Nonsaturation of the J/ψ Suppression at Large Transverse Energy in the Comovers Approach

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We show that, contrary to recent claims, the J/ψ suppression resulting from its interaction with comovers does not saturate at large transverse energy E_T. On the contrary, it shows a characteristic structure—a change of curvature near the knee of the E_T distribution—which is due to the E_T (or multiplicity) fluctuation. This change of curvature is also present in recent experimental results, although the experimetal effect is larger than in our calculation.

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An interesting result of the last run (1998 data) of the NA50 Collaboration at CERN on the transverse energy (E_T) dependence of J/ψ suppression in Pb-Pb collisions is the observation [1] of a convexity at large E_T. More precisely, for $E_T \gtrsim 100$ GeV (which corresponds to the so-called knee of the E_T distribution [1]), the slope of the ratio $R(E_T)$ of J/ψ over Drell-Yan (DY) cross sections increases with increasing E_T. In sharp contrast with this result, models [2–7] of the J/ψ suppression in non-quark-gluon plasma (QGP) scenarios—such as the one based on the interaction of the J/ψ with comovers—exhibit a clear saturation at large E_T.

In this work, we show that the above feature of the comovers model is true only up to the knee of the E_T distribution ($E_T \sim 100$ GeV). Beyond this value, we enter into the tail of the E_T distribution—where the increase in E_T is due to fluctuations. This fluctuation produces a corresponding increase in the density of comovers—which, in turn, increases the J/ψ suppression at large E_T.

In order to illustrate this phenomenon we use the model introduced in Ref. [2]. Here, as in most non-QGP models, the J/ψ suppression is due to two mechanisms: absorption of the preresonant $c\bar{c}$ pair with nucleons (the so-called nuclear absorption) and the interaction of the J/ψ with comovers. The corresponding J/ψ survival probabilities are given by [2].

$$S^{\mathrm{abs}}(b,s) = \frac{\{1 - \exp[-AT_A(s)\sigma_{\mathrm{abs}}]\}\{1 - \exp[-BT_B(b - s)\sigma_{\mathrm{abs}}]\}}{\sigma^2_{\mathrm{abs}} ABT_A(s)T_B(b - s)}, \tag{1}$$

$$S^{\mathrm{co}}(b,s) = \exp\left\{-\sigma_{\mathrm{co}} N_y^{\mathrm{co}}(b,s) \ln\left(\frac{N_y^{\mathrm{co}}(b,s)}{N_f}\right)\right\}. \tag{2}$$

The survival probability S^{co} depends on the density of comovers $N_y^{\mathrm{co}}(b,s)$ in the rapidity region of the dimuon trigger $2.9 < y_{\mathrm{lab}} < 3.9$ and $N_f = 1.15$ fm^{-2} [2,4] is the corresponding density in pp collisions. In order to compute N_y^{co}, various hadronic models have been used in the literature. For instance, in Ref. [4] it has been assumed that the hadronic multiplicity is proportional to the number of participant nucleons (the so-called wounded nucleon model), while in Ref. [2] a formula based on the dual parton model (DPM) was used—which includes an extra term proportional to the number of binary interactions. In this paper we use the DPM formula [Eq. (6) of [2]]. In both cases, the calculations do not include the fluctuations mentioned above and, therefore, cannot be applied beyond the knee of the E_T distribution—where the increase in E_T (or multiplicity) is entirely due to fluctuations. In order to introduce these fluctuations, it is convenient to recall the other formulas needed to calculate the J/ψ suppression.

At fixed impact parameter b, the J/ψ cross section is given by [2]

$$\sigma_{AB}^{\psi}(b) = \frac{\sigma_{pp}^{\psi}}{\sigma_{pp}} \int d^2s\, m(b,s) S^{\mathrm{abs}}(b,s) S^{\mathrm{co}}(b,s), \tag{3}$$

where $m(b,s) = AB\sigma_{pp}T_A(s)T_B(b - s)$. The corresponding one for DY pair production is obtained from (3) putting $\sigma_{\mathrm{abs}} = \sigma_{\mathrm{co}} = 0$ (i.e., $S^{\mathrm{abs}} = S^{\mathrm{co}} = 1$) and is proportional to AB. In this way we can compute the ratio of J/ψ over DY as a function of the impact parameter. Experimentally, however, the ratio $R(E_T)$ is given as a function of the transverse energy E_T measured by a calorimeter, in the rapidity interval $1.1 < y_{\mathrm{lab}} < 2.3$. In order to compute $R(E_T)$ we have to know the correlation $P(E_T, b)$ between E_T and the impact parameter, which is given by [2]

$$P(E_T, b) = \frac{1}{\sqrt{2\pi q a E_T^{\mathrm{NF}}(b)}} \exp\left[\frac{-[E_T - E_T^{\mathrm{NF}}(b)]^2}{2qaE_T^{\mathrm{NF}}(b)}\right]. \tag{4}$$

Here

$$E_T^{\rm NF}(b) = qN_{\rm cal}^{\rm co}(b) + k[A - m_A(b)]E_{\rm in}\,, \qquad (5)$$

$m_A(b)$ is the number of participants of A (at fixed impact parameter), $E_{\rm in} = 158$ GeV$/c$ is the beam energy, and $k = 1/4000$ [2]. In (4) and (5) $N_{\rm cal}^{\rm co}(b)$ is obtained by integrating the comover density $N_y^{\rm co}(b,s)$ over d^2s and dy (in the rapidity range of the E_T calorimeter). The second term in (5) was introduced in Ref. [2] in order to reproduce the correlation between E_T and the energy $E_{\rm ZDC}$ of the zero degree calorimeter. The value of k was determined from the best fit to this correlation. The second term of Eq. (5) was interpreted as due to intranuclear cascade—which is present here due to the location in rapidity of the NA50 calorimeter. This term is sizable for peripheral collisions, when many spectator nucleons are present, and dies away for central ones. Our results are totally insensitive to this term for $E_T > 60$ GeV—the region we are especially interested in for this work (see Fig. 2). The parameters $q = 0.56$ and $a = 0.94$ are obtained from a fit to the minimum bias E_T distribution. The parameter q gives the relation between multiplicity of comovers (positive, negative, and neutrals) and the E_T of the NA50 calorimeter (which contains only neutrals). The product qa controls the width or the standard deviation of $P(E_T, b)$. [It is not clear *a priori* whether or not the second term of Eq. (5) contributes to this standard deviation. We have checked numerically that our results are practically unchanged when $E_T^{\rm NF}(b)$ in the denominator of the bracket in Eq. (4) is replaced by the first term of Eq. (5).]

The J/ψ and DY cross sections at fixed E_T are then given by

$$\frac{d\sigma^{\psi({\rm DY})}}{dE_T} = \int d^2b\, \sigma_{AB}^{\psi({\rm DY})} P(E_T, b)\,. \qquad (6)$$

The quantity $E_T^{\rm NF}(b)$ in Eq. (5) does not contain fluctuations—hence the index NF. This is obvious from the fact that the parameter a is not present in (5). In order to see it in a more explicit way, we plot in Fig. 1 the quantity

$$F(E_T) = E_T/E_T^{\rm NF}(E_T)\,, \qquad (7)$$

where

$$E_T^{\rm NF}(E_T) = \frac{\int d^2b\, E_T^{\rm NF}(b) P(E_T, b)}{\int d^2b\, P(E_T, b)}\,. \qquad (8)$$

We see that $E_T^{\rm NF}$ coincides with E_T only up to the knee of the E_T distribution. Beyond it, $E_T^{\rm NF}$ is smaller than the true value of E_T. This difference is precisely due to the E_T fluctuation.

As discussed above, in order to compute the ratio $R(E_T)$ beyond the knee of the E_T distribution it is necessary to introduce in $N_y^{\rm co}$ the E_T (or multiplicity) fluctuations responsible for the tail of the distribution. In order to do so, we use the experimental observation that multiplicity and

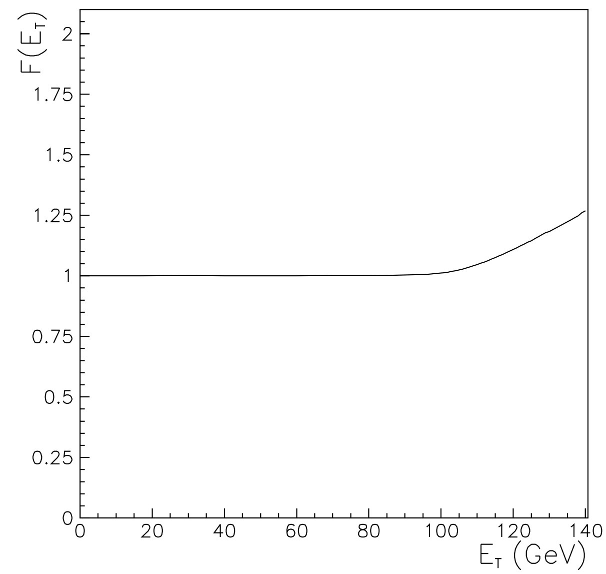

FIG. 1. The ratio $F(E_T)$ in Eqs. (7) and (8).

E_T distributions have approximately the same shape. This indicates that the fluctuations in E_T are mainly due to fluctuations in multiplicity—rather than in p_T. This leads to the following replacement in Eq. (2):

$$N_y^{\rm co}(b,s) \to N_y^{\rm co}(b,s)F(E_T)\,. \qquad (9)$$

In this way the results for the ratio $R(E_T)$ are unchanged below the knee of the distribution (see Fig. 1). Beyond it, the J/ψ suppression is increased as a result of the fluctuation. Note that we are not assuming that the multiplicity of comovers is proportional to E_T. Equation (9) implies only that the fluctuations of these two quantities, beyond the knee of the E_T distribution, are the same. Since $F(E_T)$ is different from unity only for $E_T \geq 100$ GeV, where the correlation between E_T and $E_{\rm ZDC}$ is very loose, the replacement (9) does not practically affect the $E_T - E_{\rm ZDC}$ correlation. Note also that the J/ψ suppression in S-U collisions is not affected by the E_T fluctuation in the region of E_T covered by the NA38 data—because this region does not extend beyond the knee of the E_T distribution.

Since the rapidity acceptances of the zero degree calorimeter and of the dimuon trigger do not overlap, when introducing the effect of the fluctuation via Eq. (9) we are implicitly assuming that the fluctuations in these two rapidity intervals are fully correlated [8]. In string models of the DPM type, the fluctuation in E_T is due to a fluctuation in the number of strings rather than to fluctuations within the individual strings. Because of that, fluctuations in different rapidity intervals are strongly correlated.

We turn next to the numerical results. In Ref. [2] we used for the two parameters of the model $\sigma_{\rm abs} = 6.7$ mb and $\sigma_{\rm co} = 0.6$ mb. In this case, the computed J/ψ suppression at $E_T \sim 100$ GeV is somewhat too small [2].

Clearly, we can increase it by increasing the value of $\sigma_{\rm co}$. However, we then increase the value of the suppression for peripheral collisions. This, in turn, can lead to some conflict with the S-U data (see [3] for a discussion on this point). However, recent data [9] on the J/ψ cross section in pA collisions point to a smaller value of $\sigma_{\rm abs}$—of 4 to 5 mb. With this value of $\sigma_{\rm abs}$, we can increase $\sigma_{\rm co}$ from 0.6 mb up to 1.0 mb without increasing the J/ψ suppression either for peripheral Pb-Pb or for S-U collisions. In Fig. 2 we present the result of our calculation using $\sigma_{\rm abs} = 4.5$ mb and $\sigma_{\rm co} = 1$ mb. (Theoretically, the value of $\sigma_{\rm co}$ is poorly known due to different assumptions on the scattering mechanism. If its value were substantially smaller than 1 mb, as suggested by some authors [10], the absorption by comovers would be too small to explain the data even qualitatively. However, other authors [11] find that for $\sqrt{s}$ roughly 1 GeV above threshold this cross section ranges from 1 to several mb. The uncertainty in the value of $\sigma_{\rm co}$ is a drawback of the comovers approach.) We see that the main features of the data are reproduced. In particular, our curve shows a slight change of curvature at $E_T \sim 100$ GeV, which is entirely due to the effect of fluctuations—and is seen in the 1998 NA50 data [1]. The physical origin of this change in the slope of E_T is the following: when approaching the knee of the E_T distribution from below, the number of participants approaches $2A$ and changes slowly. The latter is also true for the multiplicity of comovers. Beyond the knee, the multiplicity increases faster due to the fluctuations and produces a faster decrease of $R(E_T)$ (see Fig. 1).

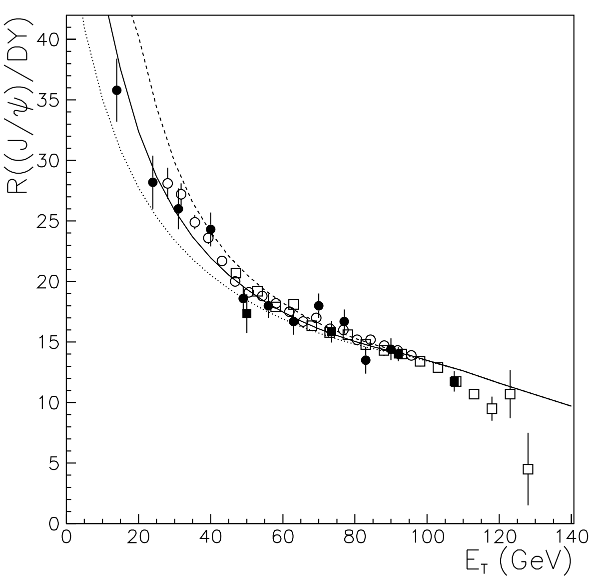

FIG. 2. The ratio $R(E_T)$ of J/ψ over DY cross sections, obtained with $\sigma_{\rm abs} = 4.5$ mb and $\sigma_{\rm co} = 1$ mb, compared to the NA50 data [1] [19]. The full curve corresponds to $k = 1/4000$ [2] in Eq. (5). The dashed curve is obtained with $k = 1/2000$ and the dotted curve with $k = 0$ (see main text). The black points correspond to 1996 Pb-Pb data, the black squares correspond to 1998 Pb-Pb data, the white points to 1996 analysis with minimum bias, and the white squares to 1998 analysis with minimum bias.

Clearly, the structure at $E_T \approx 100$ GeV in our curve is quite small. However, when comparing the curve in Fig. 2 with the corresponding theoretical curves of Fig. 6 of Ref. [1], we see that the description of the data has been considerably improved. Obviously the slope of the J/ψ suppression from the last four data points in Fig. 2 is larger than the one in the theoretical curve. However, it should be stressed that these data points are obtained from the so-called minimum bias (MB) analysis. In this analysis one measures the ratio J/ψ over MB and multiplies this ratio by the theoretical ratio MB/DY. In this theoretical calculation one assumes that the tails of the MB and DY E_T distributions are identical. It would be very important to have an experimental confirmation of this assumption.

We want to stress that the shape of our curve in the lower half of the E_T region [where the ratio $R(E_T)$ changes rather fast with E_T] is sensitive to the relation between E_T and the impact parameter. We see from Eq. (5) that this relation depends on the size of the contribution of the intranuclear cascade (parameter k). As mentioned above, the value $k = 1/4000$ used here was obtained in [2] from the best fit of the correlation between E_T and the energy $E_{\rm ZDC}$ of the zero degree calorimeter. However, since we do not have a totally reliable expression for the latter, there is an uncertainty in the value of k. In order to illustrate its effect on $R(E_T)$, we show in Fig. 2 (dashed line) the result with $k = 1/2000$, i.e., doubling the (comparatively small) contribution of the intranuclear cascade and also with $k = 0$ (dotted line). The effect is concentrated in the lower half of the E_T range. This uncertainty would not be present if the J/ψ suppression were given as a function of either E_T or charged multiplicity at midrapidities.

The mechanism of nonsaturation of the J/ψ suppression at large E_T discussed here has a solid physical basis. In principle, this mechanism should be present in nonequilibrium [6], cascade [12], and transport [13] type Monte Carlo simulations. Thus it is not clear why the J/ψ suppression calculated in [6] and the unscaled one in the second paper of Ref. [12] exhibit saturation at large E_T. On the other hand, no saturation is seen with the late comovers absorption model in the second paper of Ref. [13]—although the calculation extends only up to $E_T = 110$ GeV. Actually, in the latter model the results are similar to ours. The semianalytic discussion presented here should help to clarify the situation.

The J/ψ suppression seen by the NA50 Collaboration can also be explained in a quark-gluon plasma scenario [5,14,15]. However, the structure at $E_T \approx 100$ GeV cannot be quantitatively reproduced even assuming two deconfining phase transitions [14]. At the energies of the Relativistic Heavy Ion Collider (RHIC) at Brookhaven, it will be possible to determine whether this structure is due to QGP formation or to the multiplicity fluctuation. In Ref. [2] an estimate of the value of the J/ψ suppression for central Pb-Pb collisions at the knee of the E_T distribution has been given, using two different Ansätze for the shadowing corrections at RHIC: no shadowing [16] or strong shadowing [17]. In the latter case the central densities are reduced by a factor of 2. The values of the ratio R relative to the p-p ones are 0.03 and 0.1, respectively. The energy densities at RHIC are $\epsilon = 15\ \mathrm{GeV/fm^3}$ and $\epsilon = 7.5\ \mathrm{GeV/fm^3}$, without and with shadowing, respectively. (The effect of the shadowing corrections in our case reduces the energy density at RHIC by a factor of 2 and, therefore, is considerably larger than the 15% to 20% effect obtained by other authors [18]. Note that, in our formalism, the diagrams in the multiple-scattering picture of shadowing also reduce the soft component [17].) On the other hand, the structure at $E_T \approx 100$ GeV seen at SPS corresponds to an energy density $\epsilon = 3\ \mathrm{GeV/fm^3}$ [1]. Since this energy density is much smaller than the one at the knee of the E_T distribution at RHIC, if our interpretation is correct no structure will be present at this value of ϵ.

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