Quantum effects and color transparency in charmonium photoproduction on nuclei

Boris Z. Kopeliovich
Joint Institute for Nuclear Research, Laboratory of Nuclear Problems, Head Post Office, P.O. Box 79, 101000 Moscow, U.S.S.R.

Bronislav G. Zakharov
Moscow State University, Nuclear Physics Institute, High Energy Physics Laboratory, Moscow 177234, U.S.S.R.
(Received 1 March 1991)

We develop a rigorous quantum-mechanical treatment of color transparency effects in diffractive photoproduction of $\bar{c}c$ pairs on nuclei. The evolution of the $\bar{c}c$ wave function during propagation through a nucleus is more a considerable distortion of its form than a trivial attenuation. One of the manifestations of the quantum effects is a nuclear antishadowing of the ψ' yield, i.e., transparency above one. However, considerable nuclear shadowing is predicted for the photoproduction of J/ψ, which has a much smaller absorption cross section than ψ'.

I. INTRODUCTION. SPACE-TIME PATTERN OF CHARMONIUM PHOTOPRODUCTION

A salient prediction of QCD is a close connection between the interaction cross section of a hadron and its transverse dimension: the more compact the hadron, the more weakly it interacts [1–3]. As a result a nucleus should be transparent for a high-energy hadron participating in a process where only compact fluctuations of the hadron can survive [4,5]. It means that nuclear transparency, defined as

$$T_N=\frac{\sigma_A}{A\sigma_N}\ , \qquad (1)$$

where σ_A and σ_N are the cross sections of the process on a nuclear and nucleon targets, should be close to unity. Indeed, analyses [6,7] of experimental data from Serpukhov on quasifree charge-exchange scattering $\pi^- p\rightarrow\pi^0 n$ on bound protons at 40 GeV demonstrate a steep easing of the pion attenuation in a nucleus depending on momentum transfer. It is a clear signal of color transparency. On the contrary, measurements of quasielastic *pp* scattering at 90° in the c.m. frame, performed at BNL [8] at energies up to 13 GeV, showed an unexpected fall of the nuclear transparency. Any of the existing explanations [9,10] consider a considerable admixture of nonpointlike hadron configurations. This uncertainty obscures the situation and makes one look around for other hard processes, in which more definite information about the hadron wave function at the moment of interaction is available. The diffractive photoproduction of charmonium, considered below, is one example.

The influence of color transparency on J/ψ photoproduction on nuclei was considered on a qualitative level in Ref. [11]. The authors divided conventionally the process of J/ψ production into two stages, shown in Fig. 1. The first one is the creation of a compact $\bar{c}c$ pair, localized in a small volume with a dimension of about $1/m_c$. Because of the uncertainty principle, this stage takes in the laboratory frame a time

$$\tau_p\approx\frac{E}{(2m_c)^2}\approx 0.02\left[\frac{E}{1\ \mathrm{GeV}}\right]\mathrm{fm}\ . \qquad (2)$$

As a matter of fact, this is the lifetime of the hadronic fluctuation with $m\approx 2m_c$ in a vacuum. According to this estimate, the creation of the $\bar{c}c$ pair at energies below 100–150 GeV can be attributed to the interaction with a single nucleon in the nucleus.

The second stage is the formation of the charmonium wave function. It lasts in the laboratory frame for a period τ_F, related to the reversed distance between low-energy levels of the $\bar{c}c$ system, multiplied by the Lorentz factor:

$$\tau_F\approx\frac{2}{m_{\psi'}-m_\psi}\left[\frac{E}{2m_c}\right]\approx 0.2\left[\frac{E}{1\ \mathrm{GeV}}\right]\mathrm{fm} \qquad (3)$$

This estimate demonstrates that, starting from energies of a few tens of GeV, the formation zone of the charmonium exceeds nucleus radii. In this case one can expect that transparency should exceed the prediction of the Glauber model due to a weak attenuation of the compact $\bar{c}c$ system in a nuclear medium [11]. The Glauber model is valid only at low energies, when τ_F is much shorter than the nuclear radius.

A numerical estimation of the nuclear transparency in the J/ψ photoproduction was first performed in Ref. [12] under the assumption that the $\bar{c}c$ pair propagates along

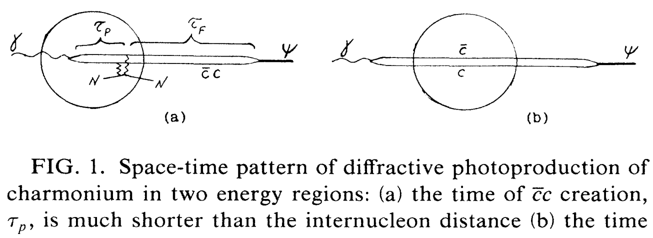

FIG. 1. Space-time pattern of diffractive photoproduction of charmonium in two energy regions: (a) the time of $\bar{c}c$ creation, τ_p, is much shorter than the internucleon distance (b) the time τ_p considerably exceeds the nuclear radius.

 ©1991 The American Physical Society

fixed trajectories starting from the pointlike configuration. It was assumed that the absorption cross section of the $\bar{c}c$ system increased proportionally to the distance covered by the quarks in the laboratory frame. The formation time τ_F was twice as small as ours. The authors of [12] found that the Glauber approach is approximately valid at SLAC energies. However their predictions overestimate the high-energy data [13].

It is worthwhile noting that, starting from the same arguments, the authors of [11] came to the opposite conclusion; they argued that the Glauber approximation should be crudely violated even at SLAC energies. From our viewpoint, the source of disagreement is the too large value of formation time used in [11], four times higher than ours, (3).

In addition, it is obvious that an approach based only on qualitative arguments and semiclassical estimations cannot produce any rigorous quantitative results. It is argued in the present paper that the inner dynamics of the $\bar{c}c$ system including quantum effects are of great importance. We approximate the wave function of charmonium with the nonrelativistic harmonic oscillator, and find an exact solution for the evolution operator of the $\bar{c}c$ system, propagating through a nuclear medium of varying density. An analogous approach was used by the authors earlier [7] for analysis of BNL data [8] on nuclear transparency in quasielastic *pp* scattering. Of course the application of the oscillator model in the latter case was rather questionable. For the $\bar{c}c$ system, it is much more justified.

From the point of view of the double-step approach [11] to the process of $\bar{c}c$ photoproduction, shown in Fig. 1, one can single out two energy regions where the treatment is most simplified. The first one corresponds to the case of $\tau_p \ll R_A$, shown in Fig. 1(a). Nuclear transparency can be written in the form

$$T_N(E)=\frac{\int d^3\mathbf{r}\,\rho_A(\mathbf{r})|\langle\Psi_f|\hat{U}|\Psi_{\rm in}\rangle|^2}{A|\langle\Psi_f|\Psi_{\rm in}\rangle|^2} \tag{4}$$

Here $\rho_A(\mathbf{r})$ is a nuclear density function; $\Psi_{\rm in}$ is a wave function of the $\bar{c}c$ system originated from the process $\gamma N\to\bar{c}cN$ [Fig. 1(a)]; Ψ_f is a wave function of the produced charmonium, J/ψ, ψ', etc.; $\hat{U}(E)$ is the evolution operator of the $\bar{c}c$ system in nuclear medium. We neglect the integration over the momentum transfer in the reaction $\gamma N\to\bar{c}cN$, because the transferred momentum cannot affect essentially the wave function $\Psi_{\rm in}$ of the compact $\bar{c}c$ system. For the same reason, i.e., due to the smallness of the radius of the $\bar{c}c$ system and its interaction cross section, we ignore in (4) any incoherent final-state interactions.

At much higher energies, when $\tau_p \gg R_A$, the pattern of charmonium photoproduction changes drastically. Now the photon converts into the $\bar{c}c$ pair long before the nucleus, as is shown in Fig. 1(b). The condition $\tau_p \gg R_A$, guarantees simultaneously a smallness of a variation of the transverse size of the $\bar{c}c$ system during propagation through the nucleus. So the influence of the nuclear medium is reduced to a simple attenuation factor, and the nuclear transparency takes the form

$$T_N(E\to\infty)=\frac{\int d^2\mathbf{b}\,T(\mathbf{b})|\langle\Psi_f|\exp[-\sigma(\rho)T(\mathbf{b})/2]|\Psi_{\rm in}\rangle|^2}{A|\langle\Psi_f|\Psi_{\rm in}\rangle|^2}\,. \tag{5}$$

Here $\mathbf{b}$ is an impact parameters of the $\bar{c}c$-pair center of mass; $T(\mathbf{b})=\int_{-\infty}^{\infty}dz\,\rho_A(\mathbf{b},z)$ is the nucleus profile function; $\sigma(\rho)$ is the interaction cross section of the $\bar{c}c$ pair, depending on its relative impact parameter ρ. The integration over ρ is assumed in (5). In analogy with formula (4) we neglect the multiple incoherent interactions of the $\bar{c}c$ pair, as well as the integration over transverse momentum.

It is worthwhile emphasizing that the color transparency phenomenon is not reduced to a simple filtering of pointlike $\bar{c}c$ pairs. Nuclear absorption distorts the form of the $\bar{c}c$ wave function. From the phenomenological point of view, this phenomenon is equivalent to the effects of the Gribov inelastic corrections [14]. The latter are known to bring about an antishadowing [15,16] in some cases, i.e., an increase of the nuclear transparency in comparison with expectations of the Glauber approximation. It is shown below that just this phenomenon takes place for ψ' photoproduction on nuclei: in spite of the attenuation in nuclear matter, the yield of ψ' per nucleon is predicted be higher than on a free-nucleon target. So nuclear transparency, formally defined in (1), is above 1 in this case.

II. EVOLUTION OF $\bar{c}c$ SYSTEM PASSING A NUCLEUS

Let us go to the c.m. of the $\bar{c}c$ pair where a nonrelativistic quantum-mechanical description is appropriate for lowest states. If one represents the evolution operator in the form of functional integral, the influence of nuclear medium will result in a supplementary attenuation factor, $\exp(-\frac{1}{2}\int dl\,\sigma(\rho)\rho_A(\mathbf{r}))$ for each virtual trajectory, where $\sigma(\rho)$ is the total cross section of the $c\bar{c}$ pair interaction with a nucleon, depending on the transverse interquark distance ρ. The integral is taken along the trajectory in the laboratory frame. So the evolution operator can be represented in the form

$$U=\int D^3\tau\exp\left[i\int dt\,L_{\rm eff}(\tau,\dot{\tau},t)\right]\,, \tag{6}$$

$$L_{\rm eff}(\tau,\dot{\tau},t)=L(\tau,\dot{\tau})+\frac{iv\gamma}{2}\sigma(\tau_T)\rho_A(\mathbf{r}(t))\,. \tag{7}$$

Here τ is an interquark radius vector; γ and v are the Lorenz factor and the velocity of the $\bar{c}c$ pair in the laboratory frame; $L(\tau,\dot{\tau})$ is the vacuum Lagrangian of the $\bar{c}c$ system. We approximate the latter with the harmonic-oscillator model

$$L(\tau,\dot{\tau})=\frac{\mu\dot{\tau}^2}{2}-\frac{\mu\omega^2\tau^2}{2}\,, \tag{8}$$

where $\mu=m_c/2$, $m_c=1.5$ GeV. The oscillatory frequency $\omega=(M_{\psi'}-M_\psi)/2$ is adjusted to the low states of charmonium.

The widespread approach to the problem of the total cross section of hadron interactions is the double-gluon approximation in QCD [1–3]. A $q\bar{q}$ pair with a relative impact parameter ρ interacts with a nucleon with the cross section

$$\sigma(\rho)=\frac{16\alpha_s^2}{3}\int d^2\mathbf{k}\frac{[1-\exp(i\mathbf{k}\cdot\rho)][1-F(\mathbf{k})]}{(k^2+m_g^2)^2} \tag{9}$$

The effective gluon mass m_g is introduced to account for the confinement. We fix it at the pion mass. The double-quark form factor $F(k)=\langle N|\exp[i\mathbf{k}(\mathbf{r}_1-\mathbf{r}_2)]|N\rangle$ is averaged over the nucleon wave function.

We take into account also the evolution of the QCD coupling, $\alpha_s(q^2)$, which is essential due to the smallness of the charmonium radius. According to the usual prescription [17], one should choose a maximal virtuality q^2 of lines on a Feynman diagram entering the vertex. So we put the product $\alpha_s(k^2)\alpha_s[\max(k^2,1/\rho^2)]$ in place of α_s^2 in (9), where $1/\rho^2$ characterizes the virtuality of the c-quark line. We use the one-loop approximation for the k^2 behavior of α_s. However at small values of k^2, perturbative QCD fails; then we fix $\alpha_s(k^2)$ at a constant value. These two regimes join at some border value of $k=k_0$:

$$\alpha_s(k)=\begin{cases}\dfrac{2\pi}{9\ln(k_0/\Lambda_{\rm QCD})} & \text{if } k<k_0\ ,\\[2mm] \dfrac{2\pi}{9\ln(k/\Lambda_{\rm QCD})} & \text{if } k>k_0\ .\end{cases}$$

Normalizing $\langle\sigma(\rho)\rangle_\pi=\sigma_{\rm tot}^{\pi N}=24$ mb, we fix $k_0=0.47$ GeV at $\Lambda_{\rm QCD}=0.2$ GeV. Computed in this way $\sigma(\rho)$ reproduces well the European Muon Collaboration (EMC) data on q^2 evolution of the nucleon structure function at small x [18].

At small values of $\rho^2\le\langle\rho^2\rangle_{\psi,\psi'}$ the cross section (9) is close to a simple behavior

$$\sigma(\rho)\approx C\rho^2\ , \tag{10}$$

used hereafter for the computing of the nuclear transparency. The factor C is fixed by the relation

$$C=\frac{\sigma_{\rm tot}(\psi N)}{\langle\rho^2\rangle_\psi}\approx\frac{\sigma_{\rm tot}(\psi' N)}{\langle\rho^2\rangle_{\psi'}}\ , \tag{11}$$

where $\sigma_{\rm tot}(\psi N)$ and $\sigma_{\rm tot}(\psi' N)$ are the average values of $\sigma(\rho)$ weighted with squares of J/ψ and ψ' wave functions, respectively. Using the oscillatory model and the double-gluon approximation we find $\sigma_{\rm tot}(\psi N)=5.75$ mb and $\sigma_{\rm tot}(\psi' N)=12.23$ mb. The mean radii are, $\langle\rho^2\rangle_\psi=2/m_c\omega$, $\langle\rho^2\rangle_{\psi'}=7\langle\rho^2\rangle_\psi/3$. Note that this value of $\langle\rho^2\rangle_\psi$ [and consequently $\sigma_{\rm tot}(\psi N)$] obtained in simplified oscillatory model is close to the result of exact calculations with the realistic wave function of J/ψ [19]. At the same time this estimate of $\sigma_{\rm tot}(\psi N)$ is considerably higher than the value extracted from photoproduction data using the vector-dominance hypothesis [20]. The latter however is known to fail crudely for J/ψ [21,22].

We compare the behavior (10) shown in Fig. 2 by a dashed curve, with the more sophisticated double-gluon approximation. One can see that both curves are very close at small $\rho^2\le\langle\rho\rangle_{\psi'}^2$.

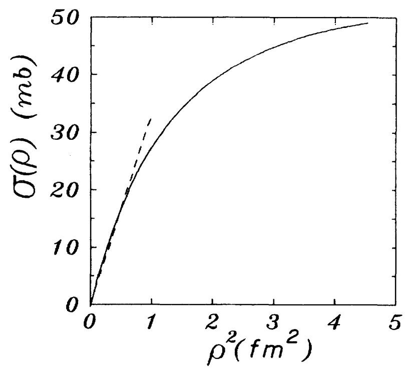

FIG. 2. Interaction cross section of a $\bar{q}q$ pair separated by a relative impact parameter ρ, on a nucleon target. The solid curve corresponds to the double-gluon-exchange approximation. The dashed curve shows a simple ρ^2 behavior (10).

The absorption term in (7) leads to a modification of the frequency ω_T of transverse oscillators [7]:

$$\omega_T=[\omega^2-i\delta(\mathbf{r})]^{1/2}\ ,$$

where

$$\delta(\mathbf{r})=\rho_A(\mathbf{r})\gamma v\omega\sigma_{\rm tot}(\psi N)\ .$$

To calculate the evolution operator of the $\bar{c}c$ system propagating through a nucleus with a varying density function we changed the latter with a multistep function. Within each slice of constant density one can use the known expression for the evolution operator for an harmonic oscillator with constant frequency [23,7]. For the one-dimensional oscillator,

$$\langle y|\hat{U}(t)|x\rangle=\left[\frac{\mu\omega}{2\pi i\sin(\omega t)}\right]^{1/2}\times\exp\left[\frac{i\mu\omega}{2\sin(\omega t)}[(y^2+x^2)\cos(\omega t)-2xy]\right], \tag{12}$$

where x and y are the initial and final coordinates of the oscillator.

The evolution operator for the multistep nuclear density can be found using the convolution relation

$$\hat{U}(t_{n+1})=\hat{U}(t_{N+1}-t_n)\otimes\hat{U}(t_n)\ . \tag{13}$$

Here t_i is the moment the border between corresponding slices passes by the $\bar{c}c$ pair. Note that the recurrent sequence (13) can be finished as soon as the nuclear density becomes sufficiently small, because the evolution operator in the vacuum provides only a phase factor, unessential for nuclear transparency (4).

After applying expression (12) and relation (13) we get the entire evolution operator, which also has a Gaussian form

$$\langle y|\hat{U}(t)|x\rangle = A(t)\exp\{i[\alpha(t)y^2+\beta(t)x^2+\gamma(t)xy]\} . \quad (14)$$

Here t is the total time of propagation of the $\bar{c}c$ pair along the given trajectory through the nucleus. Values of the factors $A(t)$, $\alpha(t)$, $\beta(t)$, and $\gamma(t)$ can be computed using the following recurrent relations following from (13):

$$\begin{aligned} A(t_{n+1}) &= A(t_{n+1}-t_n)A(t_n)\left[\frac{i\pi}{\alpha(t_n)+\beta(t_{n+1}-t_n)}\right]^{1/2} , \\ \alpha(t_{n+1}) &= \alpha(t_{n+1}-t_n)-\frac{\gamma^2(t_{n+1}-t_n)}{4[\alpha(t_n)+\beta(t_{n+1}-t_n)]} , \\ \beta(t_{n+1}) &= \beta(t_n)-\frac{\gamma^2(t_n)}{4[\alpha(t_n)+\beta(t_{n+1}-t_n)]} , \\ \gamma(t_{n+1}) &= -\frac{\gamma(t_{n+1}-t_n)\gamma(t_n)}{2[\alpha(t_n)+\beta(t_{n+1}-t_n)]} . \end{aligned} \quad (15)$$

The factors A, α, β, γ depending on the argument $t_{n+1}-t_n$ are defined according to expression (12).

Summarizing, expression (14) and relations (15) solve the problem of the determination of the entire evolution operator for a given trajectory of the $\bar{c}c$ pair.

III. CHOICE OF INITIAL WAVE FUNCTION OF THE $\bar{c}c$ PAIR.

In order to calculate the nuclear transparency using formula (4) one needs also an initial wave function $\Psi_{\rm in}$ of the diffractively produced $\bar{c}c$ pair in the process $\gamma N\to\bar{c}cN$. For the sake of simplicity we restrict ourselves to two variants of $\Psi_{\rm in}$. The first one uses the wave function of the $\bar{c}c$ pair produced in the process $\gamma N\to\bar{c}cN$ at very high energies, when the production time τ_p is considerably higher than a nucleon radius. We assume that this choice of $\Psi_{\rm in}$ can be used at maximal allowed energies (see the Introduction) about 100–150 GeV, where the ratio τ_p/r_N is of the order of 2–3.

In accordance with the diffractive mechanism of $\bar{c}c$-pair production shown in Fig. 1, the asymptotic $\bar{c}c$ wave function is a product of the quark wave function of the photon and the amplitude of $\bar{c}c$-pair interaction with a nucleon, i.e., $\sigma(\rho)$. The former should be taken for that photon component which has an helicity equal to a sum of helicities of $\bar{c}$ and c, because we are working within the nonrelativistic approach to the charmonium wave function. Using the formula of the noncovariant perturbative theory in the infinite momentum frame

$$|\Psi\rangle = \sum_n \frac{|n\rangle\langle n|\hat{V}|i\rangle}{E_i-E_n} ,$$

we get the wave function in the momentum representation

$$\Psi_\gamma(\alpha,k_T)\propto(m_c^2+k_T^2)^{-1} , \quad (16)$$

where α is the light-cone variable of the $\bar{c}c$ pair, a part of the total momentum carried by one of the quarks. Coming back to the ρ representation we get the transverse part of the photon wave function in the form of the modified Bessel function $k_0(m_c\rho)$. Gathering all parts together, the transverse part of the $\bar{c}c$ wave function $\Psi_{\rm in}$ [18] takes the form

$$\Psi_{\rm in}^T(\rho)\propto K_0(m_c\rho)\sigma(\rho) . \quad (17)$$

This behavior shown in Fig. 3, except the very far tail is quite exactly reproduced with a simple parametrization

$$\Psi_{\rm in}^T(\rho)\approx{\rm const}\times[\exp(-\rho^2/a^2)-\exp(-\rho^2/b^2)] ,$$

where $a=0.536$ fm, $b=0.11$ fm.

A longitudinal part of the $\bar{c}c$-pair wave function (16) produced in the reaction $\gamma N\to\bar{c}cN$, is independent of α. So in the c.m. of the $\bar{c}c$-pair the intrinsic momentum is distributed in a region of the order of m_c. Then the $\bar{c}c$ pair is located within a region $\Delta z\approx 1/m_c$ of the longitudinal coordinate. A specific choice of the form of the longitudinal part of $\Psi_{\rm in}$ does not play any role for the J/ψ production because of a factorization of transverse and longitudinal coordinates in the oscillatory model. It does not affect considerably the nuclear transparency for ψ' photoproduction also, in spite of the lack of the factorization. We use the Gaussian parametrization, $\Psi_{\rm in}^L\propto\exp(-z^2/d^2)$, with $d=1/2m_c$. We check below a sensitivity of the results to the parameter d.

An additional test of validity of input wave function (17) is the calculation of a ratio of yields of J/ψ to ψ' on a nucleon target:

$$R=\left|\frac{\langle J/\psi|\Psi_{\rm in}\rangle}{\langle\psi'|\Psi_{\rm in}\rangle}\right|^2 .$$

The computed value $R=6.5$ nicely agrees with the measured value [24] $R=6.8\pm2.4$. Nevertheless we should note that the input wave function (17) is approximate even at asymptotic energies. We used in (16) the free-quark approximation, neglecting the interquark interaction. The latter brings about a correction of about 30% to the mass of the $\bar{c}c$ system at a relative distance of $1/m_c$. This effect is important for the absolute value of

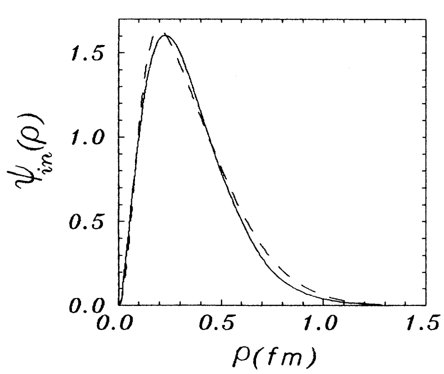

FIG. 3. The input wave function $\Psi_{\rm in}(\rho)$ of a $\bar{c}c$ pair photoproduced on a nucleon. The solid curve corresponds to the exact expression (17). The dashed curve is the result of a fit with a two-Gaussian parametrization.

the photoproduction cross section, but is not essential for the nuclear transparency.

Let us remember that the minimum in the ρ dependence of $\Psi_{\rm in}^T(\rho)$ at $\rho=0$ is the result of the modulation of the wave function with the factor $\sigma(\rho)$ in (17). It is truly only if the transverse coordinates of the $\bar{c}c$ pair are "frozen" during the interaction with a single nucleon. The latter is possible at sufficiently high energy, which provides $\tau_p \gg r_p$. At lower energies, when this condition is crudely violated, a transverse shift of quarks during the interaction with the nucleon is important. Indeed, $\Delta\rho \approx v_T \tau_p \approx 1/m_c$ where $v_T \approx 2m_c/E_\gamma$ is the velocity of transfer motion of c quarks in the laboratory frame. Thus the transverse shift of quarks during the interaction is of the order of the size of the quark localization region. As a result, the specific form of the wave function (17) is entirely wiped out. For this reason at low energies we use a simple parametrization of the initial $\bar{c}c$ wave function in the form

$$\Psi_{\rm in}^T(\rho) \propto \exp(-\rho^2/a^2) \ . \tag{18}$$

Note that the decreasing of the time of life of the $\bar{c}c$ pair, τ_p, leads to the reduction of the transverse dimension of the fluctuation, because the hadronic fluctuation of the photon starts from a point. For this reason the parameter a in (18) is not connected directly with the dimension of asymptotic distribution (17). Below we will test a few values of a.

IV. RESULTS OF CALCULATIONS

The results of a calculation of the transparency of nuclei ^{9}Be, ^{56}Fe, and ^{207}Pb with the Saxon-Woods nuclear density, for the photoproduction of J/ψ and ψ', are shown in Fig. 4. We used expression (4) and a $\bar{c}c$ initial wave function in the form of (17). Experimental data [13] on J/ψ photoproduction at an energy of $E_\gamma \approx 120$ GeV are depicted in the same picture. Though the calculations were performed in a wide energy interval, we recall that the usage of the asymptotic form of wave function (17) is questionable at low energies. In addition, the results of the Glauber model for photoproduction of J/ψ and ψ',

$$T_{N({\rm Gl})}=\frac{1}{A}\int d^2\mathbf{b}\int_{-\infty}^{\infty} dz\,\rho_A(\mathbf{b},z) \times \exp\left[-\sigma_{\rm tot}(\psi N)\int_z^{\infty} dz'\,\rho_A(\mathbf{b},z')\right] , \tag{19}$$

are also depicted in Fig. 4 with crosses.

One can see that the transparency for J/ψ photoproduction is nearly energy independent and is close to the Glauber-model prediction. At high energies, where this variant of calculations is most justified, our predictions agree well with the experimental data [13].

On the contrary, the photoproduction of ψ', displays some peculiarities. First, at high energies the transparency is considerably higher than the Glauber predictions, so inelastic corrections play an important role. Second, the transparency for ψ' is higher than for J/ψ production, in spite of the fact that $\sigma_{\rm tot}(\psi' N)$ is more than two times higher than $\sigma_{\rm tot}(\psi N)$. At last the transparency exceeds unity at energies higher than 40–60 GeV, i.e., nuclei enhance the yield of ψ'. This means that the nuclear transparency defined in (1) cannot be interpreted in accordance with an intuitive understanding, as a simple attenuation in the nuclear medium. This result shows also that the distortion of the $\bar{c}c$ wave function during propagation through the nucleus plays a more important role

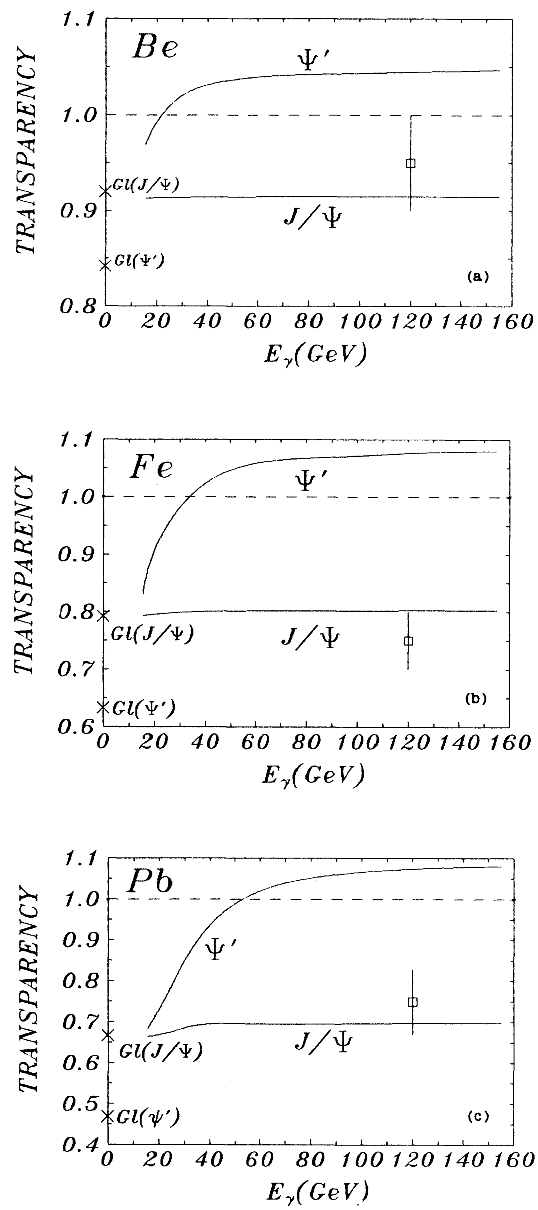

FIG. 4. Energy dependence of nuclear transparency for J/ψ and ψ' photoproduction on (a) ^{9}Be, (b) ^{56}Fe, and (c) ^{207}Pb, computed with the asymptotic input wave function (17). The crosses depicted on the y axis are the predictions of the Glauber approximation (19), independent on energy.

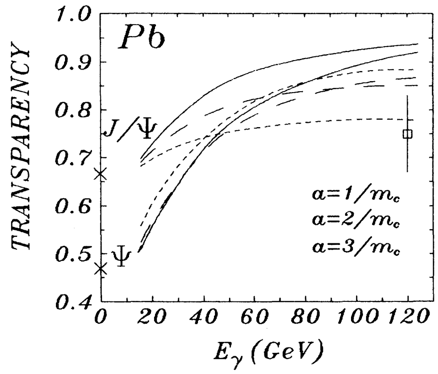

FIG. 5. Energy dependence of nuclear transparency for J/ψ and ψ' photoproduction on ^{207}Pb computed with the low-energy input wave function (18), vs value of parameter a.

than the nuclear attenuation. Of course the total yield of $\bar{c}c$ is shadowed. Note that the possibility of a positive contribution of inelastic corrections was discussed earlier in Refs. [15,16].

The results of calculations with the second variant (18) of $\psi_{\rm in}$ more appropriate at low energies are shown in Fig. 5 versus the value of the parameter a. This choice of the initial wave function essentially modifies the nuclear transparency. First, in the case of J/ψ production a strong energy dependence appears. Second, the relation between yields of J/ψ and ψ' is found to be sensitive to the size of the initial $\bar{c}c$ system: the larger the parameter a, the higher the relative yield of ψ'. Within the uncertainty of the parameter a, the results agree well with the measurements at SLAC at $E_\gamma=20$ GeV [25]. The measured value of the ratio $R=Tr(^9\text{Be})/Tr(^{180}\text{Ta})=1.21\pm0.08$ should be compared with the prediction ranging from 1.25 to 1.27 for the parameter $a=(1-3)/m_c$.

Now let us proceed to another energy region where both τ_p and τ_f are much higher than the nuclear radius R_A. Under these conditions the relative impact parameter ρ of the $\bar{c}c$ pair is "frozen" during the propagation through the nucleus, so one can use asymptotic expressions (5), for the transparency, and (17) for the $\bar{c}c$-pair wave function. The results of calculations are collected in Table I.

Note that the asymptotic values of the transparency are lower for J/ψ and a little higher for ψ' than those at intermediate energies, depicted in Fig. 4. Consequently the growth of the transparency at intermediate energies should turn to a fall at higher energies. The reason is obvious: in the latter case a path covered by the hadronic fluctuation inside a nucleus is longer. As a result the influence of the nucleus at asymptotic energies is stronger, but we have found out that it manifests itself as the shadowing for J/ψ and the antishadowing for ψ'.

TABLE I. Nuclear transparency for J/ψ and ψ' at asymptotic energies.

A	Be	Fe	Pb
J/ψ	0.85	0.72	0.5
ψ'	1.08	1.12	1.08

V. CONCLUSIONS

Let us summarize the main conclusions of the present paper.

(a) Quantum effects for a quark system propagating through a nuclear medium are very important. The nuclear absorption causes not only an attenuation but distorts also the quark wave function.

(b) Nuclear transparency essentially depends on the wave function of the quark system at the moment of its creation. Figs. 4 and 5 illustrate the sensitivity to the choice of wave function. We conclude that nuclear shadowing of charmonium production in hadron-nucleus interactions is uncertain up to the initial $\bar{c}c$ wave function which depends on the production mechanism.

(c) The relative yields of different final states considerably vary depending on their wave functions. The exciting prediction of this paper is nuclear antishadowing of ψ' photoproduction in spite of nuclear absorption.

(d) On the contrary, to the naive expectation of the Glauber approach, nuclear transparency for photoproduction of ψ' is higher than for J/ψ, in spite of the larger absorption cross section of the former. This result might explain the experimentally observed high yield of ψ' in hadron-nucleus interactions [26].

(e) At asymptotic energies nuclear effects are enhanced, both the shadowing for J/ψ and the antishadowing for ψ'.

Summarizing, the color transparency phenomenon is analyzed in the present paper in a simple and clear case of heavy-quarkonium photoproduction. The theoretical expectations agree well with the available experimental data on nuclear enhancement of the J/ψ photoproduction cross section. Nevertheless more precise measurements in a wide energy range, as well as data for ψ', are desirable to have a more definite confirmation of the color transparency phenomenon. It would be of high interest to have also data for a photoproduction of $\bar{b}b$ quarkonia in which the present approach can be used.

[1] F. Low, Phys. Rev. D **12**, 163 (1975).
[2] S. Nussinov, Phys. Rev. Lett. **34**, 1286 (1975).
[3] J. F. Gunion and D. Soper, Phys. Rev. D **15**, 2617 (1977).
[4] A. Mueller, in *Proceedings of the XVIIth Rencontre de Moriond on Elementary Particle Physics: I. Electroweak Interactions and Grand Unified Theories,* Les Arcs, France, 1982, edited by J. Tran Thanh Van (Editions Frontières, Gif-sur-Yvette, 1982), p.13.
[5] S. J. Brodsky, in *Multiparticle Dynamics 1982,* Proceedings of the XII International Symposium, Volendam,

Netherlands, 1982, edited by W. Kittel, W. Metzger, and A. Stergiou (World Scientific, Singapore, 1983), p. 963
[6] B. Z. Kopeliovich and B. G. Zakharov, Yad. Fiz. **46**, 1535 (1987) [Sov. J. Nucl. Phys. **46**, 911 (1987)].
[7] B. Z. Kopeliovich and B. G. Zakharov, Phys. Lett. B **264**, 434 (1991).
[8] A. S. Carroll *et al.*, Phys. Rev. Lett. **61**, 1698 (1988).
[9] J. P. Ralston and B. Pire, Phys. Rev. Lett. **61**, 182 (1988).
[10] S. J. Brodsky and G. F. de Teramond, Phys. Rev. Lett. **60**, 1924 (1988).
[11] S. J. Brodsky and A. H. Mueller, Phys. Lett. B **206**, 685 (1988).
[12] G. N. Farrar, H. Liu, L. L. Frankfurt, and M. I. Strikman, Phys. Rev. Lett. **61**, 686 (1988).
[13] M. D. Sokoloff *et al.*, Phys. Rev. Lett. **57**, 3003 (1986).
[14] V. N. Gribov, Zh. Eksp. Teor. Fiz **56**, 892 (1969) [Sov. Phys. JETP **29**, 483 (1969)].
[15] B. Z. Kopeliovich, L. I. Lapidus, S. V. Mukhin, and A. B. Zamolodchikov, Zh. Eksp. Teor. Fiz. **77**, 461 (1979) [Sov. Phys. JETP **50**, 434 (1979)].
[16] N. N. Nikolaev, Zh. Eksp. Teor. Fiz. **81**, 814 (1981) [Sov. Phys. JETP **54**, 434 (1981)].
[17] Yu. L. Dokshitzer, D. I. Diakonov and S. I. Troyan, Phys. Rep. **C58**, 269 (1980).
[18] N. N. Nikolaev and B. G. Zakharov, Z. Phys. C **49**, 607 (1991).
[19] M. Eichten *et al.*, Phys. Rev. D **21**, 203 (1980).
[20] M. Binkley *et al.*, Phys. Rev. Lett. **48**, 73 (1982).
[21] K. G. Boreskov and B. L. Ioffe, Yad. Fiz. **25**, 623 (1977) [Sov. J. Nucl. Phys. **25**, 331 (1977)].
[22] J. Pumplin, Phys. Rev. D **28**, 2741 (1983).
[23] R. Feynman and A. R. Hibbs, *Quantum Mechanics and Path Integrals* (McGraw-Hill, New York, 1965).
[24] U. Camerini *et al.*, Phys. Rev. Lett. **35**, 483 (1975).
[25] R. Anderson, Phys. Rev. Lett. **38**, 263 (1977).
[26] E772 Collaboration, D. M. Alde *et al.*, Phys. Rev. Lett. **66**, 133 (1991).