The data for the mean squared transverse momentum $\langle p_t^2 \rangle (E_t)$ as function of transverse energy $E_t$ of $J/\psi$ and $\psi'$ produced in Pb-Pb collisions at the CERN-SPS are analyzed and it is claimed that they contain information about the time structure of anomalous suppression. A transport equation which describes transverse motion of $J/\psi$ and $\psi'$ in the absorptive medium is proposed and solved for a QGP and a comover scenario of suppression. While the comover approach accounts for the data fairly well without adjusting any parameter, the fit to the data within the QGP scenario requires to assume anomalous suppression to become effective rather late, $3-4$ fm/c after the nuclear overlap.

The discovery in 1996 of anomalous $J/\psi$ suppression in Pb-Pb collisions at the SPS has been one of the highlights of the research with ultrarelativistic heavy ions at CERN [1]. Does it point to the discovery of the predicted quark gluon plasma (QGP)? Six years later the situation is still confused, since several models - with and without the assumption of a QGP - describe the observed suppression, after at least one parameter is adjusted. The data on the mean squared transverse momentum $\langle p_t^2 \rangle (E_t)$ [2] for the $\psi$ (this symbol stands for $J/\psi$ and $\psi'$) in the regime of anomalous suppression and as a function of transverse energy $E_t$ have received less attention - for no good reason. We claim: $\langle p_t^2 \rangle (E_t)$ contains additional information about the nature of anomalous suppression and may help to distinguish between different scenarios. In this paper we investigate how the time structure of anomalous suppression influences the values $\langle p_t^2 \rangle (E_t)$. This idea has already been considered more than 10 years ago [3] (c.f. also more recent works [4, 5]) and is based on the following observation (Fig.1): Anomalous suppression is not an instantaneous process, but takes a certain time depending on the mechanism. During this time $\psi'$s produced with high transverse momenta may leak out of the parton/hadron plasma and escape suppression. As a consequence, low $p_t \psi'$s are absorbed preferentially and the $\langle p_t^2 \rangle$ of the surviving (observed) $\psi'$s show an increase $\delta \langle p_t^2 \rangle$, which grows monotonically with the mean time $t_A$, when anomalous absorption acts [5]. In this letter we propose a general formalism of how to incorporate the effect of leakage into the various models, which have been proposed to describe
anomalous suppression and we extract information about the time $t_A$ from a comparison with experiment.

It has become customary to distinguish between normal and anomalous values of suppression $S(E_t) = \sigma^\psi(E_t)/\sigma^{DY}(E_t)$ for $\psi$’s produced in nuclear collisions, c.f. reviews [6, 7]. Here, $\sigma^\psi$ and $\sigma^{DY}$ are the production cross section for a $\psi$ and a Drell-Yan pair in an AB collision, respectively. By definition, $\psi$’s produced in pA collisions show normal suppression via inelastic $\psi N$ collisions in the final state and normal increase of $\langle p_t^2 \rangle$ (above $\langle p_t^2 \rangle_{NN}$ in $NN$ collisions) via gluon rescattering in the initial state. These normal effects are also present in nucleus-nucleus collisions and happen, while projectile and target nuclei overlap. Anomalous values of $S$ and $\langle p_t^2 \rangle$ are attributed to the action on the $\psi$ by the mostly baryon free phase of partons and/or hadrons (we call it parton/hadron plasma) which is formed after the nuclear overlap. It may lead to deconfinement of the $\psi$ via colour screening in the QGP, dissociation via gluon absorption or inelastic collisions by the comoving hadrons during the later period of the plasma evolution. In this letter we describe anomalous $\psi$ suppression within a transport theory and apply it to two rather different scenarios: (I) Absorption involving a threshold in the energy density like in a QGP scenario, (II) continuous absorption via comovers.

We denote by $d\sigma^\psi/d\vec{p}_t(E_t)$ the cross section for the production of a $\psi$ with given $p_t$ and in an event with fixed transverse energy $E_t$. It can be related to the phase space density $f^\psi$ via

$$\frac{d\sigma^\psi}{d\vec{p}_t}(\vec{p}_t, E_t) = \lim_{t \to \infty} \int d\vec{b} P(E_t; b) \int d\vec{s} f^\psi(\vec{s}, \vec{p}_t, t; \vec{b}).$$

(1)

Here, $P(E_t; b)$ describes the distribution of transverse energy $E_t$ in events with a given impact parameter $\vec{b}$ between projectile A and target B. We follow ref. [8] in notation for $P(E_t; \vec{b})$ and the values of the numerical constants. The function $f^\psi(\vec{s}, \vec{p}_t, t; \vec{b})$ is the distribution of $\psi$’s in the transverse phase space $(\vec{s}, \vec{p}_t)$ at time $t$ for given $\vec{b}$.

We define $t = 0$ as the time, when the process of normal suppression and normal generation of $\langle p_t^2 \rangle$ has ceased and denote by $f^\psi_N(\vec{s}, \vec{p}_t; \vec{b})$ the distribution of $\psi$’s at this time. $f^\psi_N$ is taken as initial condition for the motion and absorption of the $\psi$’s during the action of anomalous interactions. The evolution of the $\psi$ is described by a transport equation

$$\frac{\partial}{\partial t} f^\psi + \vec{v}_t \cdot \vec{\nabla}_s f^\psi = -\alpha f^\psi.$$  

(2)

The time dependence arises from the free streaming of the $\psi$ with transverse velocity $\vec{v}_t = \vec{p}_t/\sqrt{m_\psi^2 + p_t^2}$ (l.h.s.) and an absorptive term on the r.h.s., where the function $\alpha(\vec{s}, \vec{p}_t, t; \vec{b})$ contains all details about the surrounding matter and the absorption process. We have left out effects from a mean field, because the elastic $\psi N$ cross section is very small and have also neglected a gain term on the r.h.s., because recombination processes $c + \bar{c} + N \to \psi + X$ seem unimportant at SPS energies where at most one $c\bar{c}$ is created per event.
Figure 1: Schematic picture of the leakage phenomenon. Between the Lorentz contracted remnants “B” and “A” of the nuclei which have collided charmonia move in the created parton/hadron plasma. Those $\psi$’s with large transverse velocities $v_t$ (case (1)) may leak out and escape suppression, while low $v_t$ particles may remain (case (2)) leading to an effective increase of $\langle p_t^2 \rangle$ for the surviving (observed) $\psi$’s.

Eq. (2) can be solved analytically with the result

$$f^\psi(s, p_t, t; b) = \exp \left( -\int_0^t dt' \alpha(s - \vec{v}_t(t - t'), \vec{p}_t, t'; \vec{b}) \right) f^\psi_N(s - \vec{v}_t t, \vec{p}_t; \vec{b}),$$

which for $t = 0$ reduces to $f^\psi = f^\psi_N$. If we denote by $t_f$ the time when anomalous suppression has ceased, $\alpha(s, \vec{p}_t, t; \vec{b}) = 0$ for $t > t_f$, the limit $t \to \infty$ in eq. (1) can be replaced by setting $t = t_f$, since the distribution in $p_t$ does not change for larger $t$’s.

There is little controversy about $\psi$ production and suppression in the normal phase: The gluons, which fuse to the $c\bar{c}$, collide with nucleons before fusion and gain additional $p_t^2$. The $\psi$ on its way out is suppressed by inelastic $\psi N$ collisions without any change in $\langle p_t^2 \rangle$. Neglecting effects of formation time, one has

$$f^\psi_N(s, p_t; b) = \sigma_{NN} \int dA dB \rho_A(s, z_A) \rho_B(b - s, z_B) \cdot \exp \left( -\sigma_{abs} [T_A(s, z_A, \infty) + T_B(b - s, -\infty, z_B)] \right) \langle p_t^2 \rangle_N \exp \left( -p_t^2 / \langle p_t^2 \rangle_N \right),$$

where

$$\langle p_t^2 \rangle_N(b, \vec{s}, z_A, z_B) = \langle p_t^2 \rangle_{NN} + a_{gN} \rho_0^{-1} [T_A(\vec{s}, -\infty, z_A) + T_B(b - \vec{s}, z_B, +\infty)]$$

$$\text{(5)}$$
with the thickness function \( T(\vec{s}, z_1, z_2) = \int_{z_1}^{z_2} dz \rho(\vec{s}, z) \). All densities \( \rho_A, \rho_B \) are normalized to the number of nucleons (\( \rho_0 \) is the nuclear matter density). We shortly explain eqs. (4) and (5): For given values \( \vec{b} \) and \( \vec{s} \) in the transverse plane the \( \psi \) is produced at coordinates \( z_A \) and \( z_B \) in nuclei \( A \) and \( B \), respectively. On its way out, the \( \psi \) experiences the thickness \( T_A(\vec{s}, z_A, \infty) \) and \( T_B(\vec{b} - \vec{s}, -\infty, z_B) \) in nuclei \( A \) and \( B \), respectively and is suppressed with an effective absorption cross section \( \sigma_{\psi}^{\text{abs}} \). The two gluons which fuse carry transverse momentum from two sources: (i) Intrinsic \( p_t \), because they had been confined to a nucleon. The intrinsic part is observable in \( NN \to \psi \) collisions and leads to \( \langle p_t^2 \rangle^\psi_{NN} \) in eq. (5). (ii) Collisional contribution to \( p_t \), because in a nuclear collision, the gluons traverse thicknesses \( T_A(\vec{s}, -\infty, z_A) \) and \( T_B(\vec{b} - \vec{s}, z_B, +\infty) \) of nuclear matter in \( A \) and \( B \), respectively, and acquire additional transverse momentum via \( gN \) collisions. This is the origin of the second term in eq. (5).

The constants \( \sigma_{\psi}^{J/\psi}, \sigma_{\psi}^{\psi^\prime} \) and \( \sigma_{gN} \) are usually adjusted to the data from \( pA \) collisions, before one investigates anomalous suppression. Fig. 2 shows (dashed curves) the results for normal suppression \( S^\psi(E_t) \) and \( \langle p_t^2 \rangle^\psi(E_t) \) calculated with \( f_N^\psi \) eq. (4) in eq. (1). While the difference between calculation and data is enormous for the suppression, it is rather small for \( \langle p_t^2 \rangle^\psi(E_t) \).

Since the physical origin of anomalous suppression is not yet settled, we investigate suppression \( S^\psi(E_t) \) and \( \langle p_t^2 \rangle^\psi(E_t) \) for two models, which have rather contradictory assumptions.

I. Threshold (QGP-) scenario: \( \psi \)'s are totally and rapidly destroyed, when they are in a medium with energy density above a critical one, and nothing happens elsewhere. As a representative model we use the approach by Blaizot et al. [8].

II. Comover scenario: The plasma of comoving partons and/or hadrons leads to a continuous absorption of long duration due to inelastic collisions with the comoving particles. As a representative model, we use the approaches by Capella et al. [10] and Kharzeev et al. [12].

In this letter we study the effect of leakage on the observed values of \( \langle p_t^2 \rangle \) within two well established scenarios, using their assumptions and parameters. We do not introduce any modifications like (i) the \( p_t \) dependence of the absorption process (i.e. \( p_t \) dependence of \( \alpha(\vec{s}, \vec{p}_t, t; b) \) in eq. (2)) and (ii) expansion of the plasma during absorption. Both effects may contribute to \( \langle p_t^2 \rangle \) (we will present a qualitative discussion at the end), but both require a detailed study by themselves.

We begin with model I: In their schematic approach Blaizot et al. [8] include anomalous suppression via

\[
f^{\psi}(\vec{s}, \vec{p}_t; \vec{b}) = \theta(n_c - n_p(\vec{b}, \vec{s})) f_N^{\psi}(\vec{s}, \vec{p}_t; \vec{b}).
\]
Here \( n_c \) is a critical density and \( n_p(\vec{b}, \vec{s}) \) is the density of participant nucleons

\[
n_p(\vec{b}, \vec{s}) = T_A(\vec{s}, -\infty, +\infty)[1 - \exp \left(-\sigma^{NN}_{in}T_B(\vec{b}-\vec{s}, -\infty, +\infty)\right)] + (A \leftrightarrow B). \tag{7}
\]

According to eq. (6) all \( \psi \)'s are destroyed if the energy density (which is directly proportional to the participant density) at the location \( \vec{s} \) of the \( \psi \) is larger than the critical density. All other \( \psi \)'s survive. While the prescription eq. (6) successfully describes the data in the full \( E_t \) range of anomalous suppression after the only one free parameter, \( n_c \), is adjusted, the predictions for \( \langle p_t^2 \rangle(E_t) \) are significantly below the data, especially at large \( E_t \) (see below).

The expression eq. (6) for the phase space distribution \( f_\psi \) including anomalous suppression within the threshold model is recovered within our transport approach eq. (3) by setting

\[
\alpha(\vec{s}, \vec{p}, t; \vec{b}) = \alpha_0 \theta(n_p(\vec{b}, \vec{s}) - n_c)\delta(t) \tag{8}
\]

and taking the limit \( \alpha_0 \to \infty \). The delta function \( \delta(t) \) has to be included to recover the expression eq. (6) and may be understood by the physical picture that the energy density is highest at \( t = 0 \) and anomalous absorption most effective then.

There are various ways to introduce another time structure into the absorption term. We will try two options, one being

\[
\alpha(\vec{s}, \vec{p}, t; \vec{b}) = \alpha_0 \theta(n_p(\vec{b}, \vec{s}) - n_c)\delta(t-t_A). \tag{9}
\]

The idea of a threshold density is kept, but suppression does not act at \( t = 0 \) but at a later time \( t_A \), which time is then determined from a comparison with the data. For times \( t > t_A \), one finds from the general solution eq. (3) (and \( \alpha_0 \to \infty \))

\[
f_\psi(\vec{s}, \vec{p}, t) = \theta(n_c - n_p(\vec{b}, \vec{s}))f_N(\vec{s} - \vec{v}_tt_A, \vec{p}), \quad t > t_A \tag{10}
\]

which differs from the expression (6), by the motion in phase space of \( f_N^\psi \) during the time \( 0 \leq t \leq t_A \). For \( t > t_A \) the momentum distribution derived from \( f_\psi \) does not change any more.

The suppression \( S_\psi(E_t) \) and \( \langle p_t^2 \rangle_\psi(E_t) \) calculated with the distribution \( f_\psi \) from eq. (10) depends on the parameter \( t_A \). We use the parameters of ref. [8] where available, i.e. for \( J/\psi \):

\[
\sigma_{J/\psi}^{abs} = \frac{6.4 \text{ mb}}{f^{-1}} \quad n_c = 3.75 \text{ fm}^{-2}
\]

The parameters for the generation of \( \langle p_t^2 \rangle \) by gluon rescattering are taken from a fit to the \( pA \) data [2]:

\[
\langle p_t^2 \rangle_{NN} = 1.11 \text{ (GeV}/c)^2\quad \text{and} \quad a_{gN} = 0.081 \text{ (GeV}/c^2)^{-1}
\]

We also account for the transverse energy fluctuations [8] which have been shown to be significant for the explanation of the sharp drop of \( J/\psi \) suppression in the domain of very large \( E_t \) values, by replacing \( n_p \) by \( \frac{E_t}{\langle E_t \rangle}n_p \) where \( \langle E_t \rangle \) is the mean transverse energy at given \( \vec{b} \). We then calculate \( \sigma_{J/\psi}^{abs}/\sigma_{DY} \) as a function of \( t_A \). Since the critical density for \( J/\psi \) is quite high, the leakage affects only the very high momentum \( \psi \)'s. Since their number is small, the calculated results for the suppression \( S_\psi(J/\psi)(E_t) \) depend only little on \( t_A \) (Fig. 2).
We turn to a discussion of $\langle p_T^2 \rangle_{J/\psi}(E_t)$. Fig. 2 shows calculated curves for values of $t_A = 0$ to 4 fm/c. The dotted line ($t_A = 0$) is the result of the original threshold model with immediate anomalous suppression (and has been predicted in [13]). It fails badly at large values of $E_t$. Also no other curve with a given $t_A$ describes the data for all values of $E_t$. We have to conclude that $t_A$ depends on $E_t$: $t_A(E_t)$. The larger the values $E_t$ the later anomalous suppression acts. From a comparison with data we have $t_A \lesssim 2$ fm/c for $E_t < 80$ GeV, $t_A \simeq 2.5$ fm/c for $80 \leq E_t \leq 100$ GeV, and $t_A \simeq 3.5$ fm/c for $E_t > 100$ GeV.

Figure 2: Nuclear suppression $\sigma^{\psi}/\sigma^{\text{DY}}$ and $\langle p_T^2 \rangle_\psi$ for $J/\psi$ (above) and $\psi'$ (below) as a function of transverse energy $E_t$. Data are from [11] for $\sigma^{\psi}/\sigma^{\text{DY}}$ and from [2] for $\langle p_T^2 \rangle$. Dashed curves show the result of normal suppression alone. The dotted lines correspond to the original threshold model, which in our notation corresponds to $t_A = 0$. The other curves include anomalous suppression within the threshold model eqs. (6) and (10), where anomalous suppression is assumed to act at time $t_A > 0$. The curves are labeled by the values $t_A = 1, 2, 3, 4$ fm/c for $J/\psi$ and $t_A = 2, 4, 6$ fm/c for $\psi'$. Also the curves in $\sigma^{\psi}/\sigma^{\text{DY}}$ carry these labels, lowest curve $t_A = 0$ and monotonic increase with $t_A$.

A similar analysis is performed for the $\psi'$. The data for the suppression are taken from [14, 15], those for $\langle p_T^2 \rangle_{\psi'}$ from [2]. Since the $\psi'$ has not been treated by Blaizot et al., we fit the suppression data and find values for the absorption parameters, $\sigma^{\psi'}_{\text{abs}} = 7$ mb, $n_c^{\psi'} = 2.3$ fm$^{-2}$, leaving the parameters $\langle p_T^2 \rangle_{NN}$ and $a_{gN}$ unchanged. Since the critical density for $\psi'$ is
smaller than that for $J/\psi$, more $\psi'$s leak out of the anomalous suppression region, the change in suppression due to the increase of $t_A$ is noticeable. Fig. 2 shows the results with a good fit to the suppression data.

The data for $\langle p_t^2 \rangle_{\psi'}$ have rather large error bars. The calculated curves show a strong dependence on $t_A$, again with a trend that $\langle p_t^2 \rangle$ at larger values of $E_t$ require larger values of $t_A$. However, the numerical values $t_A(E_t)$ for $\psi'$ are above those for the $J/\psi$ by about $1 - 2\,fm/c$. This result is strange, because we expect the $\psi'$ to be destroyed more easily and therefore more rapidly. We will come back to this point in the conclusions.

The time structure introduced via eq. (9) is certainly oversimplified. Rather than having it act at one particular time $t_A$, it is more reasonable to assume that it acts during a time interval. Therefore we have also investigated the following form of the absorptive term

$$\alpha(s, p_t, t; \bar{b}) = \frac{\alpha_0}{t_1 - t_0} \theta(n_p(\bar{b}, s) - n_c) \theta(t_1 - t) \theta(t - t_0).$$

(11)

Anomalous suppression acts for times $t$ between $t_0$ and $t_1$ (which may be a function of $E_t$). Especially $t_0$ is an interesting quantity, because it gives the starting time for anomalous sup-
We calculate the suppression $S_{J/\psi}(E_t)$ and $\langle p_t^2\rangle_{J/\psi}(E_t)$ as a function of $[t_0, t_1]$ and for $\alpha_0 = 5$. All other parameters are unchanged. The suppression is well described for all intervals $[t_0, t_1]$, but the calculations for $\langle p_t^2\rangle_{J/\psi}(E_t)$ depend strongly on the choice of this interval. Fig. 3 shows some representative examples: The four windows in Fig. 3 display events with $t_0 = 0, 1, 2, 3 \text{ fm/c}$, respectively, and several values for $t_1$. It is obvious that there is not one curve, which describes all the data. Rather we find that the best curve for $E_t < 60 \text{ GeV}$ corresponds to $[0, 1]$, for $60 < E_t < 80 \text{ GeV}$ to $[1, 3]$, for $80 < E_t < 100 \text{ GeV}$ to $[2, 5]$, and for $E_t > 100 \text{ GeV}$ to $[3, 5]$. The length $t_1 - t_0$ of the interval remains approximately constant, while the beginning time $t_0$ increases with $E_t$. We will discuss the significance of these results after we have treated leakage within the comover model.

We proceed to model II, the scenario of comovers: Partons and/or hadrons which move with the $\psi$ may destroy the $\psi$ with a cross section $\sigma_{co}^{\psi}$. The comover density $n_{co}(t)$ depends on time, for which the Bjorken scenario of longitudinal expansion predicts $t^{-1}$. The absorptive term $\alpha$ in eq. (2) then takes the form

$$
\alpha(s, p_t, t; b) = \sigma_{co}^{\psi} \frac{n_{co}(b, s)}{t} \theta(\frac{n_{co}(b, s)}{n_f} t_0 - t) \theta(t - t_0),
$$

i.e. absorption by comovers starts at $t = t_0$ and ends at $t_1$, when the comover density $n_{co}(b, s) \cdot t_0/t_1$ has reached a value $n_f$ independent of $b$ and $s$. The comover approach contains a definite time structure for anomalous suppression and we have not changed it. We also account for the transverse energy fluctuations\[10\], by replacing $n_c$ by $\frac{E_t}{E_{\perp}} n_c$, and the transverse energy loss\[11\] induced by the $J/\psi$ trigger, by rescaling $< E_t >$ by a factor $\frac{n_{co}^{\psi}}{n_p}$, which have been shown to be significant for the explanation of the sharp decrease of $S_{J/\psi}(E_t)$ at $E_t > 100 \text{ GeV}$. In the choice of parameters we have followed \[10\] \[12\]: $n_{co}(b, s) = 1.5 n_p(b, s)$ with the participant density from eq. (7), $t_0 = 1 \text{ fm/c}$, $n_f = 1 \text{ fm}^{-2}$, $\sigma_{abs}^{J/\psi} = 4.5 \text{ mb}$, $\sigma_{abs}^{\psi} = 6 \text{ mb}$, $\sigma_{co}^{J/\psi} = 1 \text{ mb}$ and $\sigma_{co}^{\psi} = 3 \text{ mb}$, no additional parameter has been introduced.

The calculated $E_t$ dependence of the suppression shown in Fig. 4 fits the data acceptably well like model I. For $\langle p_t^2 \rangle$, we compare the results of two calculations with the data. The dashed lines in Fig. 4 show the calculation leaving out leakage. Formally this limit is obtained from eq. (3) by setting $\vec{v}_t = 0$ in the exponent and in $f_N^\psi$. Due to the introduction of the $E_t$ loss which is necessary to recover the $J/\psi$ suppression for large values of $E_t$, the case without leakage does not fit the data $\langle p_t^2 \rangle$ even in the domain of low $E_t$ values. Only when the leakage effect is taken into account, the calculation agrees well with the data. We stress: the calculation of $\langle p_t^2 \rangle (E_t)$ in the comover model is a true prediction in the sense that no parameter is adjusted above those which are fitted to the suppression $S_{J/\psi}(E_t)$. We have also calculated the mean time
\langle t_A^{J/\psi} \rangle \text{ for comover action by studying the suppression } S^\psi(E_t) \text{ as a function of time and taking the mean of } t \text{ with the weight } dS^\psi(E_t)/dt \text{ and find}

\begin{align*}
\langle t_A^{J/\psi} \rangle &= 3.5 \text{ fm/c} \\
\langle t_A^{\psi'} \rangle &= 3.0 \text{ fm/c},
\end{align*}

which values include the time \( t_0 \), eq. (12), between the end of normal suppression and the beginning of comover action. The values eq. (13) are found to be rather independent of \( E_t \).

![Figure 4: Nuclear suppression \( \sigma^\psi/\sigma^{DY} \) and \( \langle p_T^2 \rangle \) for \( J/\psi \) (above) and \( \psi' \) (below) as a function of transverse energy \( E_t \). Dotted and solid lines are calculated in the comover model without and with considering the leakage effect, respectively.](image)

We summarize: In this letter we have investigated the influence of leakage on the calculation of \( \langle p_T^2 \rangle(E_t) \) for \( J/\psi \) and \( \psi' \) produced in Pb-Pb collisions at SPS energies. This effect is closely related to time structure of anomalous suppression. The evolution of the \( \psi \) during anomalous suppression including leakage is described within a general transport equation. The formalism is applied to two models with rather contradictory underling physical assumptions, the threshold (QGP) and the comover models with the following results:

(i) Calculations within the original models, where leakage is left out, do not describe the data for \( \langle p_T^2 \rangle^\psi(E_t) \), the discrepancy being particularly strong at high values of \( E_t \).
Including leakage into the comover model, without changing its structure nor its parameters leads to a good agreement with the data for $\langle p_t^2 \rangle (E_t)$ for $J/\psi$ over the full range of values $E_t$.

The assumption in the threshold model that anomalous suppression acts instantaneously at $t_A = 0$, i.e. right after normal suppression is not supported by experiment. Rather for central collisions, the data of $\langle p_t^2 \rangle$ are described best, if one assumes that anomalous suppression acts at a time $t_A = 3 - 4 \text{ fm}/c$ after the nuclear overlap.

The situation for the $\psi'$ is less clear. While the data for suppression can be fitted by properly adjusting the parameters, both models underpredict the data for $\langle p_t^2 \rangle^{\psi'}(E_t)$. This is evident for the comover model. Within the threshold model the values $t_A$ required to fit the $\psi'$ data are larger than those for the $J/\psi$, which seems unreasonable to us. It could be that the error bars on the experiments are too small.

We conclude: In this letter we have investigated how leakage (escape of high $p_t \psi'$s) when introduced into exciting models of anomalous suppression influences the calculated values of $\langle p_t^2 \rangle (E_t)$. As mentioned already above, there could also be other effects which could influence $\langle p_t^2 \rangle$. We discuss them briefly: $p_t$ dependence of the mechanism responsible for anomalous suppression and transverse expansion of the plasma. Both effects can be treated within the transport approach eqs. (2), (3) by introducing an explicit dependence on $p_t$ and a modified dependence on $t$ into the function $\alpha(s, p_t, t; b)$. A first and schematic investigation on the $p_t$ has been made by Dingh [18]. One sees without calculation that if $\partial \alpha / \partial p_t < 0$, i.e. high $p_t \psi'$s are absorbed less (whatever the mechanism may be), the calculated values of $\langle p_t^2 \rangle (E_t)$ increase, thus having the same effect as an increase of the time $t_A$ eq. (9), when anomalous suppression acts. As for the transverse expansion of the plasma, its effect is qualitatively clear: It reduces the effect of leakage since it makes it harder for the $\psi'$s to escape. Although the qualitative changes on $\langle p_t^2 \rangle$ of the two effects are clear, their quantitative treatment needs a rather careful study of the underlying physics, necessitates to modify existing models, to change their input parameters and to introduce new parameters. This is beyond the scope of this letter.

Acknowledgments: We are grateful to our friends and colleagues A. Gal, J. Dolejši, Yu. Ivanov, B.Z. Kopeliovich, H.J. Pirner and C.Volpe for help and valuable comments. One of the authors (P.Z.) thanks for the hospitality at the Institute of Theoretical Physics. This work has been supported by the grant 06HD954 of the German Federal Ministry of Science and Research and by the Chinese National Science Foundation.
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