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UNRESOLVED QUESTIONS IN J/ψ PRODUCTION AND PROPAGATION IN NUCLEI

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Unresolved Questions in J/ψ Production and Propagation in Nuclei

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In order to understand the J/ψ suppression arising from the possible occurrence of the quark-gluon plasma in high-energy heavy-ion collisions, it is necessary to have a comprehensive picture how the J/ψ and its precursors are produced, what their properties after production are, and how the J/ψ and its precursors propagate inside nuclear matter. There are unresolved questions in the descriptions of J/ψ production and propagation. We outline some of these questions and discuss the approaches for their resolution.

1. Introduction

The occurrence of J/ψ suppression has been suggested as a way to probe the screening between a charm quark-antiquark pair in the quark-gluon plasma [1]. With the observation of deviation from systematics in the absorption of J/ψ in high-energy nucleus-nucleus Pb-Pb collisions [2,3], it is necessary to understand the complete process of J/ψ (and ψ') production and propagation in hadron matter in order to infer the interesting physics of their propagation in the deconfined quark-gluon plasma.

While much progress has been made in many aspects of this topic [4–15], unresolved problems remain. These problems do not exist in isolation, as the suppression of the J/ψ depends on how the produced $c\bar{c}$ object interacts with hadrons and deconfined matter. This interaction depends on the properties of the produced $c\bar{c}$, which in turn depend on the production mechanism.

Because of the limited space in this survey, we shall focus on the unresolved problems in the production and propagation of the J/ψ and its precursors in nucleon-nucleon and nucleon-nucleus collisions. J/ψ suppression in nucleus-nucleus collisions has been discussed earlier in [12–15] and will be the subject of another report.

2. The Production of $c\bar{c}$ Bound States in a Nucleon-Nucleon Collision

In a nucleon-nucleon collision, the parton of one nucleon collides with the parton of another nucleon. There is a finite probability for the collision of two gluons or a q and a $\bar{q}$ to produce a $c\bar{c}$ pair. We focus our attention on those $c\bar{c}$ pairs which have center-of-mass energies close to the bound states energies and can form various $c\bar{c}$ bound states.

The simplest theory for heavy quarkonium production is the Color-Evaporation Model [4,5]. The model is a prescription which calculates first the total cross section for $c\bar{c}$ pair production up to the $D\bar{D}$ threshold. The cross section for J/ψ production is then obtained empirically as a fraction of this total cross section. Although the Color-Evaporation Model is useful in providing a rough estimate of the quarkonium production cross section, how the colored $c\bar{c}$ system evolves into a color-singlet J/ψ state is not specified.

The Color-Singlet Model (CSM) [6] was developed to examine the production mechanism in more detail. It is assumed that the production probability amplitude is factorizable and is the product of a short-distance perturbative QCD part and a long-distance nonperturbative part. The PQCD part is given in terms of the Feynman production amplitude. The long-distance nonperturbative part is given in terms of the wave function (for J/ψ production) or its derivative (for χ production) at the origin, to take into account the interaction of the final c and $\bar{c}$ quarks. The color-singlet component of the production amplitude is projected out, to be identified as the probability amplitude for the production of the observed bound states.

As the experimental data of heavy quarkonium production accumulates, the CSM has been found to be inadequate in describing the production process (see [7–10] for reviews). First, the model predicts cross sections for high-p_T J/ψ, ψ' and χ production for $p+\bar{p}$ collisions at the Tevatron at $\sqrt{s} = 1.8$ TeV which are substantially lower than the cross sections measured by the CDF Collaboration [17]. Second, the CSM leads to a cross section ratio χ_1/χ_2 equal to 0.067 [16] for π-N collision at 300 GeV, which is too small compared to the experimental data of $\sim$0.70 [18]. Third, CSM predicts J/ψ production with a population of $J_z = \pm 1$ states substantially greater than that of the $J_z = 0$ state, leading to an anisotropic angular distribution of muons from J/ψ decay. The angular distribution is often parameterized in the form of $1 + \lambda \cos^2\theta$ in the Gottfried-Jackson frame, which is the J/ψ rest frame in which the beam goes in the z directions, while the target momentum lies in the x-z plane with momentum pointing downward ($p_x \leq 0$). The CSM gives $\lambda \sim 0.5$, but the experimental data give an isotropic distribution with $\lambda \sim 0$ for J/ψ and ψ' production in π-W collisions at 252 GeV [19] and 125 GeV [20].

The Color-Octet Model (COM) was developed to include processes which are additional to those in the color-singlet model [7]. It is assumed that besides the production of the bound state by the color-single mechanism, bound states are also produced by the color-octet formalism whereby a $c\bar{c}$ pair in a color-octet state is first formed either by gluon fragmentation or by direct parton reactions, and the octet color of the $(c\bar{c})_8$ pair is neutralized by emitting a soft gluon of low energy and momentum. The emission of the soft gluon takes place at a nonperturbative QCD time scale and is assumed to occur with a probability of unity. The cross sections for direct J/ψ and ψ' production due to the color-octet mechanism are then proportional to a lower power of α_s, leading to large color-octet contributions, in comparison with the higher-order process of color-singlet J/ψ and ψ' production by hard gluon emission. The matrix elements for the emission of the soft gluon from the color-octet $(c\bar{c})_8$ states involve nonperturbative QCD. They are left as phenomenological parameters obtained by fitting experimental data. Matrix elements have been extracted to yield good agreement with the CDF data [21].

3. Unresolved Questions in the Color-Octet Model

Although the CDF data for high p_T quarkonium production can be understood in terms of an additional color-octet mechanism, there are many unresolved questions in the Color-Octet Model. The first question is connected with J/ψ production in γp reactions at energies of 40 GeV $< \sqrt{s_{\gamma p}} <$ 140 GeV at HERA [22]. One can examine the transverse momentum p_T and the momentum fraction z of the produced J/ψ, where $z = (p \cdot p_{J/\psi})/(p \cdot p_\gamma)$ and p is the momentum of the proton. Because the photon is in a color-singlet state, the fusion of the photon with a gluon leads to a color-octet state, which according to the COM can evolve into a color-singlet state by the emission of a soft gluon of little energy and momentum. On the other hand, direct color-singlet production will involve the emission of a hard gluon through the reaction $\gamma + g \to J/\psi + g$. Because a produced hard gluon carries transverse momentum and energy while the initial colliding γ and gluon carry little transverse momentum, the kinematic region of p_T <1 GeV and another region with $z \sim 1$ are regions where the color-octet production process dominates. The color-singlet production process dominates in the other regions with $p_T > 1$ GeV and $z < 1$. However, using the color-octet matrix element as determined from the CDF measurements and assuming $\langle \mathcal{O}_8^{J/\psi}(^1S_0)\rangle = \langle \mathcal{O}_8^{J/\psi}(^3P_0)\rangle/m_c^2$, one obtains theoretical cross sections which are much larger than the experimental data for $p_T < 1$ GeV and for $z \sim 1$ [22].

The second question is connected with the color-octet matrix elements extracted from J/ψ, ψ' and χ production cross sections at fixed-target energies which give [9]

$$\langle \mathcal{O}_8^{J/\psi}(^1S_0)\rangle + \frac{7}{m_c^2}\langle \mathcal{O}_8^{J/\psi}(^3P_0)\rangle = 3.0 \times 10^{-2}\ \text{GeV}^3 \quad (1)$$

for J/ψ production, and for ψ' production

$$\langle \mathcal{O}_8^{\psi'}(^1S_0)\rangle + \frac{7}{m_c^2}\langle \mathcal{O}_8^{\psi'}(^3P_0)\rangle = 0.5 \times 10^{-2}\ \text{GeV}^3. \quad (2)$$

These need to be compared with those from high p_T CDF measurements at 1.8 TeV [21]

$$\langle \mathcal{O}_8^{J/\psi}(^1S_0)\rangle + \frac{3}{m_c^2}\langle \mathcal{O}_8^{J/\psi}(^3P_0)\rangle = 6.6 \times 10^{-2}\ \text{GeV}^3 \quad (3)$$

for J/ψ production, and for ψ' production

$$\langle \mathcal{O}_8^{\psi'}(^1S_0)\rangle + \frac{3}{m_c^2}\langle \mathcal{O}_8^{\psi'}(^3P_0)\rangle = 1.8 \times 10^{-2}\ \text{GeV}^3. \quad (4)$$

As mentioned by Beneke *et al.* [9], the fixed target energy values are a factor of 4 (7) smaller than the Tevatron values for $J/\psi(\psi')$ if $\langle \mathcal{O}_8^{\{\psi(\psi')\}}(^1S_0)\rangle = \langle \mathcal{O}_8^{\{\psi(\psi')\}}(^3P_0)\rangle/m_c^2$. The discrepancy would be lower to a factor of 2 if $\langle \mathcal{O}_8^{\{\psi(\psi')\}}(^3P_0)\rangle = 0$.

The third question which is left from the Color-Singlet Model concerns χ_1 production relative to χ_2 production. As we mentioned earlier, the color-singlet model prediction of the χ_1 cross section is much smaller than the experimental data [16] when we normalize the theoretical yield to χ_2 states. The Color-Octet Model does not improve this discrepancy on the χ_1 yield. This arises because of Yang's theorem that two on-shell gluons cannot fuse to form a spin-one state. Hence, χ_1 can be formed only by the $gg \to \chi_1 g$ process which is α_s^3 order in cross section. On the other hand, χ_2 in the color-singlet state can be formed by the fusion of two gluons with a cross section that is of order α_s^2. Thus, the predicted $\mathcal{B}_1\chi_1/\mathcal{B}_2\chi_2 = 0.13$ [9] is much smaller than the observed $\mathcal{B}_1\chi_1/\mathcal{B}_2\chi_2 \sim 1.4$ for π-N collision at 185 GeV and 300 GeV [18].

Finally, previous questions on the polarization of the produced J/ψ and ψ' for the Color-Singlet Model remains a puzzle for the Color-Octet Model. In the Color-Octet Model, the dominant production arises from the $c\bar{c}$ color-octet pair in the $^1S_0$ and the ^3P_J states whose matrix elements are in the combination of Eqs. (1) and (2). While the $^1S_0 \to {}^3S_1$ transition leads to an isotropic muon distribution, the $^3P \to {}^3S_1$ preferentially populates $J_z = \pm 1$ substates with large transverse polarization. The velocity counting rule gives $\langle \mathcal{O}_8^{\psi'}(^1S_0)\rangle$ to be of the same order as $\langle \mathcal{O}_8^{J/\psi}(^3P_0)\rangle/m_c^2$. If so, the expected angular distribution is not isotropic, with $0.31 < \lambda < 0.62$ for J/ψ, and $0.15 < \lambda < 0.4$ for ψ' [9]. The experimental polarization suggests a nearly isotropical distribution in its center-of-mass frame as would be the case for equal population of the three J_z substates.

4. Resolution of Some of the Questions in J/ψ Production

We can discuss possible approaches which may resolve some of these questions. First, we examine the question of J/ψ polarization. The polarization data can be explained as arising from the fact that in the color-octet mechanism, the soft gluon is emitted predominantly from a ${}^1S_0$ state to reach the final state J/ψ state of ${}^3S_1$ – much more prominent than from the emission from the ${}^3P_J$ state. This may appear surprising as one expects the opposite relation from the velocity counting rule, where from the comparison of the spatial transition currents, the absolute value of the ratio of the M1 amplitude to that of the E1 amplitude is of the order of v [23]. However, the transition matrix elements consist of a part due to the spatial current and a part due to the spin current. They have different dependences on the energy of the radiation. The transition from ${}^3P_0$ indeed dominates over the ${}^3S_1$ transition for high energy radiative transitions, but the situation is just reversed for soft radiative transitions when the transition due to the spin current is allowed, as was clearly demonstrated by the analogous situation in the radiative production of deuterons by the interaction of a neutron with a proton [23].

The deuteron bound state is a ${}^3S_1$ state, with the same quantum number as J/ψ and ψ'. A low energy neutron and a proton can form ${}^1S_0$ and ${}^3P_J$ states in the continuum. The cross section for a radiative transition from the ${}^1S_0$ continuum state to the bound ${}^3S_1$ deuteron state with the emission of a photon is [Eq. (XII.4.27) of [23]]

$$\sigma^{M1}({}^1S_0 \to {}^3S_1) = \pi \frac{e^2}{\hbar c}\left(\frac{\hbar}{M_N c}\right)^2 \frac{B}{Mc^2}(\mu_n - \mu_p)^2(1-\gamma a_s)\left(\frac{2B}{E_N}\right)^{1/2}, \tag{5}$$

where $B \equiv \hbar^2\gamma^2/M_N$ is the binding energy of the deuteron, $E_N \equiv 2\hbar^2k^2/M_N = 2(E_{\text{photon}} - B)$ is the asymptotic neutron kinetic energy relative to the proton, μ_n and μ_p are the magnetic moments of the neutron and proton respectively, and a_s is the scattering length between the neutron and proton. The M1 cross section has the typical $1/v = 1/\sqrt{2E_N/M_N}$ behavior, and is large for very soft radiative transitions. The magnetic dipole transition arises not from the spatial current but from the spin current which allows a spin-flip in the ${}^1S_0 \to {}^3S_1$ transition.

On the other hand, the E1 radiative transition from a ${}^3P_J$ continuum state to the bound ${}^3S_1$ deuteron state with the emission of a photon comes from the spatial transition current. The radiative transition cross section is [Eq. (XII.4.38) and (XII.4.14) of [23]]

$$\sigma^{E1}({}^3P_J \to {}^3S_1) = \frac{8\pi}{3}\frac{e^2}{\hbar c}\frac{k\gamma}{(k^2+\gamma^2)^2}(1-\gamma r_{ot})^{-1}, \tag{6}$$

where r_{ot} is the effective range of the interaction. From these results, we note that for the softest radiation near the bound state formation threshold, the magnetic dipole M1 cross section σ^{M1} varies as $1/v$ and is large, while the electric dipole $E1$ cross section σ^{E1} vanishes as $M_N v/2 = \hbar k \to 0$. Thus, radiative formation of ${}^3S_1$ by the emission of very soft radiation is dominated by the magnetic dipole transition ${}^1S_0 \to {}^3S_1$ over the electric dipole transition ${}^3P_J \to {}^3S_1$, and the velocity counting rule breaks down [23]. Experimentally, the dominance of the ${}^1S_0 \to {}^3S_1$ transition over the ${}^1P_J \to {}^3S_1$ transition in the soft photon region has been demonstrated in nuclear reactions [23].

We can examine the color-octet matrix element by making the analogy of photon radiation with gluon radiation. For c and $\bar{c}$ in color-octet states, the interaction is repulsive, and a_s is negative. It is clear from the deuteron analysis that for $J/\psi(^3S_1)$ production by very soft gluon radiation from a color-octet state, the $^1S_0 \to {}^3S_1$ transition dominates over the $^1P_J \to {}^3S_1$ transition. With the dominance of the $^1S_0 \to {}^3S_1$, the resultant population of the $^3S_1$ states is unpolarized, with isotropic distribution of muons from its decay. This explains the observed isotropic muon angular distribution.

With the dominance of $\langle\mathcal{O}_8^{J/\psi}(^1S_0)\rangle$ over $\langle\mathcal{O}_8^{J/\psi}(^3P_J)\rangle/m_c^2$, the discrepancy of the matrix elements in the Tevatron measurement and the fixed target measurements is reduced. The discrepancy can be further reduced when one takes into account the physical masses and allowing for the finite energy carried by the soft gluon. The fact that $\langle\mathcal{O}_8^{J/\psi}(^1S_0)\rangle >> \langle\mathcal{O}_8^{J/\psi}(^3P_J)\rangle/m_c^2$ also helps to reduce the discrepancy in HERA photoproduction data. The perceived discrepancy of the data arises from assuming that $\langle\mathcal{O}_8^{J/\psi}(^1S_0)\rangle = \langle\mathcal{O}_8^{J/\psi}(^3P_J)\rangle/m_c^2 = 10^{-2}$ GeV^3. If $\langle\mathcal{O}_8^{J/\psi}(^3P_J)\rangle$ is set to zero for soft gluon radiation, then the Tevatron data implies $\langle\mathcal{O}_8^{J/\psi}(^1S_0)\rangle = 3 \times 10^{-2}$ GeV^2, and the HERA data of $\gamma p \to J/\psi X$ with $p_T < 1$ GeV can be approximately consistent with theoretical predictions. The discrepancy at $z \sim 1$ remains and may suggest additional effects at $z \sim 1$ which may be beyond the scope of the COM.

Finally, the question of the ratio χ_1/χ_2 is not yet resolved. It has been suggested that such a discrepancy may be due to higher twist terms [9,24]. One can think of another intriguing possibility which may lead to an enhanced production of χ_1. Experimentally, one knows that feeding from higher states can be the source of observed bound state populations. For example, about 32.5% of the observed J/ψ comes from the feeding from the electromagnetic decay of the χ states, most notably the χ_1 and the χ_2 states, and about 7.5% from the decay of the ψ' state [18]. From the branching ratio of ψ' and the observed yield of ψ', one can also infer that about 56% of the observed ψ' comes from the feeding from higher states, perhaps from 2P states analogous to the feeding of J/ψ from 1P states. It is therefore of interest to examine whether the χ_1 states come from the feeding from higher D states. D states form the multiplet $^1D_2$, and $^3D_{\{1,2,3\}}$. Experimentally, the J=1 $^3D_1(1^{--})$ state has been observed and lies at 3.77 GeV, just above the $D\bar{D}$ threshold of 3.74 GeV. It decays predominantly by $D\bar{D}$ breakup. The $^{\{1,3\}}D_2$ and $^3D_3$ states have not been observed, but are expected to lie close to the $^3D_1(1^{--})$ state. They are likely to lie below the $\pi D\bar{D}$ threshold. While the J=3 $^3D_3(3^{--})$ state can decay by $D\bar{D}$ breakup, the J=2 $^1D_2(2^{-+})$ and $^3D_2(2^{--})$ states cannot decay by $D\bar{D}$ breakup because they are unnatural parity states with J^P=2^-. The $^3D_2(2^{--})$ state can predominantly decay into the χ_1 and χ_2 states by E1 electromagnetic transitions [25]. Such a state can feed into the population of χ_1 and χ_2. On the other hand, $^3D_2(2^{--})$ can be formed by the fusion of two gluons into a color-singlet state with a cross section of order α_s^2. The production of $^3D_2(2^{--})$ may result in an enhanced production of χ_1 and χ_2, and may explain the discrepancy with regard to the χ_1 yield in the COM.

One concludes from the above discussions that while there are questions concerning the Color-Octet Model, there may be resolution of these questions in terms of a careful refinement on the details of the model.

5. Propagation of J/ψ Precursors in Nuclear Matter

The COM gives the probability for the formation of various bound states, but gives no specific information on the type of admixture of the precursor state. However, the nature of the state vector of the precursor has great influence in its subsequent interaction with target nucleons when the precursor propagates through nuclear matter in nucleon-nucleus collisions. Therefore, nucleon-nucleus collisions provide an arena to study the nature of the state vector of the precursor. The precursor state has many degrees of freedom: color C, angular momentum and spin JLS. The precursor state can, in general, be incoherent in one degree of freedom but a coherent admixture in another degree of freedom. Theoretical and experimental investigations of the nature of the admixture in the precursor state are interesting and unresolved problems in J/ψ production.

Dynamical processes in J/ψ production are controlled by the $c\bar{c}$ pair production time $\tau_{\rm pair}$ and the evolution time $\tau_{\rm evol}$ for the produced $c\bar{c}$ pair to evolve into a bound state. The pair production is a short-distance perturbative QCD process; its time $\tau_{\rm pair}$ is about $1/2m_c$=0.07 fm/c. The evolution time is a nonperturbative QCD process, and its time $\tau_{\rm evol}$ is about $1/\Lambda_{QCD}$=0.5 fm/c measured in the J/ψ rest frame. In J/ψ production in pA collisions, the dynamics is further controlled by the next-nucleon meeting time $\tau_{\rm nn}$, the time it takes for the J/ψ precursor to meet the next target nucleon after its production. In the J/ψ rest frame, the next-nucleon meeting time $\tau_{\rm nn}$ is $d/(\gamma^2-1)^{1/2}$, where $d=2$ fm is the internucleon spacing in a nucleus at rest and $\gamma(x_F)$ is the relativistic energy/mass ratio of the moving target nucleons [15]. For $x_F>0$ at fixed-target energies of several hundred GeV, we have $\tau_{\rm nn} << \tau_{\rm evol}$. Therefore, for p-A collisions in fixed-target experiments, many of the collisions between the produced precursor and target nucleons take place before the $c\bar{c}$ pair has completed its evolution to bound color-singlet states. One can use nuclear matter in a pA collision as an arena to probe the nature of the $CJLS$ admixture of the produced precursor.

6. Color-Dependence of the Absorption Cross Section

To study the absorption of J/ψ in pA collisions, we can rely conceptually on the fact that the collision between the precursor and the target nucleon occurs at high energies and high-energy hadron-hadron cross sections are dominated by Pomeron exchange. In the Two-Gluon Model of the Pomeron (TGMP) [26–28], the color-singlet (C1) total $(c\bar{c})_1$-nucleon cross σ_1 can be expressed as T_1-T_2, where T_n is the contribution in which the two exchanged gluons interact with n quarks in the projectile. The color-singlet total $(c\bar{c})_1$-N cross sections are size-dependent. The cross section σ_1 vanishes if one of the colliding hadrons shrinks to a point, because in this limit $T_2=T_1$. A produced coherent color-singlet wave packet with a small separation between c and $\bar{c}$ will lead to small J/ψ absorption, while a large separation in the wave packet will result in a large absorption.

For $(c\bar{c})_8$-N scattering, the total color-octet (C8) cross section σ_8 is very different, as pointed out by Dolejší and Hüfner [27]. This is because the one- and two-quark contributions now add together in the form of $T_1+T_2/8$. The result is then insensitive to the c-$\bar{c}$ separation of the color-octet precursors. It is also very large, typically of the order of 30-60 mb when a perturbative propagator is used for gluons with a nonzero effective mass. Because of the insensitivity of the cross section on the color separation, the absorption cross section will be insensitive to the spatial admixture of the color-octet precursor.

For a precursor with a coherent admixture of color-singlet and color-octet states in the form (a_1 C1 + a_2 C8) the total cross section between a precursor with nucleons in the TGMP is approximately $|a_1|^2\sigma_1+|a_2|^2\sigma_8$ where $\sigma_{\{1,8\}}$ are the total cross sections evaluated with a pure color-singlet or octet state [14]. Thus, the cross section for an incoherent admixture lies in between the two limits.

7. Incoherent Admixture

The simplest description of the precursor is the model of incoherent admixture of color and spatial quantum numbers. We must now generalize this standard absorption model for precursors of one type [29] to several types of precursors [15]. The survival probability is then the weighted sum of survival probabilities of different absorption components characterized by different absorption cross sections σ_i

$$R(BA/NN,x_F) \equiv \frac{d\sigma^{AB}_{J/\psi}/dx_F}{Ad\sigma^{NN}_{J/\psi}/dx_F} = \sum_i f_i(x_F)R_i(AB)\,, \tag{7}$$

where the production fractions are normalized to $\sum_i f_i = 1$, $i = \{CJLS\}$, and

$$R_i(BA) = \int \frac{d\mathbf{b}_A}{A\sigma_{\mathrm{abs}\,i}} \frac{d\mathbf{b}_B}{B\sigma_{\mathrm{abs}\,i}} \left\{1 - \left(1 - T_B(\mathbf{b}_B)\sigma_{\mathrm{abs}\,i}\right)^B\right\}\left\{1 - \left(1 - T_A(\mathbf{b}_A)\sigma_{\mathrm{abs}\,i}\right)^A\right\}. \tag{8}$$

Each survival probability is approximately an exponential function of the average path length L. In the semilog plot of the logarithm of the total survival probability as a function of the average path length, the slope of the curve will change as a function of the average path length. The slope will be proportional to the largest absorption cross section for small L, and to the smallest absorption cross section for large L.

Ref. [15] considers the incoherent model in which the C1 and C8 cross sections can be considered to be constants, independent of the angular momentum JLS and the number of nodes n. Then there are only two components in Eq. (7). An absorption model can be constructed that respects the popular theoretical prejudices that C1 precursors are produced in point-like states and tend to be transparent with σ_1=0 in the colliding nuclear complex, where C8 precursors are strongly absorbed, with $\sigma_8 \sim$15 mb. The C8 fractions that come out of this model are quite substantial, in agreement with independent analyses of hadron production rates in free space [9,24]. These results are supported by later investigations [32]. However, the phenomenological analyses seem to show that the available data alone are not sufficiently discriminating to tell us if the C1 precursors are transparent because they are produced point-like, or if they are also absorbed because they are produced at almost full size. Better fits to these data are obtained by using an older picture in which color-singlet precursors are also absorbed by nuclei [15].

8. Coherent Admixture

Experimental quarkonium production data have been obtained only for a very limited number of nuclear mass numbers, which makes it difficult to determine with certainty whether the survival probability as a function of average path lengths has one or many absorption components. Alternatively fits by a single absorption component can be made when one makes allowance for the energy dependence of the production cross section [2,3]. Such a procedure has a high degree of uncertainty because of the uncertainty in matching data from different experimental conditions and energies. It is an unresolved problem as to whether the survival probability can be described by a single exponential component or the sum of many components.

A description of the precursor as a coherent admixture has been given previously [14]. In this description, the interaction of partons a and b in a nucleon-nucleon collision form a precursor state $|\Phi_{ab}\rangle$ with an average energy and a width of energy. As determined by the Feynman production amplitude, the precursor state is a coherent linear combination of states of various $CJLS$ states. The probability amplitude for the production of the bound $CnJLS$ state is then obtained as the projection of this precursor state to this bound state or to the combination of this bound state with a soft gluon, after taking into account the evolution of the precursor state.

We can express this mathematically as follows. The initial precursor state Φ_{ab} of the Q-$\bar{Q}$ pair from the collision at $t_i = 0$ is represented by the state vector

$$|\Phi_{ab}(t_i)>=\mathcal{M}(ab \to Q(P/2+q)\bar{Q}(P/2-q))|Q(P/2+q)\bar{Q}(P/2-q)> \tag{9}$$

where $\mathcal{M}(ab \to Q(P/2+q)\bar{Q}(P/2-q))$ is the Feynman amplitude for the $a+b \to Q+\bar{Q}$ process, and P and q are the center-of-mass and relative momentum of Q and $\bar{Q}$. One can perform a decomposition in terms of color and angular momentum states as

$$|\Phi_{ab}(t_i)>=\sum_{CJLS} \tilde{\phi}^C_{JLS}(q)|Q\bar{Q}[^SL^C_J](P)> . \tag{10}$$

A bound $Q\bar{Q}$ state with quantum numbers JLS can be written as

$$|\Psi_{nJLS};Pq>=\sqrt{\frac{2M_{nJLS}}{4m_Q m_{\bar{Q}}}}\tilde{R}_{nJLS}(q)|Q\bar{Q}[^SL^{(1)}_J](P)> . \tag{11}$$

In lowest-order perturbative QCD, the probability amplitude for the direct production of Ψ_{nJLS} is obtained by projecting $\Phi_{ab}(t_i)$ onto Ψ_{nJLS}. The projection is simplest in the Q-$\bar{Q}$ center-of-mass system where $P = (M_{nJLS}, \mathbf{0})$ and $q = (0, \mathbf{q})$, and the probability amplitude is [30,31]

$$<\Psi_{nJLS};Pq|\Phi_{ab}(t_i)>=\sqrt{\frac{2M_{nJLS}}{4m_Q m_{\bar{Q}}}}\int \frac{d^3\mathbf{q}}{(2\pi)^3}\tilde{R}_{nJLS}(\mathbf{q})\tilde{\phi}^{(1)}_{JLS}(\mathbf{q}). \tag{12}$$

Because the bound state Ψ_{nJLS} is a color-singlet state, the above projection will involve only color-singlet components of the admixture in Eq. (10).

In the next-order perturbation theory, the color-singlet bound state Ψ_{nJLS}, accompanied by a soft gluon g_s, can be produced by the color-octet component of Φ_{ab} in Eq. (10). The probability amplitude for the production of the bound state Ψ_{nJLS} accompanied by a soft gluon g_s at the hadronization time t is

$$<[\Psi_{nJLS};Pq]g_s|\Phi_{ab}(t)>=<[\Psi_{nJLS};Pq]g_s|U(t,t_i)|\Phi_{ab}(t_i)>, \tag{13}$$

where $U(t,t_i)$ is the evolution operator.

The state Ψ_{nJLS} can also be produced indirectly through the production of different bound states $\Psi_{n'J'L'S'}$ which subsequently decay into Ψ_{nJLS}. The production probability, including direct, indirect, and color-octet contributions, is then the sum of the absolute squares of various amplitudes. Heavy quarkonia can be produced by different parton combinations such as g-g, q-$\bar{q}$, and g-q collisions, which will lead to different precursor states. The total production probability will be the sum from all precursor states.

9. Propagation of a Coherent Precursor in Nuclear Matter

A coherent precursor propagates through the nuclear medium as a single object with a single absorption cross section. The time of evolution t can be represented equivalently by the corresponding path length L/v, where v is the velocity of the precursor in the medium. The state vector after propagating a distance L in the nuclear medium is related to the state vector after production by

$$|\Phi_{ab}(L)\rangle = e^{-\rho\sigma_{abs}(ab)L/2}|\Phi_{ab}(L=0)\rangle\,, \tag{14}$$

where $\sigma_{abs}(ab)$ is the precursor absorption cross section for the collision of the precursor Φ_{ab} with a nucleon, and ρ is the nuclear matter number density.

When we include precursors from different parton collisions leading to the production of the bound state Ψ_{nJLS}, there will be different factors $e^{-\rho\sigma_{abs}(ab)L/2}$ for different parton combinations a-b. At fixed-target energies, where the total yield of J/ψ in the forward direction is dominanted by contributions from gg collisions, there is essentially only a single survival factor for the total yield in forward directions. The ratio of the production of various bound states in a pA collision to a pp collision will be independent of the mass number of the nucleus, as the production of all different bound states comes from the projection of the precursor state onto the bound states after the absorption. As a consequence of Eq. (14) we have

$$\frac{\sigma(pA\to\psi'\,X)}{\sigma(pp\to\psi'\,X)} = \frac{\sigma(pA\to\psi\,X)}{\sigma(pp\to\psi\,X)} = \frac{\sigma(pA\to\chi\,X)}{\sigma(pp\to\chi\,X)}. \tag{15}$$

Experimentally, the ratio $\psi'/(J/\psi)$ is observed to be approximately a constant of the atomic numbers [33], in agreement with the present picture. The present picture predicts further that the ratio of the χ yield to the J/ψ yield will also be independent of the mass number in pA collisions. It will be of interest to test such a prediction in the future.

Recently, C. W. Wong [34] pointed out that because of the nature of the two-gluon coupling model of the Pomeron, channel coupling between the color-octet and singlet states are weak, and the propagation of an admixture of color states may nonetheless show up as two different absorption components even for a coherent admixture of color states. It is important to have accurate measurements of the absorption curve to separate out the different rates of absorption of the different components.

We conclude this section by remarking that the available data points are sparse and have large uncertainties. It is difficult to separate out the different components of the absorption process. It will be of great interest to perform accurate measurements of J/ψ production in pA collisions for a large set of nuclei so as to infer from the survival probability the number of incoherent components and their different absorption cross sections in order to obtain a better description of the J/ψ precursor.

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