Elastic and Resonance contributions to moments of the proton structure function $F_2$

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Abstract: We discuss the role of nucleon and its excited state poles in the twist expansion of the nucleon structure function moments. We find that the nucleon pole contribution was overestimated in previous analyses by a factor of two. Inclusion of this missing factor makes the duality appear for all moments at least down to $Q^2 = 1$ GeV$^2$. For resonance poles the time reversal invariance together with unitarity demand at least four of them in the forward Compton scattering amplitude. These poles as well as the elastic pole can be propagated separately through the standard Operator Product Expansion derivation for DIS. This part of the amplitude gives a coherent, positive, higher twist contribution, which can be singled out of the total and compared with the remaining one given mostly by threshold effects. A comparison of the estimated resonance contribution to the data on structure function moments allows to test large-$Q^2$ behavior of proton-resonance transition form-factors.

Keywords: Deep Inelastic Scattering, Phenomenological Models, Sum Rules.
1. Introduction

Investigations of the nucleon structure with electromagnetic probes at large momentum transfers allow for a simple treatment within the parton model. In fact, in this kinematics the inclusive lepton-nucleon scattering cross section can be related to the parton momentum distributions inside the nucleon. Latter do not follow from pQCD and they are subject to a measurement or Lattice simulations. Since structure functions themselves cannot be compared to pQCD predictions one can instead extract quantities calculable in the theory, moments. The Mellin transform of structure functions to Euclidean space offers many simplifications. First of all, it gives an access directly to the subject of pQCD, the $Q^2$-evolution, leaving only one undefined parameter to the measurement. This parameter is the absolute normalization of the moment at some arbitrary scale. This scale can be chosen in the kinematic region where approximations involved in calculations are fulfilled. The second, the calculation of higher order corrections by means of effective series summation techniques are less involved in the moment space \[1\]. Obviously, there are also difficulties: nucleon structure functions cannot be measured down to $x = 0$. Therefore, lowest moments always carry a dose of uncertainty. However, higher moments are well defined experimentally and they are subject to precise measurements \[2, 3\].

Scaling down in $Q^2$ we explore a new kinematic region, not yet understood in terms of pQCD. It is the region of bound partons, where the interaction with the probe photon likely involves more than one single parton. Correlated partons can either interact to collectively produce hadronic final state or form an excited nucleon state. These multiparton correlations are responsible for the confinement and they are the subject of intensive studies \[4, 5, 6, 7, 8, 9, 10\]. Main goal of these investigations is to estimate the contribution of multiparton correlations (higher twists) and their $Q^2$ evolution based on experimental
data. The theoretical basis of these studies [3, 11] is still poorly developed. This creates a large amount of phenomenological approaches aimed to catch main features of the problem as e.g. the parton-hadron duality phenomena [12, 1, 5, 8, 10, 13].

In this article we discuss applicability of the Operator Product Expansion (OPE) to poles of the forward Compton scattering amplitude. The separation of the pole contribution in the moments allows to put a lower limit on the higher twist contribution in the measured moments. Moreover subtraction of the pole contribution might help to the application of the constituent quark model picture as proposed in Ref. [14].

2. Forward Compton scattering amplitude

The forward Compton scattering amplitude $T_{\mu\lambda}$ of the virtual photon with four-momentum $q$ on the nucleon $N$ is given by [15]:

$$T_{\mu\lambda} = i \sum_\sigma \int d^4 \xi e^{i q \cdot \xi} \langle N, \sigma | T(J_{\mu}(\xi) J_{\lambda}(0)) | N, \sigma \rangle ,$$  

(2.1)

where the sum is running over polarization degrees of freedom $\sigma$, $J_{\mu}$ are hadronic currents, $\xi$ is the light-cone separation between currents and $T()$ stays for the time-ordered product. In case of unpolarized scattering off the nucleon the amplitude can be expressed in terms of two invariant amplitudes $T_1$ and $T_2$:

$$T_{\mu\lambda} = -T_1 \left( g_{\mu\lambda} + \frac{q_{\mu} q_{\lambda}}{Q^2} \right) + \frac{T_2}{M^2} \left( P_{\mu} + \frac{P \cdot q}{Q^2} q_{\mu} \right) \left( P_{\lambda} + \frac{P \cdot q}{Q^2} q_{\lambda} \right) ,$$  

(2.2)

where $q_{\mu} = (\nu, q)$ $(Q^2 = -q^2)$ and $P_{\mu}$ are the virtual photon and proton four-momenta, respectively.

Hadronic currents from Eq. 2.1 for particular class of processes realizing through the production of an intermediate state particle can be expressed in terms of invariant elastic and transition form-factors. The nucleon intermediate state corresponding to the elastic current according to Ref. [14] can be written as:

$$\langle N', \beta | J_{E}^{\mu}(0) | N, \alpha \rangle = \bar{u}_\beta(P') \left\{ \gamma^{\mu} F_1(Q^2) + i \sigma^{\mu\nu} \frac{q_{\nu}}{2M} F_2(Q^2) \right\} u_\alpha(P) ,$$  

(2.3)

where $F_1$ and $F_2$ are known Dirac and Pauli form-factors and $M$ is the proton mass. The intermediate particle can be also a nucleon resonance. In this case the hadronic current can be expressed as following [17]:

$$\langle N_j, \beta | e_{\mu}^{m} J_{R}^{\mu}(0) | N_2, \alpha \rangle = \frac{2}{e} \sqrt{M(M_R^2 - M^2)} A_m ,$$  

(2.4)

where $j$ is the resonance spin, $M_R$ is the resonance mass, $A_m$ is the resonance helicity amplitude, $\alpha$ and $\beta$ are helicity indices, $m$ indicates the virtual photon helicity state with respect to the proton helicity defined as in Ref. [16], and the virtual photon polarization four-vectors are given by:

$$e_{\mu}^{\pm} = \frac{1}{\sqrt{2}} (0, \pm 1, i, 0) , \quad e_{\mu}^{0} = \frac{1}{Q} (|q|, 0, 0, \nu) .$$  

(2.5)
These two contributions of the amplitude are particular because represent poles on the Riemann $\nu$-surface. The rest of the amplitude is an analytic function of $\nu$ describing the continuum of intermediate states. Therefore, the total amplitude can be described as a sum of elastic ($T^E$), resonance ($T^R$) and continuum ($T^C$):

$$T(\nu, Q^2) = T^E(\nu, Q^2) + T^R(\nu, Q^2) + T^C(\nu, Q^2).$$ (2.6)

If we plug the complex amplitude from Eq. 2.6 into the unitarity relation $SS^\dagger = 1$ (with $S = 1 + 2iT$) we obtain

$$\text{Im}T^E(\nu, Q^2) + \text{Im}T^R(\nu, Q^2) + \text{Im}T^C(\nu, Q^2) = |T^E|^2 + |T^R|^2 + |T^C|^2 + 2\text{Re}[T^R T^C^\dagger],$$ (2.7)

collecting the same singularities in r.h.s. and l.h.s. and neglecting the interference effects we obtain the optical theorem for each individual term separately.

In this article we are interested in the case of $T_2$. We assume that elastic and resonance amplitudes are simple poles and can be written in the factorized form:

$$\nu T^{E,R}_2 = R^{E,R}(Q^2) P^{E,R}(\nu),$$ (2.8)

where the residue of the elastic pole is given by

$$R^E(Q^2) = \frac{G^2_E(Q^2) + \tau G^2_M(Q^2)}{1 + \tau},$$ (2.9)

where Sachs form-factors $G_E$ and $G_M$ are related to $F_1$ and $F_2$ as following:

$$G_E = F_1 - \frac{Q^2}{4M^2} F_2, \quad G_M = F_1 + F_2.$$ (2.10)

While resonance residues can be written as

$$R^R(Q^2) = \frac{1}{2\pi\alpha} \frac{M^2_R - M^2}{4M} \frac{Q^2}{q_R^2} \left[ |A_{1/2}|^2 + |A_{3/2}|^2 + 2 \frac{Q^2}{q_R^2} |S_{1/2}|^2 \right],$$ (2.11)

here $q_R$ is the virtual photon three-momentum taken at the resonance pole.

### 3. Pole structure of the Compton amplitude

The Compton amplitude defined in the first quadrant of the complex $\nu$-plane can be analytically continued in other quadrants by means of the crossing symmetry. The analytic continuation can be performed through the following relation between crossing channels: $T_2(-\nu) = T_2(\nu)$ (time reversal) and $T_2(\nu^\ast) = T_2^\ast(\nu)$ ($\gamma\gamma \rightarrow p\bar{p}$). The Compton amplitude has a number of simple poles and branch cuts shown in Fig. 1. All these poles and cuts should obey crossing symmetry relations mentioned above. The elastic scattering amplitude has two poles situated on the real axis at $\nu = \pm \nu_E = \pm Q^2/2M$ and it has the following simple structure:

$$P^E(\nu) = \frac{\nu}{\nu_E - \nu} + \frac{\nu}{\nu_E + \nu}.$$ (3.1)

\footnote{in the following two sections we will drop Lorenz indices for simplicity}
Nucleon resonances also generate poles in the Compton amplitude. But these are located on unphysical Riemann sheets. In order to satisfy symmetry relations each resonance should have at least four poles, two on each lower and upper unphysical Riemann sheets. Nevertheless, we limit ourselves to simple poles, which do not carry information about Riemann sheet of location:

\[ P_R(\nu) = \left( \frac{\nu}{\nu_R - \nu} + \frac{\nu}{\nu_R + \nu} \right)_{+i\varepsilon} + \left( \frac{\nu^*}{\nu_R^* - \nu} + \frac{\nu^*}{\nu_R^* + \nu} \right)_{-i\varepsilon}, \]  

where \( \varepsilon \) is an infinitesimal positive constant indicating Riemann sheet of influence and the resonance is located at:

\[ \nu_R = \frac{Q^2 + (M_R^2 - M^2)}{2M} + i \frac{\Gamma_R M_R}{2M}, \]  

where \( \Gamma_R \) is its width. One can check that Eq. 3.3 yields the familiar Breit-Wigner parameterization.

Unfortunately the simple parameterization from Eq. 3.3 does not respect the second crossing symmetry property: \( T_2(\nu^*) = T_2^*(\nu) \) and the amplitude becomes imaginary on the real axis even in the region \( |\nu| < Q^2/2M \) where no branch cut is present. We can adjust this by brute force method multiplying the \( +i\varepsilon \) term in Eq. 3.3 by:

\[ G = \begin{cases} 
-1, & |\nu| > Q^2/2M \\
1, & |\nu| < Q^2/2M 
\end{cases} \]  

This picture is in contrast with Refs. [18] where only poles on the lower unphysical sheet are predicted. However, we find that the absence of poles on the upper unphysical sheet would generate a disbalance between lower and upper semi-planes of the Compton amplitude on the physical sheet. This disbalance would result in an imaginary part of the discontinuity across the branch cut and therefore would violate the optical theorem:

\[ \text{Disc}(T^R) = \frac{1}{2i} \left[ T^R(\nu + i\varepsilon) - T^R(\nu - i\varepsilon) \right] = \frac{1}{2i} T^R(\nu + i\varepsilon) \in \mathbb{C} \neq 2\pi W^R(\nu) \in \mathbb{R}, \]  

where \( W^R \) is the resonance part of hadronic tensor.

4. Low Energy Expansion

The standard OPE derivation for DIS [15] proceeds through an expansion of the dispersion integral in the series of \( \nu \) in unphysical kinematic domain \( |\nu| < Q^2/2M \), which we call here \( ^2 \) actually poles are present on many lower and many upper unphysical sheets generated by different decay channels

\( ^3 \) in principle the pole structure of the resonance amplitude is more complex: poles should “remember” the Riemann sheet which they belong to. This can be parameterized in various forms. One of the simplest solution is the following:

\[ P_R(\nu) = -i \frac{\pm \sqrt{\nu_E - \nu_R}}{\sqrt{\nu_E - \nu} - \sqrt{\nu_E - \nu_R}}, \]  

plus two conjugate poles at \( \nu_R^* \), here \( \pm \) signs refer to the principal and secondary square root values indicating therefore correct Riemann sheet. However in this study we are interested in the amplitude at small \( \nu << \nu_E \) where the difference between two cases is not so important but the simplification of using
Figure 2: The imaginary part of the resonance shape distribution: the solid line shows the shape of Eq. 3.3 (Breit-Wigner) and the dashed line represents Eq. 3.2. Note that the dotted line goes to zero at the threshold.

“low energy” regime. The Cauchy integral over the contour $C$ shown in Fig. 1 is given by:

$$T(Q^2, \nu) = \frac{1}{2\pi i} \int_C \frac{T(Q^2, \nu')d\nu'}{\nu' - \nu}. \quad (4.1)$$

Assuming that the integral over the circle of radius $R \to \infty$ vanishes and taking discontinuities across left and right branch cuts, using also the time reversal invariance $T(Q^2, -\nu) = T(Q^2, \nu)$ we obtain:

$$\int_C \frac{T(Q^2, \nu')d\nu'}{\nu' - \nu} = 2i \int_{\nu_E}^{\infty} \frac{\text{Disc}(T(Q^2, \nu'))d\nu'}{\nu' - \nu} + 2i \int_{-\nu_E}^{-\infty} \frac{\text{Disc}(T(Q^2, \nu'))d\nu'}{\nu' - \nu}$$

$$= 2i \int_{\nu_E}^{\infty} \left\{ \frac{1}{\nu' - \nu} + \frac{1}{\nu' + \nu} \right\} \text{Disc}(T(Q^2, \nu'))d\nu'. \quad (4.2)$$

From the other hand the optical theorem states that:

$$\text{Disc}(T(Q^2, \nu)) = \text{Im}(T(Q^2, \nu + i\varepsilon)) = 2\pi W(Q^2, \nu). \quad (4.3)$$

Therefore, substituting the discontinuity $\text{Disc}(T(Q^2, \nu))$ by Eq. 4.3 we can rewrite the dispersion relation in the following form:

$$T(Q^2, \nu) = 4 \int_{Q^2/2M}^{\infty} \frac{d\nu'\nu'}{\nu'^2 - \nu^2} \left\{ W(Q^2, \nu') \right\}. \quad (4.4)$$

Choosing $|\nu| < Q^2/2M$ one expands this into the geometrical series:

$$T(Q^2, \nu) = 4 \sum_{n=0, \text{even}}^{\infty} \tilde{M}_n(Q^2) \left[ \frac{1}{x} \right]^n, \quad (4.5)$$

where

$$\tilde{M}_n(Q^2) = \int_0^1 dx' x'^{n+1} W(Q^2, x'). \quad (4.6)$$
In turn, the OPE of the product of two hadronic currents separated by a small light-cone distance is given by \cite{15}:

\[ T(Q^2, \nu) = \sum_{\tau, n} C_{\tau n}(Q^2, \mu^2) O^n_\tau(\mu^2) \left( \frac{1}{x} \right)^{n} \left( \frac{1}{Q^2} \right)^{\tau/2-1}, \quad (4.7) \]

and making the standard correspondence between terms with equal powers of $1/x$ we obtain usual DIS twist expansion:

\[ \tilde{M}_n(Q^2) = \frac{1}{4} \sum_{\tau = 2}^{\infty} C_{\tau n}(Q^2, \mu^2) O^n_\tau(\mu^2) \left( \frac{1}{Q^2} \right)^{\tau/2-1}, \quad (4.8) \]

where expansion coefficients $C_{\tau, n}$ can be analytically calculated in pQCD and local operators $O^n_\tau$ can be obtained from Lattice simulations.

In order to apply this expression in practice we need to rewrite Eq. 4.8 for a measured structure function. For instance, it can be done for moments of the proton structure function $F_2 = \nu W_2$ extracted from the data in Ref. \cite{2}. $n$-th moment of the structure function $F_2$ is defined as:

\[ M_n(Q^2) = \int_0^1 dx x^n F_2(x, Q^2), \quad (4.9) \]

and therefore one obtains the following expansion:

\[ M_n(Q^2) = \sum_{\tau = 2}^{\infty} E_{n\tau}(\mu, Q^2) O^\tau_n(\mu^2) \left( \frac{\mu^2}{Q^2} \right)^{\frac{1}{2}(\tau-2)} \quad (4.10) \]

where $k = 1, 2, ..., \infty$, $n = 2, 4, ..., \infty$, $\mu$ is a reference scale, $O^\tau_n(\mu^2)$ is the reduced matrix element of the local operators with definite spin $n$ and twist $\tau$ (dimension minus spin), related to the non-perturbative structure of the target. $E_{n\tau}(\mu, Q^2)$ is a dimensionless coefficient function describing the small distance behavior, which can be perturbatively expressed as a power expansion of the running coupling constant $\alpha_s(Q^2)$.

However, we noticed that also the pole (elastic and resonance) amplitude contributing to the l.h.s of Eq. 4.13 can be expanded in the power series of $\nu$ in the region $|\nu| < Q^2/2M$:

\[ \nu T^E_2(\nu, Q^2) = 2R^E(Q^2) \sum_{n=1, odd}^{\infty} \left[ \frac{\nu}{\nu_R} \right]^n = \frac{2}{x} R^E(Q^2) \sum_{n=0, even}^{\infty} \left[ \frac{1}{x} \right]^n \quad (4.11) \]

and

\[ \nu T^R_2(\nu, Q^2) = 4R^R(Q^2) \sum_{n=1, odd}^{\infty} \left[ \frac{\nu}{\nu_R} \right]^n \cos(n\phi_R) = \frac{4}{x} |x_R| \sum_{n=0, even}^{\infty} \left[ \frac{|x_R|}{x} \right]^n \cos((n + 1)\phi_R), \quad (4.12) \]

where $\phi_R = \text{atan} \frac{\Gamma_{R M_R}}{Q^2 + (M_R^2 - M^2)}$ is the phase of $\nu_R$.

\[^4n = 0\) relation is meaningless because Regge theory predicts that the structure function behaves as $1/x$ at $x \to 0$ and therefore $n = 0$ integral in Eq. 4.1 would be diverging.\]
In case of $T_2$ amplitude the moment expansion will give

$$\nu T_2(Q^2, \nu) = \frac{4}{x} \sum_{n=0, \text{even}}^{\infty} M_n(Q^2) \left[ \frac{1}{x} \right]^n .$$

(4.13)

We can rewrite Eq. 4.13 as following:

$$\nu T_2^C(\nu, Q^2) = \frac{4}{x} \sum_{n=0, \text{even}}^{\infty} M_n(Q^2) \left[ \frac{1}{x} \right]^n - \nu T_2^E(\nu, Q^2) - \nu T_2^R(\nu, Q^2) ,$$

(4.14)

therefore the continuum part of the Compton amplitude is given by

$$\nu T_2^C(\nu, Q^2) = \frac{4}{x} \sum_{n=0, \text{even}}^{\infty} \left[ \frac{1}{x} \right]^n \left\{ M_n(Q^2) - \frac{1}{2} R^E(Q^2) - R^R(Q^2)|x_R|^{n+1} \cos ((n+1)\phi_R) \right\} ,$$

(4.15)

or in terms of measured moments\(^5\) it can be rewritten

$$M_n^C(Q^2) = M_n(Q^2) - \frac{1}{2} F^E_n(Q^2) - R^R(Q^2)|x_R|^{n-1} \cos ((n-1)\phi_R) ,$$

(4.16)

where the last two terms can be estimated using phenomenological data on elastic and resonance transition form-factors.

5. Results

First of all one can notice that the elastic contribution comes out to be a factor of two smaller than in all previous works (see for example Ref. [7]). Previously, the dispersion relation method was applied uniformly to both the elastic and inelastic channels. This led to a confusion in two aspects: 1) the discontinuity across a pole is not a well defined object from the mathematical point of view; 2) the value of the integral along the semi-circle around the pole is not vanishing. One can prove this statement by evaluating the dispersion relation for the elastic amplitude alone. If the elastic amplitude is given by:

$$T(\nu) = \frac{1}{\nu_E - \nu} + \frac{1}{\nu_E + \nu} ,$$

(5.1)

then applying the dispersion relation from Eq. 4.1 using the discontinuity across the pole one finds:

$$\frac{1}{\nu_E - \nu} + \frac{1}{\nu_E + \nu} \neq \frac{1}{\pi} \int_0^{\infty} \left[ \frac{1}{\nu - \nu'} - \frac{1}{\nu + \nu'} \right] \times \text{Im}T(\nu')d\nu' = 2 \left[ \frac{1}{\nu_E - \nu} + \frac{1}{\nu_E + \nu} \right] .$$

(5.2)

Therefore, the evaluation of the discontinuity across the elastic pole in the dispersion relation leads to the amplitude which is twice larger than the original one. The comparison of moments obtained in the present article and those from Ref. [7] is shown in the Fig. 3 together with various pQCD calculations.

\(^5\)notice that the definition in Eq. 4.13 is different from that of the Ref. [8] and we replace here $n$ with $n - 2$ for consistency
Figure 3: Total Nachtmann moments of the proton structure function $F_2$ from Ref. [2] obtained by adding the elastic contribution according to Eq. 3.16 (solid triangles) and “standard” formula (empty circles). The lines represent pQCD calculation fitted to the data at largest $Q^2$: solid line - LO, dashed line - NLO, dotted line - NNLO, dot-dashed - NLO+SGR from Ref. [1].

Using the relation from Eq. 4.16 we estimated the resonance contribution to the proton structure function moments. For the sake of simplicity we neglected the longitudinal couplings $S_{1/2}$ of nucleon resonances. For this exploratory work this is a good approximation, provided that those couplings are typically small. We used phenomenological parameterizations of the nucleon resonance helicity amplitudes from Ref. [19] and Ref. [20]. We have
Figure 4: Inelastic Nachtmann moments of the proton structure function $F_2$ from Ref. [2]: solid line represents the resonance contribution with form-factors of Ref. [19]; dashed line shows the resonance contribution using form-factors from Ref. [20].

taken into account 9 (20) well established resonances from the parameterization [20] ([19]) whose masses and widths were obtained from Ref. [21]. In order to compare this to the data extracted in Ref. [2] we had to convert Conwell-Norton moments into Nachtmann ones. In the region $M^2/Q^2 < 1$ this can be done by expanding Nachtmann moments in a power series of $M^2/Q^2$ as in the Ref. [22]. Results of the comparison are shown in Figs. 4 and Fig. 5.
Figure 5: The total higher twist contribution and the elastic+resonance contribution to the $F_2^p$ structure function moments relative to its leading twist part taken from Ref. [2]: solid line - the total higher twist contribution; dashed line - the resonance contribution using resonance form-factors from Ref. [20]; dotted line - the elastic+resonance contribution using resonance form-factors from Ref. [20].

6. Conclusions

We calculated the contribution of the nucleon elastic pole and its excited states into moments of the nucleon structure function $F_2$. Due to formation of an intermediate bound state this contribution cannot be described by an incoherent elastic scattering of the vir-
tual photon off a quark (either constituent or current). In terms of OPE this contribution belongs to higher twists. This is confirmed by the observed $Q^2$-dependence of estimated the resonance contribution shown in the Fig. 5. Vice versa, assuming that conceptually the resonance contribution as well as the elastic peak contribution should belong to higher twists one obtains a stringent test on the large-$Q^2$ behavior of resonance form-factors.

We found that the elastic peak contribution is actually twice smaller than it was known before. The reduction of the elastic contribution in moments leads to an improvement of the duality effect in the low $Q^2$ region. As you can see in Fig. 8 the second moment becomes almost constant while higher moments are much more flat after the correction and leave no room for a significant higher twist contribution (see also Fig. 5). This leads to the conclusion that the extraction of nucleon elastic form-factors from DIS pioneered in Ref. [7] is not hopeless and the large discrepancies found more recently in Ref. [23] are due to missing 0.5 factor.

In Refs. [2, 3] we have seen that different higher twist terms cancel generating the well known duality phenomenon. The resonance contribution estimated in this article yields coherent positive higher twist in the moments. Additional positive contribution to higher twists is given by the elastic scattering. While the negative higher twist which largely cancels the resonance and elastic contributions is related to number of exclusive channel thresholds (e.g. single pion production threshold). If all elastic, resonance and the threshold contributions are subtracted from the moments the remaining part should be related to the incoherent scattering off proton constituents regardless $Q^2$ value. Therefore, an estimate of the threshold contribution to the structure function moments would allow to study $Q^2$-behavior of form-factors of proton constituents.

Notice that both elastic and resonance higher twist contributions are related to nucleon and resonance bound state wave functions. Therefore these terms cannot be described by asymptotic freedom methods like Borel summation of the leading twist $\alpha_S$ series which also generates $1/Q^2$ corrections. Given large magnitude of elastic and resonance contributions in the total moments (in particular for $n > 2$) the good overlap between renormalon-based fits [24] and the data should be taken as accidental or related to the duality phenomenon.

It is clear from Fig. 5 that the resonance contribution itself does not follow DIS scaling behavior. Therefore, the observed approximate scaling of integrated resonance peaks (known as local duality) is due to compensation between rising resonances and decreasing background at low $Q^2$.

Large physics program focused on extraction of nucleon resonance form-factors from electro- and photo-production data undergoing in Jefferson Lab will allow for precise evaluation of the main higher twist contribution to the nucleon structure functions.

Acknowledgments

We are grateful to S. Simula for fruitful discussions.

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