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Suppression of heavy quarkonia in pA collisions

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In this proceeding we present results of our studies of quarkonium suppression in pA collisions. We estimated the nuclear suppression due to absorption, energy loss and gluon shadowing. In addition, we found that at high energies a sizeable additive contribution comes from multinucleon production. We demonstrate that the suggested approach can simultaneously explain a relatively small nuclear suppression of J/ψ and Υ, as well as a strong suppression of $\psi(2S)$ observed at RHIC and LHC in proton-ion collisions.

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1. Introduction

The dynamics of the heavy quarkonium $\bar{Q}Q$ is described by perturbative QCD (see [1, 2, 3] for a review of current situation), for this reason quarkonium propagation inside a medium serves as a clean probe of the gluon content of the target. This motivated extensive theoretical and experimental studies of quarkonium dynamics, which recently got a new impetus after the discovery of various $\bar{Q}Q$-containing exotic tetraquark and pentaquark candidates [4].

Inside nuclei, it is expected that in a heavy quark mass limit the nuclear suppression of quarkonia should vanish, though this behaviour might be modified at forward rapidities due to the energy loss corrections [5, 6, 7]. The small nuclear effects predicted in energy loss based models from the first sight are consistent with almost no suppression of J/ψ and Υ observed at RHIC. However, in such models it is challenging to explain why for $\psi(2S)$ a considerably stronger nuclear suppression is observed both at RHIC [8] and LHC [9]. This difference hints that the transverse size of the charmonia (formally a higher twist effect) should be taken into account. Naively, the suppression effects should grow with energy [10], reflecting the energy dependence of the underlying dipole-nucleon cross-section [11]. However, recent result from LHC [12] did not confirm this expectation and found approximately the same suppression as at RHIC.

We revisited the problem and found that the controversy is solved by a novel mechanism in which two nucleons from the heavy ion participate in quarkonium production. This additional contribution is enhanced by $A^{1/3}$ inside heavy nuclei, is growing as a function of energy faster than single-nucleon process, and compensates increasing with energy nuclear suppression expected for single-nucleon term. At the same time, this contribution becomes quite small for the heavy bottomonium, and for this reason allows us to construct a consistent description of nuclear suppression for all the aforementioned quarkonia.

The paper is structured as follows. In the Section 2 we describe briefly the framework used for evaluations (see [13] for mode details). In Section 3 we present our results and draw conclusions.

2. Quarkonium production in a dipole approach

Production of the quarkonium on single nucleons up to absorptive corrections coincides with production in pp collisions extensively studied in the literature [14, 15, 16]. The Color Singlet Model (CSM) gives a very reasonable description of data, provided k_T-dependent parton distributions are used [17]. The evaluations in the CSM model are performed in a momentum space, neglecting for simplicity the relative motion of the $\bar{Q}Q$ inside quarkonium in the heavy quark mass limit. For evaluation of the nuclear suppression, this approximation is not justified since nuclear effects are controlled by the transverse size of the dipole. For evaluations of nuclear effects it is more convinient to work in the coordinate representation and rewrite the k_T-factorization results in the framework of the color dipole approach [11, 18]. The cross-section of the quarkonia production in pp collisions in this approach is given by [13]

$$\frac{d\sigma_{pp}}{dy} = \frac{9}{8} g\left(x_1(y)\right) \int d\alpha_G d\alpha_Q^{(1)} d^2 r_Q^{(1)} d\alpha_Q^{(2)} d^2 r_Q^{(2)} d^2 r_G \Psi_M^*\left(\alpha_Q^{(1)}, r_Q^{(1)}\right) \Psi_M\left(\alpha_Q^{(2)}, r_Q^{(2)}\right) \times \quad (2.1)$$

$$\times \sum_{n,n'=1}^{6} \eta_n \eta_{n'} \mathrm{Tr}\left[\Lambda_M \Phi_{\bar{Q}Q}\left(\varepsilon_n, \vec{r}_n^{(1)}\right) \Phi_{QG}\left(\delta_n, \vec{r}_{G,n}^{(1)}\right)\right] \mathrm{Tr}\left[\Lambda_M \Phi_{\bar{Q}Q}\left(\varepsilon_{n'}, \vec{r}_{n'}^{(2)}\right) \Phi_{QG}\left(\delta_{n'}, \vec{r}_{G,n'}^{(2)}\right)\right]^*$$
$$\times\ \sigma\left(x_2, \vec{b}_n^{(1)} - \vec{b}_{n'}^{(2)}\right),$$

where we use notation $g(x)$ for the projectile gluon density and $x_{1,2}$ are related to rapidity y and transverse momenta p_T of quarkonia as $x_{1,2} \approx \sqrt{M_{\bar{Q}Q}^2 + p_\perp^2} e^{\pm y}/\sqrt{s}$. The wave function $\Psi_M(\alpha, \vec{r})$ of quarkonium is taken form the literature [20], $\Phi_{\bar{Q}Q}(\varepsilon_n, r_n)$ and $\Phi_{QG}(\delta_n, r_G)$ are the gluon splitting and emission wave functions (2×2 matrices in helicity space), $\sigma(x_2, r)$ is the color singlet dipole cross-section [18], and all the other notations are explained in [13].

The dynamics of small $\bar{Q}Q$ dipoles moving inside a nuclear matter depends crucially on the production time, $t_{\bar{Q}Q}$ and quarkonium formation time t_f [21]. For the high energy scattering at central and forward rapidities, the formation time t_f significantly exceeds the nuclear size, for this reason the $\bar{Q}Q$ pair traverses the nucleus with frozen transverse size, though might interact via soft gluon exchanges with neighboring nucleons. The mean number of such inelastic interactions for the dipole of transverse quark separation r is given by $n_{\rm coll}^{\bar{Q}Q}(r, b) = \sigma(x_2, r) T_A(b)$, where $T_A(b)$ is the nuclear thickness function. The parameter $n_{\rm coll}^{\bar{Q}Q}$ is quite small for RHIC kinematics, however grows with energy and becomes comparable to one in LHC kinematics. This implies that in addition to the CSM mechanism, there could be additional multinucleon contributions to the total pA cross-sections.

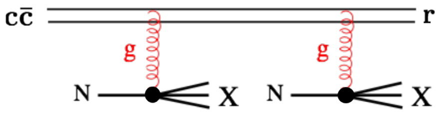

Figure 1: (Color online) Multiple color-exchange interaction of a high energy $\bar{c}c$ pair propagating through a nucleus.

For the S-wave quarkonia, the two-nucleon contribution $gNN \to J/\psi X$ shown in the Figure (1) presents a special concern, since it contributes in the same order over $\sim \mathscr{O}\left(\alpha_s(m_Q)/m_Q^2\right)$ as the single-nucleon term and requires a dedicated study. Due to difference of final states, the one- and two-nucleon contributions do not interfere and contribute to the nuclear cross-section additively. For this reason in what follows, we will refer to their contributions to suppression factor as $R_{pA}^{(1N)} = d\sigma^{(1N)}/A\, d\sigma^{(pp)}$ (single-nucleon term) and $R_{pA}^{(2N)} = d\sigma^{(2N)}/A\, d\sigma^{(pp)}$ (two-nucleon term) respectively. For the single-nucleon term the cross-section differs from (2.1) only by additional absorptive factor in the integrand,

$$\sim \int d^2b\, dz\, \rho_A(b,z) \times \tag{2.2}$$
$$\times\ \exp\left(-\frac{\sigma_4\left(x_2, \alpha_Q^{(1)}, r_Q^{(1)}, r_G\right) + \sigma_4\left(x_2, \alpha_Q^{(2)}, r_Q^{(2)}, r_G\right)}{2} T_-(b,z) - \frac{\sigma\left(x_2, r_Q^{(1)}\right) + \sigma\left(x_2, r_Q^{(2)}\right)}{2} T_+(b,z)\right),$$

where b and z are the impact parameter and longitudinal coordinate of the nucleon inside the nucleus, $\rho_A(b,z)$ is the nuclear density, and the relation of the cross-section σ_4 with color singlet dipole cross-section $\sigma(x,r)$ is explained in [13]. The cross-section of the two-nucleon mechanism is given by

$$\frac{d\sigma_{pA}^{(2N)}}{dy} = g(x_1(y)) \int d^2b \int dz_1 \int dz_2 \int d\alpha_Q^{(1)} d^2r_Q^{(1)} d\alpha_Q^{(2)} d^2r_Q^{(2)} \rho_A(b,z_1)\, \rho_A(b,z_2) \tag{2.3}$$

$$
\begin{aligned}
&\times\ \Psi^{\dagger}_{J/\psi}\left(\alpha_Q^{(1)}, r_Q^{(1)}\right)\Psi_{J/\psi}\left(\alpha_Q^{(2)}, r_Q^{(2)}\right)\Delta\Sigma_8\left(x_2, \alpha_Q^{(1)}, r_Q^{(1)}, \alpha_Q^{(2)}, r_Q^{(2)}\right)\\
&\times\ \mathrm{Tr}\left[\Lambda_M\Phi_{\bar{c}c}\left(\varepsilon_n^{(1)}, \vec{r}_n^{(1)}\right)\right]\mathrm{Tr}\left[\Lambda_M\Phi_{\bar{c}c}\left(\varepsilon_n^{(2)}, \vec{r}_n^{(2)}\right)\right]^*\Sigma^{tr}_{1^-\to 8^+}\left(x_2, \alpha_Q^{(1)}, r_Q^{(1)}, \alpha_Q^{(2)}, r_Q^{(2)}\right)\\
&\times\ \exp\left(-\frac{\sigma_3\left(x_2, \alpha_Q^{(1)}, r_Q^{(1)}\right)+\sigma_3\left(x_2, \alpha_Q^{(2)}, r_Q^{(2)}\right)}{2}T_-\left(b, z_1\right)\right.\\
&\left.-\frac{\Sigma_8\left(x_2, \alpha_Q^{(1)}, r_Q^{(1)}, \alpha_Q^{(2)}, r_Q^{(2)}\right)}{2}T_A\left(b, z_1, z_2\right)-\frac{\sigma\left(x_2, r_1\right)+\sigma\left(x_2, r_2\right)}{2}T_+\left(b, z_2\right)\right)\\
&,
\end{aligned}
$$

and the expressions for the the cross-sections σ_3, $\Delta\Sigma_8$ and $\Sigma^{tr}_{1^-\to 8^+}$, as well as description of additional corrections due to gluon shadowing and energy loss, may be found in [13].

3. Results and discussion

As was discussed in previous section, the nuclear suppression factor R_{pA} gets additive contributions from one- and two-nucleons terms $R_{pA}^{(1N,2N)}$ which we will present together with the total result. In the Figure 2 we plot the nuclear suppression factor R_{pA} for different quarkonia at RHIC and LHC energies. For J/ψ, our framework gives a reasonable description of available LHC data as well as RHIC data at forward rapidities. At central rapidities description of RHIC data in a dipole model slightly overestimates available data. This happens because typical values of x_2 in this kinematics reach $\sim 10^{-2}$, a value where an underlying eikonal approximation might be not very accurate. Although at LHC the found value of R_{pA} is close to one, we can see that this happens due to compensation of large absorption and contribution from two-nucleon term, rather than due to inherently weak nuclear effects. For $\psi(2S)$, the nuclear suppression is considerably larger, in agreement with RHIC [8] and LHC [9] data. This happens due to a node in a radial part of $\psi(2S)$ wave function and partial cancellation of contributions of small and large dipoles. For $\Upsilon(1S)$, we can see that our model agrees with available RHIC [22] data. Due to higher mass of Υ and smaller dipole sizes, all nuclear effects are smaller for this meson except energy loss corrections, which leads to sizable suppression at forward rapidities. The two-nucleon mechanism for Υ production is smaller than for J/ψ because additional gluon interactions is suppressed by smaller dipole size.

To summarize, we found that at high energies the nuclear suppression in *pA* collisions gets a sizeable contribution from the multinucleon processes. We studied in detail the largest correction of this type, which includes the two nucleons from the target, and found that it presents a substantial contribution (up to 30%) to charmonia production. This contribution explains a weak change of suppression from RHIC [23] to LHC [12], despite of significant increase of the absorptive cross-section. In the same framework we managed to describe a weak suppression of J/ψ and Υ and a large suppression of $\psi(2S)$ observed at RHIC [8] and LHC [9]. While for Υ our results may be interpreted as inherently weak nuclear effects in heavy mass limit, for J/ψ this interpretation is not accurate and happens due to partial compensation of one- and two-nucleon contributions. We expect that at even higher energies the multinucleon contributions will become more pronounced, and the cross-section will remain finite in a black disk limit, despite the fact that a contribution of single-nucleon term vanishes due to absorption.

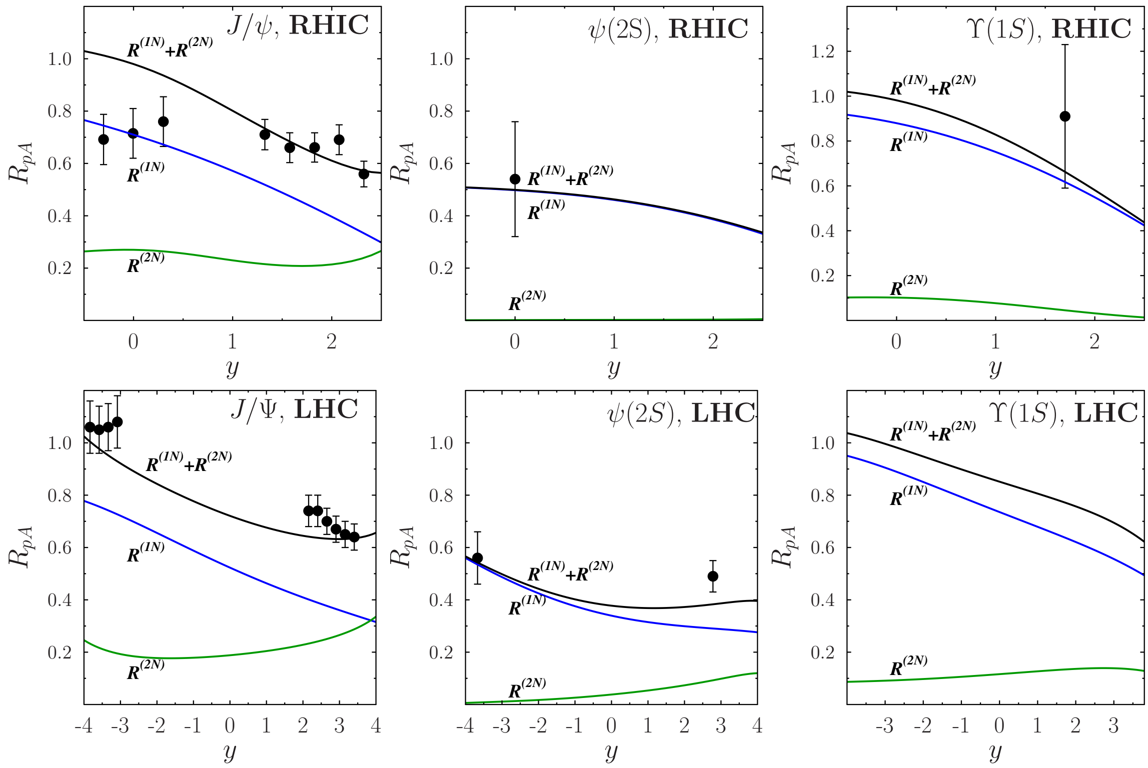

Figure 2: (Color online) Suppression of different quarkonia in heavy ion collisions for $\sqrt{s}$=200 GeV (upper row), and $\sqrt{s}$=5 TeV (lower row). Blue and green lines correspond to contribution of one- and two-nucleon terms, black line (top) is a total suppression. Data are from [12, 23, 8, 9, 22].

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