Exclusive electroproduction and the quark structure of the nucleon

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Abstract

The natural interpretation of deep inelastic scattering is in terms of hard scattering on QCD constituents of the target. We examine the relation between amplitudes measured in exclusive lepto-production and the quark content of the nucleon. We show that in the Bjorken limit, the natural interpretation of amplitudes measured in these hard exclusive processes is in terms of the quark content of the meson cloud and not the target itself. In this limit, the most efficient representation of these exclusive processes is in terms of leading Regge amplitudes.

Key words: Generalized parton distributions, inclusive reactions, exclusive reactions, Regge phenomenology
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1. Introduction. Recent interest in hard exclusive lepto-production, in particular deeply virtual Compton scattering (DVCS) and meson production, has been stimulated by the idea that these processes may give new insights into the quark structure of the nucleon [1,2,3,4,5,6,7,8,9,10]. The connection between hard exclusive amplitudes and quark distributions in the nucleon, commonly referred to as generalized parton distributions, is formally analogous to that between the deep inelastic scattering cross-section and the structure functions. As shown by Feynman [11], structure functions can be interpreted in terms of quark probability distributions in the nucleon. Duality teaches us that, at least in principle, it is possible to use any channel to describe the scattering amplitude. The parton basis of deep inelastic scattering (DIS) is an example of a process that is most efficiently interpreted in the s-channel representation. The basis of quasi-free QCD constituents is the natural choice for expressing structure functions in the Bjorken scaling regime, $Q^2 \to \infty$ and finite $x_{BJ}$. 

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In this regime the relevant matrix elements are diagonal in the parton Fock space basis. However even in the case of DIS the s-channel parton representation becomes less useful in the limit \( x_{BJ} \to 0 \). In this wee parton regime it becomes more efficient to parametrize structure functions in terms of amplitudes associated with \( t \)-channel processes. The physical interpretation of the structure functions changes in between these two regimes. As \( x_{BJ} \to 0 \) the structure function evolves to represent ladders of partons originating from \( t \)-channel meson exchanges.

As in the case of DIS, a factorization theorem in exclusive lepto-production enables one to separate the hard quark-photon (alternatively, \( W \) or \( Z \) boson) scattering from the target (nucleon) contribution \([12]\). The latter contribution is typically parametrized by the generalized parton distributions or GPDs \([1,3,5]\). In analogy with deep structure functions the GPDs are often interpreted as corresponding to some quark distributions of the nucleon \([10]\). Just as in the case of DIS, one can interpret hard exclusive lepto-production either in terms of \( s \)-channel exchanges or via \( t \)-channel exchanges.

Recently Mueller and Schaefer \([13]\) produced a conformal spin expansion of GPDs. As part of their study they investigated the extent to which the GPDs displayed the characteristics of their leading Regge trajectories. When they examined the effective slope parameters for amplitudes corresponding to \( \omega \) and \( \rho \) exchange, they found them to be extremely close to the phenomenological slopes for those trajectories, a result they called “quite astonishing.” Also, a recent analysis of \( \omega \) electroproduction with the CLAS detector at Jefferson Laboratory \([14]\) found that their data agreed quite well with standard Regge phenomenology. The purpose of this letter is to show that in the Bjorken limit exclusive lepto-production amplitudes are most naturally described in terms of \( t \)-channel processes. Our results will demonstrate why one should expect that interpretation of the quark content of exclusive lepto-production processes will be most effectively discussed in the context of the parton content of Reggeons, rather than of the nucleon.

Consider the case of exclusive vector meson production at high-\( s \) and low-\( t \). As shown by a large body of evidence \([15,16,17,18]\), such processes can be described by \( t \)-channel exchanges, where sums over exchanged mesons with all possible spins can be described by Regge trajectories \([19]\). The amplitude for a given Regge trajectory has the behavior \( s^{\alpha(t)} \). At asymptotic energies, \( W \gg 10 \) GeV Pomeron exchange dominates \([20,21]\) since it has the largest intercept \( \alpha_P(0) \sim 1 \) and the process is predominantly \( s \)-channel helicity conserving.

In this paper, we present simple arguments to justify our claim that hard exclusive processes are most naturally understood in terms of \( t \)-channel exchanges. For comparison purposes, this involves a review of some very well known material in both deep inelastic scattering and Regge phenomenology.
Such a review is necessary in order to compare and contrast the underlying mechanisms that drive inclusive and exclusive lepto-production in the Bjorken limit, and to clarify the differences between the quark/nucleon amplitudes that can be extracted from these reactions.

2. The hadronic tensor in inclusive and exclusive lepto-production

Consider a deep inclusive reaction on a nucleon, \( a^\ast(q) + N(p) \rightarrow X \). Here \( a^\ast(q) \) is a virtual photon or weak gauge boson with momentum \( q \), \( N(p) \) is a nucleon with momentum \( p \), and \( X \) is the final state. To calculate the DIS cross section, one takes the square of this amplitude and sums over all unobserved final states \( X \). As is well known \[22\], the resulting inclusive cross section can be obtained from the discontinuity across the right hand energy cut of the forward virtual Compton amplitude. This is a special case of the general exclusive amplitude

\[
a^\ast(q) + N(p) \rightarrow b(q') + N(p') ,
\]

where \( a^\ast(q) \) is a virtual boson (\( \gamma, W \) or \( Z \)) with momentum \( q \), where \( -q^2 = Q^2 \). In Eq. (1), \( N(p), N(p') \) represent the initial and final nucleons with momenta \( p \) and \( p' \), respectively (\( p^2 = p'^2 = m_N^2 << Q^2 \)).

Forward elastic amplitudes, which describe deeply inclusive processes, are characterized by \( b(q') = a^\ast(q) \) and the kinematical relations \( q' = q, p' = p \). In contrast, for hard exclusive processes \( b(q') \) is typically a real photon, meson or meson resonance; therefore the momentum \( q' \) of \( b \) satisfies \( 0 \leq q'^2 \sim m_N^2 << Q^2 \). Since our main goal is to illustrate the differences between the hadronic contribution in exclusive and inclusive processes, in the following we will ignore spin and other internal degrees of freedom and assume only scalar currents. For these processes the hadronic contribution to the cross section is determined from the hadronic tensor (a scalar function under the above approximations),

\[
T = \int d^4z e^{i \frac{q}{2} \cdot z} \langle p'|T \left[ j(z)j(-\frac{z}{2}) \right]|p \rangle .
\]

In Eq. (2) \( j(z) = \phi^\dagger(z)\phi(z) \) represents a (scalar) quark current in the Heisenberg picture which couples to the external fields representing the \( a \) and \( b \) particles. The Heisenberg nucleon states represent fully interacting nucleons; in particular, they include the meson cloud contribution. To study the valence and sea parton content of the nucleon the bare nucleon is often introduced within models that separate the QCD interactions among partons from chiral meson-nucleon interactions \[23,24\]. The \( x_{BJ} = \mathcal{O}(1) \) region is then found to be dominated by the bare nucleon and the sea contributes in the limit \( x_{BJ} \rightarrow 0 \), as expected.

From Lorentz symmetry it follows that the amplitude \( T \) in Eq. (2) is a function of four independent Lorentz scalars, \( T = T(Q^2, \nu, t, q'^2) = T(Q^2, x_{BJ}, t, q'^2) \),
with, \( \nu = \frac{p \cdot q}{m_N} = Q^2/(2x_{BJ}m_N) \), and \( t = (p' - p)^2 = (q' - q)^2 \). For inclusive processes we require the forward amplitude characterized by \( t = 0, 0 > q'^2 = q^2 = -Q^2 \), while the kinematics for exclusive processes require \( t < 0, 0 \leq q'^2 \sim m_N^2 \). Using the operator product expansion to leading order in QCD one finds that the matrix elements of the time-ordered product of the quark currents can be replaced by the product of two quark field operators and the quark propagator:

\[
T(Q^2, \nu, t, q'^2) = i \int \frac{d^4z d^4k}{(2\pi)^4} \frac{e^{-ikz}}{(z + p')^2} \langle p' | T \left[ \phi^{\dagger}(\frac{z}{2})\phi(-\frac{z}{2}) \right] | p \rangle . \tag{3}
\]

Using Wick’s theorem the normal ordered product of fields (in the interaction picture with the interaction arranged as a power series of the QCD coupling) was replaced by the time-ordered product since the c-number difference between the two types of ordering does not contribute to connected matrix elements. The integral in Eq. (3) is dominated by points on the light cone with \( z^2 \sim O(1/Q^2) \). It is convenient to use light cone coordinates, \( A^\mu = (A^+, A^-, A_\perp) \) with \( A^\pm \equiv A^0 \pm A^2 \) and to choose the frame in which \( q^+ = 0, q'^2 = Q^2, p_\perp = 0 \).

For inclusive reactions where \( q' = q \), the quark propagator in Eq. (3) becomes

\[
\left( \frac{q + q'}{2} + k \right)^2 = Q^2 \left( \frac{x}{x_{BJ}} - 1 \right) - q_\perp \cdot k_\perp + k^2 \sim Q^2 \left( \frac{x}{x_{BJ}} - 1 \right) . \tag{4}
\]

where \( x \equiv k^+/p^+ \) is the fraction of the nucleon longitudinal momentum carried by the struck quark. The approximation is based on the observation that the matrix element in Eq. (3) does not involve hard scales and thus on average \( k_\perp, k^-, k^+ << |Q| \). Under such approximations the absorptive part of the hadronic tensor, \( W \equiv T(\nu + i\epsilon) - T(\nu - i\epsilon) \) which determines the DIS cross section is given by

\[
W(Q^2, x_{BJ}) = \frac{1}{Q^2} \int dx \delta(x - x_{BJ}) F(x, Q) , \tag{5}
\]

where \( F(x, Q) \) is the structure function,

\[
F(x, Q) = \frac{1}{2} p^+ \int dz^- e^{-ixp^+z^-/2} \langle p | T \left[ \phi^{\dagger}(\frac{z}{2})\phi(-\frac{z}{2}) \right] | p \rangle_{z^+=0, |z_\perp| < \frac{1}{Q}} . \tag{6}
\]

For exclusive production with \( Q^2 >> q'^2 \geq 0, (p - p')^2 = t < 0 \), again using light cone coordinates, to leading order in \( O(Q^2) \) the quark propagator can be approximated by

\[
\left( \frac{q + q'}{2} + k \right)^2 = \frac{Q^2}{2} \left( \frac{x}{\xi} - 1 \right) . \tag{7}
\]
In the Bjorken limit $\xi = x_{BJ}/(2 - x_{BJ})$, and the longitudinal component of quark momentum in this case is $x = k^+/P^+$ with $P^+ \equiv (p^+ + p^{'+})/2$. The hadronic tensor for exclusive processes in the Bjorken limit is therefore given by,

$$T(Q^2, \nu, t, q^2) = \frac{P^+}{Q^2} \int \frac{dz^-}{(2\pi)} \frac{dx}{2} i \xi e^{-ip^+z^-} \langle p'| T \left[ \phi^\dagger \left( \frac{z}{2} \right) \phi \left( -\frac{z}{2} \right) \right] |p\rangle_{z^+=0, |z|<\frac{Q^2}{2t}}.$$

The positive energy cut contribution to the hadronic tensor, which determines the inclusive cross section, forces $x = x_{BJ} > 0$. This is not the case in exclusive processes; here the full amplitude $T$ is needed to determine the cross section, so it contains an integral over both positive and negative $x$. Defining the free quark and anti-quark creation and annihilation operators in the standard way in terms of field operators, it is possible to reinterpret the integration over the negative-$x$ region of the quark (anti-quark) operator matrix elements in terms of the positive-$x$ region of the anti-quark (quark) operator matrix elements [8]. Thus in the quark representation the matrix element in Eq. (8) receives contributions from pair creation and pair annihilation operators which mix different Fock space sectors of the nucleon wave function. Thus unlike DIS the DVCS matrix elements require nondiagonal overlaps of light front wave functions [25,26]. More detailed analysis of the correspondence between current matrix elements and the light cone wave function representation is given in [25]. We also note that calculations of exclusive cross sections based on GPD models that employ the quark handbag phenomenology also include explicit contributions from meson exchanges, most notably an elementary $t$-channel pion exchange [27,28].

When an observable becomes sensitive to mixing between elements of a particular basis, it makes it difficult to interpret the internal structure of a state. This suggests that for hard exclusive processes there may be a more efficient representation of the matrix elements defining the observable. In the following we will show that just as a hierarchy of $t$-channel processes naturally explains the low-$x$ behavior of DIS structure functions, the same is true for the amplitudes representing exclusive reactions in the entire kinematical region of Bjorken-$x_{BJ}$.

3. $t$-channel dominance of exclusive lepto-production. Duality implies that all Feynman diagrams contributing to the hadronic tensor can be classified as either $s$-channel exchanges with baryon quantum numbers, or $t$-channel exchanges with meson quantum numbers. For large $s$ and small $t$, $t$-channel exchange of a particle with spin $J$ is proportional to $\beta s^J$ with the residue $\beta$ depending on $t$ and particle masses (which in our case includes the large virtual photon mass, $Q^2$). For example in the simple model of linear meson trajectories, the spin of a particle is proportional to the square of its mass, $J(M^2) = \alpha(M^2) = \alpha_0 + \alpha' M^2$, and the sum over all exchanged mesons leads to
an amplitude proportional to $\beta_s^{\alpha(t)}$. Such a description successfully reproduces
the experimentally observed shrinkage of the forward peak with increasing en-
ergy [19]. In general in the regime where $s/|t| >> 1$ the singularity $J = \alpha(t)$ furthest to the right in the complex angular momentum plane, determines
the leading power of the energy dependence. We are concerned with reactions that obey the constraints $s/|t| >> 1$. In this kinematic region one would expect $t$-channel exchanges to accurately parametrize these amplitudes. It is
well known that in the case of DIS the Regge parametrization is relevant when
$x_{BJ} \rightarrow 0$; however, for finite $x_{BJ}$ the $t$-channel exchange description becomes
inefficient. This occurs because away from the forward region, all singularities
in the complex angular momentum plane i.e. all daughter trajectories contribute equally to the amplitude as the rightmost singularity, which defines
the leading Regge trajectory. For DIS processes as one goes away from the
region $x_{BJ} \sim 0$, it very quickly becomes more efficient to represent the am-
plitude by $s$-channel exchanges of quasi-free partons. However, we will show
that the conditions that characterize exclusive production are quite different
from the conditions governing the inclusive processes.

The contribution to the hadronic tensor from $t$-channel exchange of a spin-$J$
meson is proportional to

$$T_J = \frac{\beta_J(t)\beta_J^*(q^2, q'^2, t)}{t - M_J^2} \sum_{\lambda=J}^{\infty} \left[ \frac{(p' + p)^{\mu_1}}{2} \cdots \frac{(p' + p)^{\mu_J}}{2} \epsilon_{\mu_1 \cdots \mu_J}^{\lambda} (p' - p) \right] \times \left[ \frac{(q' + q)^{\nu_1}}{2} \cdots \frac{(q' + q)^{\nu_J}}{2} \epsilon^{\nu_1 \cdots \nu_J}_{\lambda} (p' - p) \right].$$

(9)

In Eq. (9), $\epsilon$ is the spin-$J$ polarization vector, and $\beta_J^l$ and $\beta_J^u$ are the residue
functions at the lower and upper vertex, respectively. This is shown schemati-
cally in Fig. 1. In the Bjorken limit, $s \rightarrow Q^2(1 - x_{BJ})/x_{BJ}$ and the amplitude reduces to

$$T_J = \frac{\beta_J(t)\beta_J^*(q^2, q'^2, t)}{t - M_J^2} \left( \frac{Q^2}{2x_{BJ}} \right)^J.$$

(10)

The key question is how the upper residue function depends on the large
variables ($Q^2$ and $-q'^2 = Q^2$ in the case of inclusive processes, and $Q^2$ for the exclusive amplitudes). It is well known that for kinematics that are relevant to
inclusive scattering, the upper residue function behaves as $(1/Q^2)^{J+1}$, modulo
logarithmic corrections, so that the amplitude scales, $Q^2 T_J \propto (1/x_{BJ})^J$, as expected [29,30,31,32]. Summing over all spins leads to the Regge behavior,
$Q^2 T = \sum_J T_J \propto (1/x_{BJ})^{\alpha(0)}$. The leading Regge trajectory with $\alpha(0) > 0$ will
 dominate the $x_{BJ} \rightarrow 0$ behavior of the hadronic tensor, while all daughter
trajectories with $\alpha_n(0) < \alpha(0)$ are subleading as $x_{BJ} \rightarrow 0$. For finite $x_{BJ}$,
however, daughter Regge trajectories are no longer suppressed, and as a result
Fig. 1. $t$-channel meson contribution to the hadronic tensor for exclusive lepto-production. The amplitude is summed over all spins $J$ that can contribute, and depends on the product of the residue functions $\beta$ at the upper and lower vertices.

the Regge description becomes ineffective while the $s$-channel parton model description becomes natural.

We will now show that for exclusive amplitudes the upper vertex scales with a finite power of $1/Q^2$ instead of being suppressed for high spins. Thus after summing over all spins it gives an amplitude that behaves as $T \propto (Q^2/x_{BJ})^{a(0)}$, i.e. so long as $Q^2 >> m_N^2$ is in the Bjorken regime, the amplitude is dominated by the leading Regge trajectory for all $x_{BJ}$ and not only in the limit $x_{BJ} \to 0$.

To show that, we first rewrite the contribution of a $t$-channel, spin-$J$ exchange to the matrix element in Eq. (3) in terms of the two-current correlation in the exchanged meson,

\[
\int d^4ze^{-ikz} \langle p'|T[\phi^+(\frac{z}{2}) \phi(-\frac{z}{2})] |p\rangle = \frac{\beta_{J_n}^u(t)}{t - M_{J_n}^2} \Phi_{J_n}(p - p', k) \\
\sum_{\lambda,J} \prod_{\nu=1}^{J} \epsilon_{\nu_1^{(J)}}^{J} \epsilon_{\nu_2^{(J)}}^{J} (p' - p) \left[ \frac{(p' + p)^{\mu_1}}{2} \cdots \frac{(p' + p)^{\mu_J}}{2} \right] \epsilon_{\mu_1 \cdots \mu_J}^{(J)} (p' - p)
\]

(11)

In Eq. (11), $n$ refers to other quantum numbers that distinguish between
exchanged mesons (after Reggeization $n$ distinguishes between the leading and daughter trajectories), and $\Phi_{Jn}(\Delta, k)$ is the covariant (Bethe-Salpeter) amplitude of a meson with momentum $\Delta$, where $\Delta^2 = t$. Finally $k$ is the relative momentum between the quarks, as shown in Fig. 1. Using dispersion relations, the (unnormalized) Bethe-Salpeter amplitude can be represented as [33],

$$\Phi_{Jn}(\Delta, k) = i \int_{-1}^{1} dx \int d\mu^2 \frac{g_{Jn}(x, \mu^2)}{\left[(k - \Delta)^2 - \mu^2 + i\epsilon\right]^{n+J}},$$

(12)

where the spectral density $g_{Jn}$ is related to the parton distribution amplitude in a meson and can in principle be constrained from electromagnetic data [34] and QCD asymptotics [35].

What is important for our argument are the following model independent features of the amplitude in Eq. (12). The magnitude of the relative momentum $k$ is of the order of the hadronic scale $\mu m_N$. Secondly, in the infinite momentum frame, $k \propto \Delta$, $\xi^\pm \equiv (1 \pm x)/2$ represents the fraction of the longitudinal momentum carried by the quark (antiquark), and $g$ becomes the parton distribution function. In the Bjorken limit the leading, helicity-zero component of the meson distribution amplitude has $J$-independent behavior near $\xi^\pm \to 1$ [36]. Finally the power dependence of the relative momentum is constrained by the angular momentum, i.e. the power of the denominator in Eq. (12) increases with $J$. Inserting the analytical expression for the Bethe-Salpeter amplitude of Eq. (12) into Eq. (11) and then into Eq. (3), one obtains the final expression for the contribution of spin-$J$ exchange to the hadronic tensor. It is given in terms of an integral over $k$ (see Eq. (3)) of the product of the quark propagator, the Bethe-Salpeter amplitude of Eq. (12), and a polynomial in $k$ originating from the coupling to the spin-$J$ polarization vectors (Eq. (11)). The polynomial is responsible for the $s^J \sim (Q^2)^J$ behavior of the amplitude. The integral can easily be evaluated using the Feynman parametrization which introduces an integral over the parameter $\alpha$. Ignoring terms of order $m_N^2/Q^2$ and $t/Q^2$, up to constant numerical factors one finds

$$\beta_{Jn}^{\mu} = \int_{-1}^{1} dx \int d\mu^2 g_{Jn}(x, \mu^2) \int_{0}^{1} d\alpha \frac{\alpha^J}{-\alpha \left(\frac{q^2 + q^2}{2} + x\frac{q^2 - q^2}{2} + \mu^2\right)^{n+J}},$$

(13)

For inclusive amplitudes, $q^2 = q^2 = -Q^2$ the $x$ disappears from the denominator and the integration over $\alpha$ is dominated by $\alpha \sim \mu^2/Q^2$. As a result, the entire integral is of order $(\mu^2/Q^2)^{J+1}$, as expected. However for exclusive amplitudes, $q^2 \sim 0$ the integrand is dominated by the region $1 - x = O(\mu^2/Q^2)$, and finite $\alpha$. The endpoint behavior of the distribution amplitudes $g_{nJ}$ is spin independent, and for leading-twist amplitudes $g_{Jn}(x \to 1) \sim (1-x)$. This leads
to a $J$-independent suppression of the upper vertex with $Q^2, \beta_{Jn} \sim O(\mu^4/Q^4)$ which is independent of the spin of the exchanged meson. This is our main result. As discussed above, upon summing over all spins from a single trajectory one determines that the hadronic tensor is proportional to $(Q^2/x_{BJ})^{\alpha(t)} \sim (Q^2/x_{BJ})^{\alpha(0)}$, for small $t$. Thus, in the Bjorken limit hard exclusive processes should be dominated by a single, leading Regge trajectory for all $x_{BJ}$, and not just for $x_{BJ} \to 0$. We argue that this is the most efficient way to interpret hard exclusive processes. We also note that the Regge approach to exclusive deeply virtual production was previously considered in [37], where a different $Q^2$ dependence was obtained for the full exclusive amplitude. However, those authors assumed a Regge-like amplitude with a particular $Q^2$ dependence, rather than deriving the behavior from a sum of $t$-channel poles as was done here.

As we mentioned earlier, a recent experimental analysis of $\omega$ electroproduction at Jefferson Laboratory [14] showed that their data was in good agreement with predictions from standard Regge phenomenology, while showing large uncertainties with analyses based on models of generalized parton distributions [38]. Note that our results were derived in the Bjorken limit with $s/|t| >> 1$, while the JLab results correspond to energies of a few GeV and values up to $|t| \sim 2.7$. At high energies, one expects $\omega$-photoproduction to be dominated by Pomeron exchange; however at lower energies the leading meson Regge trajectories $f_2$ and $\pi$, with intercepts $\alpha_{f_2}(0) \sim 0.5$ and $\alpha_{\pi}(0) \sim 0$ can also give sizeable contributions. This indeed was found to be the case for the CLAS data [14,39]. These exchanges are also found to be responsible for $s$-channel helicity-flip amplitudes. Extension to $|t| \gtrsim 1$ GeV$^2$ is somewhat model dependent [40,41] as one needs to extrapolate further from the physical region of the $t$-channel. In all these analyses, once $W$ is greater than a few GeV, the daughter Regge trajectories with $\alpha_n(0) < 0$ do not seem to be needed.

**Summary** We have shown that in the Bjorken limit ($s >> |t|$ and $Q^2 >> m_N^2$), the leading Regge trajectory should be expected to dominate amplitudes for exclusive lepto-production. Our arguments also imply that the generalized parton distributions can be written in terms of Reggeon-nucleon coupling and that their ‘natural’ interpretation would be in terms of the parton content of the meson cloud rather than that of the bare nucleon. The GPDs can be computed by summing amplitudes of the type given in Eq. (11) with various sum rules constraining the products of residue functions and Reggeon-quark-antiquark distribution amplitudes [42]. Our arguments are model-independent and are based on general assumptions about the analytic structure of the scattering amplitude in the complex angular momentum plane. However, analyses of hard exclusive processes, particularly those at relatively low energies and large $|t|$ values, will require detailed models that can accommodate spin-flavor dependence and build in the characteristics of the relevant Regge trajectories. The recent analysis of $\omega$-photoproduction [14] indeed suggests that Regge
phenomenology can successfully be used to describe the hadronic part of the production amplitude in exclusive lepto-production.

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