

Anomalous suppression of π^0 production at large transverse momentum in Au + Au and d + Au collisions at $\sqrt{s_{NN}} = 200$ GeV

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Abstract

We propose a model of suppression of large p_T -pions in heavy ion collisions based on the interaction of the large p_T pion with the dense medium created in the collision. The model is practically the same as the one previously introduced to describe J/ψ suppression. Both the p_T and the centrality dependence of the data are reproduced. In deuteron-gold collisions, the effect of the final state interaction with the dense medium turns out to be negligibly small. Here the main features of the data are also reproduced both at mid and at forward rapidities.

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1 Introduction

One of the most interesting results of the heavy ion program at RHIC is the so-called jet quenching [1] [2]. The yield of particles produced in AA collisions at mid-rapidities and large p_T increases with centrality much less than the number of binary collisions $n(b)$. For most central collisions the large p_T yield is suppressed by a factor 4-5 as compared to the result expected from this scaling. This phenomenon is particularly interesting since it is not observed in deuteron-gold collisions at RHIC at mid-rapidities [2].

The suppression of the yield with respect to the scaling in $n(b)$ is well known in soft collisions. In this case the phenomenon is observed in hadron-nucleus and nucleus-nucleus collisions at all energies. It is well described in the framework of string models, such as the Dual Parton Model (DPM) and the Quark Gluon String Model (QGSM), when shadowing corrections are taken into account [3]. With increasing p_T the shadowing corrections decrease [4] and the scaling with $n(b)$ is predicted in perturbative QCD. Actually, the observed increase is even faster due to initial state interactions, the so-called Cronin effect.

At very high energies the shadowing effects are very important for hA and AA -collisions (see for example [3] [4]). These nonlinear effects lead, as $s \rightarrow \infty$, to “saturation” of the distributions of partons in the colliding hadrons and nuclei [5]. Detailed calculations of shadowing effects at RHIC and LHC energies [3], show that these effects are important for the description of inclusive spectra, but the situation is still far from the “saturation” limit. This is true for particles with average momentum transfer. For particles with large momentum transfer, which we study in this paper, the situation is different. It is well known that shadowing effects for partons take place at very small x , $x \ll x_{cr} = 1/m_N R_A$ where m_N is the nucleon mass and R_A is the radius of the nucleus. On the other hand, partons which produce a state with transverse mass m_T and a given value of Feynman x_F , have $x = x_{\pm} = 1/2(\sqrt{x_F^2 + 4m_T^2/s} \pm x_F)$. Thus, at fixed initial energy (s) the condition for existence of shadowing will not be satisfied at large transverse momenta. For example in the central rapidity region ($y^* = 0$) at RHIC and for p_T of jets (particles) above 5(2) GeV/c the condition for shadowing is not satisfied and these effects are absent. It was shown in ref. [6] that in this region, where

$x \gtrsim x_{cr}$, there are in general final state interactions, which can be treated in a simple quasiclassical way. These interactions lead, in particular, to an energy loss of a parton (particle) in the dense medium produced in the collision. The situation is very similar to production of heavy quarkonia in pA and AA collisions [7], where most of the present data correspond to energies below the critical one and a simple probabilistic interpretation can be applied. Note that these “final state” interactions are absent in the shadowing region.

Hadrons loose a finite fraction of their longitudinal momentum due to secondary interactions with hadrons of the nucleus. This is a characteristic property of such theoretical models as DPM and QGSM and agrees with approximate Feynman scaling of inclusive spectra in the fragmentation regions. From this point of view, it is natural to expect that a particle scattered at some non zero angle will also loose a fraction of its transfer momentum due to final state interaction. This is a characteristic property of soft hadronic interactions. In perturbative QCD the situation is more complicated [8]. In the following we will assume that final state interactions are mostly soft ones.

The aim of the present work is to describe the suppression of the yield of pions in a framework, based on final state interactions, similar to the one used by the authors in order to describe the suppression of J/ψ [9]. In the latter case, the origin of the suppression is twofold. On the one hand, the $c\bar{c}$ pair interacts with nucleons of the nucleus (normal absorption or nuclear absorption, controlled by σ_{abs}). On the other hand, the $c\bar{c}$ pair (at times close to initial time τ_0) or the J/ψ (at larger times), interacts with the dense medium produced in the collision (anomalous absorption, controlled by $\tilde{\sigma}$). In both cases, as a result of the interaction, a $D\bar{D}$ pair is produced instead of a J/ψ . It turns out that in hadron-nucleus collisions, the density of the medium is small enough and such that the effect of the interaction with the medium is negligible. Thus, this effect is only present in nucleus-nucleus collisions – hence its qualification of anomalous.

In the case of large p_T production the particle does not disappear as a result of the interaction but its p_T is shifted to smaller values. Due to the steepness of the p_T distribution, the effect may be quite large. Moreover, in this case there is also a gain of the yield at a given p_T due to particles produced at larger p_T – which have experienced a p_T shift due to the interaction with the medium. This gain

is significantly smaller than the corresponding loss due to the steepness of the p_T distribution.

Another difference with respect to the J/ψ case is that here the suppression vanishes at low p_T . Indeed, when the p_T of the produced particle is close to $\langle p_T \rangle$, its p_T can either increase or decrease as a result of the interaction, i.e. in average the p_T shift tends to zero. Therefore, the above mechanism will not change the results obtained [3] in DPM for soft collisions.

It turns out that in our formalism most of the observed suppression takes place at very early times, where the density of the medium is higher. Since hadron formation times are longer, most of the suppression takes place at a pre-hadronic (partonic) level. Therefore, the mechanism described below is not a conventional hadronic final state interaction and, qualitatively, is expected to lead to similar results as the jet quenching – based on radiative parton energy loss. Further discussion on this point can be found in the conclusions.

2 The Model

The interaction of a large p_T particle with the soft medium is described by the gain and loss differential equations which govern final state interactions. In the following, the large p_T particle will be a π^0 and the medium will be all charged and neutral secondaries produced in $AuAu$ collisions at $\sqrt{s_{NN}} = 200$ GeV. Denoting by ρ_H and ρ_S the corresponding space-time densities, we have following [10]

$$\frac{d\rho_H(x, p_T)}{d^4x} = -\tilde{\sigma} \rho_S [\rho_H(x, p_T) - \rho_H(x, p_T + \delta p_T)] \quad (1)$$

where $\tilde{\sigma}$ is the final state interaction cross-section, averaged over the momentum distribution of the colliding particles. The first term describes the loss of π^0 's with a given p_T , due to its interaction with the medium with density ρ_S . The second term describes the (smaller) gain in the yield at a given p_T resulting from the π^0 's produced at $p_T + \delta p_T$, which have suffered a shift in p_T due to the interaction. In the conventional treatment [10] of eq. (1), one uses cylindrical space-time variables with the longitudinal proper time $\tau = \sqrt{t^2 - z^2}$, space-time rapidity $y = 1/2 \ln((t+z)/(t-z))$ – to be identified later on with the usual rapidity – and transverse coordinate s .

One also assumes longitudinal boost invariance. Therefore, the above picture is not valid in the fragmentation regions. One further assumes that the dilution in time of the densities is only due to longitudinal motion² which leads to a τ^{-1} dependence on τ .

Eq. (1) can then be written in the form [10] [11]

$$\tau \frac{N_{\pi^0}(b, s, y, p_T)}{d\tau} = -\tilde{\sigma} N(b, s, y) [N_{\pi^0}(b, s, y, p_T) - N_{\pi^0}(b, s, y, p_T + \delta p_T)] \quad (2)$$

where $N(b, s, y) \equiv dN/dy \, d^2s(y, b)$ is the density of all charged plus neutral particles per unit rapidity and per unit of transverse area at fixed impact parameter, integrated over p_T . $N_{\pi^0}(b, s, y, p_T)$ is the same quantity for π^0 's at fixed p_T .

Equation (2) has to be integrated from initial time τ_0 to freeze-out time τ_f . It is invariant under the change $\tau \rightarrow c\tau$ and, thus, the result depends only on the ratio τ_f/τ_0 . We use the inverse proportionality between proper time and densities and put $\tau_f/\tau_0 = N(b, s, y)/N_{pp}(y)$ where $N_{pp}(y) = (1/\pi R_p^2) \, dN_{pp}/dy$ is the density of charged and neutral particles per unit rapidity for minimum bias pp collisions at $\sqrt{s} = 200$ GeV. At $y^* \sim 0$, $N_{pp}(0) = 2.24 \, \text{fm}^{-2}$. This density is about 90 % larger than at SPS energies. Since the corresponding increase in the AA density is comparable, the average duration time of the interaction will be approximately the same at CERN-SPS and RHIC, about 5 to 7 fm.

Note that $N(b, s, y)$ in eq. (2) is the density at time τ_0 , i.e. the density produced in the primary collisions. It can be computed in DPM. The procedure is explained in detail in [9]. The hard density N_{π^0} in the primary collision is assumed to scale with the number of binary collisions³.

Eq. (2) can be easily integrated over τ . We obtain in this way the suppression factor $S_{\pi^0}(b, y, p_T)$ of the yield of π^0 's at given p_T and at each impact parameter, due to its interaction with the dense medium. We get

$$S_{\pi^0}(b, y, p_T) = \frac{\int d^2s \, \sigma_{AB}(b) \, n(b, s) \, \tilde{S}_{\pi^0}(b, s, y, p_T)}{\int d^2s \, \sigma_{AB}(b) \, n(b, s)}, \quad (3)$$

²Transverse expansion is neglected. The fact that HBT radii are similar at SPS and RHIC and of the order of magnitude of the nuclear radii, seems to indicate that this expansion is not large. The effect of a small transverse expansion can presumably be taken into account by a small change of the final state interaction cross-section.

³Actually, due to the p_T broadening, this scaling is not satisfied and its violation depends on the value of b (see column I of Table 1). However, if we incorporate this violation by changing $n(b, s)$ into $n(b, s)f(b)$ in eq. (3), the value of S_{π^0} will not change.

where the survival probability is given by

$$\tilde{S}_{\pi^0}(b, s, y, p_T) = \exp \left\{ -\tilde{\sigma} \left[1 - \frac{N_{\pi^0}(b, s, y, p_T + \delta p_T)}{N_{\pi^0}(b, s, y, p_T)} \right] N(b, s, y) \ell n \left(\frac{N(b, s, y)}{N_{pp}(y)} \right) \right\} . \quad (4)$$

Here $\sigma_{AB}(b) = \{1 - \exp[-\sigma_{pp} AB T_{AB}(b)]\}$, where $T_{AB}(b) = \int d^2s T_A(s) T_B(b-s)$, and $T_A(b)$ are profile functions obtained from the Woods-Saxon nuclear densities [12]. Upon integration over b we obtain the AB cross-section. $n(b, s)$ is given by

$$n(b, s) = AB \sigma_{pp} T_A(s) T_B(b-s) / \sigma_{AB}(b) . \quad (5)$$

Upon integration over s we obtain the average number of binary collisions at fixed b , $n(b)$. Note that if we neglect the second term in eqs. (1) and (2), the factor inside brackets in eq. (4) reduces to unity and we recover exactly the formula [9] [13] for the survival probability of the J/ψ .

3 Numerical Results

In order to perform numerical calculations, we need the value of $\tilde{\sigma}$ (which will be treated as a free parameter) as well as the p_T distribution of the π^0 's. Let us concentrate first on π^0 production at mid-rapidities ($|y| < 0.35$).

In pp collisions at $\sqrt{s} = 200$ GeV, the shape of the p_T distribution of π^0 can be described as $(1 + p_T/p_0)^{-n}$ with $p_0 = 1.219$ GeV/c and $n = 9.99$ [14]. The corresponding average p_T is $\langle p_T \rangle = 2p_0/(n-3) = 0.349$ GeV/c. The corresponding value in central ($n_{part} = 350$) $AuAu$ collisions at the same energy is $\langle p_T \rangle = 0.453$ GeV/c. This value is obtained from ref. [15] as an average of π^+ and π^- .

This is the well known p_T broadening, which can be described as a result of initial state interaction (see refs. [9b] [16] for the case of J/ψ). Since it is not our purpose here to describe the p_T broadening we take it from experiment. Of course, the data also contain the effect of the final state interaction. However the effect of the latter is only important at medium and large p_T and it hardly changes the value of $\langle p_T \rangle$ (from the calculation in ref. [9b] this change is of the order of 1 %). As discussed above, two mechanisms are responsible for the p_T broadening. First, the decrease of the shadowing with increasing p_T . This produces an increase of the ratio

$$R_{AA}(b, y, p_T) = \frac{dN^{AA}}{dyd^2p_T}(b)/n(b)\frac{dN^{pp}}{dyd^2p_T} \quad (6)$$

from its small p_T value (substantially lower than unity [3]) to one. This contribution to the p_T broadening can be computed [4] at each value of s , y , p_T and b without adjustable parameters in terms of the (experimentally known) diffractive cross-section. The shadowing computed in this way describes the EMC effect [4]. Moreover, incorporated in DPM, it gives a good description of the centrality dependence of the charged particle inclusive spectra in nucleus-nucleus collisions both at SPS and RHIC energies [3]. The second mechanism is the Cronin effect proper, which produces an increase of R_{AA} above unity. It turns out that at low p_T , where R_{AA} is below unity, most of the increase of R_{AA} with p_T is due to the first mechanism. However, the second is not negligible and it is difficult to compute. In view of that we proceed as follows. We assume that the p_T distribution of π^0 's in $AuAu$ at each b is obtained from the corresponding one in pp by keeping the same value of $n = 9.99$ and changing the scale p_0 into $p_0(b) = \langle p_T \rangle_b (n - 3)/2$. In this way, the average p_T of the new distribution coincides with the measured value $\langle p_T \rangle_b$ at the corresponding centrality [15]. We can thus compute the ratio R_{AA} , eq. (1), in the absence of final state interaction. The reasons for assuming that n is not changed are twofold. On the theoretical side, this is needed in order to reproduce the p_T broadening due to the variation of shadowing with p_T . Indeed, at large p_T the shadowing vanishes and the ratio R_{AA} is independent of p_T – which requires that n is constant⁴. On the experimental side, we notice that the value of n obtained in dAu is the same as in pp within errors. Note also that a substantial change in n would lead to a strong variation of R_{AuAu} and R_{dAu} with p_T at large p_T – which does not seem to be the case experimentally.

In order to fix the absolute normalization we use the value of R_{AA} obtained from DPM in the case of soft collisions (i.e. integrated over p_T) which describes well the experimental results at all centralities [3]. In this way we obtain for the 10 % most central collisions ($n_{part} = 325$) the result shown in Table 1 (column I).

To these values we apply the correction due to the suppression factor S_{π^0} in eq. (3). First, we neglect the second term in eqs. (1) and (2). The formula is then

⁴In perturbative QCD, R_{AA} should tend to unity at large p_T . However, this may occur at much larger values of p_T than present ones.

identical to the one for the J/ψ case, as discussed above, and the suppression is independent of p_T . In order to normalize our result to the experimental values of R_{AA} at large p_T , we use $\tilde{\sigma} = 1.03$ mb, which gives a suppression factor $S_{\pi^0} = 0.143$. The corresponding results are given by column II in Table 1. We see that R_{AA} increases slightly with p_T and agrees with experiment for $p_T > 5$ GeV/c. At lower p_T the result is significantly lower than the data. This is to be expected. Indeed, as discussed above, the suppression factor S_{π^0} has to vanish at small p_T – which is not the case so far. Before introducing this requirement, let us introduce the second term in eqs. (1) and (2) and let us assume that the p_T shift of the π^0 , due to its interaction with the medium, is constant. Consider two cases : $\delta p_T = 0.5$ GeV/c and $\delta p_T = 1.5$ GeV/c. Imposing in all cases the same normalization at $p_T = 7$ GeV/c ⁵ we obtain the results in columns III and IV of Table 1, respectively and in Fig. 1. We observe a slight increase of R_{AA} at large p_T , consistent with the data for $p_T > 7$ GeV/c. The results tend to those in column I with increasing value of δp_T – as it should be. The important result here is that, with constant δp_T , one obtains a small increase of R_{AA} with p_T consistent with the data at large p_T ($p_T \gtrsim 5$ GeV/c). The result is rather insensitive to the value of the shift, for any $\delta p_T \gtrsim 0.5$ GeV/c. Of course, the problem at small p_T remains. In order to cure it, we assume that $\delta p_T \propto (p_T - \langle p_T \rangle_b)$. In this case the factor S_{π^0} is 1 at $p_T = \langle p_T \rangle$ as it should be. Taking $\delta p_T = (p_T - \langle p_T \rangle_b)/20$ we obtain the values in column V of Table 1 and in Fig. 1. Note that these values are rather insensitive to the fraction (5 %) of p_T lost in each interaction. Varying this fraction between 1 % and 10 % the results do not change substantially. We see from Table 1 and Fig. 1 that the slight increase of R_{AA} obtained with constant δp_T is changed into a slight decrease, also consistent with the data, and, moreover, the agreement at small p_T is significantly improved. Actually, a better agreement in the low p_T region is obtained assuming

$$\delta p_T = (p_T - \langle p_T \rangle_b)^{1.5} / 20 . \quad (7)$$

The results are given in column VI of Table 1 and in Fig. 2.

As discussed above, we assume $\delta p_T \propto (p_T - \langle p_T \rangle)^\alpha$ in order to implement

⁵This is done by changing the only free parameter available ($\tilde{\sigma}$ in eq. (4)) in such a way that $\tilde{\sigma}[1 - N_{\pi^0}(b, s, y, p_T + \delta p_T)/N_{\pi^0}(b, s, y, p_T)] = 1.03$. This is done at $p_T = 7$ GeV/c, $\eta^* = 0$ for the 10 % most central collisions. Of course, the same value of $\tilde{\sigma}$ is then used for all values of p_T , η and b .

the condition $S_{\pi^0} \rightarrow 1$ as $p_T \rightarrow \langle p_T \rangle$. The power α controls the way in which this limit is reached. Although such a behaviour is needed at low p_T , it does not have to be the same at large p_T . (In fact a power larger than unity cannot be used at large p_T since δp_T would be larger than p_T). Actually one obtains an equally good agreement with data using $\delta p_T = (p_T - \langle p_T \rangle)^{1.5}/20$ for $p_T < 7$ GeV/c and $\delta p_T = (7 - \langle p_T \rangle)^{1.5}/20$ for $p > 7$ GeV/c. In this case the result is given in column VII of Table 1 and in Fig. 2.

The important result is that at large p_T ($p_T \gtrsim 5$ GeV/c) we obtain a slight increase of R_{AA} with p_T for constant δp_T and a slight decrease for $\delta p_T \propto p_T$ (or $p_T^{1.5}$). In both cases there is agreement with PHENIX data [1a].

The centrality dependence of $R_{AA}(p_T)$ at large p_T ($p_T > 4$ GeV/c) is reasonably well described (see Figs. 2 and 3). The constancy of $R_{AA}(p_T)$ ($p_T > 4$ GeV/c) for $N_{part} < 60$ is the result of a cancellation, in this range of N_{part} , between the increase of the p_T broadening and the increase of the suppression with increasing centrality. This centrality dependence has been reproduced in a recent work, also based on absorption in a dense medium [17].

We turn next to minimum-bias dAu collisions. Here $\langle p_T \rangle = 0.39$ GeV/c [18]. With the same value of n as above ($n = 9.99$), this corresponds to $p_0 = 1.346$. Calculating the ratio dAu to pp and fixing the normalization from the DPM value (integrated over p_T) of this ratio, we obtain the result in Table 2. These values, obtained without introducing nuclear absorption, are in agreement with experiment in the lower half of the p_T range. At large p_T , nuclear absorption is expected to be present both in dAu and $AuAu$ collisions. The dAu data at large p_T are consistent with the presence of nuclear absorption. However, the error bars are too large in order to perform a quantitative study of this question – and determine the value of σ_{abs} ⁶.

Finally, we turn to the dAu collisions at forward rapidities. Consider first the ratio R_{dAu} integrated over p_T . In our approach, R_{dAu} decreases as y^* increases. There are two effects contributing to this decrease.

The first effect is basically due to energy-momentum conservation. It has been known for a long time in hadron-nucleus collisions at SPS energies (as well as at lower ones) and it is well understood in string models such as DPM and QGSM. Recently,

⁶Introducing nuclear absorption in $AuAu$ collisions would result in a smaller value of $\tilde{\sigma}$.

it has been referred to as the low p_T “triangle” [19]. Its extreme form occurs in the hadron fragmentation region, where the yield of secondaries in collisions off a heavy nucleus is smaller than the corresponding yield in hadron-proton. This phenomenon is known as nuclear attenuation. It turns out, that, at RHIC energies, this effect produces a decrease of R_{dAu} of about 30 % between $\eta^* = 0$ and $\eta^* = 3.2$.

The second effect is the increase of the shadowing corrections in dAu with increasing y^* [4]. This produces a decrease of $R_{dAu}(p_T)$ between $\eta^* = 0$ and $\eta^* = 3.2$ of about 30 % for pions produced in minimum bias collisions⁷. This decrease is practically independent of p_T . Therefore, we expect a suppression factor of about 1.7 between $R_{dAu}(p_T)$ at $\eta^* = 0$ and at $\eta^* = 3.2$, practically independent of p_T . This is consistent with the BRAHMS results [2b].

The same result is obtained with the procedure used above in $AuAu$ collisions, if the $\langle p_T \rangle$ of pions in dAu is the same at $\eta^* = 0$ and at $\eta^* = 3.2$. This is approximately the case in $AuAu$ collisions [1b] and, in our approach, it is expected also in dAu . Indeed, as discussed above, most of the p_T broadening at low p_T is due to the variation of shadowing with p_T – and this variation is practically the same at $\eta^* = 0$ and at $\eta^* = 3.2$. On the contrary, the dependence of shadowing on centrality at fixed η is quite important and, thus, the centrality dependence of $R_{dAu}(p_T)$ is quite large. Details will be presented elsewhere⁸.

4 Conclusions

In this work the suppression of π^0 production at large p_T in $AuAu$ collisions is described in terms of final state interaction in the dense medium produced in the collision. The mechanism is similar to the one responsible for J/ψ suppression. A nice feature of our formulation is that it leads to a suppression of $R_{AA}(p_T)$ at large p_T ($p_T > 5$ GeV/c) which is rather insensitive to the size and form of the p_T shift produced by the final state interaction.

Our approach contains dynamical, non-linear, shadowing. This shadowing is

⁷The situation is different in $AuAu$. Here the shadowing decreases with increasing η , but its variation is much smaller than in dAu [4].

⁸After completion of this work, we learned of a related work, ref. [20]. In that paper only Cronin effect plus geometrical shadowing was considered in dAu collisions. So the authors are unable to describe data at large η (contrary to our approach, where dynamical shadowing is taken into account).

determined in terms of (experimentally known) diffractive cross-sections. As $s \rightarrow \infty$, it leads to saturation of the parton distributions. However, at both RHIC and LHC energies, this shadowing is significantly smaller than the one present in a saturation regime. By itself, the shadowing in our approach is not sufficient to explain the difference in R_{dAu} between $\eta^* = 0$ and $\eta^* = 3.2$ measured by BRAHMS. Indeed, it produces only about one half of the measured variation. However, when the effect of shadowing (i.e. the increase of shadowing with increasing rapidity) is combined with low p_T effects present in string models such as DPM and QGSM, agreement with the BRAHMS measurement is achieved. This also shows that decreasing x by increasing energy is not equivalent to doing so by going to forward rapidities. In the latter case the value of R_{dAu} at low p_T is substantially reduced – which obviously influences its large p_T value. This is not so in the former case.

In this paper we have restricted ourselves to the study of the π^0 large p_T suppression. Other observables have been measured such as the back-to-back and near side large p_T azimuthal correlations and large p_T azimuthal anisotropy. As discussed in [21], the present data on these observables indicate that the suppression takes place at very early times, where the density of the medium is larger. However, this does not imply that the suppression is due to non-abelian radiative parton energy loss. Indeed, as discussed in Section 1, in our mechanism the suppression also takes place at very early times, at a partonic level. In view of that, we expect that the results for these observables will be similar to the ones obtained in the radiative jet quenching scenario. In our approach the back-to-back correlation will be suppressed by the same amount as the single particle one, whereas the near-side one will remain essentially the same – since the two large p_T particles originate from the same jet. Likewise, a high p_T azimuthal anisotropy is expected in our approach due to the asymmetry of the dense medium at early times.

The large p_T suppression phenomenon is important in order to determine the properties of the dense medium in which it takes place. In our approach, using the inverse proportionality between density and interaction time, we find that in a central $Au Au$ collisions the density of the medium is about 5 times larger than in pp . This factor is predicted to be practically unchanged at LHC energies. In a radiative jet quenching scenario much larger densities have been claimed in the literature. However, in a recent paper [22], the density of the medium has been

found to increase by a factor 4 between peripheral and central $Au\ Au$ collisions at RHIC – consistent with our result.

In order to distinguish between radiative jet quenching scenario and collisional energy loss or collisional p_T -shift it will be very interesting to measure the detailed medium modifications of jet shapes and multiplicities [23]. While we do not see at present how to disentangle the different scenarios on the basis of qualitative arguments, there will hopefully be differences in their predictions at a qualitative level.

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References

- [1] a) PHENIX Collaboration, S. S. Adler et al., Phys. Rev. Lett **91**, 072301 (2003) and nucl-ex/0308006. See also D. d’Enterria in Proc. Quark Matter 2004, to be published in Nucl. Phys. A.
b) BRAHMS Collaboration, I. Arsene et al., Phys. Rev. Lett. **91**, 072305 (2003) and M. Murray in Proc. Quark Matter 2004, to be published in Nucl. Phys. A.
c) STAR Collaboration, J. Adams et al, Phys. Rev. Lett. **91**, 172302 (2003) and K. Schweda in Proc. Quark Matter 2004, to be published in Nucl. Phys. A.
d) PHOBOS Collaboration, B. B. Back et al., Phys. Rev. Lett. **88**, 022301 (2002); Phys. Lett. **B578**, 297 (2004) and P. Steinberg in Proc. Quark Matter 2004, to be published in Nucl. Phys. A.
- [2] a) PHENIX Collaboration, S. S. Adler et al., Phys. Rev. Lett. **91**, 072303 (2003).
b) BRAHMS Collaboration, see ref. [1b] and R. Debbe in Proc. Quark Matter 2004, to be published in Nucl. Phys. A.
c) STAR Collaboration, J. Adams et al, Phys. Rev. Lett. **91**, 072304 (2003).
d) PHOBOS Collaboration, B. B. Back et al., Phys. Rev. Lett. **91**, 072302 (2003).
- [3] A. Capella and D. Sousa, Phys. Lett. **B511**, 185 (2001).
A. Capella, A. Kaidalov, J. Tran Thanh Van, Heavy Ion Physics **9**, 169 (1999).
- [4] N. Armesto, A. Capella, A. B. Kaidalov, J. Lopez-Albacete and C. A. Salgado, Eur. Phys. J. **C29**, 531 (2003).
- [5] L. McLerran and Raju Venugopalan, Phys. Rev. **D49**, 2233 (1994); *ibid* 3352 (1994); *ibid* **D50**, 2225 (1994).
E. Iancu, A. Leonidov and L. McLerran, Nucl. Phys. **A692**, 583 (2001).
E. G. Ferreira, E. Iancu, A. Leonidov and L. McLerran, Nucl. Phys. **703**, 489 (2002).
For recent reviews see L. McLerran, hep-ph/0402137 and 0311028 and references therein.

- [6] K. G. Boreskov, A. B. Kaidalov, S. M. Kiselev and N. Ya. Smorodinskaya, Sov. J. Nucl. Phys. **53**, 356 (1991).
- [7] K. G. Boreskov, A. Capella, A. B. Kaidalov and J. Tran Thanh Van, Phys. Rev. **47**, 919 (1993).
N. Armesto, A. Braun, A. Capella, C. Pajares and C. A. Salgado, Nucl. Phys. **B509**, 357 (1998).
- [8] For a review see R. Baier, D. Schiff and B. G. Zakharov, Ann. Rev. Nucl. Part. Sci. **50**, 37 (2000). For more recent work see M. Gyulassy, I. Vitev, X-N Wang and B-W Zhang, nucl-th/0302077, published in Quark Gluon Plasma 3, editors: R.C. Hwa and X.N. Wang, World Scientific, Singapore; G. G. Barnaföldi, P. Lévar, G. Papp, G. Fäi and M. Gyulassy, Eur. Phys. J **C33**, S609 (2004); C. A. Salgado and U. Wiedemann, Phys. Rev. Lett. **93**, 042301 (2004); N. Armesto, C. A. Salgado and U. Wiedemann, Phys. Rev. **D69**, 114003 (2004); F. Arleo, JHEP 0211, 044 (2002).
- [9] a) N. Armesto and A. Capella, Phys. Lett. **B430**, 23 (1998).
b) N. Armesto, A. Capella and E. G. Ferreira, Phys. Rev. **C59**, 395 (1999).
c) A. Capella, E. G. Ferreira and A. Kaidalov, Phys. Rev. Lett. **85**, 2080 (2000).
d) A. Capella, A. Kaidalov and D. Sousa, Phys. Rev. **C65**, 054908 (2002).
e) A. Capella and D. Sousa, Eur. Phys. J **C30**, 117 (2003).
- [10] B. Koch, U. Heinz and J. Pitsut, Phys. Lett. **B243**, 149 (1990).
- [11] A. Capella, Phys. Lett. **B364**, 175 (1995).
- [12] C. W. Jager et al., Atomic Data and Nuclear Tables **14**, 485 (1974).
- [13] D. Kharzeev, C. Lourenço, M. Nardi and H. Satz, Z. Phys. **C74**, 307 (1997).
- [14] PHENIX Collaboration, S. S. Adler et al., Phys. Rev. Lett. **91**, 241803 (2003).
- [15] PHENIX Collaboration, S. S. Adler et al., Phys. Rev. **C69**, 034909 (2004).
- [16] D. Kharzeev, M. Nardi and H. Satz, Phys. Lett. **B403**, 362 (1997).
- [17] A. Drees, H. Feng and J. Jia, nucl-th/0310044.
W. Cassing, K. Gallmeister and C. Greiner, Nucl. Phys. **A735**, 277 (2004).

- [18] PHENIX Collaboration, D. d'Enterria, private communication.
- [19] M. Gyulassy in Proc. Quark Matter 2004, to be published in Nucl. Phys. A.
- [20] A. Accardi and M. Gyulassy, J. Phys. **G30**, S969 (2004).
- [21] X. N. Wang, Phys. Lett.**B579**, 299 (2004).
- [22] A. K. Chaudhuri, nucl-th/0409020.
- [23] C. A. Salgado and U. A. Wiedemann, Phys. Rev. Lett. **93**, 042301 (2004); I. Vitev, nucl-th/0404052.

CAPTIONS

Table 1. Values of $R_{AuAu}^{\pi^0}(p_T)$ for the 10 % most central collisions $AuAu$ collisions at mid-rapidities ($|y^*| < 0.35$). Column I is the result obtained with no final state interaction. The results in the other columns include final state interaction with several ansatzs for the p_T shift induced by this interaction (see main text).

Table 2. Values of $R_{dAu}^{\pi^0}(p_T)$ for minimum bias dAu collisions at mid rapidities ($|y^*| < 0.35$).

Fig 1. Values of $R_{AuAu}^{\pi^0}(p_T)$ for the 10 % most central collisions at mid-rapidities ($|y^*| < 0.35$), using the p_T shift given by eq. (7) $\delta p_T = (p_{T-} < p_{T>b})^{1.5}/20$ (solide line), the linear case $\delta p_T = (p_{T-} < p_{T>b})/20$ (dashed-dotted line), the quadratic case $\delta p_T = (p_{T-} < p_{T>b})^2/20$ (dashed line), and $\delta p_T = \text{constant}$ (dotted lines). See Table 1 and main text for details.

Fig 2. Values of $R_{AuAu}^{\pi^0}(p_T)$ for the 10 % most central collisions (lower line) and for peripheral (80-92 %) collisions (upper line) at mid-rapidities ($|y^*| < 0.35$), using the p_T shift given by eq. (7), $\delta p_T = (p_{T-} < p_{T>b})^{1.5}/20$, (solide line). The dashed line is obtained using eq. (7) for $p_T \leq 7$ GeV/c and $p_T = \text{constant}$ for $p_T \geq 7$ GeV/c (see main text). The data are from ref. [1a].

Fig 3. Centrality dependence of $R_{AuAu}^{\pi^0}$ for $p_T > 4$ GeV/c using the p_T shift given by eq. (7). The data are from ref. [1a].

TABLE 1

p_T (GeV/c)	I	II	III	IV	V	VI	VII
0.5	0.38	0.05			0.34	0.38	0.38
2	0.90	0.13	0.08	0.11	0.20	0.31	0.31
5	1.48	0.21	0.18	0.19	0.23	0.25	0.25
7	1.69	0.24	0.24	0.24	0.24	0.24	0.24
10	1.84	0.27	0.34	0.30	0.25	0.24	0.32

TABLE 2

p_T (GeV/c)	$R_{dAu}(p_T)$
0.39	0.63
1	0.78
2	0.92
3	1.01
5	1.10
7	1.16
10	1.21

Figure 1

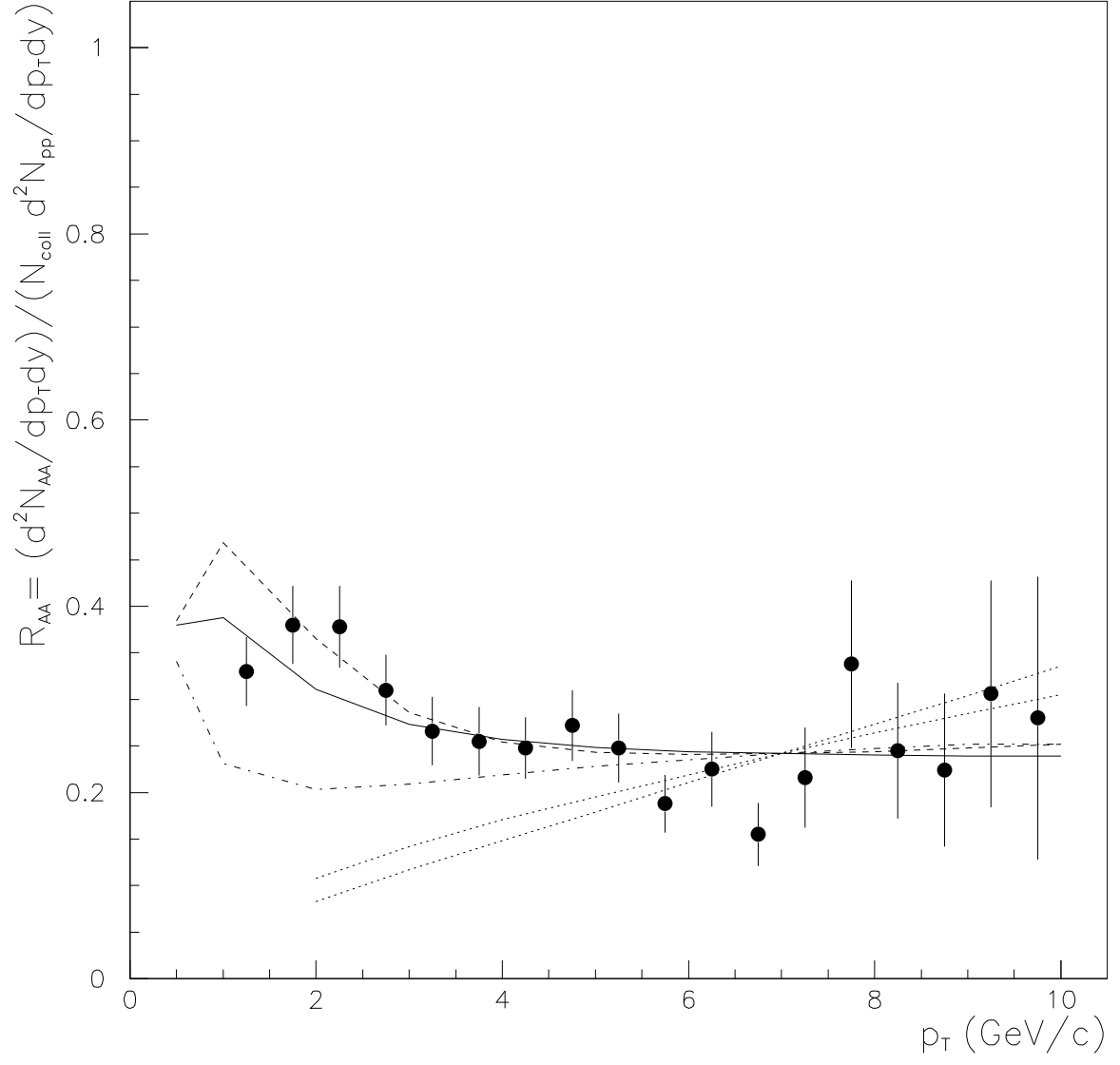


Figure 2

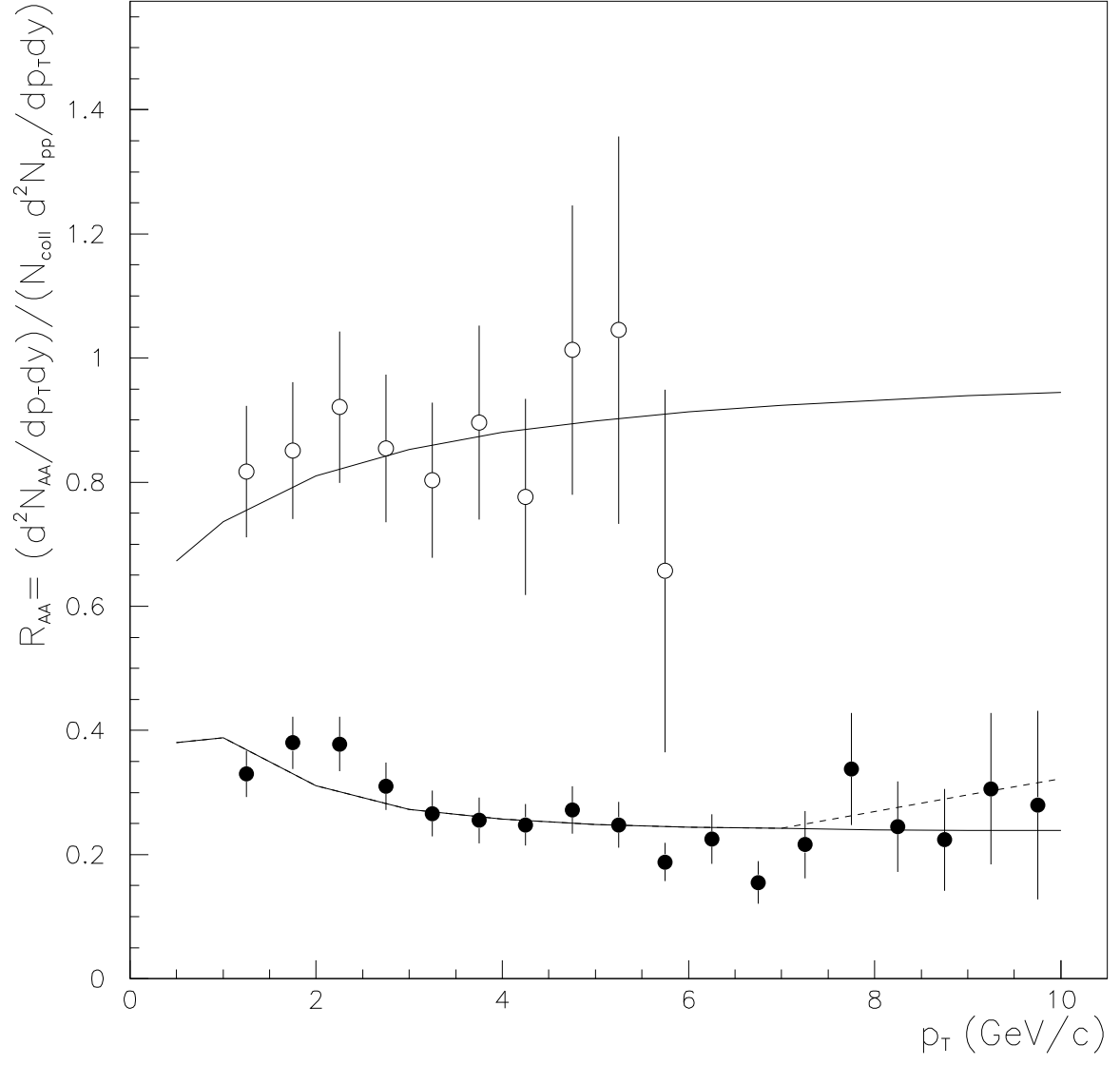


Figure 3

