J/\psi\pm g_3)=C\frac{1}{16\pi\hat{s}^2}\frac{\hat{s}^2(\hat{s}^2-M^2)^2+t^2(t-M^2)^2+u^2(u-M^2)^2}{(\hat{s}^2-M^2)^2(t-M^2)^2(u-M^2)^2}\,, \tag{2}$$

where M^2 is the mass of the J/ψ, $\hat{s}+t+u=M^2$, and the constant C is given by

$$C=5\pi^2\alpha_s^3M^3\frac{\Gamma(J/\psi\to e^+e^-)}{\alpha^2}\,,$$

where α_s is the strong coupling constant, α is the fine-structure constant, and Γ is the leptonic-decay width of the J/ψ. The thermal distribution functions $\rho^{(a,e)}$ are obtained by squaring the matrix element (appearing in the amplitude of the hard process) for a state of n gluons to make a transition into a state of $n-1$ or $n+1$ gluons via absorption or emission,

$$\langle\phi_{n-1}|a|\phi_n\rangle=\sqrt{n}\,,\quad\langle\phi_{n+1}|\hat{a}|\phi_n\rangle=(n+1)^{1/2}\,,$$

and then taking its Boltzmann average to obtain the familiar result[11]

$$\rho^a(E)=\frac{\sum ne^{-n\beta E}}{\sum e^{-n\beta E}}=\frac{1}{e^{\beta E}-1}\,,$$

$$\rho^e(E)=\frac{\sum(n+1)e^{-n\beta E}}{\sum e^{-n\beta E}}=1+\frac{1}{e^{\beta E}-1}\,.$$

It should be noted that in a related work[12] calculating J/ψ production purely from a plasma, the authors are missing the important thermal enhancement factor ρ^e in their emission graph and completely overlooked the possibility of an absorption graph.

The reason for the peculiar form of the differential cross section above (taken with respect to $\hat{s}$) is that to estimate the production of J/ψ by this mechanism one must insure that the plasma has had sufficient time to form. Examining Fig. 1, this restricts, via the uncertainty principle, how far the leg k can be off shell. Estimating the plasma formation time as being[13] $\approx 1/T$ (where T is the temperature of the plasma), we require that

$$|k^2-m_c^2|=\tfrac{1}{2}|\hat{s}-M^2|\le T^2\,, \tag{3}$$

where $m_c\approx M/2$ is the mass of the charm quark. Because the expression for $d\tilde{\sigma}/dt$ [Eq. (2)] has a well-defined (i.e., finite) infrared limit, this condition can be meaningfully enforced.

In order to judge the magnitude of this mechanism for producing J/ψ's, we must have an estimate for its expected production without a plasma. Because of condition (3), the J/ψ's are produced with low transverse energy. The best model for creating such low-$p_\perp$ J/ψ's under normal (i.e., nonplasma) conditions is the two-gluon-fusion model of Carlson and Suaya[14] in which two gluons fuse to create first a ^3P_j charmonium state which subsequently decays electromagnetically into a J/ψ. The total cross section for this process is given by[14]

$$\sigma(h_1+h_2\to J/\psi+X)=\frac{\pi^2}{8Ms_{\rm tot}}\Gamma_{\rm eff}\int\frac{dx}{x}G(x)G(\tilde{\tau}/x)\,, \tag{4}$$

where $\tilde{\tau}=M^2/s_{\rm tot}$, and $\Gamma_{\rm eff}$ is given by

$$\Gamma_{\rm eff}=\sum_{j=0}^{2}(2j+1)B(^3P_j\to J/\psi+\gamma)\Gamma(^3P_j\to 2g)\,,$$

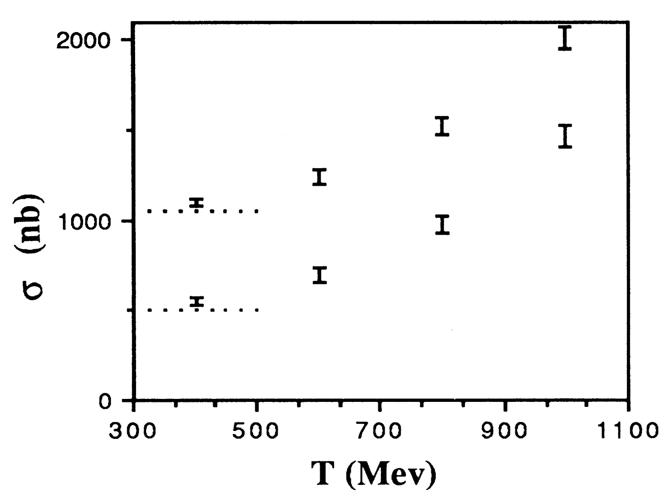

FIG. 2. The total cross section as a function of temperature. The results of the three-gluon graph have been added to the background two-gluon-fusion-model results (which appear as dotted lines). The two sets of data correspond to two different values of $\Gamma_{\rm eff}$ used in calculating the background.

where B is a branching ratio. There is, however, an uncertainty in this model in the magnitude of $\Gamma_{\rm eff}$ by over a factor of 2; 1.33 MeV $\le \Gamma_{\rm eff} \le$ 2.76 MeV [see Ref. 11 of Ref. 14(b)]. The cross section [Eq. (4)] has thus been calculated for both values of $\Gamma_{\rm eff}$, the results of which have been included in Fig. 2 where they appear as dotted lines.

The results for the total production rate [i.e., the sum over both the absorption and emission graphs of Eq. (1), and the two-gluon-fusion-model result, Eq. (4)] at RHIC energies ($P=100$ GeV) have been plotted in Fig. 2 as a function of the plasma's temperature, where the plasma has been assumed to be formed at rest in the c.m. frame of the colliding nuclei. E_3, the argument of the thermal distribution function appearing in Eq. (1), is given, under these assumptions concerning the plasma, by

$$E_3=\left|(x_1+x_2)P-\left(\frac{M^2-u}{4Px_2}+\frac{M^2-t}{4Px_1}\right)\right|.$$

The gluon structure functions that have been used are $xG(x)=3(1-x)^5$ and the strong coupling constant, $\alpha_s=0.41$, as recommended by Ref. 8. Defining $\Delta=|\hat{s}-M^2|$ (in the case of gluon bremsstrahlung $\hat{s}\ge M^2$, while in the case of absorption, $\hat{s}\le M^2$) the error bars in Fig. 2 estimate the error of integrating Eq. (1) to its kinematic limit, $\Delta=0$. At each temperature, the value of the integrand (i.e., $d\sigma/d\hat{s}$) is fairly constant over the range of integration ($2T^2\ge\Delta\ge 0$), decreasing by less than 1% at $T=600$ MeV, and by up to 15% at $T=1000$ MeV. The steep rise in $\sigma_{\rm tot}$ is due to the quadratic expansion of the integration region combined with an almost linear increase of $d\sigma/d\hat{s}$ with temperature. At all temperatures and all Δ, $d\sigma/d\hat{s}$ receives roughly equal contributions from the absorption and emission graphs, each of which, even at the lowest temperatures, are over an order of magnitude larger than the bremsstrahlung case when calculated without any plasma.

Examining Fig. 2, we see that in the case of a hot plasma a large enhancement in the total yield of J/ψ's is predicted. At $T=800$ MeV, we estimate an enhancement of $\approx 40\%$ for $\Gamma_{\rm eff}=2.76$ MeV, and one of $\approx 80\%$ if $\Gamma_{\rm eff}=1.33$ MeV. Of course this assumes that all the charmonium states survive the plasma (the effect of either a dense hadronic gas or comoving partons then has to be superimposed upon this enhancement). It should be noted that this effect disappears when increasing the mass of the heavy quarks. Comparing Eqs. (1) and (4), simple power counting gives that their ratio goes like $(T/M)^4$ (multiplied by a large coefficient arising from the Bose-Einstein distribution function). That is, the effect we are considering is higher twist. However, in the case of charmonium the charm's mass is small enough so that one should not necessarily expect these corrections to be small (and indeed they are not in our case due to the Bose-Einstein factor). In fact, one might say that the whole reason for discussing J/ψ production (as opposed to, say, Υ or toponium production) in the context of a plasma relies on the fact that the charm quark's mass is comparable to the plasma temperature.

One might wonder about the further influence of the plasma on the $c\bar{c}$ pair. The effects on their binding energy have already been discussed in the opening. These lead to a modification of the charmonium wave functions which enter Eqs. (1) and (4) from their free-space values used here. Although their importance is certainly acknowledged, no attempt has been made to incorporate these effects here, in large part due to the complexity of the problem. There still remains the possibility of further gluon attachments to Fig. 1 which could subsequently destroy the $c\bar{c}$ color-singlet state. The number of such multiple scatterings clearly depends upon the space-time evolution of the plasma as to how many interactions are physically possible. However, unlike in the one-gluon-interaction case, the $c\bar{c}$ pair, now being in a color-singlet state, will begin to interact amongst themselves in an attractive, potentially strong, manner. The work of Hansson, Lee, and Zahed[5] indicates that under such conditions the J/ψ states will not be appreciably suppressed due to further absorption of individual gluons from the plasma. It is therefore plausible that most of the J/ψ's will escape. One might likewise wonder if the ^3P_j state's formation by two-gluon fusion (which provides the dominant background of J/ψ's) might not be interfered with by one-gluon scattering off of the plasma. If this indeed happens, then one possible signal for plasma formation would be a suppression of these states accompanied by an anomalously large yield of J/ψ's. However, since the $c\bar{c}$ color-singlet pair forming a ^3P_j state has a lag time $\approx 1/T$ before the plasma appears, a significant attractive force will have probably already been formed between the pair, in which case the work of Hansson, Lee, and Zahed suggests that these states will also not be seriously suppressed.

There still remains one last source of suppression which the plasma provides. Hitherto we have focused only on the effects of the gluons in the plasma upon J/ψ production. What about the quarks? Borrowing the ideas of Brodsky and Mueller in Ref. 4, one expects that comoving plasma quarks will tend to coalesce with the heavy quarks (either before or after they are a color singlet), thus resulting in open-charm production. (It should be noted though that coalescence occurs independent of whether a plasma is or is not formed. Indeed, it is one of the nonplasma mechanisms that is perhaps responsible for the suppressed J/ψ yield observed at CERN.[1] What is being addressed here is the coalescence of the heavy quarks with light, thermalized, plasma quarks.) A full treatment of the problem should clearly include this competing effect; however, it is believed this effect is relatively small compared to the enhancement mechanism presented here as there are a lot more gluons than quarks in a plasma due to their different thermal distribution functions.

It should be finally noted that the calculation here of direct J/ψ production is at best an estimate and only a lower bound of the possible enhancement a plasma could provide by this mechanism. A static plasma, formed in the c.m. frame of the colliding nuclei, has been used. Calculating with an expanding plasma should increase the effect, as can be seen by considering the extreme case that the plasma is comoving with the produced J/ψ. Then the argument of the thermal distribution functions entering Eq. (1) would be Lorentz boosted to a lower value than that given in the static-plasma case, thus increasing the overall effect. In addition, and probably more importantly, direct production is not the only means by which three gluons can create a J/ψ state.[8] There exists a number of other charmonium states they can first form, producing a J/ψ residually via the electromagnetic decay of these excited modes (as was the case in the two-gluon-fusion model). However, unlike Eq. (2), the hard-scattering differential cross section of some of these QCD processes are infrared sick, and therefore an additional analysis of their structure, along the lines of factorization, is first required before their contributions can be estimated. All these effects are presently being studied.

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