Novel color transparency effect: 
scanning the wave function of vector 
mesons

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Abstract

We demonstrate how the virtual photoproduction of vector mesons on nuclei scans the wave function of vector mesons from the large non-perturbative transverse size $\rho \sim R_V$ down to the small perturbative size $\rho \sim 1/\sqrt{Q^2}$. Thee mechanism of scanning is based on color transparency and QCD predicted spatial wave function of quark-antiquark fluctuations of virtual photons. A rich, energy- and $Q^2$-dependent, pattern of the nuclear shadowing and antishadowing is predicted, which can be tested at the European Electron Facility and SLAC.
Introduction. In this paper we discuss the novel feature of color transparency (CT) tests of QCD: the scanning of the nonperturbative hadronic wave functions. The virtual photoproduction of vector mesons (ϒ, ϒ’, J/Ψ, Ψ’, ρ, ρ’, ...) is particularly suited for such CT tests. At high energy $\nu$, the photoproduction can be viewed [1,2] as a production of the virtual $\bar{q}q$ pair with the coherence length

$$l_c = \frac{2\nu}{Q^2 + m_V^2}. \hspace{1cm} (1)$$

The size $\rho_Q$ of the produced $\bar{q}q$ pair (the ejectile state) is controlled by the virtuality $Q^2$ of photons [3]. The $\bar{q}q$ state is projected onto (recombines into) the final-state vector meson $V$ with the formation (recombination) length

$$l_f = \frac{\nu}{m_V \Delta m}, \hspace{1cm} (2)$$

where $\Delta m$ is the typical level splitting in the quarkonium. We demonstrate, how by changing $Q^2$ and the initial size $\rho_Q$ one can scan the wave function of vector mesons starting from the non-perturbative size $\rho \sim R_V$ down to the perturbative region of $\rho \sim 1/Q$.

Besides the scanning radius $\rho_Q$, the formation length $l_f$ is a second important parameter of CT physics. Changing energy $\nu$, one can vary $l_f$ from $l_f \ll R_A$ (quasi-instantaneous formation of the final-state hadron) to $l_f > R_A$ (the frozen-size limit), and thus study the dynamical evolution of the small-sized, perturbative $q\bar{q}$ pair to the full-sized nonperturbative hadron. We predict very rich pattern of the nuclear shadowing and antishadowing phenomena, which changes with the scanning radius $\rho_Q$, i.e., with $Q^2$, and with the evolution rate, i.e., with energy $\nu$.

It is worth while to emphasize that the virtual photoproduction exemplifies in a particularly transparent way the principle idea of CT tests of QCD: (1) A small transverse-size component of the interacting hadrons, $\rho \ll R_h$, is selected by the interaction dynamics. (2) The interaction cross section $\sigma(\rho)$ of the small-sized ejectile is measured by a strength of the final state interaction (FSI) in the target nucleus [4,5]. Notice a combination of the perturbative and nonperturbative aspects of QCD: the perturbative production of the small-sized ejectile is followed by probing its size in the nonperturbative, diffractive small-angle scattering in the nuclear matter. Nevertheless, QCD as a theory of strong interactions predicts that the strength of this diffractive scattering vanishes as $\rho \rightarrow 0$: $\sigma(\rho) \propto \rho^2 [6,7]$.

The scanning mechanism. The quantum-mechanical description of the scanning goes as follows: At $l_f, l_c > R_A$ the amplitude of the forward photoproduction on the free nucleon $M_N$ and the nuclear transmission coefficient or the nuclear transparency $T_{RA} = d\sigma_A/Ad\sigma_N$ in the quasielastic photoproduction $\gamma^* A \rightarrow VA$ read [8,9]:

$$M_N = \langle V | \sigma(\rho) | \gamma^* \rangle \hspace{1cm} (3)$$

$$T_{RA} = \frac{1}{A} \int d^2b T(b) \frac{\langle V | \sigma(\rho) \exp \left[ -\frac{1}{2} \sigma(\rho) T(b) \right] | \gamma^* \rangle^2}{\langle V | \sigma(\rho) | \gamma^* \rangle^2} \hspace{1cm} (4)$$

where $T(b) = \int dz n_A(b, z)$ is the optical thickness of a nucleus, the nuclear density $n_A(b, z)$ is normalized to the nuclear mass number $A$: $\int d^3 \vec{r} n_A(\vec{r}) = A$. 

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The wave function $|\gamma^*\rangle$ of the $q\bar{q}$ fluctuations of the virtual photons was calculated in [3] in the mixed $(\vec{\rho}, z)$ representation, where $z$ is a fraction of photon’s light-cone momentum, carried by the quark, $0 < z < 1$. The most important feature of $|\gamma^*\rangle$ is an exponential decrease at large distances [3]:

$$|\gamma^*\rangle \propto \exp(-\varepsilon \rho)$$  (5)

where

$$\varepsilon^2 = m_q^2 + z(1-z)Q^2$$  (6)

Calculation of the matrix elements in (3,4) involves the $d^2\vec{\rho}dz$ integration. In the nonrelativistic quarkonium $z \approx 1/2$, so that the relevant $q\bar{q}$ fluctuations have a size $\rho \sim \rho_Q = \sqrt{m^2_q + 1/4Q^2} \approx 2\sqrt{m^2 + V^2 + Q^2}$ (7)

What enters (3,4) is a product $\sigma(\rho)|\gamma^*\rangle$. Because of CT property, $\sigma(\rho)|\gamma^*\rangle \propto \rho^2$ at small $\rho$ [6,7], this product will be sharply peaked at $\rho \approx \rho_2 = 2\rho_Q$ with the width $\Delta \rho \sim 2\rho_Q$, which leads naturally to the idea of scanning: The transition matrix elements (3,4) probe the wave function $|V\rangle$ at $\rho \sim 4\rho_Q$, and varying $\rho_Q$ by changing $Q^2$, one can scan the wave function $|V\rangle$ from large to small distances. In Fig.1 we demonstrate qualitatively how the scanning works. We also show the $z$-integrated wave functions of the ground state $|V\rangle$ and of the radial excitation $|V'\rangle$.

The nuclear matrix element in (4) can be expanded in the moments $\langle V|\sigma(\rho)^n|\gamma^*\rangle$ which are the true QCD observables of the virtual photoproduction process. These moments probe the wave function $|V\rangle$ at different values of $\rho \sim \rho_{2n}$. To the leading order in the final state interaction,

$$T_{rA} = 1 - \Sigma_V \frac{1}{A} \int d^2\vec{b} T(b)^2$$  (8)

where

$$\Sigma_V = \frac{\langle V|\sigma(\rho)^2|\gamma^*\rangle}{\langle V|\sigma(\rho)|\gamma^*\rangle}$$  (9)

This expansion works well when $1 - T_{rA} \ll 1$, and is convenient to explain how the scanning proceeds. Notice, that $\sigma(\rho)^2|\gamma^*\rangle$ peaks at $\rho \sim \rho_4 = 4\rho_Q$.

**Scanning the vector mesons: the node effect.** Start with the real photoproduction ($Q^2 = 0$) of the ground-state vector mesons ($\Upsilon, J/\Psi, \rho^0, ...$) and take the case of the charmonium. In this case $\rho_Q = 1/m_c \ll R_{J/\Psi}$, $\rho_2$ is still smaller than $R_{J/\Psi}$, but $\rho_4 \approx R_{J/\Psi}$. For this numerical reason, one finds $\Sigma_{J/\Psi} \approx \sigma_{tot}(J/\Psi N)$, i.e., the predicted nuclear shadowing will be marginally similar [8,9] to the Vector Dominance Model (VDM) prediction [10]. This holds to much extent for the $\Upsilon$ and the light vector mesons as well. At larger $Q^2$, one has $\rho_2, \rho_4 \ll R_V$ and $\Sigma_V \sim \sigma(\rho_Q) \propto \rho_Q^2$ with the calculable logarithmic corrections [3,11], so that the above marginal similarity with the VDM disappears. The nuclear transparency will tend to unity from below:

$$1 - T_{rA} \propto \frac{A}{R^2_A} \rho_Q^2$$  (10)
The case of the radial excitation $V'$ is more interesting. Radial excitations have larger radius and larger free nucleon cross section: e.g., $\sigma_{\text{tot}}(\Psi'N) \approx 2.5 \sigma_{\text{tot}}(J/\Psi N) [8]$, which classically would suggest much stronger final state interaction for $\Psi'$ than for $J/\Psi$. Presence of the node in the $V'$ wave function leads to a rather complex pattern of the shadowing and antishadowing. In the photoproduction limit, because of $\rho_2 \sim R_{\Psi'}$, there are rather strong cancellations between the contributions to the amplitude (3) from $\rho$ below and above the node (the node effect, see Fig.1). For this reason, in photoproduction on the free nucleons, one predicts the ratio of the forward production differential cross sections $r(Q^2 = 0) = d\sigma(\gamma N \to V'N)/d\sigma(\gamma N \to V N)|_{t=0} < 1$. For the charmonium the $\Psi'/(J/\Psi)$ ratio $r(0) = 0.17$. In the $\Upsilon'$ production, the initial size $\rho_Q$ scales as $1/m_q$, whereas the radius of the $\bar{b}b$ bound states decreases with $m_q$ less rapidly. For this reason for the $\Upsilon'$ the node effect is much weaker and we find the $\Upsilon'/\Upsilon$ ratio $r(0) = 0.84$. The larger is $Q^2$, the smaller size $\rho_2$ is scanned, a contribution from the region above the node becomes negligible, and $r(Q^2)$ will increase with $Q^2$ up to $r(Q^2 \gg m^2) \sim 1$ (the exact limiting ratio depends on the wave function’s at the origin).

Since $\rho_1$ is closer to the node position, the node effect is still stronger in the matrix element $\langle V'|\sigma(\rho)^2|\gamma \rangle$. For the $\Psi'$ one finds $\langle \Psi'|\sigma(\rho)^2|\gamma \rangle < 0$, so that $\Sigma_{\Psi'} < 0$ and $T_{\gamma A}(\Psi') > 1$ despite the larger free-nucleon cross section [8,9,11]. In the $\Upsilon'$ production the scanning radius compared to the bottomonium radius is relatively smaller than in the charmonium case, the node effect is weaker and FSI produces the shadowing of the $\Upsilon'$. Still, in spite of $\sigma_{\text{tot}}(\Upsilon'N) \gg \sigma_{\text{tot}}(\Upsilon N)$, shadowing of the $\Upsilon'$ is weaker than shadowing of the $\Upsilon$.

The $Q^2$-dependent scanning changes the node effect significantly. The smaller the scanning radius $\rho_Q$, the weaker are the cancellations. In the $\Psi'$ case $\langle V'|\sigma(\rho)^2|\gamma \rangle)$ becomes positive defined, so that the antishadowing changes to the shadowing, which first rises with $Q^2$, then saturates and is followed by the onset of asymptotic decrease (10). In the $\Upsilon'$ production the node effect is weaker, still it affects the $Q^2$-dependence of the shadowing making it different from that of the $\Upsilon$.

Predictions for the heavy quarkonium photoproduction are shown in Figs.2,3. Because of the small size of heavy quarkonium, the results are numerically reliable, as they are dominated by the perturbative QCD domain.

The two scenarios of scanning the light vector mesons. The case of the light vector mesons is particularly interesting, as similar (anti)shadowing effects occur in the energy and $Q^2$ range accessible at SLAC and at the planned European Electron Facility (EEF). Here at moderate $Q^2 \lesssim m^2$ the scanning radius is large, $\rho_Q \sim R_V$, so that numerically the predicted shadowing of the $\rho^0$, Fig.4, is not very accurate, although the accuracy increases gradually with $Q^2$ as the scanning radius decreases into the perturbative domain $\rho_Q \ll R_V$. Nevertheless, since the wave function of the $\rho^0$ does not have a node, we believe to describe correctly a smooth transition from the real to virtual photoproduction.

In the photoproduction limit we find a strong node effect and strong suppression of the photoproduction of the radial excitation $\rho'$ on nucleons, by more than one order of magnitude compared to the $\rho^0$ production, which broadly agrees with the scanty experimental data [10,12]. As here $\rho_Q \sim R_V$, we can not give a reliable numerical estimate of how small the $\rho'/\rho^0$ cross section ratio is. Nevertheless, we can describe the two possible
scenarios of scanning the \( \rho' \) wave function, experimental tests of which can shed light on the spectroscopy and identification of the radial excitations of the light vector mesons:

(1) **The undercompensated free-nucleon amplitude:** \( \langle \rho'|\sigma(\rho)|\gamma \rangle > 0 \).

In this scenario the \( \rho' \) case will be similar to the \( \Psi' \) case, apart from the possibility of anomalously strong nuclear enhancement \( T_{RA} \gg 1 \), which might show strong atomic number dependence. Indeed, because of the larger relevant values of \( \rho \) the cross section \( \sigma(\rho) \) is larger, and the attenuation factor \( \exp[-\frac{1}{2}\sigma(\rho)T(b)] \) in the nuclear matrix element, eq.(4), will suppress the large \( \rho \) region, effectively decreasing the scanning radius \( \rho_Q \) and diminishing the node effect. Detailed description of the atomic number dependence will be presented elsewhere.

The \( Q^2 \)-scanning too will follow the \( \Psi' \)-scenario: change from the antishadowing to the shadowing with increasing \( Q^2 \), followed by the saturation and then decrease of the shadowing according to eq.(10). The range of \( Q^2 \) at which the major effects should occur corresponds to a change of the scanning radius \( \rho_Q \) by the factor \( \sim 2 \), i.e., to \( Q^2 \sim 4m_q \sim m_{\rho'}^2 \).

(2) **The overcompensated free-nucleon amplitude:** \( \langle \rho'|\sigma(\rho)|\gamma \rangle < 0 \).

This scenario is preferred in the crude oscillator model used in [3] and fits the pattern of the node effect becoming stronger for the lighter flavours. In this case \( \rho_2 \approx R_V \) and the higher moments too will be negative valued: \( \langle \rho'\sigma(\rho)^n|\gamma \rangle < 0 \), so that in the photoproduction limit one starts with the conventional shadowing: \( T_{RA} < 1 \).

In the \( Q^2 \)-scanning process the striking effect is bringing the free-nucleon amplitude \( \langle V'\sigma(\rho)|\gamma^* \rangle \) to the exact compensation at certain moderate \( Q^2 \), when the decreasing \( \rho_2 \) intercepts \( R_V \) (strictly speaking, because of the relativistic corrections and different quark helicity states, the compensation is unlikely to be exact). As a result, one finds a spike in \( T_{RA} \), Fig.3. With the further increase of \( Q^2 \) one enters the undercompensation regime, and the further pattern of the \( Q^2 \)-scanning will be essentially the same as in the undercompensation or the \( \Psi' \) scenario.

The above described \( Q^2 \)-dependent scanning of the wave function of light vector mesons offers a unique possibility of identifying the radial excitation of the light vector mesons (for the detailed discussion of the spectroscopy of light vector mesons see [12]). The corresponding experiments could easily be performed at SLAC and EEF.

**Quantum evolution, coherence and energy dependence of FSI.**

If the coherence length \( l_c > R_A \), then amplitudes of production on different nucleons add up coherently. At moderate energy, \( l_c < R_A \), the production rates on different nucleons at the same impact parameter \( \vec{b} \) add up incoherently. In the opposite limit of \( l_c > R_A \) amplitudes of production on different nucleons add up coherently, and nuclear affects are generally weaker [8,9,11].

For the heavy quarks we have a strong inequality \( l_f \approx l_c \) [13]. The same inequality holds for the light vector mesons at \( Q^2 \gg m_V^2 \). Consequently, at moderate energy, when \( l_c < R_A \), one has still a broad energy range in which \( l_f > R_A \) and the transverse size of the \( q\bar{q} \) pair is still frozen. In this regime the nuclear transparency is given by a simple formula

\[
T_{RA} = \frac{1}{A} \int d^2b dz n_A(b,z) \frac{\langle V|\sigma(\rho) \exp[-\frac{1}{2}\sigma(\rho)T(b,z)]|\gamma^* \rangle^2}{\langle V|\sigma(\rho)|\gamma^* \rangle^2} \tag{11}
\]
where \( t(b, z) = \int_z^\infty dz'n_A(b, z') \). Notice, that compared to the high-energy limit (4), here the attenuation effect in the nuclear matrix element is weaker. The energy dependence of the nuclear transparency is given by an approximate interpolation formula

\[
Tr_A(\nu) \approx Tr(l_c << R_A) + F_{ch}(\kappa)^2[Tr_A(l_c > R_A) - Tr_A(l_c << R_A)]
\]

where \( F_{ch}(\kappa) \) is the charge form factor of the target nucleus and \( \kappa = 1/l_c = (Q^2 + m_V^2)/2\nu \).

For the general idea of derivation of (14) and successful description of the NMC data [14] of the photoproduction of the J/Ψ see Ref.9.

If \( l_f < R_A \), the spatial expansion of the \( q\bar{q} \) pair becomes important. It can be described in either the quark basis used above, or in the hadronic basis, which is nice demonstration of the quark-hadron duality.

Consider the leading term of the final state interaction in eq.(10) in the hadronic basis, i.e., in terms of Gribov’s inelastic shadowing [15]. Inserting a complete set of the intermediate states, one can write down

\[
\langle V|\sigma(\rho)^2|\gamma^* \rangle = \sum_i \langle V|\sigma(\rho)|V_i\rangle \langle V_i|\sigma(\rho)|\gamma^* \rangle
\]

In the hadronic basis the antishadowing of the Ψ′ comes from the destructive interference of the direct, VDM-like rescattering

\[
\gamma^* \rightarrow \Psi' \rightarrow \Psi'
\]

and the off-diagonal rescattering

\[
\gamma^* \rightarrow J/\Psi \rightarrow Psil'
\]

(there is a small contribution from other intermediate states too). The reason for the strong cancellation is that the \( \gamma \rightarrow \Psi' \) transition is weak compared to the \( \gamma \rightarrow J/\Psi \) transition, whereas the \( J/\Psi \rightarrow \Psi' \) transition has an amplitude of opposite sign, and smaller, that the \( \Psi' \rightarrow \Psi' \) elastic scattering amplitude. To the contrary, in the \( J/\Psi \) photoproduction, both amplitudes in the off-diagonal transitions like \( \gamma \rightarrow \Psi' \rightarrow J/\Psi \) are small, and this explains why one finds a marginal similarity to the VDM in the photoproduction.

At larger \( Q^2 \) the node effect is no longer effective in suppressing the \( \gamma^* \rightarrow \Psi' \) transition, the off-diagonal amplitudes become significant for the \( J/\Psi \) photoproduction too and one finds strong departure from the VDM for the \( J/\Psi \) too.

By the nature of CT experiments, at finite energy \( \sigma(\rho)^2 \) in the l.h.s. of eq.(13) is the nonlocal operator: the \( q\bar{q} \) state produced on one nucleon is probed by a second nucleon a distance \( \Delta z \sim R_A \) apart. As a result, the diagonal and off-diagonal rescattering amplitudes acquire the relative phase [15,16] \( \Delta \varphi_{21} = \Delta z(m_2^2 - m_1^2)/2\nu \sim \Delta z/l_f \). Upon the integration over \( \Delta z \) only the intermediate states \( |i\rangle \) for which \( \Delta \varphi_{i1} \lesssim 1 \), or

\[
|m_i^2 - m_V^2| < 2\nu/R_A
\]

will contribute to the r.h.s. of eq.(13), so that the strength of the final state interaction will change rapidly from the near-threshold energy of \( l_f \ll R_A \) to a higher energy of \( l_f > R_A \). We describe the energy-dependence in this region using the quark-basis path-integral technique suggested in [3].
At $l_f \ll R_A$ the amplitudes of transitions (16) and (17) do not interfere. However, since they are of comparable magnitude, for the $\Psi'$ the VDM prediction for the nuclear shadowing breaks down even near the production threshold: $Tr_A$ is larger than the Glauber model prediction. For the $J/\Psi$ a similar incoherent contribution is small at small $Q^2$, rises with $Q^2$ as described above, and the near-threshold value of $Tr_A$ rises too. In the $\Psi'$ case the dominant effect of the $Q^2$-dependent scanning is that the $\gamma^* \rightarrow \Psi'$ transition amplitude increases with $Q^2$ relative to the $\gamma^* \rightarrow J/\Psi$ amplitude, which enhances the diagonal (shadowing) rescattering contribution compared to the off-diagonal (anti-shadowing) rescattering contribution. As a result of this competition the near-threshold value of $Tr_A(\Psi')$ first decreases significantly with $Q^2$, followed by the $J/\Psi$-like behaviour at larger $Q^2$ (Fig.2). For the both $J/\Psi$ and $\Psi'$ the near threshold behaviour of $Tr_A$ is partially due to the kinematical rise of the threshold energy with $Q^2$. In the bottomonium the initial size $\rho_Q$ compared to the bottomonium radius is relatively smaller, than in the charmonium, and $Tr_A(\Upsilon')$ exhibits a monotonous $Q^2$ and $\nu$ dependence (Fig.2).

With the rising energy the destructive interference of transitions (14) and (15) leads to a rapid rise of the nuclear transparency $Tr_A$ and the onset of the antishadowing of the $\Psi'$ in the photoproduction limit. At the larger $Q^2$ the smaller $\rho$ are scanned, and similar rise of $Tr_A$ with energy ends up in the shadowing region.

In the photoproduction of the light vector mesons $l_f \sim l_c$ and the more elaborated technique like the effective diffraction operator technique developed in [17] is called upon. This will be a subject of the further investigation. The effect on the $Q^2$-dependence is not significant, though, and the cited results for the photoproduction of the $\rho^0$-mesons, Fig.4, were obtained still assuming $l_c > l_f$. At large $Q^2$, when for all vector mesons $l_f \gg l_c$, the nuclear transparency for $\rho^0$, $J/\Psi$, $\Upsilon$ exhibits a similar $Q^2$- and $\nu$-dependence.

**Coherence effects and scaling law for FSI.** According to eq.10, at large $Q^2$ the FSI effect scales as $A^{1/3}/Q^2$. However, this scaling law, suggested in [18], is valid only at the asymptotic energy [16].

Indeed, at moderate energy, the coherence constraint (16) limits the number of the interfering states $N_{e\bar{f}}$. To a crude approximation, in this case the effective attenuation will be controlled by not a small size $\rho_Q$ of the initial $q\bar{q}$ pair, but rather by the least possible size $\rho_{min}$ of the wave packet, which can be constructed on a truncated basis of $N_{eff}$ conspiro states [17]. (Above we already have encountered a strong effect of the coherence on the energy-dependence of $Tr_A$ at fixed $Q^2$.) Only in the very high energy limit, when $N_{eff}$ is not bounded from above, $\rho_{min} \sim \rho_Q$. The constraint (16) is less stringent for the heavy quarkonia, as the level splitting $\Delta m \ll 2m_q$, and is more noticeable for the light vector mesons, as here the $m^2$-splitting between higher excitations, which enters the coherence condition (16), rises rapidly with the mass. In Fig.5 we present our results for the $Q^2$ dependence of $1-Tr_A(\rho^0)$ at fixed energy $\nu$: it is much weaker than $\propto 1/(Q^2 + m_{\rho}^2)$, which is an appropriate variable in biew of eq.(7). At asymptotic energy the scaling law (11) works for all vector mesons, Fig.5. A somewhat late onset of the scaling law (11) for the $\rho^0$ can be understood in terms of the overcompensation scenario in the $\rho^0, \rho'$ photoproduction, as one needs a relatively larger $Q^2$ to make the node effect negligible.
Conclusions: We have shown that QCD observables of CT experiments correspond to scanning the non-perturbative hadronic wave functions with the $Q^2$-dependent scanning radius $\rho_Q$. The strength of the final state interaction in the exclusive virtual photoproduction of vector mesons is shown to depend strongly on the nodal structure of wave functions. At large energy and $Q^2$ we predict the universal pattern of shadowing for all the photoproduced vector mesons. The node effect in the virtual photoproduction can be used to identify the radial excitation states. The predicted $Q^2$ and energy dependence of virtual photoproduction on nuclei can easily be tested at SLAC and planned European Electron Facility. The dedicated experiments on virtual photoproduction of vector mesons deserve special attention, since theoretical predictions of CT signal are much more reliable, than in $(e, e'p)$ or $(p, p'p)$ reactions (for the recent review see [19]).

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References

[1] L.D.Landau and I.Ya.Pomeranchuk, ZhETF 24 (1953) 505; E.L.Feinberg and I.Ya.Pomeranchuk, Doklady AN SSSR 93 (1953) 439; I.Ya.Pomeranchuk, Doklady AN SSSR 96 (1954) 265; 96 (1954) 481; E.L.Feinberg and I.Ya.Pomeranchuk, Nuovo Cim. Suppl. 4 (1956) 652;

[2] V.N.Gribov, Sov.Phys.JETP 29 (1969) 483; 30 (1970) 709.

[3] N.N.Nikolaev and B.G.Zakharov, Z.f.Phys. C49 (1991) 607; C53 (1992) 331.

[4] A.H.Mueller, in: Proceedings of the XVII Rencontre de Moriond, Les Arcs, France. Ed. Tranh Thanh Van, Editions Frontieres, Gif-sur-Yvette, 1982, p.13.

[5] S.J.Brodsky, in: Proceedings of the XIII International Symposium on Multiparticile Dynamics, Volendam, Netherlands. Eds. E.W.Kittel, W.Metzger and A.Stergion, World Scientific, Singapore, 1982, p. 963.

[6] A.B.Zamolodchikov, B.Z.Kopeliovich and L.I.Lapidus, JETP Lett. 33, (1981) 595 33 (1981) 612.

[7] G.Bertsch, S.J.Brodsy, A.S.Goldhaber and J.R.Gunion, Phys.Rev.Lett. 47 (1981) 267.

[8] B.Z.Kopeliovich and B.G.Zakharov: Phys.Rev. D44 (1991) 3466.

[9] O. Benhar, B.Z. Kopeliovich, Ch. Mariotti, N.N. Nikolaev and B.G. Zakharov, Phys.Rev.Lett. 69 (1992) 1156.

[10] T.H.Bauer et al, Rev.Mod.Phys. 50 (1978) 261.

[11] N.N.Nikolaev, Quantum Mechanics of Color Transparency. Comments on Nuclear and Particle Physics (1992), in press. N.N.Nikolaev, Color Transparency: Facts and Fancy. IJMPE: Reports on Nuclear Physics (1992), in press.

[12] A.Donnachie and H.Mirzaie, Z.Phys. C33 (1987) 407; A.Donnachie and A.B.Clegg, Z.Phys. C34 (1987) 257; C40 (1980) 313; C42 (1989) 663; C45 (1990) 677.

[13] S.J.Brodsky and A.H.Mueller: Phys.Lett. B206 (1988) 685.

[14] NMC Collaboration, C. Mariotti, Invited Talk at Europhysics High-Energy and Lepton-Photon Conference, Geneva, August 1991, to be published in Proceedings;

[15] V.N.Gribov, Sov.Phys.JETP 29 (1969) 483; 30 (1970) 709.

[16] V.A.Karmanov and L.A.Kondratyuk, JETP Letters 18 (1973) 266.

[17] N.N.Nikolaev, A.Szczeurek, J.Speth, J.Wambach, B.G.Zakharov and B.R.Zoller, Jülich preprint KFA-IKP(Th)-1992-16 (1992), submitted to Nucl.Phys. A, and paper in preparation.
[18] J.P. Ralson and B. Pire, *Phys.Rev.Lett.* **65** (1990) 2343.

[19] N.N. Nikolaev, *High Energy Nuclear Reactions in QCD: Color Transparency Aspects*. Lecture course at RCNP Kikuchi School on Spin Physics at Intermediate Energies, 16-19 November 1992, Osaka University, to be published in Proceedings. *RCNP preprint 051*, December 1992.
Figure captions:

Fig.1. - The qualitative pattern of the $Q^2$-dependent scanning of the wave functions of the ground state $V$ and the radial excitation $V'$ of the vector meson. The scanning distributions $\sigma(\rho)\Psi_{\gamma^*}(\rho)$ shown by the solid and dashed curve have the scanning radii $\rho_Q$ differing by a factor 3. All wave functions are in arbitrary units.

Fig.2. - The predicted $Q^2$ and $\nu$-dependence of the nuclear transparency in the virtual photoproduction of the heavy quarkonia. The qualitative pattern is the same from the light to heavy nuclei.

Fig.3. - The predicted $Q^2$-dependence of the nuclear transparency in the virtual photoproduction of the $\Psi'$ (the undercompensation scenario) and $\rho'$ (the overcompensation scenario) mesons.

Fig.4. - The predicted $Q^2$ and $\nu$-dependence of the nuclear transparency in the virtual photoproduction of the $\rho^0$-mesons.

Fig.5. - Test of the scaling law (10) in the virtual photoproduction of the $J/\Psi$ and $\rho^0$ mesons:

The right box - the $\rho^0$ production at fixed energy. The dashed straight line corresponds to the $1/(Q^2 + m_{\rho}^2)$ dependence.

The left box - at the asymptotic energy, the dotted straight line corresponds to the $1/(Q^2 + m_{J/\Psi}^2)$ dependence.