Charmonium suppression at RHIC and SPS: a hadronic baseline

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Abstract

A kinetic equation approach is applied to model anomalous $J/\psi$ suppression at RHIC and SPS by absorption in a hadron resonance gas which successfully describes statistical hadron production in both experiments. The puzzling rapidity dependence of the PHENIX data is reproduced as a geometric effect due to a longer absorption path for $J/\psi$ production at forward rapidity.

Keywords: $J/\psi$ suppression, heavy-ion collisions, hadron gas

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1. Introduction

The effect of $J/\psi$ suppression \cite{1} is one of the diagnostic tools to prove quark-gluon plasma (QGP) formation in heavy-ion collisions. After a rich phenomenology provided by the CERN-SPS experiments NA3, NA38, NA50 and NA60, which nevertheless could not unambiguously prove QGP formation, the situation became even more complex due to first results from RHIC experiments PHENIX and STAR. New experiments are planned at CERN-LHC and FAIR-CBM, for a recent review, see \cite{2}. In order to make progress in the interpretation of experimental findings it is important to define benchmarks, such as a baseline from a statistical kinetic model provided in this communication, which is based on a purely hadronic description.

Measurements of $J/\psi$ suppression in Au-Au and Cu-Cu collisions by STAR and PHENIX experiments \cite{3,4,5} brought up two unexpected re-
results. First, at forward rapidity \( J/\psi \) is more suppressed than at midrapidity. Second, the dependence of \( J/\psi \) suppression at midrapidity on the number of participants \( N_{\text{part}} \) coincides with the one observed at SPS Pb-Pb collisions by the NA50 experiment \([7]\). Two different mechanisms to explain those RHIC data are usually considered in the literature: recombination \([8, 9, 10, 11]\) or statistical coalescence \([12, 13, 14, 15]\) and nuclear effects \([16, 17, 18]\).

We present a systematic step towards a general description of \( J/\psi \) absorption in the framework of a statistical analysis. The formalism allows \([19, 20]\) for a unified description of the data from the NA38 and NA50 experiments at CERN and from the PHENIX experiment at RHIC. Both PHENIX results, concerning rapidity and \( N_{\text{part}} \) dependences of \( J/\psi \) suppression in Au-Au and Cu-Cu collisions are simultaneously reproduced.

The \( J/\psi \) absorption is caused here by the effective hadronic medium consisting of a multi-component non-interacting hadron resonance gas (HRG). All hadrons from the lowest up to 2 GeV mass are taken into account as constituents of the matter \([21]\). The gas is in thermal equilibrium and expands longitudinally and transversally according to relativistic hydrodynamics \([22]\). \( J/\psi \) suppression is the result of inelastic scattering on constituents of the HRG and on nucleons in the cold nuclear matter (CNM) of the colliding ions. Both sources of \( J/\psi \) suppression, namely absorption in CNM and HRG, are considered simultaneously.

The model applied here is a straightforward generalization of the model of Refs. \([23, 19]\) to include the non-zero rapidity case. In our model, \( J/\psi \) suppression is the result of a \( J/\psi \) final state absorption in a confined medium through interactions of the type

\[
J/\psi + h \rightarrow D + \bar{D} + X,
\]

where \( h \) denotes a hadron, \( D \) is a charm meson and \( X \) stands for a particle which assure conservation laws (charge, baryon number, strangeness). According to the above assumption, charmonium states can be absorbed first in nuclear matter and soon after in the hadron gas. In our phenomenological analysis we assume universal cross sections for baryons, with the appropriate thresholds for their dissociation reactions but energy-independent, \( \sigma_{J/\psi N} = 4 \) mb. This choice corresponds to the absorption cross section \( \sigma_{\text{abs}} = 4.2 \pm 0.5 \) mb for \( J/\psi \) in cold nuclear matter measured in p-A collisions by the NA50 experiment \([24]\), shown to be compatible \([18, 25]\) with recent d-Au data from
the PHENIX experiment [26]. The absorption cross section for meson impact is taken as \( \sigma_m = 2\sigma_{J/\psi N}/3 \), according to quark counting rules.

An A-A collision at impact parameter \( b \) generates an almond shaped overlap region \( S_{\text{eff}} \) as presented in Fig. 1. The time \( t = 0 \) corresponds to the moment of the maximal overlap of the nuclei. After about half of the time the nuclei need to cross each other, matter appears in the central region. Soon thereafter matter thermalizes and this moment, \( t_0 \), is estimated at about 1 fm. In the real situation the longitudinal thickness of the matter at \( t_0 \) is also of the order of 1 fm. Then matter starts to expand and cool according to relativistic hydrodynamics (for details see [22, 19, 20]) until reaching the freeze-out temperature \( T_{f.o.} \) and subsequently freely streaming towards the detectors.

A \( c\bar{c} \) pair is created in a hard nucleon-nucleon collision during the passage of the colliding nuclei through each other. It will evolve to a charmonium eigenstate during the formation time which is of the order of \( t_0 \) while experiencing cold nuclear matter (CNM) effects. These are due to nuclear modification of the initial state gluon distributions and partial extinction while passing through the colliding nuclei and will be appropriately parametrized using data from \( pA \) collisions. We are interested here to describe the “anomalous suppression” effect by subsequent final state interactions with the expanding HRG formed between the receding nuclei. Since these nuclei which border the HRG on both sides in the longitudinal direction move almost with the speed of light, even \( J/\psi \)s which contribute to the particle spectra measured in the forward (backward) rapidity region can not escape the hadronic medium. Due to the different production process, the \( J/\psi \) velocities can be considered as independent from the velocity distribution of the HRG medium. This results in considerable impact velocities, in particular for forward (backward) produced \( J/\psi \)s traversing counter-propagating HRG flow.

The transverse expansion starts in the form of a rarefaction wave moving inward \( S_{\text{eff}} \). The evolution is decomposed into a longitudinal expansion inside a slice bordered by the front of the rarefaction wave and the transverse expansion. The rarefaction wave moves radially inward with the sound velocity \( c_s \). Since the temperature (and hadron gas density) decreases rapidly outside the wave, we shall ignore possible \( J/\psi \) scattering there. We denote by \( t_{\text{esc}} \) the moment of crossing the border of the rarefaction front.

Inside the region bordered by the front of the rarefaction wave the hydrodynamic evolution ceases when the freeze-out temperature is reached, here we take \( T_{f.o.} = 150 \text{ MeV} \) what is suggested by the statistical model analy-
sis of the PHENIX data \[27, 28\]. So $J/\psi$ is subject to absorption for the
time $t_{\text{final}} = \min\{\langle t_{\text{esc}} \rangle, t_{\text{f.o.}}\}$ where $\langle t_{\text{esc}} \rangle$ is the mean time until it escapes
the hadron gas region and $t_{\text{f.o.}}$ is the time when it passes the kinetic freeze-out surface for its absorption reactions. The rapidity dependence of $t_{\text{final}}$ is
crucial for the description of $J/\psi$ RHIC absorption processes.

For comparison with data one should estimate the number of participating
nucleons $N_{\text{part}}(b)$. The Woods-Saxon nuclear matter density distribution
for Au and Cu nuclei is assumed with parameters taken from \[29\] and the
standard expression for $N_{\text{part}}(b)$ is used \[30, 19\].

2. $J/\psi$ absorption

The $J/\psi$ absorption processes in CNM (nuclear absorption - NA) and HRG take place subsequently. Due to this separation in time the $J/\psi$ survival
probability for given rapidity in a heavy-ion collision with impact parameter $b$ assumes the form

$$R_{\text{AA}}(y, b) = S_{\text{NA}}(y, b) \cdot S_{\text{HRG}}(y, b),$$  \hspace{1cm} (2)
Figure 2: $J/\psi$ (dashed and dot-dashed lines) inside hadronic gas absorption region and the time evolution of the rarefaction front radius $s$ at different $z$ (continuous lines with labels denoting $z$ in fm). Numerical values correspond to Au-Au collisions.

where $S_{NA}(y, b)$ and $S_{HRG}(y, b)$ are $J/\psi$ survival probabilities in CNM and HRG, respectively. For $S_{NA}$ we employ the usual approximation

$$S_{NA} \approx \exp \left\{ -\sigma_{\psi N} \rho_0 \langle L \rangle \right\}, \quad (3)$$

where $\langle L \rangle$ is the mean absorption path length of the $J/\psi$ through the colliding nuclei obtained from the Glauber model and $\rho_0$ is the average nuclear matter density determined from the normalization of the Woods-Saxon distribution to the mass number $A$ for each nucleus separately. Within this approximate expression, $S_{NA}$ does not depend on rapidity.

The HRG survival probability $S_{HRG}(y, b)$ is defined as

$$S_{HRG}(y, b) = \frac{\int d^2p_T \int d^3r \mathcal{F}(\vec{r}, y, \vec{p}_T, t) \big|_{t=\infty}}{\int d^2p_T \int d^3r \mathcal{F}(\vec{r}, y, \vec{p}_T, t) \big|_{t=t_0}}, \quad (4)$$

where the initial time $t_0 = 1$ fm denotes the moment of the thermalization of the created hadronic matter and the beginning of the hydrodynamical expansion.

The rapidity-momentum $J/\psi$ distribution results from the kinematical change of variables $d^3p \to d^2p_T dy$ and is given by $\mathcal{F}(\vec{r}, y, \vec{p}_T, t) = M_T \cosh(y) \times f(\vec{r}, \vec{p}, t)$, where
\[ f(\vec{r}, \vec{p}, t) = f_0(\vec{r} - \vec{v}(t - t_0), \vec{p}) \times \exp \left\{ - \int_{t_0}^{t} dt' \sum_{i=1}^{l} \int \frac{d^3q}{(2\pi)^3} f_i(\vec{r} - \vec{v}(t - t'), \vec{q}, t') \sigma_{i} v_{i, \text{rel}} p_{i} q_{i, \nu} \right\} \]

is the formal solution of a kinetic equation (for details see [23, 19]). The sum is over all considered species of hadronic scatters with distributions \( f_i(\vec{r}, \vec{q}, t) \).

We have assumed that the initial distribution factorizes into
\[ F(\vec{r}, y, \vec{p}_T, t_0) = f_0(\vec{r}) \cdot g_0(\vec{p}_T) \cdot h_0(y) \]
with \( f_0(\vec{r}) \) and \( g_0(\vec{p}_T) \) normalized to unity. It is clear that \( h_0(y) \) has not to be specified since it cancels in the ratio (4).

Because of the arguments given in the Introduction we shall assume that both the created \( J/\psi \) and the hadronic medium are included in \( z = 0 \) plane. Thus the distribution of \( J/\psi \) at \( t_0 \) is \( f_0(\vec{s}) \) where \( \vec{s} \) belongs to the \( z = 0 \) plane. For \( g_0(\vec{p}_T) \) we take the form suggested by the initial state interactions [31, 32, 33, 34]

\[ g_0(\vec{p}_T) = \frac{2p_{T}}{\langle p_{T}^{2} \rangle_{J/\psi}} \cdot \exp \left\{ - \frac{p_{T}^{2}}{\langle p_{T}^{2} \rangle_{J/\psi}} \right\} \]

where \( \langle p_{T}^{2} \rangle_{J/\psi} \) is the mean squared transverse momentum of \( J/\psi \) measured in an Au-Au (Cu-Cu) collision [3, 4].

For the description of the evolution of the matter, relativistic hydrodynamics is employed. The longitudinal component of the solution of the hydrodynamic equations (the exact analytic solution for an (1+1)-dimensional case) reads [35]

\[ s(\tau) = \frac{s_0 \tau_0}{\tau}, \quad n_B(\tau) = \frac{n_B^0 \tau_0}{\tau}, \]

where \( \tau = \sqrt{t^2 - z^2} \) is the proper time and \( s_0 \) \( (n_B^0) \) is the initial density of entropy (baryon number).

For \( n_B = 0 \) and a uniform initial temperature distribution with a sharp edge at the border established by the nuclear surfaces, the full solution of the (3+1)-dimensional hydrodynamic equations is known [22].

The hadron distributions have the usual form:

\[ f_i(\vec{r}, \vec{q}, t) = \frac{2s_i + 1}{\exp \left\{ \frac{q_{i, \nu} (\vec{r}, t) - \mu_i (\vec{r}, t)}{T(\vec{r}, t)} \right\} + g_i}. \]
where $u^\nu(\vec{r},t)$ is the four-vector the hadron flow velocity and $m_i$, $B_i$, $S_i$, $\mu_i$, $s_i$ and $g_i$ are the mass, baryon number, strangeness, chemical potential, spin and a statistical factor of species $i$, respectively. We treat an antiparticle as a different species and introduce the chemical potential profile $\mu_i(\vec{r},t) = B_i\mu_B(\vec{r},t) + S_i\mu_S(\vec{r},t)$, where $\mu_B(\vec{r},t)$ and $\mu_S(\vec{r},t)$ are the local baryon number and strangeness chemical potentials, respectively.

Since for the flow described in [22] the radial velocity is zero inside the region bordered by the rarefaction wave and $v_z(\vec{s},z,t) = z/t$, the four-velocity $u^\nu(\vec{r},t)$ simplifies to

$$u^\nu(\vec{r},t) = 1/\tau(t,0,0,z).$$  \hspace{1cm} (9)

Note that a $J/\psi$ with longitudinal velocity $v_L$ is at $z = v_L(t - t_0)$ at the moment $t$ and the longitudinal flow at this point equals $v_z(\vec{s},z,t) = z/t = v_L(t - t_0)/t < v_L$, that is the relative longitudinal velocity of $J/\psi$ with respect to the flow is always non-zero. This reflects the fact that only the evolution of the medium is boost-invariant but not of the $J/\psi$. The latter is moving along a straight line with the constant velocity from the point of creation as it does not take part in the thermalization of the hadronic medium.

The temperature also does not depend on the radial coordinate $s$ within the region bordered by the rarefaction wave,

$$T(\vec{r},t) = T(z,t) = T(0,\tau) = T_0 \cdot \tau^{-a}$$  \hspace{1cm} (10)

and the power $a$ is the sound velocity squared at $T_0$. $a = c_s^2(T_0)$ [19, 20].

To obtain $T_0$ and $a$ the following procedure is applied. For a given centrality bin the initial energy density $\epsilon_0$ is estimated in [36] (Au-Au) and [37] (Cu-Cu). We put $n_S = 0$ for the strangeness number density and $n_{0 B} = 0$ for the initial baryon number density. It should be noted here that the $J/\psi$ suppression pattern is practically the same for values of $n_{0 B} \leq 0.65$ fm$^{-3}$ in this model as it was shown for the case of the centrality dependence in Ref. [19]. This range of $n_{0 B}$ corresponds to the possible values of $\mu_B$ up to $\sim 300$ MeV during the whole evolution. This is far above the estimated values of $\mu_B = 20 - 30$ MeV at the top RHIC energy [28, 38]. Now, expressing $\epsilon_0$, $n_{0 B}$ and $n_S$ as functions of $T$, $\mu_B$ and $\mu_S$ in the grand canonical ensemble framework, we can obtain $T_0$, $\mu_{0 B}$ and $\mu_{0 S}$. In our case $\mu_B = \mu_{0 B} = 0$ and $\mu_S = \mu_{0 S} = 0$ always, so the temperature is the only significant statistical parameter here.

For the most central collisions we have obtained $T_0 = 222.6$ MeV ($\epsilon_0 = 5.4$ GeV/fm$^3$) for Au-Au and $T_0 = 201.8$ MeV ($\epsilon_0 = 2.5$ GeV/fm$^3$) for Cu-Cu.
With \( T_0 \) known the initial entropy density \( s_0 \) is calculated. Having put \( s_0 \) into (7) and with the use of the grand canonical expressions for \( s \) we can obtain \( T(\tau) \). It turned out that \( T(\tau) \) has the form given by (10) with \( a \) varying in the range \( a = 0.148 - 0.156 \) for \( \epsilon_0 = 5.5 - 0.5 \text{ GeV/fm}^3 \).

For a \( J/\psi \) which is at \( \vec{s}_0 \in S_{\text{eff}} \) at the moment \( t_0 \) and has the velocity \( \vec{v} = \vec{v}_T + \vec{v}_L = \vec{p}_T/E + \vec{p}_L/E \) we denote by \( t_{\text{esc}} \) the moment of crossing the border of the rarefaction front.

The rarefaction front forms the ellipse in the \( z, s \) plane [22]

\[
(s - R_A)^2 + c_s^2 z^2 = c_s^2 t^2 \tag{11}
\]

The continuous lines in Fig. 2 labeled by different values of the longitudinal variable \( z \), show the time evolution of the rarefaction front radius \( s \) at the position \( z \).

The \( J/\psi \) trajectory is given by equations

\[
\vec{s} = \vec{d} + \vec{v}_T(t - t_0), \tag{12}
\]

\[
z = v_L(t - t_0). \tag{13}
\]

We have used here the notation \( \vec{d} = \vec{s}_0 - \vec{b} \) for the \( J/\psi \) which escapes via the left half of \( S_{\text{eff}} \) and \( \vec{d} = \vec{s}_0 \) in the opposite case.

The dashed and dash-dotted lines in Fig. 2 are \( (s, t) \) trajectories of \( J/\psi \) particles originating from the same point in the region \( S_{\text{eff}} \) shown in Fig. 1, but with different values of transverse velocity \( \vec{v}_T \). The dot-dashed line corresponds here to the lower value of \( v_T \).

The rarefaction front crossing time \( t_{\text{esc}} \) is a solution of the equation

\[
\sqrt{\vec{d}^2 + 2\vec{d} \cdot \vec{v}_T(t - t_0) + v_T^2(t - t_0)^2} = R_A - c_s \sqrt{t^2 - v_L^2(t - t_0)^2}. \tag{14}
\]

Having obtained \( t_{\text{esc}} \), we average it over the angle between \( \vec{s}_0 \) and \( \vec{v}_T \). Finally, to obtain \( \langle t_{\text{esc}} \rangle (b, v_T, v_L) \) - the average time it takes the \( J/\psi \) to leave the hadronic medium, when it is produced in a collision at impact parameter \( b \) with the velocity \( v_T = \sqrt{v_x^2 + v_y^2} \) and \( v_L \) - we average the result over \( S_{\text{eff}} \) with the weight given by the initial \( J/\psi \) distribution \( f_0(\vec{s}_0) \). This weight is equal to

\[
f_0(\vec{s}_0) = \frac{T_A(\vec{s}_0)T_B(\vec{s}_0 - \vec{b})}{T_{AB}(b)} \tag{15}
\]
where
\[ T_{AB}(b) = \int d^2s \ T_A(\vec{s})T_B(\vec{s} - \vec{b}). \]

\[ T_A(\vec{s}) = \int dz \rho_A(\vec{s}, z) \]

is the nuclear thickness function and \( \rho_A(\vec{s}, z) \) the Woods-Saxon nuclear density distribution.

Then the final expression for \( S_{HRG}(y, b) \) reads

\[ S_{HRG}(y, b) = \frac{1}{\int dp_T \ M_T \ g_0(p_T)} \int dp_T \ M_T \ g_0(p_T) \]

\[ \times \exp \left\{ - \int^{t_{\text{final}}}_{t_0} dt \ \sum_{i=1}^{l} \int \frac{d^3q}{(2\pi)^3} f_i(\vec{q}, t) \sigma_i v_{\text{rel},i} \frac{p_{\nu q_i^\nu}^E}{E_i} \right\}, \]

where \( t_{\text{final}} = \min\{\langle t_{\text{esc}}\rangle, t_{\text{f.o.}}\} \) and \( t_{\text{f.o.}} \) is the freeze-out moment resulting from the longitudinal Bjorken expansion.

Eq. (16) is in fact an approximation of the more involved exact formula. The approximation means that the average of integrals over trajectories has been replaced by the integration over the averaged trajectory, similarly as in Eq. (3). Our preliminary results of estimations of the exact formula suggest that the approximation influences only the normalization and in such a way that this can be compensated by the rescaling of the absorption cross-section to the lower values.

3. Results

In Fig. 3 an example of the behavior of \( t_{\text{final}} \) with rapidity is presented. This is for the case of the 0 – 20% centrality bin of Au-Au and Cu-Cu collisions and for \( p_T = 2 \text{ GeV} \), taken as a typical value for collisions with a \( \langle p_T^2 \rangle_{J/\psi} \sim 4 \text{ GeV}^2 \) \([3, 4]\).

In Figs. 4 (5) we show the \( J/\psi \) nuclear modification factor \( R_{AA} \) for different centralities versus rapidity in Au-Au (Cu-Cu) collisions vs. PHENIX data. There is an overall agreement with the data.

The main new effect of the PHENIX data, that \( J/\psi \) suppression is stronger at forward rapidity than at midrapidity is shown in Fig. 6 for the \( N_{\text{part}} \) dependence. It arises in the present approach as a result of the geometry of the kinetic freeze-out surface for the \( J/\psi \) which extends much more in the longitudinal direction for sufficiently long times, \( t \geq R_A - b/2 \). This entails a longer absorption path for \( J/\psi \) production at forward rapidity.
Figure 3: An example of $t_{\text{final}}$ in a function of rapidity for $p_T = 2$ GeV and the $0 - 20\%$ centrality bin for Au-Au (solid) and Cu-Cu (dashed) collisions.

Note that the present model simultaneously accounts for the anomalous $J/\psi$ suppression in the centrality dependence of the NA50/NA60 experiments at CERN, see [19, 20]. Can this result, obtained within a purely hadronic description, be interpreted as evidence against QGP formation in these experiments? Not at all, because of two reasons.

First, we have employed hadronic absorption cross sections which cannot be reconciled with a microscopic description in, e.g., nonrelativistic [39, 40] or relativistic [41, 42] quark models which have a fastly decreasing energy dependence and do not exceed a peak value of 2 mb, see also [43]. In order to justify the magnitude of hadronic absorption cross sections employed in the present work, a strong medium dependence is required which could for instance stem from spectral broadening of light mesons at the chiral restoration transition (mesonic Mott effect) [44].

Second, the hadronic resonance gas model gives a perfect description of the Lattice QCD thermodynamics in the vicinity of the critical temperature $T_c$ for the chiral and deconfinement transition [45, 20]. It has been demonstrated that the description of QCD thermodynamics within a hadronic basis can be extended even up to temperatures of 1.5 $T_c$, provided a suitable spectral broadening of these states above $T_c$ due to their Mott effect is taken into
Figure 4: $J/\psi$ nuclear modification factor versus rapidity in Au-Au collisions for constant $\sigma_{J/\psi N} = 4$ mb. PHENIX data are from [3]. Errors shown are the quadratic sum of statistical and uncorrelated systematic uncertainties.

For a recent development, see [47]. Such a Mott-Hagedorn resonance gas model has been used to describe anomalous $J/\psi$ suppression at SPS [48] and is in accordance with the picture of a strongly coupled QGP (sQGP) discovered at RHIC [49, 50].

Summarizing this discussion, the present model may be seen as a preparatory step towards a unified description of charmonium suppression kinetics in SPS and RHIC experiments within a quantum statistical approach to the sQGP to be developed. The geometrical effect on the rapidity dependence...
Figure 5: Same as Fig. 4 but for Cu-Cu collisions with PHENIX data from [4].

described in the present work shall be a part of such a description.

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Figure 6: $J/\psi$ nuclear modification factor versus centrality in Au-Au (left) and Cu-Cu (right) collisions for constant $\sigma_{J/\psi N} = 4$ mb. Lines are predictions of the model for midrapidity (solid) and forward rapidity (dashed). PHENIX data are from [3, 4]. Errors shown are the quadratic sum of statistical and uncorrelated systematic uncertainties.

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