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Large p_T J/ψ Production in the Central Region of AA and pp Collisions.

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ABSTRACT

We analyse large p_T J/ψ production in the central rapidity region of nucleus-nucleus (AA) and proton-proton (pp) collisions, in the framework of perturbative QCD. The nuclear effects are assumed to come entirely from the structure functions. We find qualitative similarities with the expected signature of a quark gluon plasma. The cross section for J/ψ production in AA collisions relative to that in pp collisions is suppressed more at small p_T than at large p_T . The Drell-Yan production of muon pairs, on the other hand, shows very little p_T dependence.

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The formation of a quark-gluon plasma (QGP) in relativistic heavy ion (RHI) collisions could be signalled by a suppression of the J/ψ production cross section relative to that in collisions of two protons [1]. Apart from the decrease in the total cross section, there is also expected to be a specific p_T dependence of the suppression [2]. It is expected that the ratio

$$R_\psi = \left(\frac{d\sigma}{dp_T} \right)_{AA \rightarrow J/\psi} / \left(\frac{d\sigma}{dp_T} \right)_{pp \rightarrow J/\psi} \quad (1)$$

is smallest at low p_T , and increases with p_T . This signal has been extensively studied in different hydrodynamic scenarios for RHI collisions [3]. Here we report the results of investigation of a background to this signal [4-5]. Ideally, in order to compare our results with those of the NA38 experiment at CERN [6], we should compute the ratio of $(d\sigma / dp_T)_{AA \rightarrow J/\psi}$ at high E_T and low E_T . However, such an exercise will necessarily be clouded by the assumptions (and parameters) of the model used to generate the E_T dependence, since it is experimentally known that most of the E_T in these events comes from soft QCD interactions. These are highly non-perturbative, and there are several non-equivalent ways of modelling these interactions. However, it appears plausible that a model which uses a partonic picture for producing the E_T and incorporates energy-momentum and colour conservation should yield non-trivial E_T dependence of $(d\sigma / dp_T)_{AA \rightarrow J/\psi}$. We believe that, as a result of these uncertainties, it is a good idea to examine the proposed signature of QGP in a theoretically more well-defined scheme, by restricting ourselves to R_ψ , i.e. to the original suggestions in Ref. 1 and Ref. 2. We hope to come back to the question of E_T -dependence in the near future.

The processes that create charm pairs in RHI (or even in proton-proton) collisions occur very early during the collision. The hadronisation of such a pair into charmonium is, on the other hand, a slow process. If a QGP is formed in such a collision, then the time taken for this system to thermalise is normally expected to lie between these two scales. The time scale for the production of the pair is less than $1/2m_c$ where m_c is the mass of the charm quark. Taking this to be 1.5 GeV, the time scale for production turns out to be less than 0.07 fm. The other two time scales are both hadronic and should be of the order of a few fermis.

This is the reasoning behind a perturbative QCD analysis of the process of J/ψ production. Due to the separation of the time scales for pair production and hadronisation, we can write the cross section for J/ψ production as a convolution of the perturbatively computed cross section for pair production, and a hadronisation model. The proposed signal [1-2] essentially looks at the effect of a hot environment on the hadronisation of the pair. We look at a background that arises from the well-known modification of parton densities in the nucleus—the so-called EMC-effect [7]. This affects pair production, and should be present over and above any hadronisation scheme, including those of Ref. 1 and Ref. 2. We find that this causes R_ψ to be small at small p_T , and grow with p_T . At $p_T > 2$ GeV, we find that R_ψ can become larger than unity, thus implying an enhancement in the J/ψ production cross section. We have also computed the ratio

$$R_{DY} = \left(\frac{d\sigma}{dp_T} \right)_{AA \rightarrow \mu\mu} / \left(\frac{d\sigma}{dp_T} \right)_{pp \rightarrow \mu\mu}, \quad (2)$$

under the same kinematic conditions and find that it is almost independent of p_T .

In a previous paper [4] we analysed R_ψ at small (essentially zero) p_T . The model consisted of taking the lowest-order perturbative QCD diagrams [8] for charmed pair production, and computing the corresponding hadron cross sections by convoluting over parton distributions. Since hadronisation of the pair (giving J/ψ of small p_T) can only occur through uncomputable (in perturbative QCD) soft gluon radiation, its effect is simulated by multiplying the pair production cross section by a phenomenological parameter, after constraining the invariant mass of the pair to lie below the threshold for open charm production ($4m_D^2$, where m_D is the mass of the D-meson). This approach has considerable success in describing J/ψ production in the central rapidity region of pp collisions [9], and is equivalent to semi-local duality computations of the same process.

A J/ψ at large p_T can be produced if the pair radiates a hard gluon, or an extra particle is produced in the basic reaction. Along with this hard radiation there will also be soft radiation involved in the hadronisation process of such a pair. The result is that we still have to compute the J/ψ production cross section as an open-charm production process, with the usual addition of a hadronisation model. One should also note that the soft radiation [10] can in principle carry large angular momentum, hence a helicity projection of the full perturbative QCD cross section into a spin-one channel is not called for. In

short, our model for J/ψ production remains as it was in the lowest-order computation. We compute the open charm production cross section restricted to the range $4m_c^2$ to $4m_D^2$ in the invariant mass of the pair, and multiply it by a phenomenological factor to simulate the hadronisation. Since our computation is meant to check only the initial state effects, this factor is the same in both pp and AA collisions, and hence cancels in the ratio R_ψ . This model can be used to predict the cross sections in pA collisions as well, a problem that we will discuss elsewhere.

Large p_T charm pair production involves three subprocesses—

$$q \bar{q} \rightarrow c \bar{c} g, \quad g q (\bar{q}) \rightarrow c \bar{c} q (\bar{q}), \quad \text{and} \quad g g \rightarrow c \bar{c} g. \quad (3)$$

The Feynman graphs corresponding to these are shown in Fig. 1. The corresponding amplitudes have been computed [11]. We do not reproduce them here since the expressions are rather lengthy. The phase space for these processes is 6-dimensional, since they involve 3 particles in the final state. We choose to work with the EHLQ conventions [12], modified for the case of some of the final state particles being massive.

The kinematics is defined in the centre of mass of the three final state particles. The energy in the subprocess in this system is given by the usual relation in terms of the Bjorken fractions x_1 and x_2 carried by the initial particles

$$s = x_1 x_2 S,$$

where S is the energy in the hadronic CM. The boost between these two systems can be written in terms of a boost rapidity

$$y_{boost} = \frac{1}{2} \log \left(\frac{x_1}{x_2} \right).$$

The initial partons come in with 4-momenta

$$\frac{\sqrt{s}}{2} (1, \sin \theta \cos \phi, \sin \theta \sin \phi, \cos \theta) \quad \text{and} \quad \frac{\sqrt{s}}{2} (1, -\sin \theta \cos \phi, -\sin \theta \sin \phi, -\cos \theta).$$

The xy plane contains the final particles, and the jet axis defines the +ve x-direction. The 4-momenta of c , $\bar{c}$ and the jet are respectively

$$\begin{aligned} p_3 &= \frac{\sqrt{s}}{2} x_3 (1, \beta_3 \cos \theta_{35}, \beta_3 \sin \theta_{35}, 0), & p_4 &= \frac{\sqrt{s}}{2} x_4 (1, \beta_4 \cos \theta_{45}, \beta_4 \sin \theta_{45}, 0), \\ p_5 &= \frac{\sqrt{s}}{2} x_5 (1, 1, 0, 0), & \text{where } \beta_i &= \sqrt{1 - 4m_c^2/x_i^2 s}. \end{aligned}$$

Energy conservation requires that

$$x_3 + x_4 + x_5 = 2, \quad (x_i \leq 1, i = 3, 4, 5)$$

The angles θ_{35} and θ_{45} are given by

$$\cos \theta_{35} = \frac{(\beta_4 x_4)^2 - (\beta_3 x_3)^2 - x_5^2}{2\beta_3 x_3 x_5} \quad \text{and} \quad \cos \theta_{45} = \frac{(\beta_3 x_3)^2 - (\beta_4 x_4)^2 - x_5^2}{2\beta_4 x_4 x_5}.$$

We choose as the phase-space coordinates the six quantities x_1, x_2, x_3, x_4, ϕ , and p_T of the jet (which is the same as the p_T of the J/ψ). In terms of the variables already defined, p_T is given by the relation

$$p_T = \frac{\sqrt{s}}{2} x_5 \sqrt{\cos^2 \theta + \sin^2 \theta \sin^2 \phi}.$$

The hadronic cross section can then be written as

$$\frac{d\sigma}{dx_1 dx_2 dx_3 dx_4 d\phi dp_T} = \frac{S \cos \theta \cos^2 \phi}{2p_T \sqrt{x_1 x_2}} \sum_p F_p(x_1, x_2) |A_p|^2, \quad (4)$$

where the summation is over the three sub-processes shown in Eq. (3) and A_p is the amplitude for the subprocess p . The structure function combinations $F_p(x_1, x_2)$ appropriate to each set of subprocesses are shown in Table 1, and the expressions for $|A_p|^2$ are taken from Ref. 11.

The quark distributions in nuclei are very well measured now, and there exist several models, all very well constrained by data, which give phenomenological parametrizations for these. The nuclear dependence of gluon densities is, unfortunately, not yet available. However, several experiments which have already finished running may soon provide this data [13]. In the absence of this data¹ we have to rely on a model. We choose for this purpose the gas model [14]. This is a two-parameter model of parton densities in a nucleus. Both are determined from structure function measurements, and then provide a parameter-free prediction of the gluon density. We use this model because it does well in reproducing

¹ It is of interest to note that direct photon production in nucleus-nucleus collisions measures the same gluon densities that are needed for this computation. Hence the results from experiments like WA80 can, in principle, be used to discriminate between different models of initial state effects.

the A-dependence of large p_T jet cross sections [15], a process which is gluon dominated. We emphasise that the model is used only because the data on gluon distributions are not available yet. In the near future this model dependence can be removed.

We constrain the J/ψ to lie in the central region of rapidity ($-1 \leq y_\psi \leq 1$ at CERN energies), and introduce the cut mentioned above on the invariant mass $m_{c\bar{c}}$ ($4m_c^2 \leq m_{c\bar{c}} \leq 4m_D^2$). Validity of the perturbative QCD computation requires that the scale p_T not be very small. Hence, we constrain it to be larger than 1 GeV. These cuts do not allow the value of Bjorken x to drop to low values, and we are prevented from seeing the effects of the little-understood “shadowing region” of nuclear structure functions.

We also compute, for comparison, the cross sections for dimuon production in the same kinematic region. This is important because a J/ψ is tagged by its decay into dimuons. An estimate of the A-dependence of the Drell-Yan continuum is necessary to eliminate this background. The appropriate cross section has been known for a long time [16]. The pair annihilation cross section is

$$\frac{d\sigma^A}{dM^2 dt} = \frac{8}{27\pi} \frac{\alpha^2 \alpha_s}{M^2 s^2} \frac{(t^2 + u^2 + 2M^2 s)}{tu}, \quad (5a)$$

and the QCD Compton cross section is

$$\frac{d\sigma^C}{dM^2 dt} = \frac{1}{9\pi} \frac{\alpha^2 \alpha_s}{M^2 s^2} \left[\frac{(s^2 + u^2 + 2M^2 t)}{s(-u)} + (t \leftrightarrow u) \right], \quad (5b)$$

where s , t , and u are the Mandelstam variables in the subprocess and M^2 is the invariant mass of the pair. The hadronic cross section is—

$$\frac{d\sigma^{AB}}{dM^2 dt dx_1 dx_2} = H^{AB}(x_1, x_2) \frac{d\sigma^A}{dM^2 dt} + G^{AB}(x_1, x_2) \frac{d\sigma^C}{dM^2 dt} \quad (6a)$$

where,

$$H^{AB}(x_1, x_2) = \frac{1}{x_1 x_2} \sum_f e_f^2 \left[q_f^A(x_1) \bar{q}_f^B(x_2) + \bar{q}_f^A(x_1) q_f^B(x_2) \right], \quad (6b)$$

$$G^{AB}(x_1, x_2) = \frac{1}{x_1 x_2} g^A(x_1) \sum_f e_f^2 \left[q_f^B(x_2) + \bar{q}_f^B(x_2) \right] + \frac{1}{x_1 x_2} g^B(x_2) \sum_f e_f^2 \left[q_f^A(x_1) + \bar{q}_f^A(x_1) \right] \quad (6c)$$

The cuts on this cross section were identical to those placed on the J/ψ cross section— $4m_c^2 \leq M^2 \leq 4m_D^2$, $-1 \leq y_{\mu\mu} \leq 1$, and $p_T \geq 1$ GeV.

The results of our computations of R_ψ are shown in Fig. 2. It is interesting to note that the p_T -dependence of R_ψ is similar to that predicted by the plasma models. The maximum suppression occurs at $p_T = 0$, and the ratio R_ψ increases with p_T . At a p_T of about 2 GeV, the ratio is unity, and at larger p_T there is an enhancement of the AA cross section over the pp cross section. Our results for R_{DY} are also shown in Fig. 2. This ratio has almost no p_T dependence, in sharp contrast to R_ψ .

In order to understand these trends, we recall the behaviour of the ratio of parton densities in the nucleus and a proton—

$$\rho(x) = \frac{q^A(x)}{q^p(x)}. \quad (7)$$

This is observed to be a decreasing function of x for $0.2 \leq x \leq 0.6$. It has a turnover and increases very rapidly for larger values of x . For J/ψ production at p_T of zero, the range of Bjorken- x that is relevant is $0.057 \leq x \leq 0.52$. On the other hand, at $p_T = 3$ GeV, the relevant range is $x \geq 0.10$; and weighted with the cross sections, the average value of ρ is larger than unity. In effect, as we increase p_T we probe the parton densities at successively larger values of the Bjorken variables. We also found that the subprocess whose contribution to the cross section is largest at large p_T is $qg \rightarrow c\bar{c}q$. In contrast, at smaller p_T it is the gluon-gluon subprocess that dominates. The final value of R_ψ is a payoff between these two factors.

The cause of the difference between R_ψ and R_{DY} is many-fold. One is in the range of Bjorken- x probed by this process. For the Drell-Yan cross sections at $p_T > 3$ GeV, the upper limit of x is unity, and the lower limit changes from 0.055 for $p_T=1$ GeV, to 0.08 for $p_T=3$ GeV. The mean value of the ratio ρ is smaller when weighted by the cross sections. This has to do with the fact that we observe only two hard particles here, whereas we observe three hard particles in J/ψ production. Another reason lies in the dominant subprocess, which in this case is the gluon-gluon process (contributing to more than 70% of the cross section at even the lowest p_T). Hence the structure function combination that we observe here is different from that in J/ψ production.

We do not find a strong dependence of R_ψ on the beam-target system. The numbers for oxygen-uranium collisions are similar to those for uranium-uranium collisions. This is also understandable, since the observed variation of ρ with A is weak.

In Fig. 3 we show the p_T -distribution of J/ψ . Although the figure shows only the distributions for pp and oxygen-uranium collisions, it serves to illustrate the differences between AA and pp collisions. We see a roughly exponential fall in these distributions. The behaviour of R_ψ is reflected in the feature that the AA distribution lies below the pp curve at small p_T , and crosses over at larger p_T .

In conclusion, we find that the three major qualitative aspects of the J/ψ suppression signature for QGP can be expected to arise on the basis of the structure function effects alone. Thus R_ψ , defined in Eq. (1), evaluated in the central rapidity region at the CERN energy of $\sqrt{s} = 20$ GeV/A is about 0.5–0.6 at very small p_T . It rises monotonically as p_T is increased, while the experimental background of the dimuon continuum over which the J/ψ peak is observed shows essentially no p_T dependence. We also find that for $p_T > 2$ GeV there is a possibility of enhancement of J/ψ production. Our computation yields a clear picture of these effects in terms of the underlying parton dynamics alone. It will be interesting to see whether the kinematic constraints due to a given E_T can be such that this picture remains valid for the NA38 data as well.

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TABLE

Table 1: The structure function combinations required in the hadronic cross sections for the $AB \rightarrow J/\psi + \text{jet}$. Quark number densities are denoted by q , antiquarks by $\bar{q}$, and gluons by g . The summations are over flavours of quarks. We take them over u, d, and s.

Initial State	Parton density combinations
$q \bar{q}$	$\sum_f \left[q_f^A(x_1) \bar{q}_f^B(x_2) + \bar{q}_f^A(x_1) q_f^B(x_2) \right]$
$g q$	$g^A(x_1) \sum_f \left[q_f^B(x_2) + \bar{q}_f^B(x_2) \right] + g^B(x_2) \sum_f \left[q_f^A(x_1) + \bar{q}_f^A(x_1) \right]$
$g g$	$g^A(x_1) g^B(x_2)$

FIGURE CAPTIONS

Fig. 1 Typical lowest-order QCD diagrams for the processes (a) $q\bar{q} \rightarrow c\bar{c}g$, (b) $qg \rightarrow c\bar{c}q$, (c) $gg \rightarrow c\bar{c}g$. These give rise to charmed pairs with large pair pt. The complete set of diagrams can be obtained by properly symmetrising the ones displayed.

Fig. 2 We display here (a) R_ψ as a function of p_T (the point at $p_T = 0$ is from Ref. 4), and (b) R_{DY} as a function of p_T . In both sets of computations $m_c = 1.5$ GeV. Circles denote the ratio for oxygen-uranium collisions, triangles for sulphur-uranium, and diamonds for uranium-uranium. The error bars denote the uncertainties in the Monte Carlo integrations for the cross sections.

Fig. 3 $(1/N) d\sigma/dN$ (in arbitrary units) as a function of p_T . The full line is for the pp system, and the dashed line for oxygen-uranium. The curves are normalised so that the areas under them are equal.

Fig. 2

Fig. 3