PHYSICAL REVIEW C **82**, 024901 (2010)

J/Ψ production in nuclear collisions: Theoretical approach to measuring the transport coefficient

B. Z. Kopeliovich, I. K. Potashnikova, and Iván Schmidt
Departamento de Física, Instituto de Avanzados en Ciencias en Ingenería, Universidad Técnica Federico Santa María, Centro Científico-Tecnológico de Valparaíso, Casilla 110-V, Valparaíso, Chile
(Received 17 June 2010; published 2 August 2010)

The observed p_T dependence of nuclear effects for J/Ψ produced in heavy-ion collisions at Brookhaven's Relativistic Heavy-Ion Collider (RHIC) might look puzzling, because the nuclear suppression seems to fade at large p_T. We explain this by the interplay of three mechanisms: (i) attenuation of J/Ψ in the hot medium created in the nuclear collision; (ii) initial-state shadowing of charmed quarks and attenuation of a $\bar{c}c$ dipole propagating through the colliding nuclei; (iii) a strong Cronin effect for J/Ψ caused by saturation of gluons in the colliding nuclei. All three effects are well under control and calculated in a parameter-free way, except for the transport coefficient $\hat{q}_0$ characterizing the medium. This is adjusted to the J/Ψ data and found to be in good agreement with the pQCD prediction, but more than an order of magnitude smaller than what was extracted from jet quenching data within the energy loss scenario.

DOI: 10.1103/PhysRevC.82.024901 PACS number(s): 24.85.+p, 25.75.Bh, 25.75.Cj, 14.40.Pq

I. INTRODUCTION

The recent measurements [1–3] of J/Ψ produced in heavy-ion collisions at Brookhaven's Relativistic Heavy-Ion Collider (RHIC) have revealed unusual features of the transverse momentum distribution. Although all species of hadrons measured so far demonstrate nuclear suppression, which increases with p_T and then levels off, the nuclei-to-pp ratio for J/Ψ production, plotted in Fig. 1 rises with p_T and even has a tendency to exceed one [3].

No explanation has been proposed so far, except for an exotic one [4] assuming that part of the production rate comes from accidental coalescence of c and $\bar{c}$ pairs available in the medium. Even if this might happen, one should consider first of all the conventional explanations, based on known dynamics.

We consider here three different mechanisms affecting the production rate of J/Ψ in heavy-ion collisions: (i) final-state attenuation (FSI) of J/Ψ in the dense medium; (ii) initial-state interaction (ISI), nuclear shadowing of charm quarks and the breakup of the $\bar{c}c$ dipole propagating through the colliding nuclei; (iii) ISI Cronin effect for J/Ψ caused by gluon saturation in the colliding nuclei. All three effects certainly exist and are important, and below we present their evaluation, which is performed in a parameter-free way, except the transport coefficient [5] characterizing the hot medium. This is assumed to be unknown and is adjusted to reproduce the data with the value of $\hat{q}_0 \approx 0.3 - 0.5\,\text{GeV}^2/\text{fm}$. This value is an order or two of magnitude less than what was extracted so far from high-p_T pion suppression observed in gold-gold collisions at $\sqrt{s} = 200\,\text{GeV}$, and interpreted within the energy loss scenario [6,7]. Notice that perturbatively small values of q-hat are also found for pion suppression out to pT 20 GeV using the leading order opacity Wicks-Horowitz-Dordjevic-Gyulassy (WHDG) theory [8].

Only the first of the three effects mentioned above, the J/Ψ attenuation due to FSI, was considered in the recent publication [9], but the ISI suppression was ignored. In addition, the $\bar{c}c$ separation was assumed to be fixed during propagation through the medium, while the J/Ψ wave function is fully formed within a very short distance, half a fermi (see Sec. II A). As a result, $\hat{q}_0$ was grossly (five times) overestimated in Ref. [9].

On the contrary, in Ref. [10] it was assumed that J/Ψ is suppressed only by ISI, but propagates with no attenuation through the produced dense matter. The observed nuclear effects were explained by ISI and by the suppressed feed-down from the decays of heavier states (χ, Ψ'), which can be dissolved in the hot medium. Such an approach does not look self-consistent: if J/Ψ is absorbed even in the cold nuclear matter, it should be even more suppressed propagating through a dense medium.

II. FINAL-STATE INTERACTION OF J/Ψ

A. Time evolution of a small dipole

A $\bar{c}c$ dipole is produced at $x_F = 0$ in the center of mass (c.m.) of the collisions, with a short timescale $t_p^* \sim 1/\sqrt{4m_c^2 + p_T^2}$ and with a small transverse separation $r \sim 1/m_c$. Then, it evolves its size and forms the J/Ψ wave function. The full quantum-mechanical description of this process is based on the path integral technique [11]. However, just a rough estimate of the formation time is sufficient here, since this timescale turns out to be very short.

A small-size dipole is expanding so fast that its initial size is quickly forgotten. Indeed, the speed of expansion of a dipole correlates with its size: the smaller the dipole is, the faster it is evolving. This is controlled by the uncertainty principle, $k \sim 1/r$.

$$\frac{dr}{dt} = \frac{2k}{E_c^*} \approx \frac{4}{E_{J/\Psi}^* r}, \tag{1}$$

where $E_{J/\Psi}^* = 2E_c^*$ is the J/Ψ energy in the c.m. of the collision; $k \sim 1/r$ is the transverse momentum of the c-quark relative to the J/Ψ direction. The solution of this equation reads

$$r^2(t) = \frac{8t}{E_{J/\Psi}^*} + r_0^2, \tag{2}$$

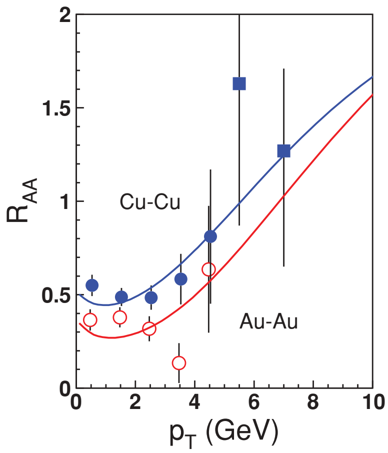

FIG. 1. (Color online) Nuclear ratio R_{AA} for central (0%–20%) copper-copper (solid circles and squares, top curve) and gold-gold (open circles, bottom curve) as a function of transverse momentum of the J/Ψ. The curves are calculated with Eq. (29) as is described in text.

where $r_0^2 \sim 1/(p_T^2 + m_c^2)$ is the initial dipole size squared, which is neglected in what follows.

According to Eq. (2), the expanding $\bar{c}c$ reaches the J/Ψ size very fast,

$$t_f^* = \frac{1}{8}\langle r_{J/\Psi}^2\rangle\sqrt{p_T^2 + m_{J/\Psi}^2} < 0.6\,\text{fm}, \tag{3}$$

for J/Ψ transverse momenta up to 5 GeV. This is about the expected time of creation of the medium.

Another estimate of the formation timescale in the c.m. of collision is [11]

$$t_f^* = \frac{2\sqrt{p_T^2 + m_{J/\Psi}^2}}{m_{\Psi^*}^2 - m_{J/\Psi}^2}, \tag{4}$$

where m_{Ψ^*} is the mass of the first radial excitation. This results in the same estimate as Eq. (3).

We conclude that what is propagating through the medium is not a small $\bar{c}c$ dipole (prehadron), but a fully formed J/Ψ.

B. J/Ψ attenuation in a dense medium

A charmonium propagates a path length L in a medium with the survival probability,

$$S(L) = \exp\left[-\int_0^L dl\sigma[r(l)]\rho(l)\right]. \tag{5}$$

Here, the path length and time are related as $l = vt$, with the J/Ψ speed $v = \sqrt{1-(2m_c/E)^2}$. The medium density is time dependent, and is assumed to dilute as $\rho(t) = \rho_0 t_0/t$ from the longitudinal expansion.

The dipole cross section for small dipoles is $\sigma(r) = Cr^2$, where r is the transverse $\bar{c}c$ separation. Correspondingly, $\sigma_{J/\Psi} = \frac{2}{3}C\langle r_{J/\Psi}^2\rangle$, and we rely on the result of the realistic model [12,13] for the mean J/Ψ radius, $\sqrt{\langle r_{J/\Psi}^2\rangle} = 0.42$ fm. The factor C for dipole-proton interactions is known from deep-inelastic scattering (DIS) data. Its value for a hot medium is unknown, as well as the medium properties. However, the factor C also controls broadening of a quark propagating through the medium [14,15],

$$\Delta p_T^2(L) = 2\frac{d\sigma(r)}{dr^2}\bigg|_{r=0}\int_0^L dl\rho(l). \tag{6}$$

Thus, the factor C is related to the transport coefficient $\hat{q}$ [5], which is in-medium broadening per unit of length,

$$C = \frac{\hat{q}}{2\rho}. \tag{7}$$

So one can represent the survival probability of J/Ψ in the medium, Eq. (5), as

$$S(L) = \exp\left[-\frac{1}{3}\langle r_{J/\Psi}^2\rangle\int_0^L dl\hat{q}(l)\right]. \tag{8}$$

The transport coefficient depends on the medium density, which is a function of impact parameter and time. We rely on the conventional form [16],

$$\hat{q}(t,\vec{b},\vec{\tau}) = \frac{\hat{q}_0 t_0}{t}\frac{n_{\text{part}}(\vec{b},\vec{\tau})}{n_{\text{part}}(0,0)}, \tag{9}$$

where $\vec{b}$ and $\vec{\tau}$ are the impact parameter of the collision and of the point where the $\hat{q}$ is defined. The transport coefficient $\hat{q}_0$ corresponds to the maximum medium density produced at impact parameter $\tau = 0$ in central collision ($b = 0$) of two nuclei, at the time $t = t_0$ after the collision. In what follows we treat the transport coefficient $\hat{q}_0$ corresponding to the medium produced in central gold-gold collision at $b = \tau = 0$, as an adjusted parameter. It is rescaled for other nuclei according to the number of participants $n_{\text{part}}(\vec{b},\vec{\tau})$ [16]. In what follows, we consider collision of identical nuclei, $A = B$, at $b = 0$.

Eventually, integrating the attenuation factor Eq. (8) over different direction of propagation of the J/Ψ produced at impact parameter $\vec{\tau}$, one gets the FSI suppression factor in the form,

$$R_{AA}^{\text{FSI}}(\vec{\tau},p_T)|_{b=0} = \int_0^\pi \frac{d\phi}{\pi}\exp\left[-\frac{1}{3}\langle r_{J/\Psi}^2\rangle\int_{l_0}^\infty dl\hat{q}(\vec{\tau}+\vec{l})\right]. \tag{10}$$

Here, $|\vec{\tau}+\vec{l}|^2 = \tau^2 + l^2 + 2\tau l\cos\phi$, and $l_0 = vt_0$. The timescale t_0 for creation and thermalization of the medium is rather uncertain, because gluons with different transverse momenta are radiated at different coherence times. We rely on the usual estimate $t_0 = 0.5$ fm.

The results are depicted by dotted curve in Fig. 2 for copper-copper collisions. We use $\hat{q}_0 = 0.45\,\text{GeV}^2/\text{fm}$, which allows one to reproduce data well, provided that other corrections, discussed in the following sections, are added.

III. INITIAL-STATE INTERACTION: SHADOWING AND ABSORPTION

A. Higher twist shadowing of charm

The same timescales, production and formation, look very different in the rest frame of one of the collision nuclei. Although a fully formed J/Ψ propagates through the hot medium, in this reference frame a $\bar{c}c$ dipole with a size

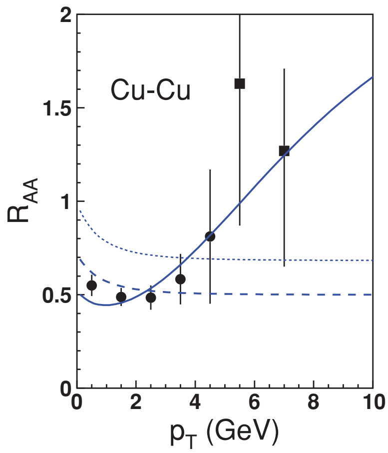

FIG. 2. (Color online) Data for the nuclear ratio R_{AA} for central copper-copper collisions as function of transverse momentum of the J/Ψ from Ref. [2] (circles) and Ref. [3] (squares). The dotted curve shows the FSI effects [Eq. (10)]. The dashed curve includes the ISI effects, charm shadowing and absorption [Eq. (19)]. The solid curve is also corrected for the Cronin effect [Eq. (28)].

"frozen" by the Lorentz time dilation propagates through the cold nuclear matter. Both the production and formation times become longer by the Lorentz factor $\gamma = 2E_\Psi/\sqrt{4m_c^2 + p_T^2}$. The coherence time of $\bar{c}c$ pair production reads,

$$t_c = \frac{E_{J/\Psi}}{(4m_c^2 + p_T^2)} = \frac{2m_c\sqrt{s}}{(4m_c^2 + p_T^2)m_N}, \qquad (11)$$

and is rather long. At $\sqrt{s} = 200$ GeV it varies from 13 to 4 fm for $0 < p_T < 5$ GeV, that is, is of the order of the nucleus size, or longer. Correspondingly, the formation time is even longer, $t_f \approx 5t_c \gg R_A$.

A full calculation of the nuclear effects for J/Ψ produced in pA collisions, including the effects of shadowing and breakup interactions of the final $\bar{c}c$, has not been done so far. Only production of the χ, the P-wave charmonium, which is a simpler case, was calculated in detail in Ref. [17]. In addition, a considerable fraction of J/Ψs are produced via decay of heavier states, Ψ', χ, etc. For our purposes it would be safer to use the experimental value of nuclear suppression observed in d-Au collisions at $\sqrt{s} = 200$ GeV [18]. Unfortunately, the experimental uncertainty is still large, so we fix $R_{dA}(x_F = 0) = 0.8$, which is about the central value.

Even if this nuclear suppression factor integrated over impact parameter is known, it is not sufficient to perform calculations for AA collisions. One has to know the b dependence of R_A. Because no relevant data are available so far, we can only rely on the theory, being constrained by the integrated value of R_A.

Because the coherence time Eq. (11) in the rest frame of one of the colliding nuclei is rather long, we assume that the $\bar{c}c$ transverse separation is "frozen" by Lorentz time delation. This grossly simplifies the calculations. In the dipole approach, the nuclear suppression factor caused by initial-state c-quark shadowing and attenuation of the $\bar{c}c$ dipole, has the form [17],

$$R_{NA}(\tau) = \frac{1}{T_A(\tau)}\left[\sum_{\lambda=\pm}\left|\int d^2r\,\Psi^*_\chi(r)(\vec{e}_\lambda\cdot\vec{r})\Psi_{g\to\bar{c}c}(r)\right|^2\right]^{-1} \times \int_{-\infty}^{\infty} dz\rho_A(\tau,z)\sum_{\lambda=\pm}\left|\int d^2r\,\Psi^*_\chi(r)(\vec{e}_\lambda\cdot\vec{r}) \times \Psi_{g\to\bar{c}c}(r)\exp\left\{-\frac{1}{2}\sigma_8(r)T_A(\tau,z) - \frac{1}{2}\sigma_{\rm dip}(r)[T_A(\tau)-T_A(\tau,z)]\right\}\right|^2. \qquad (12)$$

Here, $T_A(\tau,z) = \int_{-\infty}^{z} dz'\rho_A(\tau,z')$; $\vec{e}_\pm = (\vec{e}_x \pm \vec{e}_y)/\sqrt{2}$ is the polarization vector of the projectile gluon. The light-cone wave function of a $\bar{c}c$ in the gluon $\Psi_{g\to\bar{c}c}(r) \propto K_0(m_c r)$ [17], where K_0 is the modified Bessel function. Thus, the mean transverse size squared of a $\bar{c}c$ fluctuation of a gluon, $\langle r^2\rangle = 1/m_c^2$, is small, at least an order of magnitude smaller than that of charmonia. Apparently, in a convolution of two r distributions, narrow and wide, the mean size is controlled by the narrow one.

The cross section $\sigma_{\rm dip}(r)$ in (12) is the universal dipole-proton cross section [19], which we use in the small-r approximation $\sigma_{\rm dip}(r) \approx C(x_2)r^2$. The factor $C(x_2)$ is calculated in [20] as a function of x_2, which is the fractional light-cone momentum of the J/Ψ relative to the target.

Shadowing of the process $g \to \bar{c}c$ is controlled by the three-body g-q-$\bar{q}$ dipole cross section, which can be expressed via the conventional dipole cross sections,

$$\sigma_8(r) = \tfrac{9}{8}[\sigma_{\rm dip}(r_1) + \sigma_{\rm dip}(r_2)] - \tfrac{1}{8}\sigma_{\rm dip}(\vec{r}_1 - \vec{r}_2), \qquad (13)$$

where $\vec{r}_1$ and $\vec{r}_2$ are the transverse vectors between the gluon and the q and $\bar{q}$, respectively. We neglect the distribution of the fractional light-cone momentum of the q and $\bar{q}$, fixing it at equal shares. Then, $\vec{r}_1 = -\vec{r}_2 = \vec{r}/2$, so that

$$\sigma_8(r) = \tfrac{7}{16}\sigma_{\rm dip}(r). \qquad (14)$$

Here, we also rely on the small-r approximation.

Because we fixed the overall suppression R_{dA} at the measured value, and need to know only the impact parameter dependence, a rough estimate of Eq. (12) should be sufficient. Therefore, we approximate the result of integration over r in (12) replacing the dipole cross sections by an effective cross section, $\sigma_{\rm dip}(r) \Rightarrow \sigma_{\rm eff}$, which we can adjust to reproduce R_{dA}. Then the suppression factor Eq. (12) takes the form,

$$R_{NA}(\tau) = \frac{16}{9\sigma_{\rm eff}T_A(\tau)}\left[e^{-\frac{7}{16}\sigma_{\rm eff}T_A(\tau)} - e^{-\sigma_{\rm eff}T_A(\tau)}\right]. \qquad (15)$$

The first term in square brackets represents shadowing, the second one is related to the survival probability of the produced colorless $\bar{c}c$ dipole.

We extracted the value of $\sigma_{\rm eff}$ comparing the integrated suppression,

$$R_{NA} = \frac{1}{A}\int d^2\tau\,T_A(\tau)R_{NA}(\tau), \qquad (16)$$

with data [18] for deuteron-gold collisions at $\sqrt{s} = 200\,\text{GeV}$, $x_F = 0$. We fixed the measured ratio at $R_{dAu} = 0.8$, and found $\sigma_{\text{eff}} = 2.3\,\text{mb}$.

This value can be compared with the theoretical expectation. As was mentioned, in the convolution of the narrow distribution $\Psi_{g\to\bar{c}c}(r)$ with the large-size charmonium wave function, the latter can be fixed at $r = 0$, and the mean separation is fully controlled by the $\bar{c}c$ distribution in a gluon. Then, the mean separation squared of a produced $\bar{c}c$ pair (i.e., a fluctuation that took part in the interaction) is given by

$$\langle r^2\rangle = \frac{\int d^2r r^4 K_0^2(m_c r)}{\int d^2r r^2 K_0^2(m_c r)} = \frac{16}{m_c^2}. \quad (17)$$

Now we are in a position to evaluate the effective cross section,

$$\sigma_{\text{eff}} = C(E)\langle r^2\rangle. \quad (18)$$

The energy-dependent factor $C(E)$ is calculated in Ref. [20]. At the energy of J/Ψ $E = 300\,\text{GeV}$ ($x_F = 0$, $\sqrt{s} = 200\,\text{GeV}$) this factor varies between $C = 4.5$ in the leading order, down to $C = 3.5$ if higher order corrections are included. Correspondingly, the effective cross section Eq. (18) range is $2.5\text{mb} > \sigma_{\text{eff}} > 2\text{mb}$, which is in excellent agreement with the value extracted from the RHIC data.

We calculate the ISI suppression nucleus-nucleus collisions assuming that the suppression factors due to simultaneous propagation of the $\bar{c}c$ pair through both nuclei factorize. We ignore the possible dynamics that can breakdown this assumption [21–23], so that the ISI suppression factor for a collision of nuclei A and B with impact parameter b reads

$$R_{AB}^{\text{ISI}}(\vec{b}) = \frac{\int d^2\tau T_A(\tau) T_B(\vec{b}-\vec{\tau}) R_{NA}(\vec{\tau}) R_{NB}(\vec{b}-\vec{\tau})}{\int d^2\tau T_A(\tau) T_B(\vec{b}-\vec{\tau})}. \quad (19)$$

Thus, the initial-state interactions cause the additional suppression, Eq. (19), of J/Ψ produced in heavy-ion collision. The combined effect of ISI and FSI suppression in copper-copper central collision is shown by the dashed curve in Fig. 2. Although it agrees with the data at $p_T < 3\,\text{GeV}$, there is indication that data at higher p_T are underestimated.

B. Leading twist gluon shadowing

In addition to quark shadowing, which is a higher twist effect and scales as $1/m_c^2$, the leading twist gluon shadowing, which depends on m_c logarithmically, may be important, depending on kinematics. In terms of the Fock state decomposition, gluon shadowing is related to higher Fock components in the projectile gluon (e.g., $g \to \bar{q}qg$). Even this lowest state is heavier than just a $\bar{q}q$ and should have a shorter coherence time. In terms of Bjorken x this means that the onset of gluon shadowing is shifted toward low x_2 compared to quark shadowing. Indeed, calculations [24] show that no gluon shadowing is possible above $x_2 \approx 10^{-2}$. Moreover, it was found in Ref. [25] that the coherence length, which controls the onset of gluon shadowing, is scale independent (i.e., it the same for light and heavy quarks). This result of Ref. [25] can be understood via the energy denominator for the $g \to \bar{q}qg$ transition amplitude,

$$A(g \to \bar{q}qg) \propto \frac{1}{k^2 + \alpha_g M_{\bar{q}q}^2}, \quad (20)$$

where k and α_g are the transverse and fractional light-cone momenta of the radiated gluon, respectively. The factor α_g, which is predominantly small, suppresses the mass term in Eq. (20). In addition, the mean transverse momentum of gluons was found in Ref. [24] to be rather large, $\sqrt{\langle k^2\rangle} = 0.7\,\text{GeV}$. This is dictated by data on large mass diffraction, which is strongly suppressed compared to usual pQCD expectations. This phenomenon has been known in the Regge phenomenology as smallness of the triple Pomeron vertex. The large value of $\sqrt{\langle k^2\rangle}$, also supported by many other experimental evidences [26], leads to suppression of the gluon radiation amplitude [Eq. (20)] and weak gluon shadowing [27]. The latter is confirmed by an next-to-leading order (NLO) analysis of the DIS data [28], but contradicts the recent analysis of Ref. [29], which resulted in a very strong gluon shadowing breaking the unitarity bound [30].

The coherence length available for gluon shadowing can also be related to the value of x_2 [25],

$$l_c^g = \frac{P^g(1-x_1)}{x_2 m_N}. \quad (21)$$

Here, the denominator presents the usual Ioffe timescale, which is as long as $(0.2\,\text{fm})/x_2$, and may exceed the nuclear size at small x_2. The factor $1 - x_1$ is usually neglected, assuming that x_1 is small, which is not always the case. More important is the factor $P^g \approx 0.1$ evaluated in Ref. [25]. Its smallness is actually due to the large intrinsic transverse momentum of gluons, which we have just discussed above.

Usually, x_2 is defined as

$$x_2 = e^{-\eta}\sqrt{\frac{m_{J/\Psi}^2 + p_T^2}{s}}. \quad (22)$$

It varies with pseudorapidity and reaches a minimum at the largest measured value of η. At $\eta = 0$, with the mean value of $\langle p_T^2\rangle = 4\,\text{GeV}^2$ [18], one gets $x_2 = 0.02$, which is certainly too large for gluon shadowing. Therefore, we can safely disregard this correction in further calculations, done at $\eta = 0$.

Notice that of course x_2 decreases with η and reaches its minimal value $x_2 = 2.5 \times 10^{-3}$ at the maximal rapidity $\eta = 2$. Although this value of x_2 allows some amount of gluon shadowing, we expect a tiny correction. Indeed, within the color singlet model (CSM) [31] and its modified version [32], which provides so far the only successful description of J/Ψ production in pp collisions, the actual x_2 is considerably larger than the value given by the usual definition Eq. (22). This is because in the CSM J/Ψ is produced accompanied by a gluon, and their total invariant mass $M_{gJ/\Psi}$ is considerably larger than $m_{J/\Psi}$. With the mass distribution, $d\sigma/dM_{gJ/\Psi}^2 \propto M_{gJ/\Psi}^{-6}$, one gets the mean invariant mass squared,

$$\langle M_{gJ/\Psi}^2\rangle = 2m_{J/\Psi}^2, \quad (23)$$

which leads to a new more correct value $\tilde{x}_2 \approx 2x_2$. With the corrected minimal value $\tilde{x}_2(\eta = 2) = 0.005$ gluon shadowing correction is tiny, just a few percent [24,33].

IV. BROADENING OF GLUONS: CRONIN EFFECT

In pA collisions, projectile gluons propagating through the nucleus experience transverse momentum broadening from multiple collisions. As a result, the mean transverse momentum of produced charmonia is larger than in pp collisions. The dipole approach [14,15] is rather successful predicting broadening for heavy quarkonia [20] and heavy Drell-Yan dileptons [34], $\Delta pA = \langle p_T^2\rangle_{pA} - \langle p_T^2\rangle_{pp}$ in a parameter-free way, relying on the phenomenological cross section [24] fitted to photoproduction and DIS data. Broadening for a gluon of energy E propagating a nuclear thickness T_A reads [20]

$$\Delta_{pA}(E) = \frac{9}{16}\langle T_A\rangle\sigma_{\rm tot}^{\pi p}(E)\left[Q_{qN}^2(E) + \frac{3}{2\langle r_{ch}^2\rangle_\pi}\right], \quad (24)$$

where the proton saturation scale is

$$Q_{qN}(E) = 0.19\,\text{GeV} \times \left(\frac{E}{1\ \text{GeV}}\right)^{0.14}. \quad (25)$$

In fact, the broadening Eq. (24) is the saturation scale in the nucleus calculated in the leading order (i.e., without corrections for gluon saturation in the medium). Those corrections lead to about 20% reduction of Δ_q [20].

Remarkably, broadening does not alter the shape of the p_T distribution of produced J/Ψ. Indeed, data on pp, pA, and even AA collisions, at the energies of fixed target experiments [35] and at RHIC [18], are described well by the simple parametrization,

$$\frac{d\sigma}{dp_T^2} \propto \left(1 + \frac{p_T^2}{6\langle p_T^2\rangle}\right)^{-6}, \quad (26)$$

where $\langle p_T^2\rangle$ is the mean transverse momentum squared, which varies dependent on the process. Therefore, the simplest way to calculate the p_T dependence of the nuclear cross section would be just making a shift Δ in the mean value $\langle p_T^2\rangle$ for pA compared to pp, where Δ is broadening given by Eq. (24). Then, the nuclear ratio as a function of p_T reads

$$R_{pA}(p_T) = \frac{\langle p_T^2\rangle R_{pA}}{\langle p_T^2\rangle + \Delta_{pA}}\left(1 + \frac{p_T^2}{6\langle p_T^2\rangle}\right)^6 \times \left(1 + \frac{p_T^2}{6[\langle p_T^2\rangle + \Delta_{pA}]}\right)^{-6}, \quad (27)$$

where R_{pA} is the pA over pp ratio of p_T-integrated cross sections [Eq. (16)].

This simple procedure looks natural, although is not really proven. We can test it with the precise data from the E866 experiment at Fermilab at $\sqrt{s} = 39$ GeV. All the input parameters in Eq. (27) are known from the same measurements [36,37] and other experiments at the same energy [38–40] and also from our calculations [Eq. (24)]. $\langle p_T^2\rangle = 1.5\,\text{GeV}^2$; $\Delta = 0.08\,\text{GeV}^2 \times A^{1/3}$; $R_{pA} = A^{-0.05}$. The A dependence of the nuclear ratio calculated with Eq. (27) as a function of p_T is compared with E866 data in Fig. 3 for the exponent characterizing the A dependence, $\alpha = 1 + \ln(R_{pA})/\ln(A)$.This comparison confirms the validity of the chosen procedure, at least within the measure interval of p_T.

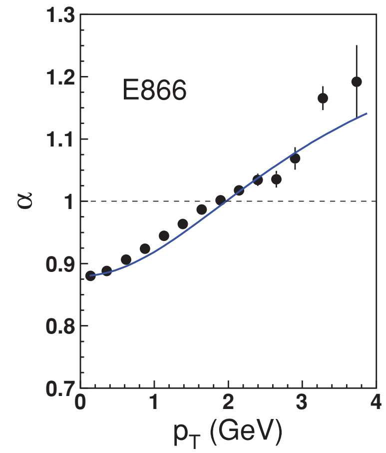

FIG. 3. (Color online) The exponent $\alpha = 1 + \ln(R_{pA})/\ln(A)$ as a function of p_T calculated with Eq. (27) in comparison with data from the E866 experiment [36].

Inspired by the good agreement with pA data, we apply the same procedure to evaluation of the Cronin effect for J/Ψ production in central AA collisions. One can use for $R_{AA}(p_T)$ the same Eq. (27), which should be modified replacing Δ and $\langle p_T^2\rangle$ by corresponding values for AA collisions at $\sqrt{s} = 200$ GeV. Also, one should replace the ratio of the p_T-integrated cross sections $R_{pA} \Rightarrow R_{AA}^{\rm ISI}(b=0)$, which is calculated with Eq. (19).

$$R_{AA}^{\rm ISI}(b=0,\tau,p_T) = \frac{\langle p_T^2\rangle R_{AA}^{\rm ISI}(b=0,\tau)}{\langle p_T^2\rangle + \Delta_{AA}(\tau)}\left(1 + \frac{p_T^2}{6\langle p_T^2\rangle}\right)^6 \times \left(1 + \frac{p_T^2}{6[\langle p_T^2\rangle + \Delta_{AA}(\tau)]}\right)^{-6}, \quad (28)$$

where the ratio of p_T-integrated cross sections, $R_{AA}^{\rm ISI}(b=0,\tau)$, is given by Eq. (19) without integration over τ.

According to Eqs. (24) and (25) broadening slowly rises with energy. However, the J/Ψ energy in the nuclear rest frame is about the same in the E866 experiment ($\langle E\rangle = 230$ GeV) and in Phenix data at $x_F = 0$ and $\sqrt{s} = 200$ GeV ($E = 330$ GeV), so we neglect the difference. Thus, in AA collisions Δ simply doubles compared to pA, and we get $\Delta = 0.64\,\text{GeV}^2$ for copper-copper and $\Delta = 0.93\,\text{GeV}^2$ for gold-gold collisions. Notice that the mean value of transverse momentum squared in pp collisions at $\sqrt{s} = 200$ GeV, $\langle p_T^2\rangle = 4\,\text{GeV}^2$, is considerably larger than at $\sqrt{s} = 39$ GeV.

Eventually, we are in a position to combine all three effects and calculate the p_T dependence of the nuclear ratio in central AA collisions,

$$R_{AA}^{J/\Psi}(b=0,p_T) = \frac{\int_0^\infty d^2\tau T_A^2(\tau) R_{AA}^{\rm FSI}(\tau,p_T) R_{AA}^{\rm ISI}(\tau,p_T)}{\int_0^\infty d^2\tau T_A^2(\tau)}. \quad (29)$$

The result is depicted by the solid curve in Fig. 2 in comparison with data for copper-copper collision. Calculations and data for gold-gold collisions are also shown in Fig. 1.

Notice that the procedure Eq. (27) has not been confronted with data above $p_T = 4$ GeV, so our extrapolation and

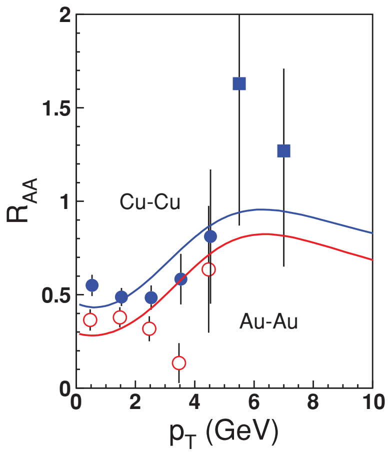

FIG. 4. (Color online) The same as in Fig. 1, but the Cronin effect is calculated differently, with Eq. (30).

predicted steep rise of the ratio at p_T, which might create problems with k_T factorization, is not well justified. For this reason we tried another way to implement broadening into the p_T distribution. An alternative way would be a simple convolution of broadening, which we take in the Gaussian form, with the p_T distribution in pp collisions. Then the nuclear modification factor $R_{AA}^{\rm ISI}(\tau, p_T)$ gets the form,

$$R_{AA}^{\rm ISI}(\tau, p_T) = \frac{R_{AA}^{\rm ISI}}{\pi\,\Delta(\tau)}\left(1 + \frac{p_T^2}{6\langle p_T^2\rangle}\right)^6 \int d^2k e^{-k^2/\Delta(\tau)} \times \left(1 + \frac{(\vec{p}_T - \vec{k})^2}{6\langle p_T^2\rangle}\right)^{-6}. \tag{30}$$

The results for copper-copper and lead-lead collisions are plotted in Fig. 4. We see that the description of data at $p_T < 5$ GeV is unchanged compared to what was depicted in Fig. 1. This means that our determination of $\hat{q}_0$ from the data is stable against the choice of the way how the p_T broadening is included. Only the behavior at larger p_T, where available data have poor accuracy, is altered, showing a weaker Cronin enhancement.

V. PROBING DENSE MATTER AT SUPER PROTON SYNCHROTRON (SPS)

The nuclear suppression caused by FSI of the J/Ψ with the produced medium was determined in the NA50 and NA60 experiments comparing the measured nuclear suppression R_{AA} with what one could expect as the cold nuclear effects in initial state extrapolating from pA data. The latest results from the NA60 experiment [41] for maximal number of participants corresponding to central collisions show that the FSI suppression factor is 0.75 ± 0.7. This experimental uncertainty is shown by the horizontal stripe in Fig. 5. The curve shows dependence of the FSI suppression factor $R_{AA}^{\rm FSI}(b=0)$, calculated with Eq. (29), on the transport coefficient. The factor $R_{AA}^{\rm ISI}$ was excluded from Eq. (29) and integration over p_T performed. Together with the experimental uncertainty, this curve provides an interval of values of $\hat{q}_0 = 0.23 \pm 0.07$

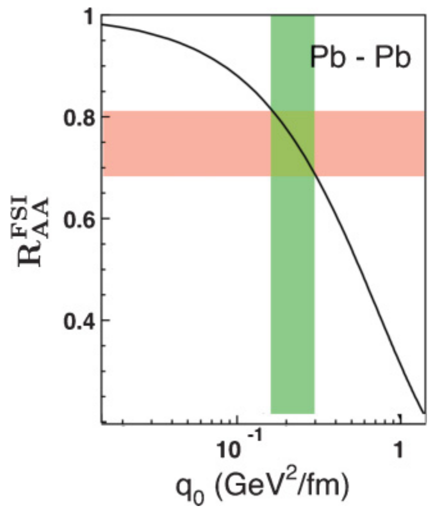

FIG. 5. (Color online) The FSI attenuation factor for J/Ψ produced in central lead-lead collisions at SPS. The curve corresponds to suppression versus $\hat{q}_0$ calculated with Eq. (29) with excluded ISI factor and integrated over p_T. The horizontal stripe shows the magnitude and uncertainty of suppression reported by the NA60 experiment [41]. The vertical stripe shows the interval of values of $\hat{q}_0$ that allow one to describe the observed suppression.

(vertical stripe), which allows one to describe the observed "anomalous" suppression. It is about twice as small as we got from RHIC data.

VI. SUMMARY AND OUTLOOK

We performed an analysis of data for p_T-dependent nuclear effects in J/Ψ production in central copper-copper and gold-gold collisions observed at RHIC. These data, looking puzzling at first glance, have not received a proper interpretation so far. We evaluated the final-state attenuation of the produced J/Ψ in the created dense medium relating it to the transport coefficient (i.e., broadening of partons propagating through the medium). The key point, which allows one to establish this relation, is the dipole description of broadening [14,15].

The observed nuclear effects in J/Ψ production in AA collisions is interpreted as a combination of FSI of the fully formed J/Ψ in the dense medium, and the ISI effects in production of J/Ψ caused by multiple interactions of the colliding nuclei. The latter includes attenuation of the produced $\bar{c}c$ dipole propagating through both nuclei, higher twist shadowing of charm quarks, and leading twist gluon shadowing. In addition, gluon saturation in nuclei leads to a considerable broadening of gluons, which causes a strong Cronin effect for J/Ψ. This explains the observed unusual rise of the nuclear ratio with p_T.

All effects are evaluated in a parameter-free way, except for the unknown properties of the produced hot medium. We employed the conventional model for the space-time development of the produced matter relating it to the number of participants. The only parameter adjusted to data, $\hat{q}_0$, is the transport coefficient corresponding to a maximal density of the matter produced in central gold-gold collision. We found that the J/Ψ data from RHIC are well reproduced with $\hat{q}_0 \approx 0.3 - 0.5\,\text{GeV}^2/\,\text{fm}$. This is close to the expected value $\hat{q}_0 = 0.5\,\text{GeV}^2/\,\text{fm}$ [5], and more than an order of magnitude less than was found from jet quenching data within the energy loss scenario [42].

We also examined the J/Ψ data from the NA60 experiments at SPS, which are available for p_T-integrated cross sections, and with already separated ISI effects. From the observed suppression in central lead-lead collisions we found $\hat{q}_0 \approx 0.23 \pm 0.07\,\text{GeV}^2/\,\text{fm}$.

We performed the calculations assuming direct J/Ψ production, but it is known that about 40% comes from the feed-down by decays of heavier charmonium states, χ and Ψ' [43]. Those states are about twice as big as the J/Ψ [13], and correspondingly should have a larger absorption cross section. Therefore, adding these channels of J/Ψ production can only reduce the value of $\hat{q}_0$ (i.e., the above values should be treated as an upper bound). The bottom bound can be estimated assuming that the r^2 approximation is valid up to the size of these excitations (which is certainly an exaggeration). Then the bottom bound for $\hat{q}_0$ is the above values times a factor 0.7.

We conclude that p_T-dependent nuclear effects for J/Ψ production in heavy-ion collisions provide a sensitive probe for the dense medium produced in these collisions. Both experimental and theoretical developments need further progress. More accurate pA (or dA) data are needed for a better control of the ISI effects in AA collisions. Data for other heavy quarkonia (Ψ', χ, Υ) would bring new precious information.

On the theoretical side, more calculations are required to describe the observed centrality and rapidity dependence of nuclear effects. The azimuthal asymmetry can also be calculated. The small-r approximation for the dipole cross section used here may not be sufficiently accurate for a very dense matter. One should rely on a more elaborated r dependence. A more realistic model for the space-time development of the dense medium, including transverse expansion, should be considered.

Notice that the observed strong suppression of open heavy flavors remains a challenge. Although charm and light quarks are suppressed similarly [8] because of a strong dead-cone effect imposed by the intensive vacuum gluon radiation [44], a strong suppression of open bottom is difficult to explain with a transport coefficient of a small magnitude.

ACKNOWLEDGMENTS

We are grateful to Dima Kharzeev for informative discussions. This work was supported in part by Fondecyt (Chile) Grant Nos. 1090236, 1090291, and 1100287, by DFG (Germany) Grant No. PI182/3-1, and by Conicyt-DFG Grant No. 084-2009.

[1] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. Lett. **98**, 232301 (2007).
[2] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. Lett. **101**, 122301 (2008).
[3] B. I. Abelev *et al.* (STAR Collaboration), Phys. Rev. C **80**, 041902 (2009).
[4] X. Zhao and R. Rapp, arXiv:0806.1239 [nucl-th].
[5] R. Baier, Yu. L. Dokshitzer, A. H. Mueller, S. Peigne, and D. Schiff, Nucl. Phys. B **484**, 265 (1997).
[6] D. d'Enterria, arXiv:0902.2011 [nucl-ex].
[7] U. A. Wiedemann, arXiv:0908.2306 [hep-ph].
[8] S. Wicks, W. Horowitz, M. Djordjevic, and M. Gyulassy, Nucl. Phys. A **784**, 426 (2007).
[9] B. Wu and B. Q. Ma, arXiv:1003.1692 [hep-ph].
[10] F. Karsch, D. Kharzeev, and H. Satz, Phys. Lett. B **637**, 75 (2006).
[11] B. Z. Kopeliovich and B. G. Zakharov, Phys. Rev. D **44**, 3466 (1991).
[12] W. Buchmüller and S. H. H. Tye, Phys. Rev. D **24**, 132 (1981).
[13] J. Hüfner, Yu. P. Ivanov, B. Z. Kopeliovich, and A. V. Tarasov, Phys. Rev. D **62**, 094022 (2000).
[14] M. B. Johnson, B. Z. Kopeliovich, and A. V. Tarasov, Phys. Rev. C **63**, 035203 (2001).
[15] J. Dolejsi, J. Hufner, and B. Z. Kopeliovich, Phys. Lett. B **312**, 235 (1993).
[16] X. F. Chen, C. Greiner, E. Wang, X. N. Wang, and Z. Xu, arXiv:1002.1165 [nucl-th].
[17] B. Kopeliovich, A. Tarasov, and J. Hufner, Nucl. Phys. A **696**, 669 (2001).
[18] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. C **77**, 024912 (2008); **79**, 059901(E) (2009).
[19] B. Z. Kopeliovich, L. I. Lapidus, and A. B. Zamolodchikov, JETP Lett. **33**, 595 (1981) [Pisma Zh. Eksp. Teor. Fiz. **33**, 612 (1981)].
[20] B. Z. Kopeliovich, I. K. Potashnikova, and I. Schmidt, Phys. Rev. C **81**, 035204 (2010).
[21] J. Hüfner and B. Z. Kopeliovich, Phys. Lett. B **445**, 223 (1998).
[22] J. Hüfner, Y. B. He, and B. Z. Kopeliovich, Eur. Phys. J. A **7**, 239 (2000).
[23] J. Hüfner, B. Z. Kopeliovich, and A. Polleri, Eur. Phys. J. A **11**, 457 (2001).
[24] B. Z. Kopeliovich, A. Schäfer, and A. V. Tarasov, Phys. Rev. D **62**, 054022 (2000).
[25] B. Z. Kopeliovich, J. Raufeisen, and A. V. Tarasov, Phys. Rev. C **62**, 035204 (2000).
[26] B. Z. Kopeliovich, I. K. Potashnikova, B. Povh, and I. Schmidt, Phys. Rev. D **76**, 094020 (2007).
[27] B. Z. Kopeliovich, J. Nemchik, A. Schäfer, and A. V. Tarasov, Phys. Rev. Lett. **88**, 232303 (2002).
[28] D. de Florian and R. Sassot, Phys. Rev. D **69**, 074028 (2004).
[29] K. J. Eskola, H. Paukkunen, and C. A. Salgado, J. High Energy Phys. 07 (2008) 102.
[30] B. Z. Kopeliovich, E. Levin, I. K. Potashnikova, and I. Schmidt, Phys. Rev. C **79**, 064906 (2009).
[31] E. L. Berger and D. Jones, Phys. Rev. D **23**, 1521 (1981).
[32] V. A. Khoze, A. D. Martin, M. G. Ryskin, and W. J. Stirling, Eur. Phys. J. C **39**, 163 (2005).

[33] B. Z. Kopeliovich, J. Raufeisen, A. V. Tarasov, and M. B. Johnson, Phys. Rev. C **67**, 014903 (2003).

[34] M. B. Johnson, B. Z. Kopeliovich, M. J. Leitch, P. L. McGaughey, J. M. Moss, I. K. Potashnikova, and I. Schmidt, Phys. Rev. C **75**, 035206 (2007).

[35] M. H. Schub *et al.* (E789 Collaboration), Phys. Rev. D **52**, 1307 (1995); **53**, 570(E) (1996).

[36] M. J. Leitch *et al.* (FNAL E866/NuSea Collaboration), Phys. Rev. Lett. **84**, 3256 (2000).

[37] M. A. Vasiliev *et al.*, Phys. Rev. Lett. **83**, 2304 (1999).

[38] D. M. Alde *et al.* (E772 Collaboration), Phys. Rev. Lett. **66**, 2285 (1991).

[39] M. J. Leitch *et al.* (E789 Collaboration), Phys. Rev. D **52**, 4251 (1995).

[40] P. L. McGaughey, J. M. Moss, and J.-Ch. Peng, NUCOLEX 99, Wako, Japan, 1999, arXiv:hep-ph/9905447; Annu. Rev. Nucl. Part. Sci. **49**, 217 (1999).

[41] E. Scomparin (NA60 Collaboration), Nucl. Phys. A **830**, 239C (2009).

[42] A. Adare *et al.* (PHENIX Collaboration), Phys. Rev. C **77**, 064907 (2008).

[43] L. Antoniazzi *et al.* (E705 Collaboration), Phys. Rev. Lett. **70**, 383 (1993).

[44] B. Z. Kopeliovich, H. J. Pirner, I. K. Potashnikova, and I. Schmidt, arXiv:1007.1913v1 [hep-ph].