Abstract

We suggest that the new regime of $J/\psi$ suppression in Pb-Pb collisions found by the NA50 experiment at CERN is the result of non-trivial space-time evolution due to specific behavior of the Equation of State (EOS) near the QCD phase transition. We also study another suppression channel, the conversion of $J/\psi$ into $\eta_c$ during the late cool hadronic stage, and find it rather inefficient.
1. Recent studies by the NA50 collaboration \cite{1} in Pb(158 GeV-A)-Pb collisions show significant $J/\psi$ suppression compared to an extrapolation of known trends in lighter beams and p-A collisions. The data show a new suppression mechanism for impact parameters $b < 8\,fm$. Surprisingly, this new suppression regime seems to set in very sharply. Such abrupt behavior may be a manifestation of dramatic collective effects, related to qualitative changes in hot/dense hadronic matter due to the QCD phase transition. \footnote{We stress however, that we discuss the new trend over the whole region and not just its onset.}

Logically speaking, additional $J/\psi$ suppression can be explained either by (i) increasing absorption rates or by (ii) increasing the time that $J/\psi$ spends in the hot/dense medium. In contrast to the available literature devoted to the first scenario we study the second one. An increase in the Quark-Gluon Plasma(QGP) lifetime was found in the hydrodynamic framework in \cite{2}, as a result of a discontinuity in the Equation of State (EOS) at the QCD phase transition.

2. The oldest and the simplest argument for suppression \cite{3} is that once QGP is formed, free gluons can rather easily “ionize” $J/\psi$ in an elementary process $gJ/\psi \to \bar{c}c$ similar to the photo-effect \cite{8,4,5}. Later \cite{6,7} it was argued that successive charmonium levels $\psi'$, $\chi$, $J/\psi$ “melt” at high temperatures and densities due to Debye screening in the QGP, and ionizing gluons are superfluous.

There is a significant phenomenological difference between these two mechanisms. “Melting” implies the existence of certain thresholds as a function of matter energy density and can therefore generate a rapid onset of suppression. Furthermore, it naturally explains why different charmonium states are suppressed at different densities. So, although there are currently no realistic models of this type, this scenario may be successful.

Gluonic excitation on the other hand, is simply a probabilistic process, and like any other absorption processes (“co-movers” for example) cannot create discontinuous behavior by itself. Therefore a recent paper \cite{5} proposed that the formation of QGP is discontinuous; QGP is produced suddenly after a certain critical energy density $\epsilon_c$ is reached. This work was motivated by the classical nucleation theory of “bubbles” at a first-order phase transition and has a number of theoretical and phenomenological problems. First, the thermodynamics of the transition tells us that plasma formation is discon-
tinuous in temperature but not in energy density. A mixed phase linearly interpolates between the energy densities of the two phases, $\epsilon_{\text{min}}$ and $\epsilon_{\text{max}}$. Second, this scenario has a phenomenological problem: a jump into the QGP phase at some $\epsilon_c$ changes the equation of state discontinuously and implies a jump in entropy (or multiplicity) [21]. Preliminary data [1] on $<n_{\text{ch}}(E_T)>$ shows no discontinuity.

3. As deduced from lattice simulations at finite temperature, the QCD phase transition leads to a very non-trivial EOS. The ratio of pressure to energy density, $p/\epsilon$ has a deep minimum at the end of the phase transition, $\epsilon_{\text{max}} \sim 1 - 2 GeV/fm^3$. For the resonance gas and the high density QGP this ratio is relatively large (.2 and 1/3 respectively) but is only about .05 at this minimum. This minimum is referred to as the softest point of the EOS and is associated with a slow hydrodynamic expansion. When the initial energy density was close to the softest point, the QGP was found to live 2-3 times longer than the more energetic and more central collisions at the SPS [2].

In this letter we study the extent to which the non-trivial features of $J/\psi$ suppression can be explained by such an increase in the QGP lifetime. The hydro solution in [2] was made only for azimuthally symmetric central collisions. Assuming that about half of the collision energy is contained in the fully stopped part of the matter, the above authors expect $\epsilon_{\text{max}}$ to be reached at $P_{\text{lab}} \sim 40 GeV/A$ for central AuAu collisions. (As this lies between the nominal AGS and SPS energies, they proposed making additional low energy runs at the SPS.) Of course, the initial energy density decreases with impact parameter and since all values of $b$ are triggered on anyway, scanning $b$ is easier than scanning the collision energy in the experiment. Examining the dependence of energy density on impact parameter, one can see that indeed an impact parameter $b \approx 8 fm$ in Pb-Pb collision roughly corresponds to the energy density of the observed change of trend in $J/\psi$ suppression.

Before turning to specifics, we point out a qualitative difference between our scenario and the one based on “melting”. In the latter case the various charmonium states are completely suppressed for all $\epsilon > \epsilon_c$, while in our

\[ ^2 \text{The arguments of [3] rely on the non-equilibrium kinetics of bubble formation. Note however, that a very small surface tension found by lattice studies, about } 0.01 T_f^3, \text{ makes the formation of small bubbles quite probable [11].} \]

\[ ^3 \text{Roughly, since } p \text{ is the moving force and } \epsilon \text{ is the mass to be moved, acceleration is proportional to this ratio.} \]
case the suppression is centered around $\epsilon \approx \epsilon_{max}$ and the survival probability increases for $\epsilon > \epsilon_{max}$. This difference is manifest in the specific predictions discussed below.

Hydro evolution for non-central collisions is much more complicated because of a directed flow (see e.g. [12, 13]). Fortunately, we avoid these effects by restricting the discussion to $J/\psi$ suppression. For the first few fermi/c expansion is predominantly longitudinal, and thus similar for central and non-central collisions.

4. It is clear that significant changes in the QGP lifetime should affect $J/\psi$ suppression, even if absorption rates are unchanged. However, $J/\psi$ may escape from the system, which obviously limits the sensitivity to QGP lifetime.

In order to reproduce many known observables, such as the distribution in the transverse energy $E_T$ and its correlation with $b$, we have constructed a small event generator to model our scenario. For Pb(158-GeV)-Pb collisions at a given impact parameter, we generate charmonium events which have survived nuclear absorption and are now in a hot de-confined medium. The survival probability for a given event is $e^{-\Gamma t_p}$ where $t_p$ is the time spent transversing this medium and $\Gamma$ is the ionization rate. $\Gamma$ is chosen to match the observed 25% suppression immediately following the jump. Our goal is to provide a reasonable explanation for the region near the discontinuity without too much regard for the most central points.

In our model $J/\psi$ interacts with the hot medium only when the wounded nucleon density, $n_w$, is greater than $3.1 \text{ fm}^{-2}$. This condition is similar to [14] and determines the transverse size of the region. We estimate the longitudinal size from the Bjorken formula. The average energy density reached in a collision at impact parameter $b$ is

$$<\epsilon> = \frac{3}{\Delta y c \tau} \frac{dE_T}{dA} = \frac{3 q}{\Delta y c \tau} <n_w(b)>$$

where $<n_w(b)>$ is the average wounded nucleon density, and $q = .4 \text{ GeV}$ is a calorimeter dependent constant as defined in [15]. Since the jump is observed for $b \approx 8.5 \text{ fm}$ and $<n_w(b = 8.5)> \approx 2.6 \text{ fm}^{-2}$, we choose $\tau = 1.9 \text{ fm}$ to match the energy density of the jump to the softest point, $\epsilon_{max} = 1.6 \text{ GeV/fm}^3$. The thickness is then $2c\tau$, or the distance the colliding pancakes have receded during the equilibration time.
In this model \( J/\psi \) suppression begins only for wounded nucleon densities that are 20% larger than the softest point average, i.e., significant suppression begins only above the mixed phase, \( \epsilon_{\text{max}} \). For the lifetime of this hot region which destroys \( J/\psi \) we take the forms shown in Fig. 1a and calculate the corresponding suppression patterns shown in Fig. 1b. To calculate the time inside plasma we need to distribute these events in coordinate and momentum space. The distribution in the transverse plane is given by Glauber theory, like in [15] and the distribution in the longitudinal direction is presumed uniform. We have calculated transverse momentum distribution of \( J/\psi \) within Glauber theory (as in [14]). The longitudinal momentum or \( x_f \) distribution is known from p-A data and in the plasma frame follows the form \( \sim (1 - |x_f|)^3 \). [17, 18]

We can then determine the time \( t_p \) that \( J/\psi \) remains in plasma and the corresponding survival probability, \( e^{-\Gamma t_p} \), and find the \( J/\psi \) yield

\[
\frac{d^2\sigma^{J/\psi}}{db^2} = \int_{\vec{s},z,\vec{p}_T,x_f} \frac{d^2\sigma^{J/\psi}_{\text{Glauber}}}{db^2} f_b(\vec{s}) f(z) f_{b,s}(\vec{p}_T) f(x_f) e^{-\Gamma t_p}
\]

Here \( \vec{s} \) is the coordinate in the transverse plane and each \( f \) is a distribution functions described above. Such a \( b \) dependent cross section may be converted to an \( E_T \)-dependent cross section using the \( E_T - b \) assignment of [15]. Dividing this cross section by the Glauber value we compare our survival probability with preliminary NA50 experimental results in Fig. 1b.

The usual Glauber-type \( J/\psi \) absorption on nucleons describes well p-A, lighter ion, and peripheral Pb-Pb data and and therefore we plot only deviations from this trend in Fig. 1b. The increased QGP lifetime leads to additional suppression for intermediate \( E_T \) between 50 – 100 GeV. Since \( J/\psi \) may leak out of the plasma region the discontinuity due to the lifetime is softened and we do not reproduce the very sharp jump seen in the data. However, agreement with most of the points in the intermediate region is significantly improved. Furthermore, our curves reproduce the vanishingly small slope of the suppression in this region.

5. The \( p_T \)-dependence of the suppression may separate the different suppression mechanisms discussed above. We have therefore calculated \( <p_T^2> \) vs. \( E_T \). As a function of \( b \), \( <p_T^2> \) is given by

\[
<p_T^2>(\vec{b}) = <p_T^2> \frac{d^2\sigma^{J/\psi}}{d^2b} > \frac{d^2\sigma^{J/\psi}}{d^2b}
\]
Converting the numerator and denominator to $E_T$ dependent cross sections we plot $<p_T^2>$ vs. $E_T$ in Fig. 1c. As a result of the usual initial state parton re-scattering, $<p_T^2>$ increases with the centrality of the collision. If $J/\psi$ “melt” in the central region then only peripheral $J/\psi$ survive, and $<p_T^2>$ decreases for the most central collisions, as suggested by Kharzeev et al. In our scenario however, some $J/\psi$ may still “leak” even from the central regions, and this flattens out the $<p_T^2>$ vs. $E_T$ dependence. This dependence was indeed observed by the NA50 collaboration [1].

Figure 1: For the different lifetimes of the plasma region shown in (a) the corresponding suppression pattern is shown in (b). (c) The $J/\psi$ $p_T$ distribution of this work (curve C in (a)) compared to [16].

6. This part is rather independent from the rest of the paper; Here we discuss the conversion of $J/\psi$ into its spin partner $\eta_c$. This channel is one more potential “sink” for $J/\psi$, which (to our knowledge) has not been investigated previously.

These reactions may be important in the later stages of the collision. The “photo-effect” above requires $\sim 1$ GeV gluons to ionize $J/\psi$ and therefore can only proceed during the hot plasma stage of the collision. In contrast, since $J/\psi$ and $\eta_c$ are nearly degenerate, spin exchange reactions require little input energy and may proceed during the cool hadronic stage, which lasts for 10-20 fm/c in Pb-Pb collisions. If these reactions have significant rates, some fraction of $J/\psi$ can be converted into $\eta_c$. The inverse reaction is of course also possible, but the rather short natural lifetime $\tau_{\eta_c} \simeq 15 fm/c$ prevents...
significant feedback.

The two “spin exchange” reactions we consider here are \( \psi + \pi \to \eta_c + \rho \) and \( \psi + \rho \to \eta_c + \pi \). Below we calculate these reaction rates in a hadronic gas at \( T=150 \) MeV using a non-relativistic quark model.

First we consider the reaction \( \psi + \pi \to \eta_c + \rho \). In a non-relativistic quark model this reaction proceeds via a spin exchange collision between the c quark in \( \psi \) and the u quark in \( \pi \). We use a point like spin-spin interaction between the two particles in the center of mass frame of the pion and \( J/\psi \).

\[
H_{\text{int}} = C\vec{\sigma}_c \cdot \vec{\sigma}_d \delta^3(\vec{r}_\pi - \vec{r}_\psi)
\]

From D-D* meson splitting we take the value \( C = 4.4 \) (GeV \(^{-2}\)). Using this interaction Hamiltonian and assuming identical spatial wave functions for \( \pi, \rho \) and \( \psi, \eta_c \), one obtains the following cross section

\[
\sigma_{s \to s'} = \frac{C^2}{v_\pi + v_\psi} |\langle s' | \vec{\sigma}_c \cdot \vec{\sigma}_\pi | s \rangle|^2 \int \frac{d^3p_\eta_c}{(2\pi)^3} \frac{d^3p_\rho}{(2\pi)^3} (2\pi)^4 \delta^4(p_{\text{tot}})
\]

The unpolarized cross section is found by summing the cross section over final spins and averaging over initial spins

\[
\sigma(E_\pi) = \frac{C^2}{\pi} \left[ \frac{E_\rho p_\rho}{v_\pi + v_\psi} \right]_{\text{cm}} \langle \langle \sigma \cdot \sigma \rangle \rangle
\]

where \( \langle \langle \sigma \cdot \sigma \rangle \rangle = \frac{1}{N_s} \sum_{ss'} |\langle s' | \vec{\sigma}_c \cdot \vec{\sigma}_\pi | s \rangle|^2 \)

for \( J/\psi + \pi \to \eta_c + \rho \) we have \( \langle \langle \sigma \cdot \sigma \rangle \rangle = 1 \) and for \( J/\psi + \rho \to \eta_c + \pi \) we have \( \langle \langle \sigma \cdot \sigma \rangle \rangle = 1/3 \). The value of the cross section(3) is determined from the incident energy of the pion and the kinematics of the collision. For a stationary \( J/\psi \) the threshold for this process occurs when \( E_\pi \simeq 710 \) MeV. The cross section is about 1.2mb at 800 MeV.

Now we calculate the reaction rate in an ideal pion gas for the process \( J/\psi + \pi \to \eta_c + \rho \). The \( J/\psi \) is at rest in the pion gas. In this energy range the

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\(^4\)The exact physical nature of the spin-spin forces between light and heavy quarks is not important for our results. Note however that it may be either due to the one-gluon exchange (as assumed in the works mentioned above) or due to small size instantons, see S. Chernyshev, M.A. Nowak, I. Zahed, Phys. Rev. D53, 5176 (1996).
pion mass is neglected and Boltzmann statistics are reasonable. The reaction rate is given by (flux) x (cross section)

\[ W_{\psi \rightarrow \eta_c} = \frac{g_\pi}{2\pi^2} \int_{\text{threshold}}^{\infty} e^{-E_\pi/v_\pi} v_\pi E_\pi^2 \sigma(E_\pi) dE_\pi \]  

(4)

where \( g_\pi = 3 \) is the isospin of the pion. Using the cross section above (3) we have calculated the reaction rates numerically

\[ W_{J/\psi + \pi \rightarrow \eta_c + \rho} = 0.29 \text{ MeV} \]  

(5)

A very similar calculation gives the reaction rate for \( J/\psi + \rho \rightarrow \eta_c + \pi \). For example the unpolarized cross section is given by \( C^2/\pi^2 [p_\pi E_\pi/(v_\rho + v_{J/\psi})]_{cm} \langle \langle \sigma \cdot \sigma \rangle \rangle \). For this process we find

\[ W_{J/\psi + \rho \rightarrow \eta_c + \pi} = 0.30 \text{ MeV} \]  

(6)

It has been suggested that the mass of the \( \rho \) shifts downward in medium, and the CERES dilepton data seem to support this claim. For reference we calculate the above rates for \( m_\rho = 385 \text{ MeV} \)

\[ W_{J/\psi + \pi \rightarrow \eta_c + \rho} \approx \]  

\[ W_{J/\psi + \rho \rightarrow \eta_c + \pi} \approx 1.1 \text{ MeV} \]  

(7)

Even these rates are rather small, and we conclude that it is unlikely that this channel can kill more than few percent of \( J/\psi \).

Finally, we consider experimental checks of \( \eta_c \) production. Its \( \gamma \gamma \) branching is too small to be useful. Among the known modes the most promising seems to be a \( \phi \phi \) decay with branching ratio of .7\%. As the \( \phi \phi \) background is not large, \( \phi \) identification/mass resolution in the KK channel is generally good, and the \( \eta_c \) peak is rather narrow, \( \eta_c \) identification may be feasible.

7. Discussion and Outlook. Recent findings by the NA50 collaboration indicate that an additional \( J/\psi \) suppression mechanism sets in for \( b < 8 \text{ fm} \). This suppression can be created by a change in the absorption of \( J/\psi \), or by a change in the space-time evolution, or both. “Melting” may increase the absorption of \( J/\psi \) in plasma, while a change in the EOS may increase the lifetime of the plasma itself.

Two qualitative differences separate these scenarios experimentally. First, “melting” is not directly related to the phase transition, but is a subtle phenomenon where basically \( \chi \) states no longer exist. The space-time evolution
on the other hand, is directly related to the EOS and hence to the transition point itself. Experimentally therefore, a change of behavior at the corresponding impact parameter in the hadronic observables sensitive to the transverse expansion, e.g. the $p_t$ slopes (especially of heavy secondaries N, d), HBT transverse sizes etc., would signal a change in the space-time evolution of the system.

A second difference, is that while “melting” is presumed to destroy completely the various charmonium states, the absorption in our scenario comes from the usual probabilistic rates, and therefore $J/\psi$ may escape from the central plasma region. This produces a qualitative difference in the mean transverse momentum vs. centrality, as discussed above. Much more detailed studies of the $p_t$ dependence of the suppression are now underway and should clarify this issue.

There remain some unresolved experimental issues with the preliminary data we use. More statistics and further analyses are needed to explain some fluctuations in the data points and most importantly to verify the shape of the transition to the new regime, which now appears as a discontinuous jump. Finally, it is important to test experimentally an assumption, common to all explanations given above, i.e. that the phenomenon is induced at some value of energy density. This can be done by lowering the collision energy of the Pb beam.

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