Color transparency
after the NE18 and E665 experiments:
Outlook and perspectives at CEBAF.*

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ABSTRACT

CEBAF is a high-luminosity factory of virtual photons with variable virtuality \(Q^2\) and transverse size. This makes CEBAF, in particular after the energy upgrade to (8-12)GeV, an ideal facility for uncovering new phenomena, and opening new windows, at the interface of the perturbative and nonperturbative QCD. We discuss color transparency as the case for a broad program on electroproduction of vector mesons \(\rho^0\), \(\omega^0\), \(\phi^0\) and their radial excitations \(\rho', \omega', \phi'\) at CEBAF. We also comment on the second generation of experiments on color transparency in \(^4\text{He}(e,e'p)\) scattering, which are also feasible at CEBAF.

In 1994, we can make more reliable projections into future because our understanding of the onset of color transparency has greatly been augmented by two experiments completed in 1993:

i) no effect of CT was seen in the SLAC NE18 experiment on \(A(e,e'p)\) scattering at virtualities of the exchanged photon \(Q^2 \lesssim 7\) GeV\(^2\),

ii) strong signal of CT was observed in the FNAL E665 experiment on exclusive \(\rho^0\)-meson production in deep inelastic scattering in the same range of \(Q^2\).

We discuss the impact of these observations on the CEBAF experimental program. We argue they both are good news, both were anticipated theoretically, and both rule in the correct QCD mechanism of the onset of CT.

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1 Introduction

The fundamental prediction of QCD is that the quark configurations with small transverse size $\vec{r}$ have small interaction cross section [1], which was dubbed color transparency (CT) [2]. Looking for CT is long discussed as the case for the high-luminosity, high-duty cycle, (10-20) GeV electron facility, which is well documented in the ELFE project [3] (ELFE = European Laboratory for Electrons). In the meantime, the good news from CEBAF is a possibility of the (8-12)GeV upgrade, which opens exciting possibilities of doing CT physics at CEBAF.

CEBAF is a high-luminosity factory of virtual photons. Higher energy means a higher virtuality $Q^2$ of photons, and higher $Q^2$ means that smaller sizes are becoming accessible reaching eventually into the perturbative QCD region. Higher energy also means longer lifetime of these small-size states. However, from the very start, we must emphasize that as far as CT physics is concerned, the purely perturbative region lies well beyond the kinematical range of CEBAF experiments on exclusive processes, even after the (8-12)GeV energy upgrade. Testing the purely perturbative QCD to few decimal places is a task of inclusive experiments at superhigh energy facilities like LEP or HERA. Even at LEP and HERA, the predictive power of the purely perturbative QCD rapidly deteriorates when the exclusive processes are considered. The real task of the CEBAF experiments is to uncover new QCD phenomena in exclusive reactions at the interface of the perturbative and nonperturbative QCD. Very conservative conclusion of this overview is that the energy-upgraded CEBAF shall do the job.

Before jumping into conclusions on the feasibility of CT physics at CEBAF, one must recall and critically summarize the results of the two CT experiments completed in 1993:

- The $A(e, e'p)$ reaction on the $D$, $C$, $Fe$ and $Au$ targets was studied by the SLAC NE18 collaboration with the negative result: no CT effects are seen at $Q^2 \leq 7$ GeV$^2$ [4].

- The FNAL E665 experiment [5] on exclusive production of the $\rho^0$ mesons in deep inelastic scattering of muons on nuclei produced a solid evidence for CT in precisely the same range of $Q^2$ as explored in the NE18 experiment.

The early history of CT focused on the quasielastic $A(e, e'p)$ scattering of electrons on nuclei. A number of predictions of precocious CT at low $Q^2$ were published (for the review and references see [6]), and the failure to confirm this precocious CT in the NE18 experiment considerably dampened the whole subject of CT. Fortunately, more consistent treatment rather predicted a very slow onset of CT in the $A(e, e'p)$ scattering [7,8]. As a matter of fact, the NE18 results do perfectly confirm the correct theory and rule in the mechanism of CT, which is alive and well, and we can joyfully recite Mark Twain’s telegram to the Associated Press: "The reports of my death were an exaggeration".

The parallel development was a theory of CT in (virtual) photoproduction of vector mesons $\gamma^*N \rightarrow VN$. From the theoretical point of view, this is a much cleaner case, with a well understood shrinkage of the transverse size of the virtual photon with the increase of $Q^2$ [9,10]. The prediction [9], not a postdiction, of the precocious onset of CT was confirmed by the FNAL E665 experiment [5], which put the CT physics in the right perspective.

The strong point which we wish to make in this review is that after the energy upgrade, CEBAF experiments on exclusive electroproduction of vector mesons can significantly contribute to our understanding of the onset of CT. Furthermore, the experiments on production of the radially excited vector mesons will open an entirely new window not only on the
mechanism of CT, but also on the quark structure and poorly known spectroscopy of radial excitations. In the $A(e,e'p)$ sector, we comment on the potential of experiments on the $^4He$ target, in which the onset of CT is sooner than for any other nucleus owing to the small size of the $^4He$ nucleus.

In this contribution to the Workshop on CEBAF at Higher Energies we concentrate on the recent experimental and theoretical developments, for the earlier reviews on the subject see [3,11-14].

2 CT in exclusive production of vector mesons

2.1 CT and dipole cross section

In order to be quantitative, let us set up the theoretical framework, which is the lightcone dipole-cross section representation [15]. Mesons can be viewed as color dipoles. The distribution of the transverse size $\vec{r}$ of color dipoles in the meson is given be the $q\bar{q}$ wave function $\Psi(z, \vec{r})$, where $z$ is the fraction of meson’s momentum carried by the quark. This mixed $(z, \vec{r})$ lightcone representation is custom tailored for description of CT. By the Lorentz-dilation, in the high-energy scattering the dipole size $\vec{r}$ becomes as good a conserved quantum number as an angular momentum.

The fundamental quantity which describes all the scattering processes, is the dipole cross section $\sigma(\nu, r)$ for interaction of the color dipole of size $r$ with the target nucleon. Of course, apart from the $q\bar{q}$ Fock state, the lightcone hadrons contain the higher $q\bar{q}g$... Fock states. The effect of gluons in the projectile dipole brings in the dependence of $\sigma(\nu, r)$ on energy $\nu$, which can be related to the gluon structure function $G(x, q^2)$ of the target nucleon [16,17]. Specifically, at small $r$, the dipole cross section is $\propto r^2$,

$$\sigma(\nu, r) = \frac{\pi^2}{3} r^2 \alpha_S(r) G(x, q^2)$$

where $q^2 = A_0/r^2$ with $A_0 \approx 10$ [18], and $x = q^2/2m_p\nu$. This CT property $\propto r^2$ derives from the decoupling of gluons from a very small color dipole. At large $r \gtrsim 1f$, the dipole cross section saturates because color forces do not propagate beyond the confinement radius. At high energy, the dipole cross section is a solution of the generalized BFKL equation [19] and describes, for instance, structure functions of deep inelastic scattering at HERA [20]. For any $q\bar{q}$ state, the total cross section of interaction with the target nucleon equals

$$\sigma_{tot} = \int_0^1 dz \int d^2\vec{r} \sigma(\nu, r)|\Psi(r, z)|^2.$$  

The dipole cross section is a universal function of $r$, the dependence on the process is only contained in the wave function of the $q\bar{q}$ state. Fig. 1 shows qualitatively, how the dipole cross section is probed in different processes:

- The total pion-nucleon cross section $\sigma_{tot}(\pi N)$ and the real photoabsorption cross section probe the region $r \gtrsim 1f$.
- The $\sigma_{tot}(J/\Psi N) \approx (5-6)\text{mb}$ comes from $r \sim R_{J/\Psi} \approx 0.45f$.
- The real photoproduction of the open charm probes the region of $r \sim 1/m_c \approx 0.15f$. 

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The proton structure function $F_2(x, Q^2)$ receives contribution from $1/\sqrt{Q^2} \lesssim r \lesssim 1f$, and CT property of the dipole cross section is crucial for the Bjorken scaling [12,15,17].

The scaling violations in $F_2(x, Q^2)$ are dominated by the contribution from $r \sim 1/\sqrt{Q^2}$.

In CT experiments one looks for a weak intranuclear final (and initial) state interactions, which will be the case if the nuclear production amplitude is dominated by the dipole cross section at small $r$ such that $\sigma(\nu, r)$ is much smaller than the free-nucleon cross section. Whether the particular exclusive reaction is selective of such a small $r$ or not, requires the case-by-case study.

### 2.2 How CT is tested in leptonproduction of vector mesons?

The exclusive (elastic) production $\gamma^* p \rightarrow V p$ ($V = \rho^0, \phi^0, J/\Psi,...$) is an ideal laboratory for testing CT ideas [9,10,12,21-24]. The forward production amplitude equals

$$M = \langle V|\sigma(\nu, r)|\gamma^* \rangle = \int_0^1 dz \int d^2 \vec{r} \sigma(\nu, r) \Psi_V^*(r, z) \Psi_{\gamma^*}(r, z)$$

(3)

Here $\Psi_{\gamma^*}(r, z)$ is the wave function of the $q\bar{q}$ fluctuation of the photon, which at high energy $\nu$ is formed at a large distance (the coherence length)

$$l_c = 2\nu/(Q^2 + m_V^2)$$

(4)

in front of the target nucleon. The most important feature of $\Psi_{\gamma^*}(r, z)$ as derived in [15] is an exponential decrease at large size

$$\Psi_{\gamma^*}(r, z) \propto \exp(-\varepsilon r),$$

(5)

where

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

(6)

and $m_q$ is the quark mass. Therefore, the amplitude (3) receives the dominant contribution from $r \sim r_S \approx 3/\varepsilon$. In the nonrelativistic quarkonium $z \approx 1/2$, $m_V \approx 2m_q$, and we conclude that the vector meson production amplitude probes the dipole cross section at the scanning radius $r_S$ given by [9,10,23]

$$r_S = \frac{6}{\sqrt{m_V^2 + Q^2}}.$$  

(7)

The scanning phenomenon is qualitatively illustrated in Fig. 2. If the scanning radius $r_S \lesssim R_V$, where $R_V$ is the radius of the vector meson, then for the (T) transverse and (L) longitudinal mesons

$$M_T \propto r_S^2 \sigma(r_S) \propto \frac{1}{(Q^2 + m_V^2)^2}$$

(8)

$$M_L \approx \sqrt{Q^2/m_V} M_T \propto \sqrt{Q^2/m_V} \frac{1}{(Q^2 + m_V^2)^2}$$

(9)

This prediction [10] of the dominance of the longitudinal cross section agrees with the E665 [5] and the NMC [25] data. Compared to the vector dominance model, the predicted amplitudes contain extra $(Q^2 + m_V^2)$ in the denominator, the origin of which is precisely in CT property $\sigma(r_S) \propto r_S^2$. 

The scanning radius $r_S$ decreases, and the virtual photon shrinks, with increasing $Q^2$. Notice, however, the large numerical factor in Eq. (7), for which the onset of small-size dominance requires very large $Q^2$. Remarkably, this large factor derives precisely from CT property of the dipole cross section. Because of this large scanning radius, the vector meson production probes the gluon structure function $G(x, q^2)$ at $q^2 \sim (10-20)\text{GeV}^2$ studied so far.

Because of a large scanning radius, the simple nonrelativistic approximation remains viable in quite a broad range of $Q^2$, and the production amplitude can be calculated using the constituent quark wave functions of vector mesons. For the same reason of large $r_S$, the asymptotic predictions \[ \sigma_T \propto 1/Q^8 \text{ and } \sigma_L \propto 1/Q^6, \] can not readily be tested at $Q^2 \lesssim (10-20)\text{GeV}^2$ studied so far.

In the conventional quark model, the radii of the $\rho$ and $\pi$ mesons are about identical. Once the radius of the $\rho^0$ is specified, further predictions \cite{23,27} for the $\rho^0$ production cross section do not contain any adjustable parameters. They are presented in Fig. 3. We use the low-energy dipole cross section of Ref. \cite{15}. Shown is the combination of the longitudinal and transverse cross sections as measured by the NMC collaboration. The agreement with the recent NMC data \cite{25} is excellent. This agreement in a broad range of $Q^2$ shows we have a good understanding of the dipole cross section from the hadronic scale $r \sim (1-2)f$ down to the smallest dipole size $r \sim 0.3f$ achieved in the NMC experiment at $Q^2 \sim 20\text{GeV}^2$. Over this range of radii $r$, the dipole cross section drops by approximately one order in magnitude, and this agreement of the total production rate with experiment is by itself a very important test of CT. Similar calculations \cite{23} give an excellent description of the $J/\Psi$ production and of the real photoproduction of the $\rho^0$ at HERA.

2.3 CT in exclusive production on nuclei

Having tested CT property of the dipole cross section in the production on free nucleons, now we turn our attention to the production on nuclei. The $q\bar{q}$ pair produced on the target nucleon, recombines into the observed vector meson with the recombination (formation) length

\[ l_f = \frac{\nu}{m_V\Delta m}, \]  

where $\Delta m$ is the typical level splitting in the quarkonium. At high energy $\nu$ both $l_c$ and $l_f$ exceed the radius $R_A$ of the target nucleus, which greatly simplifies the theoretical analysis. The E665 data \cite{5} correspond to this situation. In this high-energy limit, nuclear transparency for the incoherent (quasielastic) production on a nucleus $\gamma^* A \rightarrow V A^*$ equals \cite{12,22} (we suppress the energy-dependence of $\sigma(\nu, r)$)

\[ T_A = \frac{\sigma_A}{A\sigma_p} = \frac{1}{A} \int d^2\vec{b}T(b) \langle V|\sigma(r)\exp\left[-\frac{1}{2}\sigma(r)T(b)\right]|\gamma^*\rangle^2 \]  

\[ = 1 - \Sigma_V \frac{1}{A} \int d^2\vec{b}T(b)^2 + ... , \]

where

\[ T(b) = \int dzn_A(b, z) \]
is the optical thickness of a nucleus at the impact parameter \(b\) and \(n_A(b, z)\) is the nuclear matter density. In the quasielastic production one sums over all excitations and breakup of the target nucleus. The total cross section of the coherent (elastic) production \(\gamma^* A \rightarrow V A\), when the target nucleus remains in the ground state, equals [22]

\[
\sigma_{coh}(V A) = 4 \int d^2\bar{b} \left| \langle V|1 - \exp \left[ -\frac{1}{2}\sigma(r)T(b) \right] |\gamma^*\rangle \right|^2 \\
= 16\pi \frac{d\sigma(\gamma^* N \rightarrow VN)}{dt} \left|_{t=0} \int d^2\bar{b}T(b)^2 \left[ 1 - \frac{1}{2}\Sigma_V T(b) + ... \right] \right.
\]

(14)

In Eqs. (12,14) we have explicitly shown the leading terms of final state interaction (FSI). The strength of FSI is measured by the observable [12]

\[
\Sigma_V = \frac{\langle V|\sigma(r)^2|\gamma^*\rangle}{\langle V|\sigma(r)|\gamma^*\rangle}.
\]

(15)

Evidently, the matrix element in the numerator of (14) is dominated by

\[
r \sim r_{FSI} = \frac{5}{3}r_S = 10\sqrt{Q^2 + m_V^2},
\]

(16)

which gives an estimate [10,23]

\[
\Sigma_V \approx \sigma(r_{FSI})
\]

(17)

CT and/or weak FSI set in when \(r_{FSI} \ll R_V\), i.e., when \(\Sigma_V \ll \sigma_{tot}(VN)\). In this regime of CT, the \(\Sigma_V\) is insensitive to the wave function of the vector meson, so that predictions of FSI effects are less model independent than predictions for the total production cross section. The large value of \(r_{FSI}\) Eq. (16) implies that FSI only slowly vanishes with the increase of \(Q^2\), and this slow onset of CT is driven by the very mechanism of CT.

### 2.4 The E665 experiment [5]: the decisive proof of CT

Our predictions [9,10] for nuclear effects in the coherent and incoherent exclusive \(\rho^0\) production are compared with the E665 data in Figs. 4-7. In Fig. 4 we show nuclear transparency for the incoherent production. Nuclear attenuation is very strong at small \(Q^2\) and gradually decreases with \(Q^2\). The effect is particularly dramatic for the heavy nuclei (\(Ca, Pb\), and leaves no doubts the E665 observed the onset of CT. The predicted \(Q^2\) dependence of nuclear transparency for the forward coherent production on nuclei \(T_A^{coh}(coh) = (d\sigma_A^{coh}/A^2d\sigma_N)|_{t=0}\) is shown in Fig. 5. We predict a rise of \(T_A^{coh}\) with \(Q^2\) towards \(T_A^{coh} = 1\) for vanishing FSI. The predicted \(Q^2\) dependence of the coherent production cross section relative to the cross section for the carbon nucleus is presented in Fig. 6. In the regime of vanishing FSI,

\[
R_{coh}^{(CT)}(A/C) = \frac{12\sigma_A}{A\sigma_C} \approx \frac{AR_{ch}(C)^2}{12AR_{ch}(A)^2},
\]

(18)

which gives \(R_{coh}^{(CT)}(Ca/C) = 1.56\) and \(R_{coh}^{(CT)}(Pb/C) = 3.25\). Here \(R_{ch}(A)\) is the charge radius of a nucleus. The observed growth of the \(Pb/C\) ratio with increasing \(Q^2\) gives a solid evidence for the onset of CT.

The (approximate) \(A^\alpha\) parametrization is a convenient short-hand representation of the \(A\)-dependence of nuclear cross sections, although the so-defined exponent \(\alpha\) slightly depends
on the range of the mass number $A$ used in the fit. Then, Eq. (12) predicts that $\alpha_{inc}(Q^2)$ tends to 1 from below, as $Q^2$ increases. In the limit of vanishing FSI Eqs. (14,18) predict $\sigma_{coh} \sim A^{4/3}$, so that $\alpha_{coh}(Q^2)$ tends to $\approx \frac{4}{3}$ from below as $Q^2$ increases (more accurate analysis shows that the no-FSI cross section in the $C-Pb$ range of nuclei has the exponent $\alpha_{coh} \approx 1.39$). The agreement between the theory and the E665 fits is good (Fig. 7). Both the $\alpha_{coh}(Q^2)$ and $\alpha_{inc}(Q^2)$ rise with $Q^2$, which is still another way of stating that the E665 data confirm the onset of CT. Even at the highest $Q^2 \sim 10$ GeV$^2$ of the E665 experiment, the residual nuclear attenuation is still strong, confirming our prediction of large value of $r_{FSI}$ Eq. (16).

### 2.5 Scaling properties of nuclear attenuation: electroproduction of the $\rho^0$ vs. the real photoproduction of the $J/\Psi$

Eq. (15) in conjunction with Eq. (7) suggests that nuclear attenuation approximately scales with $(Q^2 + m_V^2)$. Specifically, we predict identical nuclear transparency in the real photoproduction ($Q^2 = 0$) of the $J/\Psi$ and electroproduction of the $\rho^0$ at $Q^2 \approx m_{J/\Psi}^2 - m_{\rho^0}^2$, which nicely agrees with the experiment:

For the $\rho^0$ production at $Q^2 = 7$ GeV$^2$ the E665 experiment gives $T_{pb}/T_C = 0.6 \pm 0.25$. This can be compared with the NMC result $T_{Sn}/T_C = 0.7 \pm 0.1$ for the real photoproduction of $J/\Psi$ in the similar energy range [28]. Similar scaling law holds for the coherent production of the $\rho^0$ and the $J/\Psi$. In the regime of vanishing FSI Eq. (12) gives [10,22] $R_{coh}(Sn/C) = 2.76$, $R_{coh}(Fe/Be) = 2.82$, $R_{coh}(Pb/Be) = 4.79$. The experimental data on the real photoproduction of $J/\Psi$ give $R_{coh}(Sn/C) = 2.15 \pm 0.10$ in the NMC experiment [29] and $R_{coh}(Fe/Be) = 2.28 \pm 0.32$, $R_{coh}(Pb/Be) = 3.47 \pm 0.50$ in the Fermilab E691 experiment [30]. These data have successfully been described [22] in the discussed framework using the dipole cross section [15]. In all cases the observed $\sim 30\%$ departure of the observed ratios for the $J/\Psi$ from predictions for vanishing FSI is of the same magnitude as in the highest $Q^2$ bin of the E665 data on the $\rho^0$ production (Fig. 6). This scaling relationship between production of different vector mesons [10] provides a very important cross check of the mechanism of CT.

Regarding the statistical accuracy of the data, the potential of the FNAL and CERN muon scattering experiments is nearly exhausted. Now we shall discuss what new can be done at CEBAF assuming the energy range $\nu \lesssim (8-10)$ GeV and $Q^2 \lesssim (4-6)$ GeV$^2$.

### 3 What next: Vector mesons at CEBAF

#### 3.1 Two scales: the coherence length and the formation length

The coherence length $l_c$ and the formation length $l_f$ control the two different aspects of nuclear attenuation. The formation length tells how rapidly the electroproduced $q\bar{q}$ state evolves into the observed hadron. For the $\rho^0$ production,

$$l_f \sim 0.4f \cdot \left( \frac{\nu}{1\text{GeV}} \right). \quad (19)$$

In the low energy limit of $l_f \ll R_A$ we have an instantaneous formation of the final-state hadron, which then attenuates with the free-nucleon cross section and CT effects are absent.
At high energies $l_f > R_A$, the formation of the observed hadron takes place behind the nucleus, and CT becomes possible. For the observation of the fully developed CT one needs
\[ \nu \gtrsim (3 - 4) \cdot A^{1/3} \text{GeV}. \] (20)

CT effects already start showing up, though, if $l_f \gtrsim l_{int}$, where the interaction length, or the mean free path, equals
\[ l_{int} = \frac{1}{n_A \sigma_{tot}(VN)} \approx 2f \cdot \left( \frac{30mb}{\sigma_{tot}(VN)} \right). \] (21)

Therefore, purely kinematically, CT effects are within the reach of CEBAF experiments after the 8 – 12GeV energy upgrade. Notice, that the formation length $l_f$ does not depend on the photon’s virtuality $Q^2$.

On the other hand, the coherence length $l_c$ tells at which distance from the absorption point the pointlike photon becomes the hadronlike $q \bar{q}$ pair. If $l_c \gtrsim R_A$, then the whole thickness of the nucleus contributes to attenuation of the $q \bar{q}$ pair. Changing the virtuality of the photon $Q^2$ and/or reducing the photon’s energy $\nu$, one can make $l_c \ll R_A$. In this case, the incoming photon is absorbed approximately uniformly over the volume of the target nucleus, and attenuation of the produced $q \bar{q}$ pair will take place over half of the total thickness of the nucleus. In the practically interesting cases of $Q^2 \gg m_V^2$ for the light mesons, or in the real and virtual photoproduction of heavy quarkonia $J/\Psi, \Upsilon$ we have $l_c \ll l_f$. If $l_f \gtrsim R_A$, but $l_c \ll R_A$, then nuclear transparency equals
\[ T_A(l_c \gtrsim R_A) < T_A(l_c \ll R_A) \] (24)

The energy dependence of the nuclear transparency is given by an approximate interpolation formula
\[ T_A(\nu) \approx T_A(l_c << R_A) + G_A(\kappa)^2 [T_{RA}(l_c > R_A) - T_{RA}(l_c << R_A)] \] (25)

where $G_A(\kappa)$ is the charge form factor of the target nucleus and $\kappa$ is the longitudinal momentum transfer in the transition $\gamma^* N \rightarrow VN$:
\[ \kappa = \frac{1}{l_c} = \frac{Q^2 + m_V^2}{2\nu}. \] (26)

The predicted energy dependence of nuclear transparency [22] is in excellent agreement with the NMC data on the real photoproduction of the $J/\Psi$ [28] shown in Fig. 8. The NMC data rule out the widely discussed models of ”quantum” and ”classical” diffusion by Farrar, Frankfurt and Strikman [31], which has also predicted the precocious CT in the $A(e, e'p)$ scattering (for the review see [6]).
3.2 Quantum evolution and energy dependence of FSI.

If \( l_f < R_A \), the frozen-size approximation is no longer applicable, and spatial expansion of the \( q\bar{q} \) pair during its propagation inside a nucleus becomes important. The consistent path-integral description of the spatial expansion effects was developed in [21]. Interaction of the \( q\bar{q} \) pair with the nuclear matter can be described by the optical potential \( V_{opt}(r) \propto \sigma(r)n_A(b,z) \), which adds to the confining \( q\bar{q} \) potential \( U(r) \). The longitudinal distance \( z \) plays the role of time. The path integral technique provides a systematic procedure for calculating the evolution kernel \( K(r',r,t) \) in the combined potential \( U(r) + V_{opt}(r) \).

The particularly elegant exact solution for the evolution kernel is found for the harmonic oscillator \( U(r) \propto r^2 \) and \( \sigma(r) \propto r^2 \) [21].

The resulting predictions for the energy dependence of nuclear transparency \( T_A \) for the production of different vector mesons are shown in Figs. 9,10. The salient features of \( T_A \) are [9,21,22]:

- At small \( Q^2 \), nuclear transparency for the \( \rho^0 \) and the \( J/\Psi \) decreases with energy \( \nu \), starting at the value given by Eq. (22) and levelling off at the value given by Eq. (12). This decrease is due to the increase of the coherence length \( l_c \), discussed in section 3.1.

- At larger \( Q^2 \), the trend changes: \( T_A \) first increases with \( Q^2 \), and then decreases for the same reason of the rise of the coherence length \( l_c \). In agreement with Eqs. (25,26), the larger is \( Q^2 \), the higher is the energy \( \nu \) at which the levelling off of \( T_A \) takes place.

- For the \( \Upsilon \) production, nuclear transparency \( T_A \) starts increasing with energy at all values of \( Q^2 \), in close similarity to the \( J/\Psi \) production at large \( Q^2 \), and in agreement with our conclusion that nuclear attenuation scales with \( (Q^2 + m_V^2) \), see section 2.5. For instance, we predict \( T_A(J/\Psi, Q^2 = 100\text{GeV}^2) \approx T_A(\Upsilon, Q^2 = 0) \).

- The radial excitations \( \Psi' \), \( \Upsilon' \) have larger size, and larger free-nucleon cross section thereof, than the ground states \( J/\Psi \) and \( \Upsilon \), respectively. Nonetheless, the radial excitations are predicted to have weaker nuclear attenuation, which is a completely counterintuitive result.

- At last but not the least, for the \( \rho^0 \) production we predict very rich pattern of the \( \nu \) and \( Q^2 \) dependence precisely in the kinematical range of CEBAF. CT effects are large and can readily be observed at CEBAF. The \( \rho' \) production will be treated in section 3.7.

These properties of nuclear transparency can best be understood in terms of the interplay of CT with the node effect.

3.3 CT and the node effect: antishadowing phenomenon.

The wave function of the radial excitation \( V'(2S) \) has a node. For this reason, in the \( V' \) production amplitude there is the node effect - cancellations between the contributions from \( r \) below, and above, the node. The product \( \sigma(r)\Psi_{\gamma^*}(z,r) \) acts as a distribution, which probes the wave function of the \( V(1S) \) and \( V(2S) \) states at the scanning radius \( \sim r_S \) [9,10,12] and the node effect evidently depends on the scanning radius \( r_S \), see Fig. 2.

If the node effect is strong, even the slight variations of \( r_S \) lead to an anomalously rapid variation of the \( V'(2S) \) production amplitude, which must be contrasted to the smooth \( Q^2 \)
and $r_S$ dependence of the $V(1S)$ production amplitude. Evidently, the stronger is the node effect and the smaller is the $V'(2S)$ production amplitude, the higher is the sensitivity to the model for the $V'(2S)$ wave function, and in some cases only firm statement will be the fact of the strong suppression of the $V'(2S)$ production.

For the real photoproduction of the $\Psi'$, the calculations in [21] gave $\sigma(\gamma N \to \Psi' N)/\sigma(\gamma N \to J/\Psi N) = 0.17$, which is in excellent agreement with the NMC result $0.20 \pm 0.05 (\text{stat}) \pm 0.07 (\text{syst})$ for this ratio [29]. In this case the node effect is already rather strong for the fact that the scanning radius $r_S(Q^2 = 0)$ is rather close to the $J/\Psi$ radius $R_{J/\Psi}$. For the $\Upsilon'$, the scanning radius is substantially smaller than $R_{\Upsilon'}$, which is relatively large for the small strong coupling $\alpha_S(R_{\Upsilon'})$. For the light mesons, the scanning radius $r_S$ is larger and, at small $Q^2$, the node effect is much stronger, see section 3.7.

Because $r_{FSI}$ is larger than $r_S$, the node effect in the strength of FSI given by Eq. (13) becomes stronger. For the $J/\Psi$, one finds the overcompensation: $\langle V|\sigma(r)^2|\gamma\rangle < 0$, which leads to $\Sigma_V < 0$ and to the antishadowing phenomenon $T_A > 1$ shown in Fig. 10. For the $\Upsilon'$, we find the undercompensation: $\langle V|\sigma(r)^2|\gamma\rangle > 0$, which leads to $\Sigma_V > 0$ and to the shadowing $T_A < 1$. None the less, the node effect shows up: for the $\Upsilon'$ with its large radius, nuclear attenuation is weaker than for the $\Upsilon$. With increasing $Q^2$, when $r_{FSI} \ll R_{\Upsilon'}$, the node effect becomes negligible, and the $V(1S)$ and $V'(2S)$ states will have identical nuclear attenuation, see Fig. 10.

### 3.4 The interplay of CT, of the node effect and of quantum evolution

Of course, the $\Psi'$ has a larger radius and larger free-nucleon cross section $\sigma_{tot}(\Psi' N) \sim (2.5-3)\sigma_{tot}(J/\Psi N)$. How come, then, that in the real photoproduction $\Sigma_{\Psi'} < 0$ and we find the antishadowing of the strongly interacting $\Psi'$ alongside with shadowing for the $J/\Psi$?

Although the above derivation of antishadowing was (deceptively) simple, still another look at antishadowing and the variation of energy dependence of nuclear transparency with $Q^2$ is in order [9]. Let us consider for simplicity the $J/\Psi$, $\Psi'$ system. The numerator of the strength of FSI $\Sigma_V$ can be expanded in terms of the complete set of intermediate states $|i\rangle$ of charmonium

$$\langle V|\sigma(r)^2|\gamma^*\rangle = \sum_i \langle V|\sigma(r)|\gamma_i\rangle \langle \gamma_i|\sigma(r)|\gamma^*\rangle.$$  \hspace{1cm} (27)

In terms of this expansion, antishadowing of the photoproduction of the $\Psi'$ comes from the destructive interference of the two dominant intermediate states: the direct, VDM-like rescattering

$$\gamma^* \to \Psi' \to \Psi'$$  \hspace{1cm} (28)

and the off-diagonal rescattering

$$\gamma^* \to J/\Psi \to \Psi'$$  \hspace{1cm} (29)

(there is a small contribution from other intermediate states too).

Then, for the $\Psi'$ production, the strength of FSI is given by

$$\Sigma_{\Psi'} = \sigma_{tot}(\Psi' N) + M(J/\Psi N \to \Psi' N) \cdot \frac{M(\gamma^* N \to J/\Psi N)}{M(\gamma^* N \to \Psi' N)}.$$  \hspace{1cm} (30)
Because of the interplay of CT and the node effect, we have \(M(\gamma^* N \to J/\Psi N)/M(\gamma^* N \to \Psi' N) \gg 1\). For the same interplay of CT and the node effect, there is an overcompensation in the amplitude of the off-diagonal transition

\[
M(J/\Psi N \to \Psi' N) < 0 ,
\] (31)

and numerically this amplitude is not very small, \(M(J/\Psi N \to \Psi' N) \sim -\sigma_{tot}(J/\Psi N)\). Consequently, the second, negative valued, off-diagonal term in \(\Sigma_{\Psi'}\) takes over the \(\sigma_{tot}(\Psi' N)\) (the higher intermediate states also slightly contribute to the antishadowing effect).

What happens with increasing \(Q^2\) is very simple: The scanning radius \(r_S\) decreases with \(Q^2\) and the node effect become weaker, the ratio of amplitudes \(M(\gamma^* N \to J/\Psi N)/M(\gamma^* N \to \Psi' N)\) decreases with \(Q^2\) and tends to approximately unity, whereas \(\sigma_{tot}(\Psi' N)\) and \(M(J/\Psi N \to \Psi' N)\) do not depend on \(Q^2\). As a result, the off-diagonal contribution in Eq. (30) becomes small, and at large \(Q^2\) the antishadowing of \(\Psi'\) changes to the shadowing.

Similarly, for the \(J/\Psi\) production

\[
\Sigma_{J/\Psi} = \sigma_{tot}(J/\Psi N) + M(J/\Psi N \to \Psi' N) \cdot \frac{M(\gamma^* N \to \Psi' N)}{M(\gamma^* N \to J/\Psi N)}. \] (32)

In this case, for the real photoproduction the off-diagonal contribution is small because of \(M(\gamma^* N \to J/\Psi N)/M(\gamma^* N \to \Psi' N) \gg 1\), which is a consequence of CT. Therefore, for this numerical reason, for the real photoproduction we find

\[
\Sigma_V \approx \sigma_{tot}(J/\Psi N) \] (33)

which explains why the predicted nuclear attenuation of the \(J/\Psi\) is marginaly similar [21,22] to the vector dominance model estimates. With increasing \(Q^2\), the off-diagonal term in \(\Sigma_{\Psi'}\) increases and starts cancelling the term \(\sigma_{tot}(J/\Psi N)\), thus depleting \(\Sigma_{J/\Psi}\) and leading to CT effect.

Consider now the \(J/\Psi\) and \(\Psi'\) photoproduction at a finite energy. The intermediate \(J/\Psi\) and \(\Psi'\) have different masses \(m_1\) and \(m_2\) and propagate with momenta which differ by

\[
\kappa_{21} = \frac{m_2^2 - m_1^2}{2\nu} \approx \frac{1}{l_f} \] (34)

If the production and rescattering points are a distance \(\Delta z = z_2 - z_1\) apart, then the off-diagonal contribution to the rescattering amplitude acquires the phase factor \(\exp[i\kappa_{21}(z_2 - z_1)]\) with respect to the elastic rescattering contribution [32,33]. Upon the integration over \(z_{1,2}\), the effect of this phase factor is

\[
\langle \exp[i\kappa_{21}(z_2 - z_1)] \rangle = G_A^2(\kappa_{21}) , \] (35)

and the off-diagonal contributions to \(\Sigma_V\) will enter with the suppression factor \(G_A(\kappa_{11})^2\) (for the detailed analysis see [7,8,9,13])

\[
\Sigma_{\Psi'} = \sigma_{tot}(\Psi' N) + M(J/\Psi N \to \Psi' N) \cdot \frac{M(\gamma^* N \to J/\Psi N)}{M(\gamma^* N \to \Psi' N)} G_A^2(\kappa_{12}) . \] (36)

\[
\Sigma_{J/\Psi} = \sigma_{tot}(J/\Psi N) + M(J/\Psi N \to \Psi' N) \cdot \frac{M(\gamma^* N \to \Psi' N)}{M(\gamma^* N \to J/\Psi N)} G_A^2(\kappa_{12}) . \] (37)
Evidently, at low energy such that $\kappa_{12}R_A \gtrsim 1$, i.e., at $l_f \ll R_A$, the nuclear form factor vanishes $G_A(\kappa_{12})^2 \ll 1$. Only the diagonal contributions to $\Sigma_V$ survive, and CT effects which come from the off-diagonal terms in $\Sigma_V$, vanish at low energy. For instance, for the $\Psi'$ one starts with the shadowing $T_A < 1$, which with increasing energy and the opening of the off-diagonal channels, changes to the antishadowing. For the $J/\Psi$ and $\rho^0$, at small energy there is a competition of CT effect which rises with energy, and of the effect of growth of the coherence length. The latter takes over at small $Q^2$, whereas at larger $Q^2$ the opening of the off-diagonal transitions leads to a rapid near-threshold rise of nuclear transparency. For the $\rho^0$ production, nuclear transparency $T_A$ is a lively function of energy $\nu$ and $Q^2$ in precisely the kinematical region accessible at CEBAF.

3.5 Measuring the $J/\Psi$-nucleon cross section at CEBAF

The smallness of the off-diagonal rescattering in the real photoproduction of the $J/\Psi$, see Eqs. (33), (37), leads to an important prediction [9,22] that nuclear attenuation allows to evaluate $\sigma_{tot}(J/\Psi N)$ using the conventional VDM formulas. This suggestion was carried over in an analysis [34] of the data [29,30] on the coherent photoproduction on nuclei with the result $\sigma_{tot}(J/\Psi N) \sim (5-7)\text{mb}$. The CEBAF experiments will allow measurement of this cross section at low energy, although being very close to the threshold requires a good understanding of the Fermi-smearing effects.

For the $\Psi'$ photoproduction at low energies, the off-diagonal transitions (29) are non-negligible even close to the threshold, and for the $\Psi'$ the VDM prediction for nuclear shadowing breaks down completely: $T_A$ is always larger than the VDM prediction.

3.6 The coherence and formation lengths revisited: Vanishing nuclear shadowing in inclusive DIS at CEBAF coexists with lots of CT in exclusive production at CEBAF

Above we have repeatedly emphasized that the onset of CT is entirely controlled by the formation length $l_f$, which does not depend on $Q^2$. The onset of CT is quantified by Eqs. (35,37), in which the formation length enters via the nuclear form factor $G_A(\kappa_{12})$, where $\kappa_{12}$ is the longitudinal momentum transfer in the off-diagonal transition $V_1N \rightarrow V_2N$.

At large $Q^2$ we have $l_c \ll l_f$ and much larger longitudinal momentum transfer $\kappa$ Eq. (26) in the transition $\gamma^*N \rightarrow VN$. However, because $\kappa \gg \kappa_{12}$, this momentum transfer is approximately the same for all intermediate states. At $l_c \ll R_A$, the corresponding overall phase factor simply drops out from the incoherent production cross section. Neither does this phase factor affect $T_A$ dramatically in the transient regime of $l_c \sim R_A$, see Eq. (23).

In the coherent production on nuclei, the major effect of the momentum transfer $\kappa$ is that the nuclear production amplitude acquires the overall factor $G_A(\kappa)$, which significantly suppresses the coherent production amplitude but does not affect the nuclear attenuation properties.

As we discussed in section 2.2, the vector meson production on the free nucleon probes the gluon distribution in the target proton. Stretching the so-called factorization theorems, one can be tempted to conclude that nuclear attenuation in the production of vector mesons is given by the nuclear shadowing of gluon structure function [26]. Indeed, exclusive electroproduction of vector mesons is the typical diffraction dissociation of the photon $\gamma^*N \rightarrow XN$, and virtual diffractive transitions $\gamma^* \rightarrow X \rightarrow$ in the nuclear forward Compton scattering
amplitude are precisely the source of nuclear shadowing [15]. But, subtle is the nuclear shadowing!

The mere definition of the shadowed nuclear parton distributions is only useful provided that the shadowing term by itself satisfies the conventional evolution equations. To a certain approximation, this is the case [17]. However, there are no theorems on the universality of these shadowing corrections in all hard scattering processes and shadowing corrections may defy the factorization theorems [16]. Here we present simple arguments, essentially of kinematical origin, why the factorization theorems must be taken with the grain of salt.

The contribution of the intermediate state $X$ to the nuclear Compton scattering amplitude contains the excitation $\gamma^* N \rightarrow X N$ on one nucleon and the de-excitation $X N \rightarrow \gamma^* N$ on another nucleon. In both transitions there is a longitudinal momentum transfer $\kappa$ Eq. (20). Consequently, the corresponding contribution to nuclear shadowing in the structure function will enter with the suppression factor $G_A(\kappa)^2$. The onset of nuclear shadowing requires $l_c \gtrsim R_A$, so that the larger is $Q^2$, the higher energy is required for the onset of nuclear shadowing. In the opposite to that, nuclear attenuation and CT effects in the nuclear electroproduction of vector mesons only requires $l_f \gtrsim l_{int}, R_A$, and this condition does not depend on $Q^2$ and does not require $l_c \gtrsim R_A$. For instance, there will be no nuclear shadowing in the inclusive electroproduction on nuclei in the kinematical range of the CEBAF experiments, but lots of CT effects in the exclusive production of vector mesons on nuclei at CEBAF. We wish to emphasize this simple, but important, point in view of the opposite claims made to this effect by Frankfurt and Strikman at this Workshop.

### 3.7 Anomalous electroproduction of radial excitations $\rho', \phi'$: CEBAF’s new window at CT

For the light vector mesons, at small $Q^2$ the scanning radius $r_S \sim R_V$, and there is an exciting, and most likely, possibility of the overcompensation already in the free-nucleon production amplitude: $M = \langle V|\sigma(r)|\gamma^* \rangle < 0$. The $\rho'$ production on nuclei is indispensable for testing the node effect and $Q^2$ dependence of the scanning radius $r_S$, because nuclear attenuation gives still another handle on the scanning radius [9,13,27]. For the sake of simplicity, we discuss the quasielastic (incoherent) $\rho'$ production on nuclei assuming that $l_f \gtrsim R_A$. Extension to lower energies and to the coherent production is straightforward and the prediction [13,27] of the anomalous $Q^2$ and $A$ dependence persists in these cases too.

The $A$-dependence of the node effect comes from the nuclear attenuation $\exp[-\frac{1}{2}\sigma(r)T(b)]$ in the nuclear matrix element $M_A(T) = \langle V|\sigma(r)\exp[-\frac{1}{2}\sigma(r)T(b)]|\gamma^* \rangle$. The possibility of the $A$-dependent node effect in hadronic diffraction production on nuclei $hA \rightarrow h^*A$ was pointed out in [35].

Firstly, consider the $Q^2$ dependence of the $\rho'/\rho^0$ ratio on the free nucleon. Increasing $Q^2$ and decreasing the scanning radius $r_S$, one will bring the $\rho'$ production on the free nucleon to the exact node effect, and the $\rho'/\rho^0$ ratio takes the minimum value at a certain finite $Q^2$, see Fig. 11. Because of the $r$-dependence of the attenuation factor, in the nuclear amplitude the node effect will be incomplete. Consequently, as a function of $Q^2$, nuclear transparency $T_A$ will have a spike $T_A \gg 1$ at a finite value of $Q^2$ [9].

Secondly, consider the $\rho'$ production on nuclei at a fixed value of $Q^2$ such that the free nucleon amplitude is still in the overcompensation regime. Increasing $A$ and enhancing the importance of the attenuation factor $\exp[-\frac{1}{2}\sigma(r)T(b)]$, we shall bring the nuclear amplitude
to the nearly exact compensation regime. Therefore, the $\rho'/\rho^0$ production ratio, as well as nuclear transparency for the $\rho'$ production, will decrease with $A$ and take a minimum value at a certain finite $A$. With the further increase of $A$, the undercompensation regime takes over, and we encounter very counterintuitive situation: nuclear transparency for the $\rho'$ is larger for heavier, more strongly absorbing nuclei! This situation is illustrated in Fig. 12a and must be contrasted with a smooth and uneventful decrease of transparency for the $\rho^0$ production on heavy nuclei.

With the further increase of $Q^2$ one enters the pure undercompensation regime for all the targets. Nuclear undoing of the node effect enhances $M_A$ and nuclear transparency $T_A$, whereas the overall attenuation factor $\exp[-\frac{1}{2}\sigma(r)T(b)]$ decreases $M_A$ and $T_A$. Of these two competing effects, the former remains stronger and we find antishadowing of the $\rho'$ production in a broad range of $A$ and $Q^2$, see Figs. 12c-12e. Typically, we find a nuclear enhancement of the $\rho'/\rho^0$ production ratio on heavy targets by one order in the magnitude with respect to the free nucleon target. This makes leptoproduction on nuclei the $\rho'$ factory, and the $\rho'$ production experiments at CEBAF can contribute much to the poorly understood spectroscopy of the radially excited vector mesons. Only at a relatively large $Q^2 \gtrsim 2\text{GeV}^2$, the attenuation effect takes over, and nuclear transparency for the $\rho'$ production will start decreasing monotonically with $A$ (Fig. 12f). Still, this decrease is much weaker than for the $\rho^0$ meson. At very large $Q^2$, when the node effect disappears because of the small scanning radius $r_S$, nuclear transparency for the $\rho^0$ and the $\rho'$ production will become identical. This pattern repeats qualitatively the one studied for the $J/\Psi$ and $\Psi'$ mesons in sections 3.3 and 3.4.

The above presented results refer to the production of the transversely polarized $\rho^0$ and $\rho'$ mesons. Accurate separation of the transverse and longitudinal cross section can easily be done in the high statistics CEBAF experiments. Here we only wish to mention the interesting possibility that for the longitudinally polarized $\rho'$ mesons, the exact node effect is likely to take place at a value of $Q^2$ larger than for the transverse $\rho'$, so that polarization of the produced $\rho'$ can exhibit very rapid change with $Q^2$.

The numerical predictions are very sensitive to the position of the node in the wave function of the $\rho'(2S)$. It is quite possible that the dip of nuclear transparency $T_A$ will take place for targets much heavier than in Figs. 12a,12b, and disappearance of the node effect and the onset of the more conventional nuclear shadowing $T_A < 1$ for the $\rho'$ production like in Fig. 12f only will take place at much larger $Q^2$. Also, the possibility of the undercompensation at $Q^2 = 0$ can not be excluded. However, the strikingly different $A$-dependence of the incoherent $\rho^0$ and $\rho'$ production on nuclei persists in such a broad range of $Q^2$ and of the scanning radius $r_S$, that the existence of the phenomenon of anomalous $A$ and $Q^2$ dependence of the $\rho'$ production is not negotiable. It is a direct manifestation of the color-transparency driven $Q^2$ dependence of the scanning radius and, as such, it deserves a dedicated experimental study.

The case of the $\phi^0$ and $\phi'(1680)$ production was studied in [24] using the path integral technique [21] modified to include corrections for the departure from the $\propto r^2$ dependence of the dipole cross section $\sigma(r)$. Corrections to the $\phi^0$ production are negligible, corrections to the $\phi'$ production do not exceed 10%. Because of the smaller radius of the $\phi^0$, $\phi'$, the numerical predictions are more reliable than for the $\rho'$. On the free nucleon, we find undercompensation for the transverse $\phi'$ production. There is much similarity to the $J/\Psi$, $\Psi'$ system, but the node effect is stronger than for the charmonium. In the nuclear production of the $\phi'$, the antishadowing effect shown in Fig. 13 gradually builds up with increasing en-
enery, and numerically is much stronger than for the $\Psi'$. Nuclear attenuation of the $\phi'(1680)$ is weaker than for the $\phi^0$ starting already at low energy (Fig. 14). Because of the large mass and small radius of the $\phi^0$, nuclear transparency for the $\phi^0$ production only weakly depends on $Q^2$ at CEBAF, but for the $\phi'(1680)$ production the predicted $Q^2$ dependence is quite strong (Fig. 15) and can easily be measured at CEBAF. Here the larger $Q^2$ and smaller scanning radius $r_S$ predict increasing attenuation of the $\phi'(1680)$. However, for the lead target we expect the onset of increasing nuclear transparency at already moderate $Q^2 > 1 GeV^2$.

Few more comments about the possibilities of CEBAF are worth while. Because of the strong suppression of the $\rho'/\rho^0$ and $\phi'/\phi^0$ production ratio by the CT and node effects, the high luminosity of CEBAF is absolutely crucial for high-statistics experiments on the $\rho'$, $\phi'$ production. Notice, that the most interesting anomalies in the $A$ and $Q^2$ dependence take place near the minimum of the $\rho'$ production cross section. Furthermore, the observation of the $\rho'$ production requires detection of its 4-pion decays, and here one can take advantage of the CLAS multiparticle spectrometer available at CEBAF.

4 FSI and nuclear transparency in $A(e, e'p)$ scattering

4.1 Multiple-scattering expansion for the nuclear spectral function

We are interested in $A(e, e'p)$ scattering at large $Q^2 > (1-2)GeV^2$, when the struck proton has the kinetic energy $T_{kin} \approx Q^2/2m_p \approx (0.5-1)GeV$. Such a proton has the free nucleon total cross section $\sigma_{tot}(pN) \approx 40 mb$. The corresponding mean free path in the nuclear medium $l_{int} \approx 1.5f$ is short and implies strong FSI and strong nuclear attenuation of struck protons. At large $Q^2$ this FSI is expected to vanish by virtue of CT. One needs first a reliable formalism for description of FSI of the struck proton, and the coupled-channel generalization [7,8,36,37] of the Glauber’s multiple scattering theory [38] provides the necessary framework.

The quantity measured in the ideal $A(e, e'p)$ scattering is the spectral function $S(E_m, \vec{p}_m)$ as a function of the missing energy $E_m$ and missing momentum $\vec{p}_m = (p_{m,z}, \vec{p}_\perp)$, for the kinematics see Fig. 16. The discussion greatly simplifies if the measured cross section is integrated over the sufficiently broad range of the missing energy $E_m$, when the closure can be applied [7,8,36]. We assume this is the case. If the plane-wave impulse approximation (PWIA) were applicable, then one would have have measured the single-particle momentum distribution $n_F(\vec{p}_m)$ ([39], for the recent review see [40]), which is related to the one-body nuclear density matrix $\rho_1(\vec{r}, \vec{r}')$ as

$$d\sigma_A \propto n_F(\vec{p}_m) = \frac{1}{Z} \int dE_m S_{PWIA}(E_m, \vec{p}_m) = \int d\vec{r}'d\vec{r} \rho_1(\vec{r}, \vec{r}') \exp[i\vec{p}_m(\vec{r}' - \vec{r})].$$

(38)

In this case the missing momentum $\vec{p}_m$ is precisely the intranuclear momentum of the struck proton.

PWIA is the theorists dream. The PWIA spectral function is what theorists do calculate in their models, in the real life the outgoing proton can not be described by the plane wave. At large $Q^2 > 1 GeV^2$, FSI of the struck proton can be treated in the Glauber approximation. The FSI modifies Eq. (38) [36,37]:

$$d\sigma_A \propto \frac{1}{Z} \int dE_m S(E_m, \vec{p}_m) =$$
where \( \vec{r} = (b, z), \vec{r}' = (\tilde{b}, \tilde{z}), \Delta = \vec{r} - \vec{r}' \), \( \alpha_{pN} \) denotes the ratio of the real to imaginary parts of the forward proton-nucleon scattering amplitude and

\[
\xi(\Delta) = \int d^2\vec{q} \, \frac{d\sigma_{el}(pN)}{d^2\vec{q}} \exp(i\vec{q}\Delta) .
\]  

The derivation of attenuation factors in [33] uses the independent particle model. The effect of two-nucleon correlations in nuclear attenuation was studied in detail in [41] and found to be negligible.

FSI leads to the four important effects [13,36,37]

- Firstly, as a result of the phase factor in the integrand of (33) which is of the form \( \exp[i\frac{1}{2}\sigma_{tot}(pN)\alpha_{pN}n_A(b, z)(z-z')] \) the spectral function \( S(k_\perp, k_z) \) is probed at a shifted value of the longitudinal momentum with

\[
k_z - p_{m,z} = \Delta p_{m,z} \approx \frac{1}{2} \sigma_{tot}(pN) \alpha_{pN} n_A \approx \alpha_{pN} \cdot 35 \text{ (MeV/c)} .
\]

In the kinematical range of the NE18 experiment \( \alpha_{pN} \approx -0.5 \) [42] and \( \Delta p_{m,z} \approx -20 \text{ MeV/c} \) is quite large. This shift makes the measured \( p_{m,z} \) distribution strongly asymmetric about \( p_{m,z} = 0 \).

- Secondly, the factor \( \exp[t(\tilde{b}, z)\xi(\Delta)] \) in Eq. (33) leads to a broadening of the \( p_\perp \) distribution. The multiple elastic-rescattering expansion for the \((E_m, p_{m,z})\)-integrated spectral function reads

\[
f_A(\vec{p}_\perp) = \frac{1}{Z} \int dE_m \, dp_{m,z} \, S(E_m, p_{m,z}, p_\perp) = \sum_{\nu=0}^{\infty} W^{(\nu)} n^{(\nu)}(\vec{p}_\perp) .
\]  

Here the \( p_\perp \)-distribution in the \( \nu \)-fold rescattering \( n^{(\nu)}(\vec{p}_\perp) \) equals

\[
n^{(\nu)}(\vec{p}_\perp) = \int d^2\vec{s} \, \frac{B}{\nu \pi} \exp\left(-\frac{B}{\nu} s^2\right) n_F(\vec{p}_\perp - \vec{s})
\]

where \( B \) denotes the diffraction slope for elastic \( pN \) scattering, \( d\sigma_{el}(pN)/dt \propto \exp(-B|t|) \), and

\[
W^{(\nu)} = \frac{1}{A} \int dz d^2\tilde{b} \, n_A(\tilde{b}, z) \exp\left[-\sigma_{tot}(pN)t(\tilde{b}, z)\right] \left[\frac{t(\tilde{b}, z)\sigma_{el}(pN)}{\nu!}\right]^{\nu}
\]

gives the probability of having \( \nu \) elastic rescatterings.

- Thirdly, FSI introduces the attenuation effect. Namely, whereas in the PWIA one has \( \int d^3\vec{p}_m \, n_F(p_m) = 1 \), with allowance for FSI

\[
T_A = \frac{1}{Z} \int dE_m \, d^3\vec{p}_m \, S(E_m, \vec{p}_m) = \int d^3\vec{p}_\perp \, f_A(p_\perp)
\]

\[
= \sum_{\nu=0}^{\infty} W^{(\nu)} = \frac{1}{A} \int dz d^2\tilde{b} \, n_A(\tilde{b}, z) \exp\left[-\sigma_{in}(pN)t(\tilde{b}, z)\right] < 1 .
\]
In the completely integrated nuclear spectral function \( [15] \), nuclear attenuation is given by \( \sigma_{in}(pN) \) [7]. This result is self-obvious: elastic rescatterings only deflect, but do not absorb, the struck proton, and the effect of deflection is not relevant for the full \( 4\pi \) acceptance.

- On the other hand, the forward peak of \( f_A(p_{\perp}) \) at \( \vec{p}_{\perp} = 0 \) is dominated by \( W^{(0)} \), which is also given by Eq. \( (15) \) but with \( \sigma_{in}(pN) \) substituted by \( \sigma_{tot}(pN) \). In this case elastic rescatterings also contribute to the observed attenuation.

The further discussion will be centered on: (1) how CT affects the integrated \( T_A \) and forward (the in-parallel kinematics) \( W^{(0)} \) nuclear transparency; (2) why the onset of CT in \( A(e, e'p) \) scattering is so slow; (3) the theoretical interpretation of the nonobservation of CT in the NE18 experiment; (4) the discussion of feasibility of CT studies at CEBAF.

### 4.2 FSI effect dominates at large transverse missing momenta

Measuring the spectral function at large missing momenta \( p_m \) is of great interest, because large \( p_m \) are expected to give a direct handle on the short-range correlations (SRC) of nucleons, which is widely discussed for 30 years since seminal works by Gottfried and Srivastava [43]. Our important finding [36] is that in the transverse kinematics, the FSI effect completely takes over the SRC effect.

The global effects of rescatterings are summarized in Table 1. Here we present, for different values of \( Q^2 \), the fractions \( P^{(\nu)} = W^{(\nu)}/T_A \) of the \( \nu \)-fold rescatterings, the nuclear transparency \( W^{(0)} \) for quasifree knockout in parallel kinematics \( p_{\perp} = 0 \), the total transparency \( T_A \) and the average number of secondary rescatterings \( \langle \nu \rangle \). The \( Q^2 \) dependence of \( 1 - W^{(0)} \) is reminiscent of the energy dependence of \( \sigma_{tot}(pN) \), which is nearly flat at the kinetic energy \( T_{kin} \approx Q^2/2m_p \gtrsim 0.5 \text{ GeV} \) [42]. The \( Q^2 \) dependence of \( 1 - T_A \) repeats the energy dependence of \( \sigma_{in}(pN) \), which rises rapidly up to \( T_{kin} \sim 1.5 \text{ GeV} \) and then approximately levels off [42]. The \( Q^2 \) dependence of the difference of the total and forward transparency \( T_A - W^{(0)} \) and of the multiplicity of secondary rescatterings \( \langle \nu \rangle \) repeats the energy dependence of the elastic cross section \( \sigma_{el}(pN) = \sigma_{tot}(pN) - \sigma_{tot}(pN) \), which steadily decreases with \( T_{kin} \) [42]. The difference between \( W^{(0)} \) and \( T_A \) is a convenient measure of the strength of rescatterings, and in the CEBAF range of \( Q^2 \) this difference is very large.

For our numerical estimates of the \( p_{\perp} \)-distribution we use a simple, yet realistic, parameterization of the SPMD [44]

\[
\begin{align*}
  n_F(\vec{k}) & \propto \exp \left( -\frac{5}{2} k^2/k_F^2 \right) + \epsilon_0 \exp \left( -\frac{5}{6} k^2/k_F^2 \right),
\end{align*}
\]

where \( \epsilon = 0.03 \) and the value of the Fermi-momentum \( k_F \) has been taken from Ref. [39]: \( k_F(C) = 221 \text{ MeV}/c \) and \( k_F(Pb) = 265 \text{ MeV}/c \). This parameterization is consistent with the results from the y-scaling analysis [45]. The steeply decreasing first term corresponds to the mean-field component of the SPMD, while the second term describes the SRC tail. In Fig. 17 we present our predictions for \( f_A(\vec{p}_{\perp})/T_A \) for the \(^{12}C(e, e'p)\) and \(^{208}Pb(e, e'p)\) reactions at \( Q^2 = 1, 2, 4 \) (GeV/c)^2, respectively. The FSI contribution to \( f_A(\vec{p}_{\perp}) \) starts to dominate over the PWIA component \( P^{(0)}n_F(p_{\perp}) \) already at \( p_{\perp} \gtrsim 350 \text{ MeV}/c \) for carbon and \( p_{\perp} \gtrsim 300 \text{ MeV}/c \) for lead, which is precisely the region thought of being dominated by the SRC tail in the single-particle momentum distribution.
An obvious signal of rescattering is the multinucleon emission (MNE). Rescatterings are not imperative for MNE, but rescatterings leading to \( p_\perp \gtrsim k_F \) are necessarily followed by MNE. The contribution of rescatterings to MNE is characterized by the average number of secondary rescatterings \( \langle \nu(p_\perp) \rangle \), which is very large (Table 1). Evidently, the strength of MNE coming from the rescattering mechanism must be the same in the longitudinal and transverse cross sections, which is a strong prediction. Presently, it can not be tested for the lack of the experimental data on large transverse missing momenta \( p_\perp \gtrsim k_F \) taken at \( Q^2 \gtrsim (1-2)\text{GeV}^2 \), which is the domain of the forthcoming CEBAF experiments. This FSI driven MNE contributes significantly to the spectral function at large missing energy \( E_m \).

Namely, the contribution \( \Delta E_m \) to the missing energy \( E_m \) from the kinetic energy of the recoil nucleons can be estimated as

\[
\Delta E_m \approx \left\langle \frac{s^2}{2m_p} \right\rangle = (1 + \epsilon_0) \frac{k_F^2}{5m_p} + \frac{p_\perp^2}{2m_p}. \quad (47)
\]

Incidentally, the SRC mechanism leads to a similar relationship between \( \Delta E_m \) and \( p_\perp \). The rescattering effect dominates the tail of the missing-energy distribution

\[
\frac{df_A}{dE_m} \propto \exp \left( -\frac{E_m}{E_0} \right), \quad (48)
\]

where the slope

\[
E_0 \approx \frac{1}{2Bm_p} \approx \frac{8\pi \sigma_{el}(pN)}{m_p \sigma_{tot}(pN)^2}. \quad (49)
\]

The definitive signature of the FSI mechanism is a strong \( Q^2 \) dependence of the slope \( E_0 \), which decreases from \( E_0 \sim 250 \text{ MeV} \) at \( Q^2 \sim 1 \text{ (GeV/c)}^2 \) down to \( E_0 \sim 90 \text{ MeV} \) at \( Q^2 \sim 6 \text{ (GeV/c)}^2 \). Such a large value of \( E_0 \) shows that in order to exhaust the closure, one has to integrate up to very large \( p_\perp \) and very large missing energies \( E_m \gtrsim E_0 \).

4.3 Rescattering effect in \( d(e, e'p) \)

The rescattering effect is quite strong even in such a diluted target as the deuteron. The momentum distribution \( f_d(p) \) of the observed protons equals [37]

\[
f_d(p_{m,z}, \vec{p}_\perp) = \left| \phi_d(p_{m,z}, \vec{p}_\perp) - \frac{\sigma_{tot}(pn)}{16\pi^2} \int d^2\vec{k} \phi_d(p_{m,z}, \vec{p}_\perp - \vec{k}) \exp \left( -\frac{1}{2} B \vec{k}^2 \right) \right|^2. \quad (50)
\]

Here \( \phi_d(\vec{k}) \) is the momentum-space wave function of the deuteron, and gives the undistorted single-particle momentum distribution, the second term describes the rescattering effect. The measured transparency \( T_d \) depends on the acceptance \( p_{max} \):

\[
T_d = \frac{\int_{p_{max}} p_\perp^2 f_d(\vec{p}_\perp) d^2\vec{p}_\perp \phi_d(\vec{p}_\perp)^2}{\int_{p_{max}} p_\perp^2 \phi_d(\vec{p}_\perp)^2}. \quad (51)
\]

The attenuation effect \( T_d < 1 \) comes from the interference of the undistorted and rescattering terms in (50). If \( R_d p_{max}^2 \gg 1 \), but \( B p_{max}^2 \ll 1 \), which is the case for the NE18, then

\[
1 - T_d \sim \frac{\sigma_{tot}(pn)}{2\pi R_d^2} \sim 0.07, \quad (52)
\]
which is about twice the Glauber shadowing effect in the $\sigma_{tot}(Nd)$. Our prediction for $T_d$ is shown in Fig. 18 and the agreement with the NE18 data [4] is good.

At still larger values of $p_\perp$, the spectral function will be entirely dominated by square of the rescattering term in (50), cf. Eqs. (42,43),

$$f_d(p_{m,z}, p_\perp) \sim \frac{\sigma_{el}(pN)}{4\pi R_d^2} |\phi_d(p_{m,z})|^2 B \exp(-Bp_\perp^2)$$

It exhibits features of the rescattering contribution to the spectral function which are common to all the nuclei:

- The $p_\perp$ distribution has the broad tail with the small slope, which is the diffraction slope of the $pN$ elastic scattering.

- The longitudinal momentum distribution is much narrower and is given by the PWIA single-particle momentum distribution.

In the opposite to that, the SRC effect in the single-particle momentum distribution is an isotropic function of the missing momentum $\vec{p}_m$, which must enable the experimental separation of the isotropic SRC effect from the sideways contribution from FSI.

### 4.4 The theoretical interpretation of the NE18 data $A(e, e'p)$ scattering

The most important conclusion from the above discussion is that the distortion effects make absolutely illegitimate a factorization of the measured spectral function $S(E_m, \vec{p}_m)$ into the PWIA spectral function $S_{PWIA}(E_m, \vec{p}_m)$ and an attenuation factor which is independent of the missing energy and momentum. In order to disentangle the genuine attenuation effect from the distortion effects, one can define nuclear transparency $T_A$ as the ratio of the experimentally measured, and the theoretically calculated PWIA, cross sections integrated over the experimental acceptance domain $D(NE18)$ in the $(E_m, p_{m,z}, p_\perp)$ space:

$$T_A(NE18) = \frac{\int_{D(NE18)} dE_m dp_{m,z} dp_\perp S(E_m, p_{m,z}, p_\perp)}{\int_{D(NE18)} dE_m dp_{m,z} dp_\perp S_{PWIA}(E_m, p_{m,z} + \Delta p_{m,z}, p_\perp)}$$

In order to eliminate the spurious effect in $T_A(NE18)$ due to the FSI generated asymmetry of the nuclear spectral function, we strongly advocate including the effect of the shift (41) in the calculation of the PWIA cross section. It was not included in the NE18 analysis. The effect of the shift (41) on $T_A$ depends on the $p_{m,z}$-acceptance. It vanishes for the wide $p_{m,z}$-acceptance. For the narrow acceptance centered at $p_{m,z} = p^*$, the longitudinal asymmetry effect can be evaluated as

$$\frac{T_A(\Delta p_{m,z})}{T_A(\Delta p_{m,z} = 0)} \approx 1 + \frac{5}{2} \cdot \frac{(\Delta p_{m,z} + p^*)^2 - (p^*)^2}{k_f^2}$$

and can be quite strong at large $p^*$. At $p^* = 0$, which is approximately the case in the NE18 experiment, the asymmetry effect enhances $T_A$ by $\lesssim 2\%$ and is still within the error bars of the NE18 experiment.

The difference between $\sigma_{tot}(pN)$ and $\sigma_{in}(pN)$, and between $W^{(0)}$ and $T_A$ thereof, is very large at moderate energies. $T_A(NE18)$ includes partly the elastically rescattered struck
protons. Since the diffraction slope $B$ rises with the proton energy, the higher is $Q^2$ the larger is the fraction of the elastically rescattered struck protons included in $T_A(NE18)$. In Fig. 18 we present our predictions [38] for $T_A$, $W(0)$ and $T_A(NE18)$ in which the $p_{\perp}$ integration for heavy nuclei is extended up to $p_{\text{max}} = 250\,\text{MeV/c}$ as relevant to the NE18 situation [4]. We find very good quantitative agreement with the NE18 data. The principle conclusion from this comparison is that there is no large signal of CT in the $A(e, e'p)$ scattering at $Q^2 \leq 7\,\text{GeV}^2$.

For the $^{12}\text{C}$ target we also show the effect of CT which is still small even at the largest $Q^2$ of the NE18 experiment (Fig. 18b), which must be contrasted with large CT effect in the E665 experiment. What makes the onset of CT in exclusive vector meson production and $(e, e'p)$ scattering so much different?

5 The onset of CT in $A(e, e'p)$ scattering

5.1 The ejectile state has a large size

The $(e, e'p)$ scattering can be viewed as an absorption of the virtual photon by the target proton, which leads to the formation of the ejectile state $|E\rangle$, which then is projected onto the observed final state proton. In the $A(e, e'p)$ scattering, this ejectile state propagates in the nuclear medium before evolving into the observed proton. The strength of FSI depends on what is the transverse size and the interaction cross section of the ejectile state.

In the exclusive production of vector mesons, the wave function of the ejectile state equals [9,21]

$$\Psi_E(r) \propto \sigma(r)|E\rangle \Psi_{\gamma*}(r)$$  \hspace{1cm} (56)

Because of CT, in this case the ejectile wave function $\Psi_E(r)$ has a hole at $r = 0$, but at large $Q^2$ the decrease of the wave function of the photon (5) takes over and $\Psi_E(r)$ has a small size $\sim r_s$. Notice, that in terms of the ejectile state $|E\rangle$ the strength of FSI can be rewritten as $\Sigma_V = \langle V|\sigma(r)|E\rangle/\langle V|E\rangle$.

The wave function of the ejectile state in the quasielastic scattering of electrons is much simpler. It can be found in any textbook in quantum mechanics and/or nuclear/particle physics. Indeed, the charge form factor $G_{em}(Q)$ of the $q\bar{q}$ ”proton” can be written as (here the $\vec{r}$-plane is normal to the momentum transfer $\vec{Q}$):

$$G_{em}(Q) = \langle p|E\rangle = \int dzd^2r\Psi_p^*(z, r)\Psi_E(\vec{r}, z) = \int dzd^2r\Psi_p^*(z, r)\exp\left(\frac{i}{2}Qz\right)\Psi_p(\vec{r}, z).$$  \hspace{1cm} (57)

Therefore, the ejectile wave function

$$\Psi_E(\vec{r}, z) = \exp\left(\frac{i}{2}Qz\right)\Psi_p(\vec{r}, z).$$  \hspace{1cm} (58)

Because $|\Psi_E(\vec{r}, z)|^2 = |\Psi_p(\vec{r}, z)|^2$, the ejectile wave packet has the transverse size identical to the size of the proton [46,13]. This result is readily generalized to the three-quark proton and to the relativistic lightcone formalism. The often made statement that the ejectile has a small size $\sim 1/\sqrt{Q^2}$ is quite misleading. The small size only is believed to gradually appear when the ejectile wave function is projected at very large $Q^2$ onto the observed proton [46,47], but one must bear in mind the well known warning by Isgur and Llewellyn Smith [48] that the onset of the small-size dominance in the charge form factor is very slow.
Even provided that the proton form factor is dominated by the small size contribution, it still would be erroneous to identify the size of the ejectile state with the region of \( r \) which contribute to the form factor of the proton. In the problem of interest it is the large-size ejectile state with the initial wave function \( |E⟩ \) which starts propagating and having FSI in the nuclear medium, evaluations of CT effects in models which start with the ejectile of vanishing size must be taken with grain at salt.

5.2 Color transparency sum rules and vanishing FSI: conspiracy of hard and soft scattering

The ejectile state can be expanded in terms of the electroproduced states \( |i⟩ \) and the form factors \( G_{ip}(Q) \) of the \( ep \to ei \) transitions:

\[
|E⟩ = J_{em}(Q)|p⟩ = \sum_i |i⟩\langle i|J_{em}(Q)|p⟩ = \sum_i G_{ip}(Q)|i⟩
\]

where \( G_{ip}(Q) \) are the \( ep \to ei \) transition form factors. Detailed description of how CT emerges in the coupled-channel multiple scattering theory is given in [7,8,13,46]. There is much similarity to the treatment of CT in the production of vector mesons. To the leading order in FSI,

\[
T_A = 1 - \Sigma_{ep}(Q) \frac{1}{2A} \int d^2b T(b)^2 + ..., \tag{60}
\]

where the strength of FSI is given by

\[
\Sigma_{ep}(Q) = \frac{⟨p|\hat{σ}|E⟩}{⟨p|E⟩} = \frac{1}{G_{em}(Q)} \sum_i σ_{pi}G_{ip}(Q) = σ_{tot}(pp) + \frac{1}{G_{em}(Q)} \sum_{i≠p} σ_{pi}G_{ip}(Q). \tag{61}
\]

Here \( G_{em}(Q) = G_{pp}(Q) \) is the form factor of the elastic \( ep \) scattering, \( \hat{σ} \) is the cross section or diffraction operator, which gives the forward diffraction scattering amplitudes \( M(jp \to kp) = i⟨k|\hat{σ}|j⟩ = iσ_{jk} \). The normalization is such that \( σ_{tot}(jN) = σ_{jj} = \text{Im}M(jN \to jN) \).

Weak FSI and/or CT effect only is possible if the off-diagonal contributions in Eq. (60) cancel the \( σ_{tot}(pN) \), which comes from the conventional Glauber rescattering of the struck proton. Evidently, such a cancellation requires a very special conspiracy between the electromagnetic form factors of hard \( ep \) interaction, and amplitudes of soft, forward, diffractive hadronic transitions \( ip \to pp \). Now we shall demonstrate that such a conspiracy takes place.

In QCD the (anti)quarks in the hadron always have the one-gluon exchange Coulomb interaction at short distances. This Coulomb interaction leads to a very special asymptotics of the form factor at \( Q^2 >> R_p^2 \) [46,47] (for the sake of simplicity we discuss the two-quark state in the nonrelativistic approximation)

\[
G_{ik}(Q) \propto \frac{V(\vec{Q})}{Q^2}Ψ_i^*(0)Ψ_k(0) \tag{62}
\]

Here \( Ψ_i(0) \) are the wave functions at the origin, and \( V(Q) \) is the one-gluon exchange quark-quark hard scattering amplitude. Making use of the QCD asymptotics (62) we find [13]

\[
Σ_{ep}(Q) = \frac{1}{G_{em}(Q)} \sum_i σ_{pi}G_{ip}(Q) \propto \frac{1}{G_{em}(Q)} \frac{V(\vec{Q})}{Q^2} \sum_i σ_{pi}Ψ_i^*(0)Ψ_p(0) \propto \sum_i σ_{pi}Ψ_i^*(0). \tag{63}
\]
Remarkably, the r.h.s. of Eq. (63) vanishes by virtue of CT sum rules [7]. The proof goes as follows: The state of transverse size $\vec{r}$ can be expanded in the hadronic basis as $|\vec{r}\rangle = \sum_i \Psi_i(\vec{r})^* |i\rangle$ and the hadronic-basis expansion for $\sigma(r)$ reads as

$$\sigma(r) = \langle \vec{r}|\hat{\sigma}|\vec{r}\rangle = \sum_{i,k} \langle \vec{r}|k\rangle \langle k|i\rangle \langle i|\vec{r}\rangle = \sum_{i,k} \Psi_i^*(\vec{r})\Psi_k(\vec{r})\sigma_{ki}$$

(64)

By virtue of CT $\sigma(r = 0) = 0$, and we obtain the "CT sum rule" [7]

$$\sum_{i,k} \Psi_k(0)^*\Psi_i(0)\sigma_{ki} = 0.$$  

(65)

Considering the matrix elements $\langle \vec{r}|\hat{\sigma}|k\rangle$ at $\vec{r} \to 0$, one readily finds a whole family of CT sum rules [7,13]

$$\sum_i \Psi_i(0)\sigma_{ik} = 0.$$  

(66)

It is precisely the sum rule (66), which in conjunction with the QCD asymptotics of the electromagnetic form factors (62) ensures that the strength of FSI (61) vanishes and CT indeed takes place.

### 5.3 The onset of CT and the coherency constraint

As we learned in section 3.4, at finite energy only few intermediate states contribute coherently into $\Sigma_V$. The case of the $\Sigma_{ep}$ is quite similar. From the kinematics of deep inelastic scattering

$$m_i^2 - m_p^2 = 2m_p\nu - Q^2 - 2\nu k_z,$$

(67)

where $k_z$ is the longitudinal momentum of the target nucleon. Consequently, different components $|i\rangle$ of the ejectile wave packet (59) are produced with the longitudinal momenta differing by [7]

$$\kappa_{ip} = \frac{m_i^2 - m_p^2}{Q^2} = \frac{m_i^2 - m_p^2}{2\nu},$$

(68)

and acquire the relative phase $\kappa_{ip}(z_2 - z_1)$ during propagation from the $(z_1)$ production to the $(z_2)$ rescattering point. Then, repeating considerations which lead to Eqs. (36, 37), we find [7,8,13]

$$\langle p|\hat{\sigma}J_{em}|p\rangle = \sum_i \langle p|\hat{\sigma}|i\rangle \langle i|J_{em}|p\rangle$$

$$\Rightarrow \sum_i \langle p|\hat{\sigma}|i\rangle \langle i|J_{em}|p\rangle \exp[i\kappa_{ip}(z_2 - z_1)] \Rightarrow \sum_i \langle p|\hat{\sigma}|i\rangle \langle i|J_{em}|p\rangle G_A^2(\kappa_{ip})$$

(69)

and

$$\Sigma_{ep}(Q) = \sigma_{tot}(pN) + \sum_{ip\neq p} \frac{G_{ip}(Q)}{G_{em}(Q)}\sigma_{ip}G_A^2(\kappa_{ip})$$

(70)

At small $Q^2$ such that

$$Q^2 \ll 2R_A m_p \Delta m \sim (3 - 5) \cdot A^{1/3} \text{GeV}^2,$$

(71)

for all inelastic channels $G_A(\kappa_{ip}) \ll 1$, so that $\Sigma_{ep} = \sigma_{tot}(pN)$, and FSI measures the free-nucleon cross section. With increasing $Q^2$, for larger and larger number of inelastic intermediate states channels $G_A(\kappa_{ip}) \approx 1$, the CT sum rule (66) will be better saturated and $\Sigma_{ep}(Q) \to 0$.  

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5.4 What makes the onset of CT in \( A(e, e'p) \) and vector meson production different?

In the vector meson production, it is the energy \( \nu \) which controls the number of intermediate states which contribute coherently to \( \Sigma_V \) and the signal of CT. The virtuality \( Q^2 \) is an independent parameter, which controls the strength of destructive interference between the diagonal (elastic) and off-diagonal (inelastic) contributions to \( \Sigma_V \). In the quasielastic scattering of electrons \( Q^2 \approx 2m_p\nu \), and one can increase the energy \( \nu \) only at the expense of very high \( Q^2 \).

There is one more fundamental difference between the two processes. Compare in more detail the \( \Sigma_{J/\Psi} \) Eq. (36) and \( \Sigma_{ep}(Q) \) Eq. (70). Notice, that

\[
M(J/\Psi N \rightarrow \Psi'N) = \sigma_{J/\Psi \Psi'}
\]

and, in the generic case, the off-diagonal matrix elements \( \sigma_{ip} \) are typically smaller than the diagonal \( \sigma_{ii} = \sigma_{tot}(i N) \). The difference comes from the two ratios

\[
\frac{M(\gamma^* N \rightarrow \Psi'N)}{M(\gamma^* N \rightarrow \Psi N)} \quad \text{and} \quad \frac{G_{ip}(Q)}{G_{em}(Q)},
\]

which play the identical role in Eqs. (36), (70), respectively. In the vector meson production, the ratio (72) increases with \( Q^2 \), which further enhances the off-diagonal contribution to \( \Sigma_{J/\Psi} \) and the signal of CT thereof. In the case of \( A(e, e'p) \) at large \( Q^2 \), the ratio \( G_{ip}(Q)/G_{em}(Q) \) does not change with \( Q^2 \), see Eq. (62) and Stoler’s review [49].

5.5 The realistic estimates of the signal of CT in \( A(e, e'p) \) scattering

One needs to evaluate the off-diagonal contributions to \( \Sigma_{ep} \). Let us accept the optimistic scenario that the form factor ratios \( G_{ip}(Q)/G_{em}(Q) \) are precociously close to QCD asymptotic predictions (62), and concentrate on the off-diagonal diffraction production amplitudes \( \sigma_{ip} \). The experimentally observed mass spectrum of the diffraction excitation of protons is shown in Fig. 19 (for the review see [50]). Because of the nuclear form factor \( G_A(\kappa_{ip}) \), the onset of CT signal starts with the contribution to \( \Sigma_{ep} \) from the low-mass states.

In the diffractive mass spectrum of Fig. 19 there is a prominent signal of excitation \( pp \rightarrow (\pi N)p \) of the low-mass \( \pi N \) state, which dominates at \( M^2 \approx 2\text{GeV}^2 \). Its origin is in the admixture of the \( \pi(3q) \) Fock state in the lightcone proton [13,51]

\[
|N\rangle = \cos \theta |3q\rangle + \sin \theta |3q + \pi\rangle,
\]

where the mixing angle \( \theta \) is related to the number of pions \( n_\pi \) in the nucleon as

\[
\sin^2 \theta = \frac{n_\pi}{1 + n_\pi}
\]

There is a mounting evidence for such a hybrid quark core-pion structure of the nucleon, coming from the violation of the Gottfried Sum Rule and the recent NA51 observation [52] of large \( \bar{u}/d \) asymmetry in the proton sea in the Drell-Yan production [53]. The diffractive excitation of the \( N + \pi \) Fock component of the nucleon (known since long also as the Drell-Hiida-Deck process [50]), gives an excellent description of the \( \pi N \) mass spectrum. For the purposes of the present analysis it is important, that the \( \pi N \) state contributes to \( \Sigma_{ep} \) a term [13]

\[
\Delta \Sigma_{ep}(\pi N) \approx -\sigma_{tot}(\pi N) \frac{n_\pi}{1 + n_\pi} \approx -8\text{mb}
\]
This result can be interpreted as follows. The pion-baryon component of the nucleon has a large size and a rapidly decreasing form factor, so that only the baryonic $3q$ core of the nucleon contributes to the ejectile state (the detailed analysis of electromagnetic form factor of the hybrid proton is given in [51]). The interaction cross section of the $3q$ core stripped off of pions, is smaller that $\sigma_{\text{tot}}(pN)$ by precisely the amount (75). Notice, how the nonperturbative quark core-pion Fock component of the nucleon gives a unified description of

- the diffraction dissociation of nucleons in the soft diffractive scattering,
- the asymmetry of the quark-antiquark sea in protons as probed in deep inelastic scattering,
- the onset of color transparency in quasielastic scattering of electrons on nuclei.

Diffraction excitation of higher mass states comes from the excitation of the $3q$ core of the nucleon. In [7] we described interaction of the $3q$ core in a QCD approach based on the dipole cross section [15], qualitatively shown in Fig. 1. The resulting diffraction matrix $\sigma_{ik}$ has the built-in CT property and satisfies CT sum rules. The proton and its excitations are described by the diquark-quark harmonic oscillator model. Within the model, diffraction dissociation is dominated by excitation of $N^*$ states of the positive parity shells. We find a good parameter-free description [54] of the diffractive mass spectrum shown in Fig. 19, and rightfully expect the realistic estimate of CT effects in $\Sigma_{ep}(Q^2)$. The excess of the observed mass spectrum in the so-called triple-pomeron mass region $M^2 \gtrsim 10\text{GeV}^2$ is an effect of excitation of higher $qqqg_{1..g_n}$ Fock states [17], which was not included in [7]. However, in the practically accessible range of $Q^2$, CT effect is completely dominated by excitation of the $\pi N$ states and resonances of the first excited positive-parity shell $L = 1$, with little contribution of states of higher shells $L \geq 2$.

The resulting predictions [54] for $\Sigma_{ep}(Q)$ are shown in Fig. 20. With increasing $Q^2$, the strength of FSI decreases much slower than $\propto 1/Q^2$: the increase of $Q^2$ by one order in magnitude from 5 to 50 GeV$^2$ is followed by decrease of $\Sigma_{ep}(Q)$ only by the factor 3 from 36mb to 12 mb, respectively. The CT effect in $T_A(Q^2)$ follows closely the variation of $\Sigma_{ep}(Q^2)$. The heavier is the nucleus, the larger $Q^2$ is needed for the onset of CT. The values of $Q^2$ studied by NE18 correspond to the threshold of CT for the carbon nucleus, already at $Q^2 \sim 20\text{ GeV}^2$ we predict a substantial CT effect.

6 CT in $A(e, e'p)$ scattering at CEBAF

6.1 $^4He(e, e'p)$ scattering: CEBAF’s choice

The principal inhibitor of the precocious CT is the coherency constraint, quantified by the factor $G_A(\kappa_{ip})^2$ in the off-diagonal contribution to $\Sigma_{ep}(Q)$. In order to have large CT effect one must excite intermediate states of high mass $m_i$. The condition $\kappa_{ip}R_A \lesssim 1$ implies, that excitation of the same mass on different nuclei requires

$$Q^2 \gtrsim (m_i^2 - m_p^2)R_A m_p$$

Consequently, one can hope to relax the coherency constraint and enhance the CT effect going to lighter nuclei which have smaller radius and less steeply decreasing charge form factor $G_A(\kappa)$. 

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With allowance for the center-of-mass motion, for the light nuclei \[13,54\]
\[\Sigma_{ep}(Q^2) = \sigma_{tot}(pN) + \sum_{i \neq p} \sigma_{ip} \frac{G_{ip}(Q^2)}{G_{pp}(Q^2)} G_A^2 \left( \frac{A}{A-1} \kappa^2_{ip} \right). \]  
(77)

The factor \(A/(A-1)\) in the nuclear form factor, which comes from the center-of-mass motion important in light nuclei, is the most unwelcome news. For the deuteron target, it makes the onset of CT slower than for the \(^{12}\)C target. The \(^3\)He nucleus has a size smaller than the \(^{12}\)C nucleus, but the onset of CT in both cases will be equally slow, which makes the \(^3\)He(e, e\(^\prime\)p) reaction ill-suited for the search of CT effects at CEBAF. Only for the \(^4\)He nucleus, the onset of CT will be sooner than for the \(^{12}\)C nucleus, which makes \(^4\)He the CEBAF’s choice target.

In the \(^4\)He(e, e\(^\prime\)p) scattering, nuclear attenuation effect is typically \(1 - T_A \sim 0.25\) [55] and with the onset of CT it will decrease \(1 - T_A \propto \eta = \frac{\Sigma_{ep}(Q)}{\sigma_{tot}(pN)}. \)  
(78)

Our predictions [54] of \(\eta\) for \(^4\)He(e, e\(^\prime\)p) scattering are shown in Fig. 21. We also show the decomposition of the CT signal \(1 - \eta\) in terms of excitation of inelastic intermediate states. At \(Q^2 \leq (5-6)\text{GeV}^2\) the signal of CT comes primarily from the \(\pi N\) intermediate states. In this case the reduction of \(\Sigma_{ep}(Q)\) from \(\sigma_{tot}(pN) = 40\text{mb}\) measures how much the size of the quark core of the nucleon stripped off the pionic cloud, is smaller than the size of the physical proton. With the increasing \(Q^2\), excitation of the intermediate nucleonic resonances starts contributing to the signal of CT and takes over at \(Q^2 \gtrsim 7\text{GeV}^2\). The experimental observation of the CT predicted \(\sim 20\%\) decrease of \(\eta\) at \(Q^2 = 6\text{GeV}^2\), requires the measurements of \(T_A\) with \(\lesssim 1\%\) statistical accuracy. This challenge can be met at CEBAF thanks to its high luminocity.

6.2 The signal of CT at large transverse missing momenta

The broadening of the \(p_\perp\) distribution discussed in sections 4.1-4.3 is generated by FSI and must vanish in the CT regime. With allowance for inelastic intermediate states, the probability of \(\nu\)-fold rescattering \(W^{(\nu)}\) will still be given by Eq. (44) subject to the substitution [13,14]
\[\sigma_{el}(pN)^\nu \rightarrow \sigma_{el}(pN)^\nu \left[ \frac{\langle p|\hat{\sigma}^\nu E \rangle}{\langle p|E\sigma_{tot}(pN)^\nu \rangle} \right]^2. \]  
(79)

In the case of \(^4\)He(e, e\(^\prime\)p) scattering, the tail of the \(p_\perp\)-distribution \(f_A(p_\perp)\) will be dominated by single scattering term, which will be proportional to the factor
\[W_1 \propto \eta^2 = \left( \frac{\Sigma_{ep}(Q^2)}{\sigma_{tot}(pN)} \right)^2. \]  
(80)

Our prediction [54] for the rescattering suppression factor \(\eta^2\) is shown in Fig. 21. Here the CT effect is twice as large compared to nuclear transparency \(T_A\). One has to pay a heavy price, though: very accurate theoretical calculation of the rescattering effect in the absence of CT is a must before one can claim the reliable experimental determination of this suppression factor and separation of the CT effect. Such a calculation must be very
accurate, indeed, because the magnitude and the $p_{\perp}$-dependence of the differential cross section of elastic $pN$ scattering changes very rapidly over the range of $Q^2$ of the interest.

Furthermore, the rescattering component of $f_A(p_{\perp})$ appears on the background from the short-range correlation effect, see Fig. 17, and very good understanding of this background is necessary for reliable determination of a small decrease of $Q^2$ with the increase of $Q^2$.

The coincidence detection of the recoil nucleon little helps in this respect: once the struck proton emerges with large $p_{\perp} \gg k_F$, it will always recoil against the second nucleon emitted with the momentum $\vec{p}_2 \approx -\vec{p}_m$, and one can not tell whether the large $p_{\perp}$ originated from the correlation of two nucleons at short distances, or from the rescattering process. For this reason, the theoretical interpretation of high-precision measurements of the $Q^2$ dependence of nuclear transparency $T_A$ will be much easier, compared to the $Q^2$ dependence of the large-$p_{\perp}$ production and/or the $^4He(e,e'pp)$ reaction.

6.3 Fermi-motion and asymmetry effect in nuclear transparency

In the $(e,e'p)$ scattering, the longitudinal momentum $k_z$ of the target nucleon is uniquely determined by the electron scattering kinematics:

$$x = \frac{Q^2}{2mp\nu} = 1 + \frac{k_z}{m_p}. \quad (81)$$

By kinematics of the electroproduction Eq. (67), the different components of the ejectile wave packet are produced on the target nucleon having different longitudinal Fermi motion, see Eq. (68). Consequently, varying the Bjorken variable $x$, one can change the composition of the ejectile wave packet [8,56,57], which gives a new handle on the CT effect. A detailed theory of the effect is given in [8].

Following West, it is convenient to introduce the PWIA structure function

$$F_A(x) = \int d^3k z \frac{dn}{dk z} \delta(1 + \frac{k_z}{m_p} - x) = \frac{1}{2\pi} \int dk z \delta(1 + \frac{k_z}{m_p} - x) \int dz \rho(0,z) \exp(ik_z z). \quad (82)$$

Then, the Fermi-motion effect is quantified by the $x$-dependent decomposition of the ejectile state

$$|E\rangle_{\text{eff}} = \sum_i G_{ip}(Q) |i\rangle \frac{F_A(x + \frac{1}{2}\Delta x_{ip})}{F_A(x)} G_A(\kappa_{ip}) \quad (83)$$

and the $x$-dependent strength of FSI

$$\Sigma_{ep}(x, Q^2) = \sigma_{tot}(pN) + \sum_{i \neq p} \frac{\sigma_{pi} G_{ip}(Q)}{G_{em}(Q)} \frac{F_A(x + \frac{1}{2}\Delta x_{ip})}{F_A(x)} G_A(\kappa_{ip})^2, \quad (84)$$

where

$$\Delta x_{ip} = \frac{\kappa_{ip}}{m_p} = \frac{m_p^2 - m_{ip}^2}{Q^2}. \quad (85)$$

In their counterpart of (83), instead of $F_A(x_{Bj} + \frac{1}{2}\Delta x_{ip})/F_A(x)$ Jennings and Kopeliovich [56] give $\sqrt{F_A(x + \Delta x_{ip})}/F_A(x)$, which is incorrect.

The $x$-dependence of $\Sigma(x, Q^2)$ and of the transparency $T_A(x, Q^2)$ comes from the fact that the ratio $F_A(x + \frac{1}{2}\Delta x_{ip})/F_A(x)$ is asymmetric around $x = 1$ [8,56,57]. It is smaller than unity at $x > 1$, which suppresses the inelastic contribution to $\Sigma_{ep}$ and hence reduces the
color transparency signal. For $x < 1$, on the other hand, transparency is enhanced. With $F_A(x)$ evaluated using the single-particle momentum distribution (10), we find the results shown in Fig. 22. At small $Q^2$ the Fermi motion effect is numerically small because the contributions from inelastic channels to $\Sigma_{ep}(x, Q^2)$ are suppressed by $G_A(\kappa_{ip})^2$. At higher $Q^2$ these contributions increase and the $x$-dependence of $T_A(x, Q^2)$ becomes significant. The SRC component of $F_A(x)$ has a significant influence on the Fermi motion effect reducing the departure of the ratio $F_A(x + \frac{1}{2}\Delta x_{ip})/F_A(x)$ from unity in the region of $x$ where $F_A(x)$ is dominated by this component. This is demonstrated by the dashed curves in Fig. 22, where $T_A(x, Q^2)$ is evaluated with $F_A(x)$ which includes only the mean-field component. $T_A(x, Q^2)$ has a peak and the signal of CT is maximized at $x \sim 0.8$.

We emphasize that the asymmetry of nuclear transparency about $x = 1$ is the CT effect. It is large and persists at large $Q^2$. Zooming at $x \sim 0.8$ will significantly enhance the chances of discovering CT at CEBAF. However, there is a strong background to this CT induced asymmetry, which comes from the effective shift of the missing momentum (11), which also generates the asymmetry in nuclear transparency. Namely, $x \sim 0.8$ corresponds to $p^* \sim k_F$, and the FSI generated asymmetry (55) is as large as the CT effect shown in Fig. 22.

7 Conclusions

The E665 observation [5] of strong signal of CT in exclusive production of vector mesons is a major breakthrough in the subject of CT and paves the way for dedicated experiments on CT. The E665 effect was predicted [9] (not postdicted!), and the agreement between the theory and experiment is very good. The nonobservation of CT in the NE18 experiment [4] was also predicted [7,8] and is a good confirmation of theory. We have a good understanding of why the onset of CT in $A(e, e'p)$ and $\gamma^* A \rightarrow VA, VA^*$ reactions is so strikingly different.

The purpose of this presentation was to make the strong case for a broad experimental program on electroproduction of light vector mesons $\rho^0, \omega^0, \phi^0, \rho', \omega', \phi'$ at CEBAF. Here the principal benefit is the theoretically well understood shrinkage of the virtual photon which elastically produces the vector meson. Variations of the scanning radius are very strong already at $Q^2 \lesssim (1-3) \text{GeV}^2$, and the first-class CT experiments with large expected signal of CT can be performed at the energy-upgraded CEBAF [58]. CT in conjunction with the node structure of wave functions of radially excited vector mesons $V' = \rho', \omega', \phi'$ leads to a very rich pattern of anomalous $A$ and $Q^2$ dependence of the $V'$ production. This unique possibility of directly scanning the wave function of radial excitations at CEBAF must not be overlooked, and CEBAF is the unique facility which can do the job. Furthermore, the CEBAF experiments on radially excited vector mesons can contribute much to the spectroscopy of vector mesons, which presently leaves much to be desired [59]. Besides the above CT aspects, there is much interest in studies of nuclear modifications of properties of vector mesons in electroproduction on nuclei at CEBAF [60].

$^4$He emerges as CEBAF’s choice target to search for CT in $A(e, e'p)$ scattering. The realistic, and conservative, estimates show that CT effect in the conventional transmission experiment can be observed at CEBAF. CT effect can be enhanced studying the asymmetry of nuclear transparency (the Fermi-motion effect) in the parallel kinematics. Here more theoretical work is needed to reliably calculate the background asymmetry originating from the more conventional final state interaction. CT effect also can be enhanced looking at large missing momentum in the transverse kinematics. Here a very accurate calculation of
the rescattering effect and of the background short-range correlation effect is needed for a reliable normalization of the cross section and extraction of CT effect.

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References

[1] A.B.Zamolodchikov, B.Z.Kopeliovich and L.I.Lapidus, JETP Lett. 33 (1981) 595; G.Bertsch, S.J.Brodsky, A.S.Goldhaber and J.R.Gunion, Phys.Rev.Lett. 47 (1981) 267.

[2] 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; S.J.Brodsky, in: Proceedings of the XIII International Symposium on Multiparticle Dynamics, Volendam, Netherlands. Eds. E.W.Kittel, W.Metzger and A.Stergion, World Scientific, Singapore, 1982, p. 963; S.J.Brodsky and A.H.Mueller: Phys.Lett. B206 (1988) 685.

[3] N.N.Nikolaev and J.Speth, The ELFE Project: an Electron Laboratory for Europe, edited by J.Arvieux and E.De Sanctis, Editrice Compositori, Bologna (1993) pp.179-212.

[4] NE18 Collaboration: N.C.R.Makins et al., MIT preprint LNS-93/11.

[5] E665 Collaboration: G.Fang, Talk at the PANIC, Perugia, July 1993, and private communication.

[6] L.L.Frankfurt, G.A.Miller and M.I.Strikman, “Colour Transparency and Nuclear Phenomena”, Comments in Particle and Nuclear Physics 21 (1992) 1 and references therein; B.Jennings and G.A.Miller, Phys.Lett. B236 (1990) 209; Phys.Rev. D44 (1991) 692.

[7] N.N.Nikolaev, A.Szczurek, J.Speth, J.Wambach, B.G.Zakharov and V.R.Zoller, Nucl.Phys. A567 (1994) 781.

[8] N.N.Nikolaev, A.Szczurek, J.Speth, J.Wambach, B.G.Zakharov and V.R.Zoller, Phys.Lett. B317 (1993) 287.

[9] B.Z.Kopeliovich, J.Nemchik, N.N.Nikolaev and B.G.Zakharov, Phys. Lett. B309 (1993) 179.

[10] B.Z.Kopeliovich, J.Nemchik, N.N.Nikolaev and B.G.Zakharov, Phys. Lett. B324 (1993) 469.

[11] B.Z.Kopeliovich, Sov.J.Part.Nucl. 21 (1990) 117.

[12] N.N.Nikolaev, Quantum Mechanics of Color Transparency. Comments on Nuclear and Particle Physics 21 (1992) 41; N.N. Nikolaev, High Energy Nuclear Reactions in QCD: Colour Transparency Aspects, in Lecture Notes of 1992 RCNP Kikuchi School on Spin Physics at Intermediate Energies, 16-19 November 1992, edited by M.Fujiwara and M.Kondo, RCNP-P-128, Osaka University.

[13] N.N.Nikolaev, Color transparency: Novel test of QCD in nuclear interactions. Surveys in High Energy Physics, 1994, in press.
[14] N.N. Nikolaev and B.G. Zakharov, *Proceedings of the Workshop on Electron Nucleus Scattering*, Marciana Marina, Elba, 5-10 July 1993, edited by O. Benhar, A. Fabrocini and R. Schiavilla, World Scientific (1994) p.323.

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

[16] V. Barone, M. Genovese, N.N. Nikolaev, E. Predazzi and B.G. Zakharov, Z. Phys. **C58** (1993) 541; Int. J. Mod. Phys. A, **8** (1993) 2779; Phys. Lett. **B326** (1994) 161.

[17] N. Nikolaev and B.G. Zakharov, The triple-pomeron regime and the structure function of the pomeron in the diffractive deep inelastic scattering at very small $x$, *Landau Inst. preprint Landau-16/93* and *Jülich preprint KFA-IKP(Th)-1993-17*, June 1993, to appear in Z. Phys. C (1994); The pomeron in the diffractive deep inelastic scattering, to appear in JETP **78(5)** (1994).

[18] N.N. Nikolaev and B.G. Zakharov, On determination of the large-$\frac{1}{x}$ gluon distribution at HERA, *Jülich preprint KFA-IKP(TH)-1994-12*, March 1994, to appear in Phys. Lett. B (1994).

[19] N.N. Nikolaev, B.G. Zakharov and V.R. Zoller, JETP Letters **59** (1994) 8; The BFKL and GLDAP regimes for the perturbative QCD pomeron, *Jülich preprint KFA-IKP(TH)-1994-02*, January 1994; JETP **78(6)** (1994); The spectrum and solutions of the generalized BFKL equation for total cross sections, *Jülich preprint KFA-IKP(TH)-1994-01*, January 1994, to appear in Phys. Lett. B (1994);

[20] N.N. Nikolaev and B.G. Zakharov, Phys. Lett. B**327** (1994) 149.

[21] B.Z. Kopeliovich and B.G. Zakharov: Phys.Rev. D**44** (1991) 3466.

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

[23] J. Nemchik, N.N. Nikolaev and B.G. Zakharov, Scanning the BFKL pomeron in elastic production of vector mesons at HERA, *Jülich preprint KFA-IKP(Th)-1994-17*, May 1994, submitted to Phys. Lett. B.

[24] O. Benhar, S. Fantoni, N.N. Nikolaev and B.G. Zakharov, Workshop at CEBAF at High Energies, 14-16 April 1994, CEBAF, and paper in preparation.

[25] NMC collaboration: A. Sandacz, private communication.

[26] S.J. Brodsky et al., SLAC-PUB-6412R (1994).

[27] J. Nemchik, N.N. Nikolaev and B.G. Zakharov, Anomalous $A$-dependence of diffractive leptoproduction of radial excitation $\rho'(2S)$, *Jülich preprint KFA-IKP(Th)-1994-18*, May 1994, submitted to Phys. Lett. B.

[28] NMC Collaboration: C. Mariotti, *Proceedings of the Joint International Lepton-Photon Symposium and Europhysics Conference on High-Energy Physics*, 25 July - 1 August 1991, Geneva, editors S. Hegarty, K. Potter and E. Quercigh, World Scientific, v.1, p.154;
[29] P.Amaudruz et al., Nucl. Phys. B371 (1992) 533.
[30] M.D.Sokoloff et al., Phys. Rev. Lett. 57 (1986) 3003.
[31] G.N.Farrar, L.L.Frankfurt, K.F.Liu and M.Strikman, Phys.Rev.lett. 64 (1990) 2996; Phys.Rev.Lett. 61 (1988) 686.
[32] V.N.Gribov, Sov.Phys.JETP 29 (1969) 483; 30 (1970) 709.
[33] V.A.Karmanov and L.A.Kondratyuk, JETP Letters 18 (1973) 266.
[34] L.Lesniak, Phys. Lett. B302 (1993) 140.
[35] N.N.Nikolaev, Color transparency: facts and fancy. Int. J. Mod. Phys. E: Reports on Nuclear Physics 3(2) (1994).
[36] N.N.Nikolaev, A.Szczurek, J.Speth, J.Wambach, B.G.Zakharov and V.R.Zoller, Multiple scattering effects in the transverse-momentum distributions from (ee’p) reactions, Juelich preprint KFA-IKP(Th)-1993-29, submitted to Nucl. Phys. A.
[37] N.N.Nikolaev, A.Szczurek, J.Speth, J.Wambach, B.G.Zakharov and V.R.Zoller, Theoretical interpretation of the NE18 experiment on nuclear transparency in A(e,e’p) scattering, submitted to Phys. Rev. C.
[38] R.J.Glauber, in: Lectures in Theoretical Physics, v.1, ed. W.Brittain and L.G.Dunham. Interscience Publ., N.Y., 1959; R.J.Glauber and G.Matthiae, Nucl. Phys. B21 (1970) 135.
[39] E.J.Moniz et al., Phys.Rev.Lett. 26 (1971) 445.
[40] S.Boffi, C.Giusti and F.D.Pacati, Phys.Rep. 226 (1993) 1.
[41] N.N.Nikolaev, A.Szczurek, J.Speth, J.Wambach, B.G.Zakharov and V.R.Zoller, Phys.Lett. B317 (1993) 281.
[42] C.Lechanoine-LeLuc and F.Lehar, Rev. Mod. Phys. 65 (1993) 47.
[43] K.Gottfried, Ann.Phys. 21 (1963) 29; W.Czyz and K.Gottfried, Nucl.Phys. 21 (1961) 676; Ann.Phys. 21 (1963) 47; N.Srivastava, Phys.Rev. 135B (1964) 612.
[44] H.Araseki and T.Fujita, Nucl.Phys. A439 (1985) 681; K.Saito and T.Uchiyama, Z.Phys. A322 (1985) 299.
[45] D.B.Day, J.S.McCarthy, T.W.Donnelly and I.Sick, Ann.Rev.Nucl. Sci. 40 (1990) 357; C.Cioffi degli Atti, E.Pace and G.Salme, Phys.Rev. C43 (1991) 1155.
[46] N.N.Nikolaev, JETP Lett. 57 (1993) 85.
[47] I.M.Narodetsky, Yu.A.Simonov and F.Palumbo, Phys.Lett. B58 (1975) 125; S.J.Brodsky and G.P.Lepage, in Perturbative Quantum Chromodynamics, editor A.H.Mueller, World Scientific, Singapore, 1989.
[48] N.Isgur and C.H.Llewellyn Smith, Nucl. Phys. B317 (1989) 527; Phys. Rev. Lett. 52 (1984) 1080.
[49] P. Stoler, Phys. Rev. D44 (1991) 1973; Phys. Rep. C226 (1993) 103.

[50] N.N. Zotov and A.V. Tzarev, Sov. Phys. Uspekhi 51 (1988) 121; G. Alberi and G. Goggi, Phys. Rep. 74 (1981) 1.

[51] N.N. Nikolaev, A. Szczurek, J. Speth, J. Wambach, B.G. Zakharov and V.R. Zoller, Pions in the lightcone nucleon and electromagnetic form factors, Jülich preprint KFA-IKP(TH)-1993-35, to appear in Z. Phys. A (1994).

[52] NA51 Collaboration: L. Kluberg, Workshop on Deep Inelastic Scattering and related problems, Eilat, February 6-9, 1994.

[53] N.N. Nikolaev, A. Szczurek, J. Speth, J. Wambach, B.G. Zakharov and V.R. Zoller, Pion content of the proton as seen in the Derell-Yan process, Jülich preprint KFA-IKP(TH)-1993-35, submitted to Phys. Rev. Lett. (1994).

[54] N.N. Nikolaev, A. Szczurek, J. Speth, J. Wambach, B.G. Zakharov and V.R. Zoller, Quark-hadron duality and color transparency sum rules for $A(e, e'p)$ scattering, paper in preparation.

[55] L.G. Dakhno and N.N. Nikolaev, Nucl. Phys. A436 (1985) 653.

[56] B. Jennings and B.Z. Kopeliovich, Phys. Rev. Lett. 70 (1993) 3384.

[57] A. Bianconi, S. Boffi and D.E. Kharzeev, Phys. Lett. B305 (1993) 3384; Nucl. Phys. A565 (1993) 767.

[58] G. Fang, talk at the Workshop on CEBAF at Higher Energies, 14-16 April 1994, CEBAF.

[59] T.H. Bauer et al., Rev. Mod. Phys. 50 (1978) 261; A. Donnachie and H. Mirzae, Z. Phys. C33 (1987) 407; A. Donnachie and A.B. Clegg, Z. Phys. C45 (1990) 313, and references therein.

[60] M. Kossov, talk at the Workshop on CEBAF at Higher Energies, 14-16 April 1994, CEBAF.
Table 1. Probabilities $P^{(\nu)}$ for $\nu$-fold rescatterings (in per cent) for the $(e, e'p)$ reaction on $^{12}$C (upper block) and $^{208}$Pb (lower block). In addition we show the nuclear transparency $W^{(0)}$, for quasifree knockout in the parallel kinematics, the nuclear transparency for the semi-exclusive reaction rate $T_A$ and the average number of rescattering $\langle \nu \rangle$.

| $Q^2$ | 0   | 1   | 2   | 3   | 4   | $W^{(0)}$ | $T_A$ | $\langle \nu \rangle$ |
|-------|-----|-----|-----|-----|-----|-----------|-------|---------------------|
| 1     | 65.15 | 23.40 | 8.09 | 2.49 | 0.67 | 0.63 | 0.97 | 0.51            |
| 1.2   | 68.37 | 22.06 | 6.96 | 1.98 | 0.50 | 0.60 | 0.87 | 0.45            |
| 1.5   | 71.94 | 20.39 | 5.76 | 1.48 | 0.34 | 0.56 | 0.77 | 0.38            |
| 2     | 74.18 | 19.26 | 5.03 | 1.21 | 0.26 | 0.54 | 0.72 | 0.34            |
| 4     | 81.31 | 15.24 | 2.87 | 0.49 | 0.08 | 0.54 | 0.67 | 0.23            |
| 6     | 83.40 | 13.89 | 2.31 | 0.35 | 0.05 | 0.55 | 0.66 | 0.19            |
| 8     | 85.65 | 12.33 | 1.76 | 0.23 | 0.03 | 0.56 | 0.65 | 0.17            |
| 1     | 30.13 | 21.60 | 16.09 | 11.69 | 8.12 | 0.25 | 0.84 | 1.98            |
| 1.2   | 39.32 | 23.70 | 15.00 | 9.39 | 5.70 | 0.23 | 0.57 | 1.46            |
| 1.5   | 49.36 | 24.44 | 12.81 | 6.72 | 3.46 | 0.20 | 0.41 | 1.03            |
| 2     | 55.08 | 24.12 | 11.22 | 5.25 | 2.43 | 0.19 | 0.34 | 0.83            |
| 4     | 67.45 | 21.54 | 7.30 | 2.48 | 0.83 | 0.19 | 0.28 | 0.49            |
| 6     | 70.66 | 20.39 | 6.23 | 1.91 | 0.57 | 0.20 | 0.28 | 0.42            |
| 8     | 74.37 | 18.79 | 5.02 | 1.34 | 0.35 | 0.20 | 0.27 | 0.35            |
Figure captions:

Fig. 1 - Qualitative dependence of the universal dipole cross section $\sigma(r)$ on the dipole size $r$ for scattering of the color dipole on the nucleon target. Different processes probe this universal dipole cross section at different $r$ [12].

Fig. 2 - The qualitative pattern of $Q^2$-dependent scanning of the wave functions of the ground state $V$ and the radial excitation $V'$ of vector meson [9,27]. The scanning distributions $\sigma(r)\Psi_{r'}(r)$ shown by the solid and dashed curve have the scanning radii $\rho_S$ differing by a factor 3. All wave functions are in arbitrary units.

Fig. 3 - Predictions [23] for the exclusive production of vector mesons $\rho^0, \phi^0$ vs. the NMC data [25].

Fig. 4 - Predictions [10] of nuclear transparency $T_A = \sigma_A/A\sigma_N$ for the incoherent exclusive production of $\rho^0$ mesons vs. the E665 data [5].

Fig. 5 - Predictions [10] of nuclear transparency $T_A^{(coh)}/T_C^{(coh)} = [144 d\sigma_A/A^2 d\sigma_C]_{q^2=0}$ for the forward coherent production of the $\rho^0$ mesons.

Fig. 6 - Predictions [10] of the $Q^2$ dependence of the ratio of cross sections $R_{coh}(A/C) = 12\sigma_A/A\sigma_C$ for coherent production of the $\rho^0$ meson vs. the E665 data [5]. The arrows indicate predictions for the complete CT (vanishing FSI).

Fig. 7 - Predictions [10] of the $Q^2$ dependence of exponents of parametrizations $\sigma_A (inc) \propto A^{\alpha_{inc}}, [d\sigma_A (coh)/dq^2]_{q^2=0} \propto A^{\alpha_{coh}}(q^2=0)$ and $\sigma_A (coh) \propto A^{\alpha_{coh}}$ vs. the E665 data [5].

Fig. 8 - Energy dependence of the ratio of transparency in quasielastic $J/\Psi$ photoproduction on tin and carbon predicted [22] from Eq. (25) (shown by solid line) versus the NMC data [28] for the 200 and 280 GeV primary muons [34]. The dotted line shows the prediction of the classical expansion model by Farrar et al. [31].

Fig. 9 - The predicted $Q^2$ and $\nu-$dependence of the nuclear transparency in the virtual photoproduction of the $\rho^0$-mesons [9].

Fig. 10 - The predicted $Q^2$ and $\nu-$dependence of the nuclear transparency in the virtual photoproduction of the heavy quarkonia [9]. The qualitative pattern is the same from the light to heavy nuclei.

Fig. 11 - The $Q^2$ dependence of the $\rho'(2S)/\rho^0(1S)$ ratio of forward production cross sections, which exhibits a dip because of the exact node effect in the $\rho'(2S)$ production.

Fig. 12 - The $Q^2$ and $A$ dependence of nuclear transparency for the $\rho^0(1S)$ and $\rho'(2S)$ electroproduction on nuclei.

Fig. 13 - Predictions [24] for the onset of antishadowing in the incoherent real photoproduction of radial excitation $\phi'(2S)$ on nuclear targets.

Fig. 14 - Predictions [24] for nuclear transparency $T_A = \sigma_A/A\sigma_N$ in electroproduction of the $\phi^0(1S)$ and $\phi'(2S)$ on nuclei.
Fig. 15 - Predictions [24] for the CT and node-effect driven steep $Q^2$ dependence of nuclear transparency in the electroproduction of the $\phi'(2S)$ on nuclei.

Fig. 16 - The final-state interaction in quasielastic $(e,e'p)$ scattering. FSI regenerates the observed proton $|p\rangle$ from electroproduced intermediate states $|i\rangle$.

Fig. 17 - Predictions [36] for the transverse missing momentum distribution of struck protons from $^{12}C(e,e'p)$ (left hand side) and $^{208}Pb(e,e'p)$ reactions (right hand side) for different values of $Q^2$. The solid lines correspond to the full momentum distributions. The PWIA component of the missing momentum distribution ($\nu = 0$) is shown by the long-dashed line. The $\nu$-fold rescattering contributions are represented by the short-dashed lines ($\nu = 1$), the dash-dotted lines ($\nu = 2$) and the dotted lines ($\nu = 3$) respectively.

Fig. 18 - Predictions [37] for the $Q^2$ dependence of nuclear transparency are compared with the NE18 data [4]. The dashed curve denotes the $p_\perp$-integrated nuclear transparency $T^A_\nu$; the dotted curve gives the nuclear transparency $W^{(0)}$ measured in parallel kinematics at $p_\perp = 0$; the solid curve represents the nuclear transparency $T^A_{\nu}(NE18)$ including the NE18 acceptance cuts [4]; the dot-dashed curve in panel b) shows the predicted [37] effect of the onset of CT.

Fig. 19 - The predicted [54] mass spectrum (the solid line) in the diffraction dissociation of protons $pp \rightarrow pX$ into state of invariant mass $M$. The data are taken from the compilation [50]. Contribution to the mass spectrum from excitation of the $\pi N$ continuum is shown by the dotted line; contributions from excitation of resonances corresponding to the positive parity excited shells is shown by dashed lines.

Fig. 20 - Predictions [54] for the $Q^2$-dependence of the strength of FSI $\Sigma^{ep}(Q^2)$ in $^{12}C(e,e'p)$ reaction (left box) and of nuclear transparency $T^A_\nu$ for different nuclear targets.

Fig. 21 - The top box: Predictions [54] for the $Q^2$ dependence of the scaled strength of FSI $\eta = \Sigma^{ep}(Q^2)/\sigma_{tot}(pN)$ in $^4He(e,e'p)$ reaction (the solid line). Shown separately are the contributions to $1 - \eta$ which is the signal of CT, from excitation of intermediate $\pi N$ state (the dotted line) and from excitations of nucleonic resonances (the dashed line).

The bottom box: Predictions [54] for the CT effect in the rate of $^4He(e,e'p)$ reaction with high transverse missing momentum. The quantity shown is the relative strength of the elastic-rescattering tail in the $p_\perp$-distribution which is given by $\eta^2 = (\Sigma^{ep}(Q^2)/\sigma_{tot}(pN))^2$.

Fig. 22 - The predicted [8] dependence of nuclear transparency $T^A_\nu$ (upper part) and of the observable $\Sigma^{ep}(x,Q^2)$ (lower part) on the Bjorken variable $x$ for different values of $Q^2$ (in $(GeV/c)^2$) listed in the figure. The dashed curves are obtained if only the mean-field component of the single-particle momentum distribution is retained in evaluation of $F_A^\nu(x)$ in Eq. (82). top results are for $^{12}C(e,e'p)$ scattering.
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