Exclusion high-$Q^2$ electroproduction:
light-cone wave functions
and electromagnetic form factors of mesons.

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Abstract

In the light-cone technique, one can describe exclusive electroproduction treating mesons
and baryons as (non-pointlike) partons of the nucleon. In this technique, the non-
perturbative light-cone wave function of the meson-baryon Fock state of the physical
nucleon defines the universal, projectile independent, density (flux) of on-mass-shell
mesons in the proton. We apply the light-cone technique to electroproduction reaction
$ep \rightarrow e\pi^+n$ long considered as a means of measuring the electromagnetic form factor
of the pion. We show that the interpretation of the long-standing puzzle of large trans-
verse cross section ($\sigma_T$) in terms of the $\gamma^*\rho \rightarrow \pi$ transition on the $\rho$ mesons in the
light-cone proton is possible, but requires quite a slow decrease of the $F_{\rho\pi}(Q^2)$ form
factor. This interpretation can be tested in the related $ep \rightarrow e\pi^0p$ reaction.

Corrections which are due to the final-state meson-baryon interactions (FSI) are
evaluated and are shown to amount to a 25% effect at moderately large $Q^2$. Vanishing
FSI with increasing $Q^2$ - the color transparency phenomenon - is shown to be very
strong.
1 Introduction

The data on pion electroproduction above the resonance region are analyzed in terms of nucleon and pion pole diagrams, shown in Fig.1. If the kinematics are chosen appropriately one argues that the pion pole diagram dominates [1]. This observation is used in determining the pion electromagnetic form factor [2,3]. Also it has been predicted [2,3,4], that at high $Q^2$ the total cross section $\sigma(ep \rightarrow \pi^+ n)$ is dominated by the longitudinal (scalar) component ($\sigma_L$) and the ratio $\sigma_L/\sigma_T$ grows rapidly at high $Q^2$ reaching the value $\sigma_L/\sigma_T \approx 15$ already at $Q^2 \approx 3 GeV^2$. The analysis of data taken at small and large values of the polarization parameter made it possible to separate the longitudinal and transverse components of the electroproduction cross section for $Q^2 = 1.2 - 3.3 GeV^2$ [2]. Surprisingly, the measured ratio $\sigma_L/\sigma_T$ appeared to be decreasing function of $Q^2$ which takes the value $\sigma_L/\sigma_T \sim 1$ at $Q^2 = 3.3 GeV^2$ in strong contradiction with the theoretical estimates [2,3,4].

Anticipating the results, we explain that the effect of strong absorption of the transverse photons is entirely due to the admixture of the transversely polarized $\rho$ mesons in the light-cone proton and the final state pion is generated in the magnetic dipole transition $\gamma^* \rho \rightarrow \pi$. The longitudinal component of total cross section at high $Q^2$ is determined by the light-cone density of the $\pi N$ Fock states in the proton.

The observables in exclusive electroproduction are very sensitive to the dynamics of the momentum transfer in the $\pi NN$ vertex [5]. In the traditional approach the uncorrelated monopole (dipole) $t$- and $u$-channel pion-nucleon form factors are introduced. Such a description has many flaws. In particular it violates the electric charge conservation and do not satisfy the energy-momentum sum rules. An attempt to minimize the momentum non-conservation effects by the appropriate choice of the $\pi NN$ form factor has been undertaken in [6].

On the way toward the consistent relativistic description of the meson-baryon component of the nucleon the Llewellyn Smith’s (LS) paper [7] was of the particular importance. In [7] the role of the energy-momentum sum rules has been emphasized and the LS-ansatz (see below) has been formulated as a constraint on the functional form of the flux of pions and nucleons. In [8,9] the light-cone wave function (LCWF) technique [10] was extended to the meson/baryon exchange processes and the light-cone densities of $\pi N$, $\pi \Delta$, $K \Lambda$, $\eta N$ Fock states have been derived. The corollary of the LCWF is that the LS-ansatz and the local gauge invariance, and momentum conservation thereof, are satisfied by construction.

This approach has been applied recently to the processes involving vector mesons. In this case, as well as in the case of spin-3/2, (corresponding Lagrangians involve the couplings with derivatives), some care is needed to avoid the conflict with gauge invariance, as it has been emphasized in [9]. The results consistent with gauge invariance were obtained in ref. [11].

The light-cone approach has one more important virtue evidently related to those mentioned above. Namely, it allows one to interpret the densities of mesons and baryons as the (non-perturbative) partons of the physical nucleon and, what is much more substantial, to study them like parton densities in inclusive DIS. The only difference is, however, that the exclusive final states have to be analysed in the one-meson electro-
production with the partonic kinematics. Indeed, in this case the photon of energy $\nu$ and virtuality $Q^2$, greatly exceeding the parton (meson/baryon) virtuality $k^2$,

$$Q^2 \gg m^2, k^2, \ (1)$$

probes the density of mesons and/or baryons in the proton at the light-cone Sudakov variable

$$\alpha = \frac{Q^2}{2m_p\nu} \equiv x. \ (2)$$

It is important that the cross section of the reaction $ep \rightarrow eMB$ does factorize

$$\frac{d\sigma}{dx dQ^2} \sim f_{\pi N}(x) F_{\pi(N)}^2(Q^2). \ (3)$$

Here, $f_{\pi N}(\alpha)$ is the parton (meson/baryon) density, which scales, and $F_{\pi(N)}(Q^2)$ is the on-mass-shell electromagnetic form factor of the struck parton.

Factorization implies a possibility of separate analysis of the light-cone meson-baryon density functions and the electromagnetic form factors of mesons and baryons. Consideration of the problems related to the practical realization of such a program is the subject of the present paper.

Previously, the pion electroproduction at small $Q^2$ has been studied in terms of a pion distribution function in [12]. For the perturbative QCD consideration of the exclusive processes see ref.[13].

2 Single pion electroproduction: light-cone parton densities and form factors

We find it convenient to use a light-cone momentum notation where a vector is denoted by

$$k_\mu = (k_-, k_+, k_\perp) = (\beta p_-, \alpha p_+, k_\perp)$$

with

$$k_\perp = (k_0 \pm k_3)/\sqrt{2}.$$ 

For the process (4) we choose a frame where

$$p_\mu = (p_-, p_+, 0, 0)$$

and

$$q_\mu = (q_-, q_+, 0, 0)$$

with the Bjorken variable

$$x = -q^2/2pq = \alpha = -q+/p_+ \left(1 + O\left(m_N^2/Q^2\right)\right).$$

That is we choose $p_+q_- \gg p_-q_+$. Still, it is assumed that the invariant mass, $W$, of the final hadronic state is large enough, $W^2 \approx Q^2(1 - x)/x > 5 GeV^2$, so that we do not encounter the problems of the resonance physics [14].
The differential cross section of the exclusive reaction (Fig. 1a)

\[ \sigma(ep \rightarrow e\pi^+ n) \]

takes the form

\[
\frac{d\sigma(ep \rightarrow e\pi^+ n)}{dx dQ^2 dk_\perp^2} = 2K_L(x, Q^2)f_{\pi N}(1 - x, k_\perp^2) F_\pi^2(Q^2) \\
+ 2K_T(x, Q^2)f_{\rho N}(1 - x, k_\perp^2) \frac{1}{8}Q^2 F_\rho^2(Q^2)
\]

(5)

where \( f_{\pi N}(\alpha, k_\perp^2) \) is the density of nucleons with the transverse momentum \( k_\perp \) and the fraction \( \alpha \) of the proton’s light-cone momentum. While \( f_{\pi N}(1 - \alpha, k_\perp^2) \) represents the density of pions with the momentum fraction \( 1 - \alpha \), in agreement with the LS-ansatz.

In other words \( f_{\pi N} \), being normalized to the number of nucleons (pions) \( n_{N(\pi)} \), represents the square of the entire LCWF of the \( \pi N \) Fock state

\[ f_{\pi N}(\alpha, k_\perp^2) = |\phi_{\pi N}(\alpha, \vec{k}_\perp)|^2 \]

(6)

which involves not only radial but also the spin-orbit degrees of freedom. With the \( \pi NN \) vertex of the form \( g_{\pi NN} \bar{u}\gamma_5 u \phi_\pi \Gamma \) the density of \( \pi^0 p \) states reads [8,9]

\[ f_{\pi N}(\alpha, k_\perp^2) = \frac{g_{\pi NN}^2 [m_N^2 (1 - \alpha)^2 + k_\perp^2]}{16\pi^2 \alpha^2 (1 - \alpha)} |\phi(M_{\pi N}^2(\alpha, k_\perp^2))|^2. \]

(7)

Hence, the isospin factors 2 in rhs of eq.(5)

The radial part of the LCWF,

\[ \varphi(M^2) = \frac{\Gamma(M^2)}{M^2 - m_\rho^2}, \]

(8)

depends only on the invariant mass squared of the two-body meson-baryon (MB) state,

\[ M_{MB}^2(\alpha, k_\perp^2) = \frac{m_B^2 + k_\perp^2}{\alpha} + \frac{m_M^2 + k_\perp^2}{1 - \alpha}. \]

(9)

The vertex function \( \Gamma(M^2) \) is parametrized as

\[ \Gamma(M_{MB}^2) = \exp \left[ -\frac{1}{2} R_{MB}^2 \left( M_{MB}^2 - M_0^2 \right) \right]. \]

(10)

Originally [8,9], the \( MNB \) couplings, \( g_{MNB} \), as well as still another nonperturbative parameter - the radii \( R_{MB} \) of the meson baryon Fock state were inferred from an analysis of the experimental data on the fragmentation of high-energy proton into nucleons and hyperons - the process dominated by stripping off the mesons of the meson-baryon Fock states [15]. The normalization to the point \( M_0^2 = (m_B + m_M)^2 \) of the \( MB \)-system threshold was accepted. The continuation into the unphysical region of the meson pole, \( M_0^2 = m_N^2 \), accepted in many of the recent papers, calls for knowledge of the final state interaction effects which are discussed below.
The kinematical factors $K_L$ and $K_T$ in (5) are as follows

$$K_T(x, Q^2) = \frac{4\pi \alpha_{em}^2}{Q^4}(1 - y + \frac{1}{2}y^2)$$  \hspace{1cm} (11)

$$K_L(x, Q^2) = \frac{4\pi \alpha_{em}^2}{Q^4}(1 - y), \quad y = \frac{Q^2}{xs}, \quad s = 2pp_e,$$  \hspace{1cm} (12)

where $p_e$ and $p$ are the four-momenta of the beam-electron and the target-proton, respectively.

In (5) $F_\pi(Q^2)$ is the on-shell charge form factor of the pion. It is worth recalling that in the light-cone parton model the condition (1) guarantees the on-shellness of partons [16]. Still, the dispersion integral representation [17] which deals only with the on-mass-shell internal particles in corresponding unitarity diagrams reproduces the above factorization. Note in this connection that the vertex function $\Gamma(M^2)$ (10) cuts off the large transverse momenta at $k_\perp^2 \sim \alpha(1 - \alpha)/R^2$. The typical $\alpha$'s we are dealing with are

$$0.2 \lesssim \alpha \lesssim 0.8$$  \hspace{1cm} (13)

In this specific region the reggeization effects can be neglected. Then, for $R^2 \approx 1 GeV^2$ the main contribution to the cross section (5) comes from rather small transverse momenta $k_\perp^2 \lesssim 0.3 GeV^2$.

The longitudinal momentum partition in the $\pi N -$ system, as it is prescribed by the $\pi N$ LCWF, is very far from uniformity. The nucleon carries the momentum fraction

$$\alpha \sim m_N/\left(m_N + \sqrt{m_N^2 + k_\perp^2}\right),$$  \hspace{1cm} (14)

only the rest, $1 - \alpha$, falls to the pion. Such a disbalance leads to a suppression of the pion-pole contribution to the cross section (5) for $\alpha \gtrsim 0.7$. By the same reasons the nucleon-pole contribution vanishes for $\alpha \lesssim 0.3$.

The second term in (5) is due to the $\rho^+ - \pi^+$ transition which is under the control of the form factor $F_{\rho\pi}(Q^2)$

$$\langle \pi(k')|J_\mu^{\rho\pi}(0)|\rho(k)\rangle = F_{\rho\pi}(Q^2)\epsilon^{\mu\nu\lambda\sigma}k'_\nu k_\lambda \rho_\sigma,$$  \hspace{1cm} (15)

where $\rho_\sigma$ is the $\rho$-meson polarization vector. Since the amplitude (15) is purely transversal, it has been predicted [18,19] that at asymptotically high $Q^2$

$$F_{\rho\pi}(Q^2) \sim Q^{-4}.$$  \hspace{1cm} (16)

At the same time, the magnetic dipole (M1) transition (15) generates an extra power of $Q^2$ in the electroproduction cross section (5).

In the region of $Q^2 \approx 5 - 10 GeV^2$, which we are interested in, the form factor $F_{\rho\pi}$ is not calculable within pQCD (for more discussion on the slow onset of the hard scattering regime in the form factors see [20]). This is the reason why we resort to the parametrization of $F_{\rho\pi}$ as

$$F_{\rho\pi}(Q^2) = g_{\rho\pi\gamma} \left(1 + Q^2/\Lambda_{\rho\pi}^2\right)^{-2},$$  \hspace{1cm} (17)
where $\Lambda_{\rho\pi}$ should be chosen to match the experimental data. The normalization is such that $F_{\rho\pi}(0) = g_{\rho\pi\gamma}$. The SU(3) quark model prediction \[21,22\] $g_{\rho\pi\gamma} = 2/3 \mu_p$, $\mu_p = 2.79/m_p$, which we rely upon, does not contradict the measured $\rho^\pm$ radiative decay width \[23\].

The density of the transversely polarized $\rho^0$-mesons in the proton, $f_{\rho N}^T$, can easily be obtained

$$f_{\rho N}^T(\alpha, k^2_\perp) = \frac{g_{\rho pN}^2}{16\pi^2} \frac{m_N^2}{\alpha^2(1-\alpha)} \times \left\{ 2 \left[ (1-\alpha)^2 + \tau(1+\alpha^2) \right] + 4(f/g)(1-\alpha)^2 + (f/g)^2(1-\alpha)^2(\tau+2)(\tau+1) \right\} \varphi(M_{\rho N}^2)^2. \tag{18}$$

Here $g$ and $f$ are the vector and tensor $\rho NN$ couplings and the dimensionless variable $\tau$ is introduced

$$\tau = \frac{k^2_\perp}{m_N^2(1-\alpha)^2}. \tag{19}$$

Our definition of the meson-baryon coupling constants as well as of the interaction Lagrangians corresponds to that accepted in \[24\], $g^2/4\pi = 0.84$, $f/g = 6.1$. Note that $\alpha$ in (18) is the light-cone fraction of the proton momentum carried by the nucleon in the $\rho N$ Fock state. The applicability region of the above formula for $f_{\rho N}^T(\alpha, k^2_\perp)$ is not much broader than $0.3 \lesssim \alpha \lesssim 0.7$.

The data \[2,3\] on electroproduction of single charged pions from hydrogen targets were taken for $W > 2.15$ GeV and $Q^2 = 0.6 - 9.8$ GeV$^2$ with values of the polarization parameter

$$\epsilon = \frac{1 - y - m_N^2 Q^2/s^2}{1 - y + \frac{1}{2} y^2 + m_N^2 Q^2/s^2}$$

in the range $0.35 < \epsilon < 0.45$. Combination with data taken at $\epsilon$ near 1 allowed the authors \[2,3\] to separate the contribution from transversely polarized and longitudinal photons in the range $1.2$ GeV$^2 < Q^2 < 3.3$ GeV$^2$.

Figure 2 shows the measured angular dependence of the transverse cross section $d\sigma_T/d\Omega_\pi$ and longitudinal $d\sigma_L/d\Omega_\pi$ components of the virtual photoproduction reaction $\gamma^* p \rightarrow \pi^+ n$ for two $(W, Q^2)$ points. The angle $\theta$ is the centre-of-mass angle between the pion and the virtual photon.

For the electromagnetic pion form factor we use the expression $F_\pi = (1 + Q^2/\Lambda_\pi^2)^{-1}$ with $\Lambda_\pi^2 = m_\rho^2$. The so-obtained longitudinal (scalar) cross section $d\sigma_L/d\Omega_\pi$ is shown in Fig.2. We conclude that there is no dramatic difference between the data and the pion-pole dominance model suggested long ago \[1\]. It is clear, however, that data leave enough room for the variations of the cut-off parameter $\Lambda_\pi$ as well as for the extra non-$\pi$-pole contributions to $\sigma_L$ \[5\].

The $\rho$-dominated transverse cross section is expected to be an increasing function of $\theta$ in the range $\theta = 0^0 - 20^0$. This effect is due to the strong $k_\perp$-dependence of the tensor term

$$\frac{f}{4m_N} \bar{\psi}\sigma_{\mu\nu}\psi(\partial^\mu \rho^\nu - \partial^\nu \rho^\mu)$$
in the $\rho NN$ interaction Lagrangian (terms $\propto f/g$, $\propto f^2/g^2$ in eq.(18)). Remind that the ratio of the tensor to vector coupling constant, $f/g$, in eq.(18) is large.

The measured angular dependence of the transverse component of the cross section is in quantitative agreement with our estimates at $Q^2 = 3.3\, GeV^2$.

However, the above comparison suffers from some uncertainties. First of all, the absolute normalization depends strongly on the accuracy of determination of the values $Q^2$ and $W$ since the $\rho N$ light-cone density (18) is a sharp function of the variable $x \simeq Q^2/(Q^2 + W^2)$ in the range of $x \simeq 0.2 - 0.3$ corresponding to the above values of $Q^2$ and $W$. The resolutions obtained in [2,3] are of about $0.5\, GeV^2$ in $Q^2$ and $1\, GeV$ in $W$, not better.

Being taken at the face value, the data would imply the substantial excess of $\rho$-mesons above the pions at $x \sim 0.2 - 0.3$. However, we keep the ratio of the $\rho N$ and $\pi N$ light-cone densities at approximately the same level as in [11] to avoid contradiction with data on the proton-neutron charge exchange inclusive reactions, the universal cut-off parameter in eq.(10) is put equal to $R^2 = 0.8\, GeV^{-2}$. Then, to reproduce the observed $Q^2$-dependence of the transverse cross section we have to assume that the form factor $F_{\rho \pi}(Q^2)$ has a very slow preasymptotics and matches the pQCD behaviour (16) only at very high $Q^2 \gg \Lambda_{\rho \pi}^2$. Our estimates presented in Fig.2 correspond to $\Lambda_{\rho \pi} = 3m_\rho$. Evidently, the accurate measurements of the $x$-dependence of the transverse component of the cross section must precede any determinations of the $\rho \pi \gamma$ form factor. Such an analysis is also of interest in its own right, as it provides direct information on the $\rho N$ light-cone density function.

The possibility to detect neutral mesons [25] makes the reaction

$$ep \rightarrow e\pi^0 p \quad (20)$$

particularly interesting as it is free of the pion-pole term and at high $Q^2$ is dominated by the transverse $\omega$-pole contribution. Its differential cross section reads

$$\frac{d\sigma(ep \rightarrow e\pi^0 p)}{dx dQ^2 dk^2_\perp} = K_T(x, Q^2) f_{\omega\pi}^T(1 - x, k^2_\perp) \frac{1}{8} Q^2 F_{\omega\pi}(Q^2), \quad (21)$$

where the density function $f_{\omega\pi}^T$ coincides with $f_{\rho N}^T$ at $f/g = 0$ and $g_{pp\rho\pi} \rightarrow g_{pp\omega\pi}$ [24].

The coupling constant $g_{\omega\pi\gamma}$ is large. Experimentally [26],

$$\frac{g_{\omega\pi\gamma}^2}{g_{\rho\pi\gamma}^2} \approx 10., \quad (22)$$

which is in agreement with the early $SU(3)$ quark model predictions [21,22]

$$F_{\rho\pm\pi\pm} = F_{\rho\pi^0} = \frac{1}{3} F_{\omega\pi}. \quad (23)$$

Due to the extra power of $Q^2$ typical of the magnetic dipole interaction the cross section of the single $\pi^0$ production at high $Q^2$ is dominated by the $\omega - \pi^0$ radiative transition. Evidently, it holds true even at the "normal" value of the parameter $\Lambda_{\omega\pi} \simeq m_\rho$. This observation makes the idea of the precision measurements of the $Q^2$-dependence of the cross section of the reaction $p(e, e'\pi^0)p$ [25] especially appealing.
The appearance of the pQCD-generated longitudinal component of the cross section (21) is discussed in [5].

Note, the term corresponding to the contribution of the diagram of Fig.1b

\[ K_T(x,Q^2)f_{\pi N}(x,k^2_{\perp})F_p^2(Q^2), \]

where \( F_p(Q^2) = F_1(Q^2) \) is the Dirac form factor of the proton is of the next order in powers of \( Q^{-2} \) and can be neglected. There is also the kinematical suppression of the small-\( x \) region, since the function \( f_{\pi N}(x,Q^2) \) peaks at \( x \approx 0.7 \) and steeply falls at \( x \to 0 \).

Note, the non-zero contribution to the cross section of the reaction (20) could come from the interference diagram of Fig. 1c. This is typically a non-partonic contribution. In our particular case the disbalanced kinematics of the \( \pi N \) system (see above) generates large invariant masses of the intermediate \( \pi N \) state which are about

\[ \sim \frac{m_N}{m_\pi} \]

times as large as the values of \( M^2_{\pi N} \) typical of the diagrams 1a,b. This is the reason why the interference diagram (Fig.1c) loses the competition with both the diagrams 1a and 1b for all \( x \)'s, independently of the specific form of the vertex function.

3 Final-state interactions and the color transparency phenomenon

In the reaction \( p(e,e'\pi^+)n \) at \( Q^2 \sim 5 - 10\, GeV^2 \) the cms energy squared of the pion-nucleon final state rescattering is large,

\[ s_{\pi N} \simeq \frac{1-x}{x}Q^2 \quad (24) \]

and we can use the Glauber approximation [27,28,29] for the final-state interaction (FSI) of the struck pion with the spectator neutron. Then the pion-nucleon density, \( f_{\pi N}(x,k^2_{\perp}) \), which plays the role of the momentum distribution of the observed pion (nucleon) can be rewritten as

\[ f^{FSI}_{\pi N}(x,k^2_{\perp}) = \left| \phi_{\pi N}(x,k_{\perp}) - \frac{1}{4\pi} \int d^2k_{\perp} \phi_{\pi N}(x,k_{\perp} - k_{\perp})f_{el}(k_{\perp}) \right|^2, \quad (25) \]

where

\[ f_{el}(k_{\perp}) = \frac{\sigma_{tot}^{\pi N}}{4\pi} \exp\left( -\frac{1}{2}b_{el}k^2_{\perp} \right) \quad (26) \]

is the \( \pi n \) elastic scattering amplitude with the slope parameter \( b_{el} = 10\, GeV^{-2} \) [30].

To integrate over \( k_{\perp} \) in (25) it is convenient to rewrite \( \phi_{\pi N} \) as a product

\[ \phi_{\pi N} = \chi \cdot \Gamma, \quad (27) \]

where \( \chi(x,k_{\perp}) \) is a smooth function of \( k_{\perp} \), while the most singular part of \( \phi_{\pi N} \) is the vertex function \( \Gamma(M^2_{\pi N}) \). Then factoring out of the integral the slowly varying function
\( \chi(x, \vec{k}_\perp) \) enables one to write the momentum distribution, corrected for FSI effects, \( f_{\pi N}^{FSI} \), in a factorized form:

\[
f_{\pi N}^{FSI}(x, k_\perp^2) = f_{\pi N}(x, k_\perp^2) \left[ 1 - \eta(x, k_\perp^2) \right]^2,
\]

where the attenuation function, \( \eta(x, k_\perp^2) \), equals

\[
\eta(x, k_\perp^2) = \frac{\sigma_{\pi N}^{tot}}{16\pi} \exp \left[ \frac{B^2(x)k_\perp^2}{B(x) + \frac{1}{2}b_{el}} \right]
\]

and

\[
B(x) = \frac{1}{2x(1-x)}R_{\pi N}^2.
\]

A quick estimate yields for the attenuation effect

\[
1 - 2\eta \approx 0.75.
\]

Still, the FSI effects distort the transverse momentum distribution. Thus, an accurate evaluation of the FSI effects is necessary for the quantitative interpretation of the data.

Eq.(25) implies that the attenuation is due to the interference of the elastic rescattering amplitudes. It has been shown in [31] that at high \( Q^2 \) the interference of elastic (\( \pi \), in our case) and inelastic (\( A_1 \), for instance) intermediate states gives rise to delicate cancellations of the contributions from the ground state and higher excitations thus resulting in CT phenomenon and/or weak final state interaction.

The patterns of the multiple scattering on both the nucleon and the nucleus are similar ones. But there is one important difference. A slow onset of CT on nuclei predicted in [31] and confirmed by the recently completed NE18 experiment [32] is due to the nuclear form factor suppression of the high mass intermediate states:

\[
M_{max}^2 \simeq \frac{s_{\pi N}}{m_NR_A}.
\]

Then to get the \( Q^2 \)-dependence of the strength of FSI denoted as \( \Sigma_{\pi N} \), it suffices to rescale the function \( \Sigma_{\pi N}(Q^2) \), already calculated in [34] making use of the diffraction operator technique, as follows

\[
\Sigma_{\pi N}(Q^2, x) = \Sigma_{\pi A} \left( \frac{R_A}{R_N}s_{\pi N} \right).
\]

Then, a quick estimate of the CT effects comes from substituting \( \sigma_{\pi N}^{tot} \) in eq.(29) by \( \Sigma_{\pi N}(Q^2, x) \). It can easily be seen that already at \( Q^2 = 10 \text{ GeV}^2 \) and \( x \sim 0.5 \) the cross section \( \Sigma_{\pi N}(Q^2) \) is approximately half the "normal" \( \pi N \) cross section at the energy \( \simeq 10 \text{ GeV} \) and vanishes at \( Q^2 \simeq 30 \text{ GeV}^2 \). Note that the rescaling (33) can not be quite exact at \( Q^2 \sim 3 \text{ GeV}^2 \), near the threshold of the first radial excitation (\( A_1 \)-meson). Due to the large \( \pi - A_1 \) mass splitting, there is a region of very slow variations of \( \Sigma_{\pi N}(s_{\pi N}) \) which is followed by a sharp fall off.

This observation is important for the correct continuation of the vertex function \( \Gamma(M^2) \), determined in the physical region of \( M^2 \geq (m_B + m_M)^2 \), to the meson-pole.
at \( M^2 = m_N^2 \). Indeed, any meson-baryon LCWF having a claim on description of the hadronic processes involves implicitly the FSI effect. Disappearance of the latter at high \( Q^2 \) implies that at moderate \( Q^2 \) the effective \( \pi NN \) coupling constant does not equal its standard pole value and must be renormalized. The renormalization amounts to the 25\% effect.

### 4 Conclusions

We analyze the high-\( Q^2 \) exclusive electroproduction in terms of the light-cone parton densities with mesons and baryons as the (non-perturbative) partons. In the partonic kinematics the observable cross sections are expressible directly in terms of the on-shell electromagnetic form factors. The transverse component of the cross section \( ep \rightarrow e\pi^+ n \) is dominated by the \( \rho \)-pole contribution. As a hint to the relevance of such an interpretation of data we consider the observed angular distribution of charged pions.

We consider the opportunity of measuring the magnetic dipole form factors \( \rho \pi \gamma, \omega \pi \gamma \) in the reaction with the charged and neutral pions as a very promising one [25]. In particular, it has been emphasized above that at high \( Q^2 \) the helicity non conserving \( \omega - \pi \) transition prevails in the \( p(e, e'\pi^0)p \) reaction due to both the specific \( Q^2 \)-enhancement of the \( M1 \) radiative transition and large value of the \( \omega \pi \gamma \) coupling.

Note, the early onset of the parton model regime enables to study the form factors \( F_{\rho \pi}(Q^2), F_{\omega \pi}(Q^2) \) in the substantially non-perturbative region and, what is more important, to retrace the onset of the pQCD regime at a very high \( Q^2 \) as well.

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Figure captions:

Fig. 1 - The one-meson electroproduction on the proton. The virtual photons, mesons and baryons are represented by wavy, dashed and solid lines, respectively.

Fig. 2 - The observed angular dependence of the transverse and longitudinal components of the cross section for the reaction $\gamma^* p \rightarrow \pi^+ n$ for the two $(W, Q^2)$ points. The data points are from [3]. Our estimates are shown by the solid curves.
$W = 2.65 \text{ GeV}$

$Q^2 = 2.0 \text{ GeV}^2$
$W = 2.65 \text{ GeV}$

$Q^2 = 3.3 \text{ GeV}^2$
$W = 2.65 \text{ GeV}$

$Q^2 = 2.0 \text{ GeV}^2$
$W = 2.65 \text{ GeV}$
$Q^2 = 3.3 \text{ GeV}^2$