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

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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 \( J/\psi \) 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 $s \to \infty$, 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 $x$, $x \ll x_{cr} = 1/m_NR_A$ 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 $x = x_{\pm} = 1/2(\sqrt{x^2_F + 4m^2_T/s} \pm x_F)$. 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...
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 $J/\psi$ [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 $\sigma_{abs}$). On the other hand, the $c\bar{c}$ pair (at times close to initial time $\tau_0$) or the $J/\psi$ (at larger times), interacts with the dense medium produced in the collision (anomalous absorption, controlled by $\bar{\sigma}$). In both cases, as a result of the interaction, a $D\bar{D}$ pair is produced instead of a $J/\psi$. 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 $J/\psi$ case is that here the suppression vanishes at low $p_T$. Indeed, when the $p_T$ of the produced particle is close to $<p_T>$, 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 $\pi^0$ and the medium will be all charged and neutral secondaries produced in $AuAu$ collisions at $\sqrt{s_{NN}} = 200$ GeV. Denoting by $\rho_H$ and $\rho_S$ the corresponding space-time densities, we have following [10]

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

(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 $\pi^0$'s with a given $p_T$, due to its interaction with the medium with density $\rho_S$. The second term describes the (smaller) gain in the yield at a given $p_T$ resulting from the $\pi^0$'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 $s$. 4
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\(^2\) which leads to a \(\tau^{-1}\) dependence on \(\tau\).

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

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

(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 \(\pi^0\)'s at fixed \(p_T\).

Equation (2) has to be integrated from initial time \(\tau_0\) to freeze-out time \(\tau_f\). It is invariant under the change \(\tau \to c\tau\) and, thus, the result depends only on the ratio \(\tau_f/\tau_0\). 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 \(y^* \sim 0\), \(N_{pp}(0) = 2.24\) 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 \(\tau_0\), 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 \(N_{\pi^0}\) in the primary collision is assumed to scale with the number of binary collisions\(^3\).

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

\[
S_{\pi^0}(b, y, p_T) = \frac{\int d^2s \sigma_{AB}(b) n(b, s) S_{\pi^0}(b, s, y, p_T) \sigma_{AB}(b) n(b, s)}{\int d^2s \sigma_{AB}(b) n(b, s)},
\]

(3)

\(^2\)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.

\(^3\)Actually, 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 \(S_{\pi^0}\) will not change.
where the survival probability is given by

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

(4)

Here \( \sigma_{AB}(b) = \{1 - \exp[-\sigma_{pp} A B T_{AB}(b)]\} \), where \( T_{AB}(b) = \int d^2sT_A(s)T_B(b - s) \), 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

\[
n(b, s) = AB \sigma_{pp} T_A(s) T_B(b - s)/\sigma_{AB}(b) .
\]  

(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 \( J/\psi \).

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 \( \pi^0 \)'s. Let us concentrate first on \( \pi^0 \) production at mid-rapidities (\( |y| < 0.35 \)).

In \( pp \) collisions at \( \sqrt{s} = 200 \) GeV, the shape of the \( p_T \) distribution of \( \pi^0 \) 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 \( < p_T >= 2p_0/(n-3) = 0.349 \) GeV/c. The corresponding value in central (\( n_{part} = 350 \)) \( AuAu \) collisions at the same energy is \( < p_T >= 0.453 \) GeV/c. This value is obtained from ref. [15] as an average of \( \pi^+ \) and \( \pi^- \).

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 \( J/\psi \)). 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 \( < p_T > \) (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
\[ R_{AA}(b, y, p_T) = \frac{dN^{AA}}{dyd^2p_T}(b)/n(b) \frac{dN^{pp}}{dyd^2p_T} \] (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 \( \pi^0 \)'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 constant\(^4\). 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 \( S_{\pi^0} \) in eq. (3). First, we neglect the second term in eqs. (1) and (2). The formula is then

\(^4\)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 present ones.
identical to the one for the $J/\psi$ 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 $S_{\pi^0}$ 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 $\pi^0$, 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 $\delta p_T$ – as it should be. The important result here is that, with constant $\delta p_T$, 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 $\delta p_T \propto (p_T - < p_T >_b)$. In this case the factor $S_{\pi^0}$ is 1 at $p_T = < p_T >$ as it should be. Taking $\delta p_T = (p_T - < p_T >_b)/20$ 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 $\delta p_T$ 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

$$ \delta p_T = (p_T - < p_T >_b)^{1.5}/20 .$$  \hspace{1cm} (7)

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

As discussed above, we assume $\delta p_T \propto (p_T - < p_T >)^{\alpha}$ in order to implement\(^5\)This 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 $p_T$, $\eta$ and $b$.\(^6\)
the condition $S_{x^0} \to 1$ as $p_T \to < p_T >$. The power $\alpha$ 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 $\delta p_T$ would be larger than $p_T$). Actually one obtains an equally good agreement with data using $\delta p_T = (p_T - < p_T >_b)^{1.5}/20$ for $p_T < 7\text{ GeV/c}$ and $\delta p_T = (7 - < p_T >_b)^{1.5}/20$ for $p > 7\text{ GeV/c}$. In this case the result is given in column VII of Table 1 and in Fig. 2.

The important result is that at large $p_T$ ($p_T \gtrsim 5\text{ GeV/c}$) we obtain a slight increase of $R_{AA}$ with $p_T$ for constant $\delta p_T$ 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 $p_T$ ($p_T > 4\text{ GeV/c}$) is reasonably well described (see Figs. 2 and 3). The constancy of $R_{AA}(p_T)$ ($p_T > 4\text{ GeV/c}$) 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 $< p_T > = 0.39\text{ GeV/c}$ [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 $\sigma_{abs}$\textsuperscript{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,

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

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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 collisions\(^7\). 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 $<p_T>$ 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 $\eta$ is quite important and, thus, the centrality dependence of $R_{dAu}(p_T)$ is quite large. Details with be presented elsewhere\(^8\).

\section{Conclusions}

In this work the suppression of $\pi^0$ 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 $J/\psi$ 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

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

\(^8\)After 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 $\eta$ (contrary to our approach, where dynamical shadowing is taken into account).
determined in terms of (experimentally known) diffractive cross-sections. As \( s \to \infty \),
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 \( \pi^0 \) large \( p_T \) suppres-
sion. 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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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 cuadratic 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

![Graph showing $R_{AA}$ (p_T > 4.0 GeV/c) vs. N_{part}.](image-url)

The graph illustrates the behavior of $R_{AA}$ for different values of $N_{part}$, demonstrating a decrease as $N_{part}$ increases.