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

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Abstract. We propose a model of the 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.

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 non-linear effects lead, as $s \to \infty$, 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, $x \ll x_{\mathrm{cr}} = 1/m_N R_A$ 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 the Feynman variable x_{F} have $x = x_{\pm} = 1/2(\sqrt{x_{\mathrm{F}}^2 + 4m_{\mathrm{T}}^2/s} \pm x_{\mathrm{F}})$. Thus, at fixed initial energy (s) the condition for the 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 [6] that in this region, where $x \gtrsim x_{\mathrm{cr}}$, 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 lose 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 lose a fraction of its transfer momentum due to final state interaction. This is a characteristic property of soft hadronic

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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 $\sigma_{\rm abs}$). On the other hand, the c–$\bar{c}$ pair (at times close to the initial time τ_0) or the J/ψ (at larger times), interacts with the dense medium produced in the collision (anomalous absorption, controlled by $\widetilde{\sigma}$). 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 $p_{\rm T}$ production the particle does not disappear as a result of the interaction but its $p_{\rm T}$ is shifted to smaller values. Due to the steepness of the $p_{\rm T}$ distribution, the effect may be quite large. Moreover, in this case there is also a gain of the yield at a given $p_{\rm T}$ due to particles produced at larger $p_{\rm T}$ – which have experienced a $p_{\rm T}$ shift due to the interaction with the medium. This gain is significantly smaller than the corresponding loss due to the steepness of the $p_{\rm T}$ distribution.

Another difference with respect to the J/ψ case is that here the suppression vanishes at low $p_{\rm T}$. Indeed, when the $p_{\rm T}$ of the produced particle is close to $\langle p_{\rm T}\rangle$, its $p_{\rm T}$ can either increase or decrease as a result of the interaction, i.e. in average the $p_{\rm T}$ 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 $p_{\rm T}$ particle with the soft medium is described by the gain and loss differential equations which govern final state interactions. In the following, the large $p_{\rm T}$ particle will be a π^0 and the medium will be all charged and neutral secondaries produced in AuAu collisions at $\sqrt{s_{NN}} = 200\,\text{GeV}$. Denoting by ρ_H and ρ_S the corresponding space-time densities, we have the following [10]:

$$\frac{\mathrm{d}\rho_H(x,p_{\rm T})}{\mathrm{d}^4x} = -\widetilde{\sigma}\rho_S\left[\rho_H(x,p_{\rm T}) - \rho_H(x,p_{\rm T}+\delta p_{\rm T})\right] \quad (1)$$

where $\widetilde{\sigma}$ is the final state interaction cross-section, averaged over the momentum distribution of the colliding particles. The first term describes the loss of the π^0's with a given $p_{\rm T}$, due to its interaction with the medium with density ρ_S. The second term describes the (smaller) gain in the yield at a given $p_{\rm T}$ resulting from the π^0's produced at $p_{\rm T}+\delta p_{\rm T}$, which have suffered a shift in $p_{\rm T}$ due to the interaction. In the conventional treatment [10] of (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 motion[1] which leads to a τ^{-1} dependence on τ.

Equation (1) can then be written in the form [10, 11]

$$\tau\frac{N_{\pi^0}(b,s,y,p_{\rm T})}{\mathrm{d}\tau} = -\widetilde{\sigma}N(b,s,y) \times\left[N_{\pi^0}(b,s,y,p_{\rm T}) - N_{\pi^0}(b,s,y,p_{\rm T}+\delta p_{\rm T})\right], \quad (2)$$

where $N(b,s,y) \equiv \mathrm{d}N/\mathrm{d}y\ \mathrm{d}^2s(y,b)$ is the density of all charged plus neutral particles per unit rapidity and per unit of transverse area at fixed impact parameter, integrated over $p_{\rm T}$. $N_{\pi^0}(b,s,y,p_{\rm T})$ is the same quantity for π^0's at fixed $p_{\rm T}$.

Equation (2) has to be integrated from the initial time τ_0 to freeze-out time $\tau_{\rm f}$. It is invariant under the change $\tau \to c\tau$ and, thus, the result depends only on the ratio $\tau_{\rm f}/\tau_0$. We use the inverse proportionality between proper time and densities and put $\tau_{\rm f}/\tau_0 = N(b,s,y)/N_{pp}(y)$ where $N_{pp}(y) = (1/\pi R_p^2)\ \mathrm{d}N_{pp}/\mathrm{d}y$ is the density of charged and neutral particles per unit rapidity for minimum bias pp collisions at $\sqrt{s} = 200\,\text{GeV}$. 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 (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 collisions[2].

Equation (2) can easily be integrated over τ. We obtain in this way the suppression factor $S_{\pi^0}(b,y,p_{\rm T})$ of the yield of π^0's at given $p_{\rm T}$ and at each impact parameter, due to

[1] Transverse 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.

[2] Actually, due to the $p_{\rm T}$ 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 (3), the value of S_{π^0} will not change.

its interaction with the dense medium. We get

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

where the survival probability is given by

$$\widetilde{S}_{\pi^0}(b,s,y,p_\mathrm{T}) = \exp\left\{ -\widetilde{\sigma}\left[1 - \frac{N_{\pi^0}(b,s,y,p_\mathrm{T}+\delta p_\mathrm{T})}{N_{\pi^0}(b,s,y,p_\mathrm{T})}\right] \times N(b,s,y) \ln\left(\frac{N(b,s,y)}{N_{pp}(y)}\right)\right\}. \quad (4)$$

Here $\sigma_{AB}(b) = \{1 - \exp[-\sigma_{pp} AB T_{AB}(b)]\}$, where $T_{AB}(b) = \int \mathrm{d}^2 s T_A(s) T_B(b-s)$, 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 (1) and (2), the factor inside brackets in (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 $\widetilde{\sigma}$ (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 $\sqrt{s} = 200\,\mathrm{GeV}$, the shape of the p_T distribution of π^0 can be described by $(1 + p_\mathrm{T}/p_0)^{-n}$ with $p_0 = 1.219\,\mathrm{GeV}/c$ and $n = 9.99$ [14]. The corresponding average p_T is $\langle p_\mathrm{T}\rangle = 2p_0/(n-3) = 0.349\,\mathrm{GeV}/c$. The corresponding value in central ($n_\mathrm{part} = 350$) AuAu collisions at the same energy is $\langle p_\mathrm{T}\rangle = 0.453\,\mathrm{GeV}/c$. This value is obtained from [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 (b) of [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 $\langle p_\mathrm{T}\rangle$ (from the calculation in (b) of [16] this change is of the order of 1%). As discussed above, two mechanisms are responsible for the p_T broadening. First, let us take the decrease of the shadowing with increasing p_T. This produces an increase of the ratio

$$R_{AA}(b,y,p_\mathrm{T}) = \frac{\mathrm{d}N^{AA}}{\mathrm{d}y\mathrm{d}^2 p_\mathrm{T}}(b)/n(b)\frac{\mathrm{d}N^{pp}}{\mathrm{d}y\mathrm{d}^2 p_\mathrm{T}} \quad (6)$$

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 one 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 $p_0(b) = \langle p_\mathrm{T}\rangle_b\ (n-3)/2$. In this way, the average p_T of the new distribution coincides with the measured value $\langle p_\mathrm{T}\rangle_b$ at the corresponding centrality [15]. We can thus compute the ratio R_{AA}, (1), in the absence of final state interactions. 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 constant[3]. 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 $R_{d\mathrm{Au}}$ 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 ($n_\mathrm{part} = 325$) the result shown in Table 1 (column I).

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

Table 1. Values of $R^{\pi^0}_{\mathrm{AuAu}}(p_\mathrm{T})$ 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 the main text)

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

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

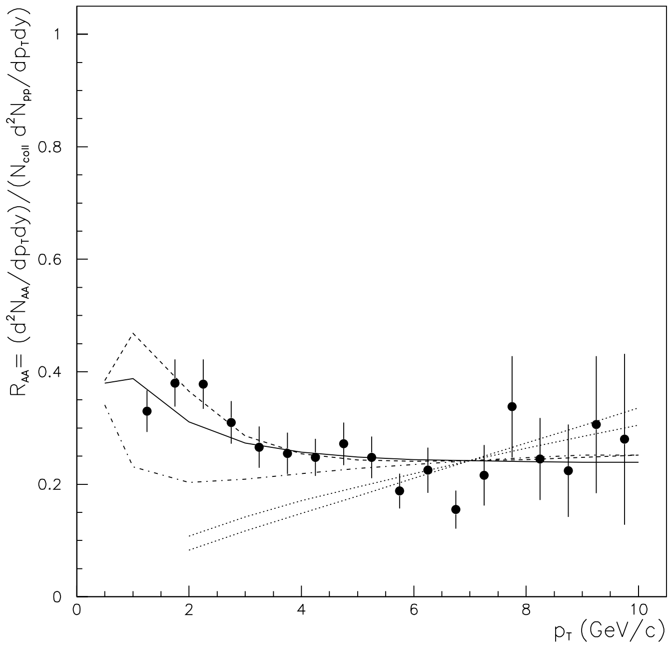

Fig. 1. Values of $R^{\pi^0}_{\mathrm{AuAu}}(p_{\mathrm{T}})$ for the 10% most central collisions at mid rapidities ($|y^*| < 0.35$), using the p_{T} shift given by (7) $\delta p_{\mathrm{T}} = (p_{\mathrm{T}} - \langle p_{\mathrm{T}}\rangle_b)^{1.5}/20$ (solide line), the linear case $\delta p_{\mathrm{T}}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

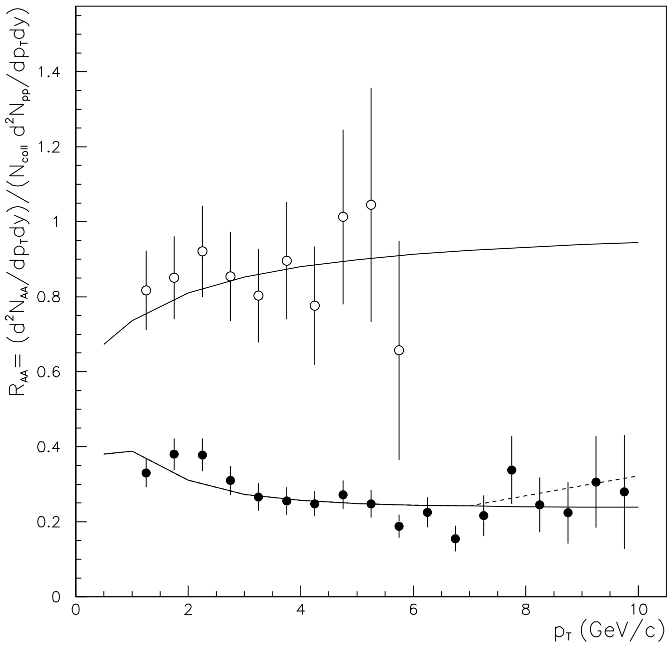

Fig. 2. Values of $R^{\pi^0}_{\mathrm{AuAu}}(p_{\mathrm{T}})$ for the 10% most central collisions (lower line) and for peripheral (80–92%) collisions (upper line) at mid rapidities ($|y^*| < 0.35$), using the p_{T} shift given by (7), $\delta p_{\mathrm{T}} = (p_{\mathrm{T}} - \langle p_{\mathrm{T}}\rangle_b)^{1.5}/20$,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

$\widetilde{\sigma} = 1.03$ 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 $p_{\mathrm{T}} > 5\,\mathrm{GeV}/c$. 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 (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: $\delta p_{\mathrm{T}} = 0.5\,\mathrm{GeV}/c$ and $\delta p_{\mathrm{T}} = 1.5\,\mathrm{GeV}/c$. Imposing in all cases the same normalization at $p_{\mathrm{T}} = 7\,\mathrm{GeV}/c$[4] 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 $p_{\mathrm{T}} > 7\,\mathrm{GeV}/c$. 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} ($p_{\mathrm{T}} \gtrsim 5\,\mathrm{GeV}/c$). The result is rather insensitive to the value of the shift, for any $\delta p_{\mathrm{T}} \gtrsim 0.5\,\mathrm{GeV}/c$. Of course, the problem at small p_{T} remains. In order to cure it, we assume that $\delta p_{\mathrm{T}} \propto (p_{\mathrm{T}} - \langle p_{\mathrm{T}}\rangle_b)$. In this case the factor S_{π^0} is 1 at $p_{\mathrm{T}} = \langle p_{\mathrm{T}}\rangle$ as it should be. Taking $\delta p_{\mathrm{T}} = (p_{\mathrm{T}} - \langle p_{\mathrm{T}}\rangle_b)/20$ 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

$$\delta p_{\mathrm{T}} = (p_{\mathrm{T}} - \langle p_{\mathrm{T}}\rangle_b)^{1.5}/20. \qquad (7)$$

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

As discussed above, we assume $\delta p_{\mathrm{T}} \propto (p_{\mathrm{T}} - \langle p_{\mathrm{T}}\rangle)^{\alpha}$ in order to implement the condition $S_{\pi^0} \to 1$ as $p_{\mathrm{T}} \to \langle p_{\mathrm{T}}\rangle$. 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 the data using $\delta p_{\mathrm{T}} = (p_{\mathrm{T}} - \langle p_{\mathrm{T}}\rangle_b)^{1.5}/20$ for $p_{\mathrm{T}} < 7\,\mathrm{GeV}/c$ and $\delta p_{\mathrm{T}} = (7 - \langle p_{\mathrm{T}}\rangle_b)^{1.5}/20$ for $p > 7\,\mathrm{GeV}/c$. 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} ($p_{\mathrm{T}} \gtrsim 5\,\mathrm{GeV}/c$) we obtain a slight increase of R_{AA} with p_{T} for constant δp_{T} and a slight decrease for $\delta p_{\mathrm{T}} \propto p_{\mathrm{T}}$ (or $p_{\mathrm{T}}^{1.5}$). In both cases there is agreement with PHENIX data; see (a) of [1].

The centrality dependence of $R_{AA}(p_{\mathrm{T}})$ at large p_{T} ($p_{\mathrm{T}} > 4\,\mathrm{GeV}/c$) is reasonably well described (see Figs. 2 and 3). The constancy of $R_{AA}(p_{\mathrm{T}})$ ($p_{\mathrm{T}} > 4\,\mathrm{GeV}/c$) for $N_{\mathrm{part}} <$

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

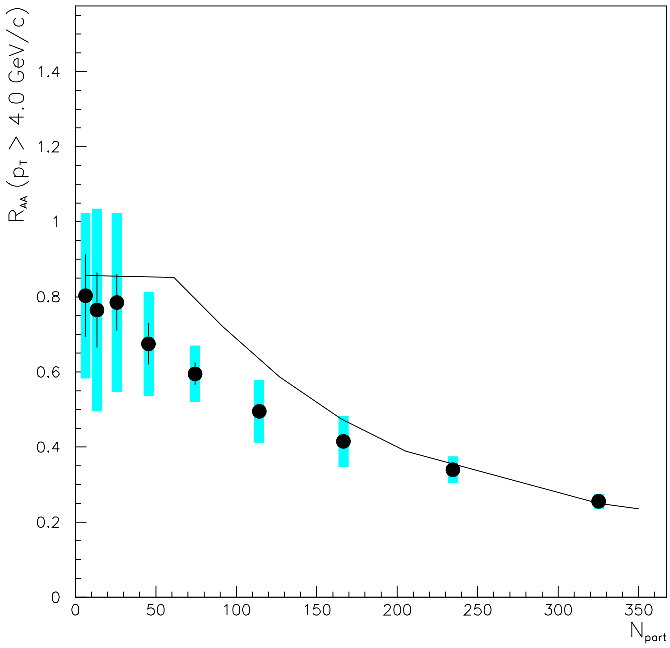

Fig. 3. Centrality dependence of $R^{\pi^0}_{\mathrm{AuAu}}$ for $p_\mathrm{T} > 4\,\mathrm{GeV}/c$ using the p_T shift given by (7). The data are from (a) of [1]

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 $\langle p_\mathrm{T}\rangle = 0.39\,\mathrm{GeV}/c$ [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[5].

Table 2. Values of $R^{\pi^0}_{d\mathrm{Au}}(p_\mathrm{T})$ for minimum bias dAu collisions at mid rapidities ($|y^*| < 0.35$)

p_T (GeV/c)	$R_{d\mathrm{Au}}(p_\mathrm{T})$
0.39	0.63
1	0.78
2	0.92
3	1.01
5	1.10
7	1.16
10	1.21

[5] Introducing nuclear absorption in AuAu collisions would result in a smaller value of $\widetilde{\sigma}$.

Finally, we turn to the dAu collisions at forward rapidities. Consider first the ratio $R_{d\mathrm{Au}}$ integrated over p_T. In our approach, $R_{d\mathrm{Au}}$ 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, 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 hadron–proton yield. This phenomenon is known as nuclear attenuation. It turns out, that, at RHIC energies, this effect produces a decrease of $R_{d\mathrm{Au}}$ 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 $R_{d\mathrm{Au}}(p_\mathrm{T})$ between $\eta^* = 0$ and $\eta^* = 3.2$ of about 30% for pions produced in minimum bias collisions[6]. This decrease is practically independent of p_T. Therefore, we expect a suppression factor of about 1.7 between $R_{d\mathrm{Au}}(p_\mathrm{T})$ at $\eta^* = 0$ and at $\eta^* = 3.2$, practically independent of p_T. This is consistent with the BRAHMS results; see (b) of [2].

The same result is obtained with the procedure used above in AuAu collisions, if the $\langle p_\mathrm{T}\rangle$ of pions in dAu is the same at $\eta^* = 0$ and at $\eta^* = 3.2$. This is approximately the case in AuAu collisions (see (b) of [1]) 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 $R_{d\mathrm{Au}}(p_\mathrm{T})$ is quite large. Details with be presented elsewhere[7].

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 $R_{AA}(p_\mathrm{T})$ at large p_T ($p_\mathrm{T} > 5\,\mathrm{GeV}/c$) 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 determined in terms of (experimentally known) diffractive cross-sections. As $s \to \infty$, it leads to saturation of the parton distributions. However, at both

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

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

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 $R_{d\mathrm{Au}}$ 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 $R_{d\mathrm{Au}}$ 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 Sect. 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 AuAu collision the density of the medium is about five 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 four between peripheral and central AuAu 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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