Deeply Virtual Compton Scattering and Experiments at the Jefferson Lab

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Abstract

Experimental data on deep-inelastic electron-nucleon scattering in a wide range of x and Q^2 are analyzed in a simplified analytic model realazing parton-hadron duality. Special emphasis is paid to the treatment of the background. The role of the spin selection rules is discussed.

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1 Introduction

Generalized Parton Distributins (GPD) [1, 2, 3] play an important role in the strong interaction theory since they provide a unifying framework for the interpretation of an entire set of fundamental quantities of hadronic structure, such as, the vector and axial vector nucleon form factors, the polarized and unpolarized parton distributions, and the spin components of the nucleon due to orbital excitations. Deeply Virtual Compton Scattering (DVCS) is one of the key reactions to determine the GPDs experimentally, and it is the simplest process that can be described in terms of GPDs.

One of the first experimental observations of DVCS was based on the recent analysis of the JLAb data from the CLAS collaboration with a 4.2 GeV polarized electron beam in a kinematical regime near $Q^2 = 1.5\ GeV^2$ and $x = 0.22$, where Q^2 is minus the photon virtuality and x is the Bjorken variable. New measurements at at higher energies are curently being analyzed, and dedicated experiments are planned [4]. The high luminocity available for these measurements will make it possible to determine details of the Q^2, x and t dependences of GPDs.

On the theoretical side, much progress has been achieved [1, 2, 3] in treating GPD in the framework of the quantum chromodynamics (QCD) with the light-cone technique. On the other hand, the prevailing non-nonpertubative effects (resonance production, the background, low-Q^2 effects) dominating the kinematical region of present measurement and the underlying dynamics still leave much ambiguity in the above-mentioned field-theoretical approach. Therofore, as an alternative or complementary approach we have suggested in a number of papers [5, 6, 7, 8] to use dual amplitudes with Mandelstam analyticity (DAMA) as a model for GPD in general and DVCS in particular. We remind that DAMA realizes duality between direct-channel resonances and high-energy Regge behaviore ("Veneziono-duality"). By introduning Q^2-dependence in DAMA, we have extended the model off-shell and have shown [5, 6] how parton-hadron (or "Bloom-Gilman") duality is realized in this way. With the above specification, DAMA can serve as and explicite model for all values of the Mandelstam variables s, t and u as well as Q^2, thus realizing the ideas of DVCS and related GPDs.

We study inclusive electron-nucleon scattering shown in Fig. 1 with usual notations (see ref. [8] for more details). The photoabsorption cross section in the resonance region has been studied in a large number of papers [10, 11, 12, 13, 14] (for a comprehensive review see Ref. [15]). Most of the data come from SLAC and have been compiled by Stoler [16]. There are nearly 20 resonances in the $\gamma^* p$ system in the region between the pion-nucleon threshold and below 2 GeV, but only a few of them can be identified more or less unambiguously. One reason is that they overlap and compete with changing Q^2 and the other is the uncertainty due to the background. Therefore, instead of identifying each resonance, one considers three maxima above the elastic scattering peak, corresponding to some "effective" resonance contributions.

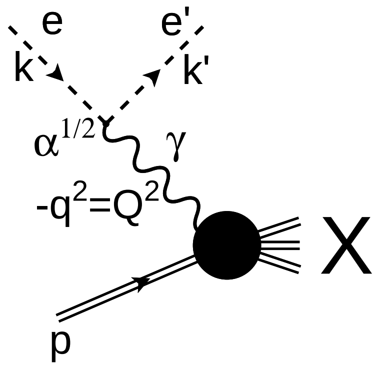

Figure 1: Kinematic of deep inelastic scattering.

In this work we present a simplified phenomenological analytical model, realizing parton-hadron duality, which is motivated by a termination of the real part of nonlinear complex Regge trajectories [7, 8]. These trajectories play a crucial role in the dynamics of the strong interactions. Actually, the trajectories can be considered as the basic dynamical variables, replacing the usual Mandelstam variables s, t and u (which enters only through the trajectories). The parameters of the trajectories can be fitted independently of the masses and widths of the known resonances, therefore, in principle, they reflect more adequately the position of the peaks in ep scattering, formed by the interplay of different resonances.

In concentrating on this aspect of the dynamics, we leave more freedom to the choice of the Q^2-dependent form factors. We start with a simplified model, disregarding the helicity structure of the amplitude and relevant selection rules, concentrating on the role of the Regge trajectories, analyticity and duality. Their role was treated in a number of papers [17] (recently in Ref. [20]). We ignore the relatively small (and poorly known) contribution from the cross sections involving longitudinally polarized photons, σ_L . In doing so, we anticipate the connection [5] with the small-x (high-energy) domain, where these simplifications are commonly accepted.

The central object of the present study is the nucleon SF, uniquely related to the photoproduction cross section by

$$F_2(x,Q^2) = \frac{Q^2(1-x)}{4\pi\alpha(1+\frac{4m^2x^2}{Q^2})}\sigma_t^{\gamma^*p}(s,Q^2)\ , \tag{1}$$

where total cross section, $\sigma_t^{\gamma^*p}$, is the imaginary part of the forward Compton scattering amplitude, $A(s,Q^2)$,

$$\sigma_t^{\gamma^*p}(s) = \mathcal{I}m\ A(s,Q^2)\ . \tag{2}$$

The center of mass energy of the γ^*p system, the negative squared photon virtuality Q^2 and the Bjorken variable x are related by

$$s = Q^2\frac{(1-x)}{x} + m^2, \tag{3}$$

We adopt the two-component picture of strong interactions [18], according to which direct-channel resonances are dual to cross-channel Regge exchanges and the smooth background in the $s-$channel is dual to the Pomeron exchange in the $t-$channel. As explained in Ref. [5], the background corresponds in a dual model to a pole term with an exotic trajectory that does not produce any resonance.

In the dual-Regge approach [5, 6, 7, 8] the Compton scattering can be viewed as an off-mass shell continuation of a hadronic reaction, dominated in the resonance region by non-strange (N and Δ) baryon trajectories. The scattering amplitude follows from the pole decomposition of a dual amplitude [5]

$$A(s,Q^2)\Big|_{t=0} = norm \sum_{i=N_1^*,N_2^*,\Delta,E} A_i \sum_{n=n_i^{min}}^{n_i^{max}} \frac{f_i(Q^2)^{2\left(n-n_i^{min}+1\right)}}{n-\alpha_i(s)}\ , \tag{4}$$

where i runs over all the trajectories allowed by quantum number exchange, $norm$ and A_i's are constants, $f_i(Q^2)$'s are the form factors. These form factors generalize the concept of inelastic (transition) form factors to the case of continuous spin, represented by the direct-channel trajectories. The n_i^{min} refers to the spin of the first resonance on the corresponding trajectory i (it is convenient to shift the trajectories by 1/2, therefore we use $\alpha_i = \alpha_i^{phys} - 1/2$, which due to the semi-integer values of the baryon spin leaves n in Eq. (4) integer). The sum over n goes with step 2 (in order to conserve parity).

It follows from Eq. (4) that

$$\mathcal{I}m\ A(s,Q^2) = norm \sum_{i=N_1^*,N_2^*,\Delta,E} A_i \sum_{n=n_i^{min}}^{n_i^{max}} \frac{[f_i(Q^2)]^{2\left(n-n_i^{min}+1\right)}\mathcal{I}m\ \alpha_i(s)}{(n-\mathcal{R}e\ \alpha_i(s))^2 + (\mathcal{I}m\ \alpha_i(s))^2}\ . \tag{5}$$

The first three terms in (5) are the non-singlet, or Reggeon contributions with the N^* and Δ trajectories in the s-channel, dual to the exchange of an effective bosonic trajectory (essentially, f) in the t-channel, and the fourth term is the contribution from the smooth background, modeled by a non-resonance pole term with an exotic trajectory $\alpha_E(s)$, dual to the Pomeron (see Ref. [5]). As argued in Ref. [5], only a limited number, $\mathcal{N}$, of resonances appear on the trajectories, for which reason we tentatively set $\mathcal{N} = 3$ - one resonance on each trajectories (N_1^*, N_2^*, Δ), i.e. $n_i^{max} = n_i^{min}$. We tried also with higher values of $\mathcal{N}$, but our analyses shows that $\mathcal{N} = 3$ is a reasonable approximation – even if additional peaks appear, they are suppressed with respect to the dominant one (first on each trajectory), because of the Q^2-behaviour of the form factors. Thus, the limited (small) number of resonances contributing to the cross section results not only from the termination of resonances on a trajectory but even more due to the strong suppression coming from the numerator (increasing powers of the form factors). Thus, for practical resonance we have replaced the formal condition $\mathcal{R}e\ \alpha(s) < const$ by a finite sum in Eq. (5).

In this work we use Regge trajectories with threshold singularities and nonvanishing imaginary part in the form:

$$\alpha(s) = \alpha_0 + \alpha_1 s + \alpha_2(\sqrt{s_0} - \sqrt{s_0 - s}), \tag{6}$$

where s_0 is the lightest threshold, $s_0 = (m_\pi + m_p)^2 = 1.14$ GeV^2 in our case, and linear term approximates the contribution from heavy thresholds [5, 6, 7, 8].

We have fitted the parameters of the baryon trajectories, given by Eq. (6), such as to reproduce the experimental masses and widths of the $\Delta(1236)$, $N^*(1520)$ and $N^*(1680)$ resonances (see [8] for more details) - the corresponding trajectory parameters are marked by † in the Table 10 (see Ref. [8] for more details).

Since, by definition, the smooth background does not show any resonance, here we keep only one term in the sum [1]. n_E^{min} is the first integer larger then $Max(\mathcal{R}e\ \alpha_E)$ – to make sure there are no resonances on the exotic trajectory. We take the exotic trajectory in the form

$$\alpha_E(s) = \alpha_E(0) + \alpha_{1E}(\sqrt{s_E} - \sqrt{s_E - s}), \tag{7}$$

where the intercept $\alpha_E(0)$, α_{1E} and the effective exotic threshold s_E are free parameters. As a first approximation we can assume the following expression for the exotic trajectory [5]:

$$\alpha_E(s) = 0.5 + 0.12(\sqrt{s_E} - \sqrt{s_E - s})\ , \tag{8}$$

where $s_E = 1.145^2$ GeV^2; in this case $n_E^{min} = 1$.

To start with we use the simplest, dipole model for the form factors, disregarding the spin structure of the amplitude and the difference between electric and magnetic form factors:

$$f_i(Q^2) = \frac{1}{(1 + \frac{Q^2}{Q_{0,i}^2})^2}\ . \tag{9}$$

where $Q_{0,i}^2$ are scaling parameters. The relative growth of the three resonance peaks and background will depend on the scaling factor $Q_{0,i}^2$. Therefore we choose $Q_{0,E}^2 > Q_{0,N_2^*}^2 > Q_{0,N_1^*}^2 > Q_{0,\Delta}^2$ in order to satisfy the experimentally observed behaviour of these terms, for example, the rise of the background contribution with respect to the resonance one with increasing Q^2; the relative growth of the N_1^* and N_2^* peaks with respect to the Δ peak.

[1]In Ref. [9] the whole DAMA integral was calculated numerical and it has been shown that in the resonance region the direct-channel exotic trajectory gives non-neglactable contribution, reaching up to 10-12%.

2 Fits to the SLAC and JLab data

In this Section we present a numerical analysis of our model based on the experimental data from SLAC [16] and JLAB [21][2]. This set of experimental data is not homogeneous, i.e. points at low s (high x) are given with very small experimental errors, thus "weighting" the fitting procedure not uniformly. This forced us to make a preselection for the fitting procedure, although all the experimental points are presented in the Figures (see ref. [8] for more details).

Table 1: Parameters of the fit. In the first column we show the result of the fit when the parameters of the baryonic trajectories are fixed. The second column contains the result of the fit when the parameters of the trajectories are varied. † denotes parameters of the physical baryon trajectories from ref. [8]. * Coefficient $norm$ is chosen in such a way as to keep $A_{N_1^*} = 1$ in order to see the interplay between different resonances.

N_1^*	α_0	-0.8377 (fixed)†	-0.8070
	α_1	0.95 (fixed)†	0.9632
	α_2	0.1473 (fixed)†	0.1387
	$A_{N_1^*}$	1 (fixed)*	1 (fixed)*
	$Q^2_{N_1^*}$, GeV^2	2.4617	2.6066
N_2^*	α_0	-0.37(fixed)†	-0.3640
	α_1	0.95 (fixed)†	0.9531
	α_2	0.1471 (fixed)†	0.1239
	$A_{N_2^*}$	0.5399	0.6086
	$Q^2_{N_2^*}$, GeV^2	2.9727	2.6614
Δ	α_0	0.0038 (fixed)†	-0.0065
	α_1	0.85 (fixed)†	0.8355
	α_2	0.1969 (fixed)†	0.2320
	A_Δ	4.2225	4.7279
	Q^2_Δ, GeV^2	1.5722	1.4828
	s_0, GeV^2	1.14 (fixed)†	1.2871
E	α_0	0.5645	0.5484
	α_2	0.1126	0.1373
	s_E, GeV^2	1.3086	1.3139
	A_{exot}	19.2694	14.7267
	Q^2_{exot}, GeV^2	4.5259	4.6041
	$norm$	0.021	0.0207
$\chi^2_{d.o.f.}$		28.29	11.60

The first approach to the fitting procedure consists in fixing the parameters of the baryon trajectories such as order to reproduce the correct masses and widths, leaving the four scaling constants Q_i^2, four factors A_i and the parameters of the exotic trajectories to be fitted to

[2] We are grateful to M.I. Niculescu for making her data compilation available to us.

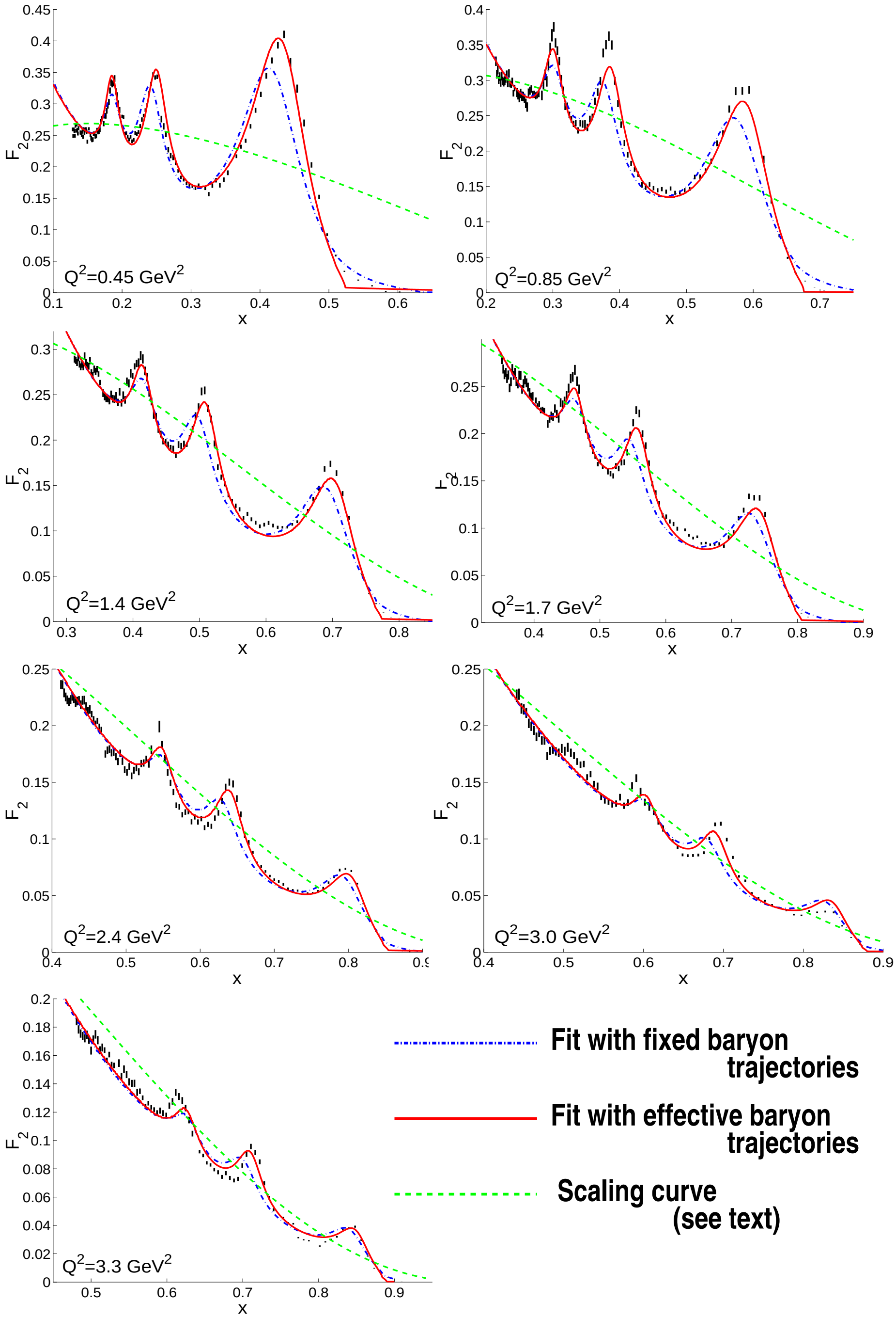

Figure 2: F_2 as a function of x for $Q^2 = 0.45 - 3.3$ GeV^2.

the data. The results are shown in Table 1 (first column) and the plots of the SF against x are presented in Fig. 2 (dashed-dotted lines). One can see that the agreement with the experimental data is poor ($\chi^2_{red} = 28.29$). To improve the fit we proceed as follows.

We try to account for the large number of resonances (about 20) present in the energy range under investigation, which overlap, as noted in above. To do this we consider the dominant resonances (N_1^*, N_2^* and Δ) as "effective" contributions to the SF. In other words we require that they mimic the contribution of the dominant resonances plus the large number of subleading contributions, which, together, fully describe the real physical system. Therefore, we have refitted the data, allowing the baryon trajectories parameters to vary.

The resulting parameters of such a fit are reported in Table 1 (second coloumn). It is worth noting that although the range of variation was not restricted, the new parameters of the trajectories stay close to their physical values, showing stability of the fit and thus reinforcing our previous considerations. From the relevant plots, shown in Fig. 2 with full lines, one can see that the improvement is significant, although agreement is still far from being perfect ($\chi^2_{d.o.f.} = 11.6$).

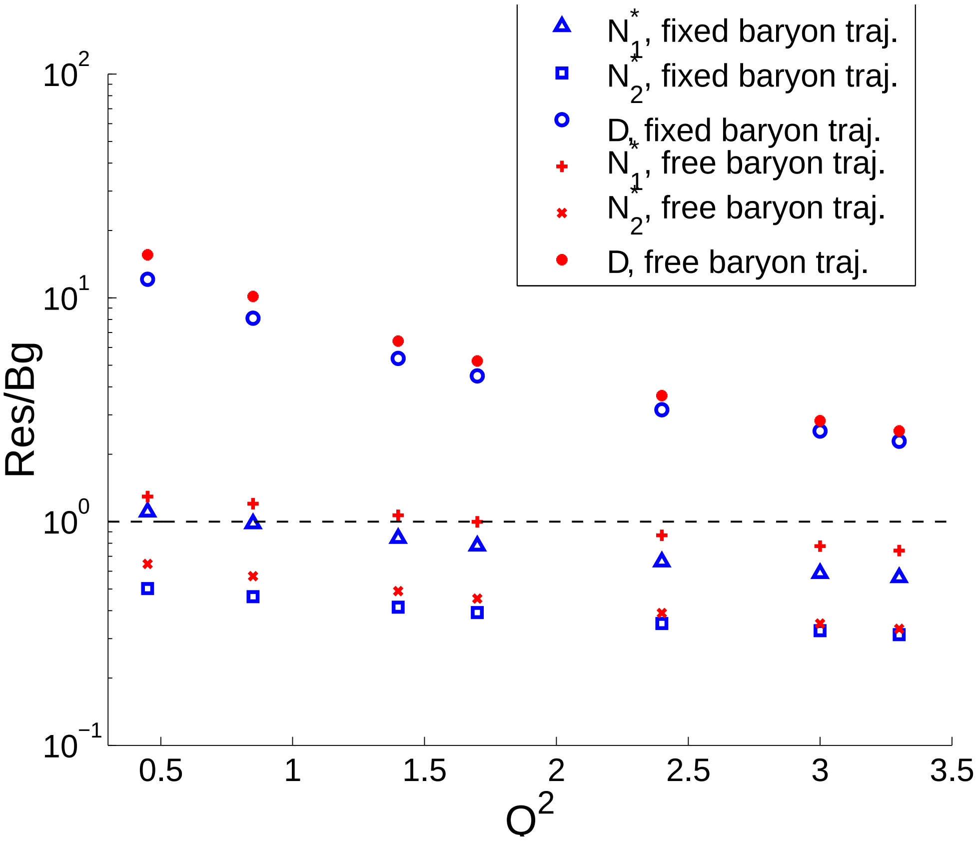

Figure 3: The ratio of the resonance to background components of the SF at the resonance peaks. See text for more details.

The dashed lines in Fig. (2) present a "scaling curves", i.e. a phenomenological parameterizations of the SF exhibiting Bjorken scaling and fitting the data. We have used the parameterizations studied in Ref. [22].

We have calculated also the Q^2-dependent ratio of the resonance to background components of the SF at three fixed values x, namely at three physical resonance peaks, $s_{N_1^*}$, $s_{N_2^*}$, s_Δ, for the fit with fixed physical baryon trajectories and at effective resonance peaks, for the fit with free baryon trajectories. On this plot the "background" for the selected resonance consists of three parts, i.e. the contribution from the exotic trajectory (usual background term) and the contributions from the two other resonances. The results are shown in Fig. 3. One may see that for N_2^* and for N_1^* for $Q^2 > 1.5$ GeV^2 the "background" contributes more than the resonant term itself. It was suggested in ref. [20] to study the Δ peak only in order to be able to neglect the contribution from the background. Our analyses shows that even for the Δ peak the background can not be neglected for Q^2 larger than $1.5 - 2$ GeV^2.

One may also ask the question of how good the dipole expression for the form factors, Eq. (9) does work. To answer this question we performed the fit letting the powers of the $1/(1+Q^2/Q^2_{0,i})$ in Eq. (9) free to vary. The results show that second power is a good approximation - powers change only by about 5%. As we shall see below the dipole approximation may deteriorate towards large values of Q^2 due to the spin effects, ignored in the present model.

Consider now the behaviour of $F_2(x,Q^2)$ at large x when s is kept in the resonance region. Let us remind the reader that x, Q^2 and s are related by Eq. (3) with $m = m_p$. Thus, each term in the rhs of Eq. (5), using Eqs. (1, 2) looks like

$$F_2(x,Q^2)_{i,n} = \frac{Q^2(1-x)}{4\pi\alpha(1+\frac{4m_p^2x^2}{Q^2})} \frac{norm\ A_i}{(1+\frac{Q^2}{Q^2_{0,i}})^{4(n-n_i^{min}+1)}}$$

$$\cdot \frac{\mathcal{I}m\ \alpha_i(s)}{(n-\mathcal{R}e\ \alpha_i(s))^2 + (\mathcal{I}m\ \alpha_i(s))^2}\ . \tag{10}$$

In the limit of x going to 1 and s in the resonance region ($1-4$ GeV^2), $Q^2 = x(s-m_p^2)/(1-x)$ is much larger than s and $Q^2_{0,i}$, which are of the same order. Thus we end up with

$$F_2(x,Q^2) \approx norm \sum_{i=N_1^*,N_2^*,\Delta,E} A_i \sum_{n=n_i^{min}}^{n_i^{max}} (1-x)^4$$

$$\cdot M_{i,n}(x,Q^2)\left(1 - 4\frac{Q^2_{0,i}+s-m_p^2}{Q^2} + O\left(\frac{1}{Q^4}\right)\right), \tag{11}$$

where

$$M_{i,n}(x,Q^2) = \frac{(s-m_p^2)x}{4\pi\alpha(1+\frac{4m_p^2x^2}{Q^2})}\left(\frac{Q^2_{0,i}}{s-m_p^2}\right)^4 \frac{\mathcal{I}m\ \alpha_i(s)}{(n-\mathcal{R}e\ \alpha_i(s))^2+(\mathcal{I}m\ \alpha_i(s))^2}\ . \tag{12}$$

In our range of interest $M_{i,n}$ is a slowly varying function of both x and Q^2. For each (i,n) the term proportional to $(1-x)^{4(n-n_i^{max}+1)}$ shows the main tendency of $F_2(x,Q^2)_{i,n}$, while $M_{i,n}$ is responsible for the "fine structure" - resonances at large x. Of course, for each trajectory i the main contribution comes from the first resonance - $(1-x)^4$.

The important ingredient neglected so far in the model is spin, i.e. helicity structure of the scattering amplitude and relevant selection rules. These change the form factors in a non-trivial way, thus complicating the Q^2-dependence of the SF's (see Ref. [20] for a recent treatment of the problem). These corrections have not yet been included in our study and might be responsible for relatively poor agreement with data. We hope to address this problem in a forthcoming work.

At this point it might be interesting to see the effect of spin corrections. As it has been shown in Ref. [20], if one explicitly takes into account the spin structure of the F_2, the main contribution from each resonance in the limit $x \to 1$ ($Q^2 \to \infty$) is proportional to $(1-x)^3$. Thus our model, neglected spin effects, strongly underestimate the physical SF.

3 Conclusions

The idea of the present paper is that deep inelastic scattering can be described by a sum of direct channel resonances lying on Regge trajectories. The form of these trajectories is crucial for the dynamics. It is constrained by analyticity, unitarity and by the experimental data. The use of baryon trajectories instead of individual resonances not only makes the model economic (several resonances are replaced by one trajectory) but also helps in classifying the resonances, by including the "right" ones and eliminating those nonexistent.

To fix the ideas and make a rough fit to the data, we constructed a simplified model with just 3 baryon trajectories, in which heavy thresholds have been replaced for simplicity by a linear term, and with the lowest-lying resonances. In fact, apart from the "prominent" three resonances many more should be included by means of relevant baryon trajectories. To this end an independent study of baryon trajectories and updated fits to dozens of existing resonances should be done. We intend to continue working in this direction.

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References

[1] D. Müller *et al.*, Fortschr. Phys. **42** (1994) 2.

[2] X. Ji, Phys. Rev. Lett. **78** (1997) 610; Phys. Rev. D **56** (1997) 7114.

[3] A.V. Radyushkin, Phys. Lett. B **380** (1996) 417; Phys. Rev. D **56** (1997) 5524.

[4] L. Elounadrhiri, hep-ph/0210341,

[5] L.L. Jenkovszky, V.K. Magas, E. Predazzi, Eur. Phys. J. **A12** 361 (2001).

[6] L.L. Jenkovszky, V.K. Magas, E. Predazzi, nucl-th/0110085; L.L. Jenkovszky, V.K. Magas, hep-ph/0111398; L.L. Jenkovszky, T. Korzhinskaja , V.I. Kuvshinov and V.K. Magas, Proceedings of the New Trend in High-Energy Physics, Yalta, Crimea, Ukraine, September 22-29, 2001, edited by P.N. Bogolyubov and L.L. Jenkovszky (Bogolyubov Institute for Theoretical Physics, Kiev, 2001), p. 178.

[7] L.L. Jenkovszky, V.K. Magas and T. Szabo, Proceedings of the International Worshop on Very High Mulitiplicities, Dubna, Russia, 7-9 April 2001, p. 185.

[8] R. Fiore, A. Flachi, L.L. Jenkovszky, A. Lengyel, V.K. Magas, hep-ph/0206027, to appear in Eur. Phys. J. A.

[9] L.L. Jenkovszky, S. Yu. Kononenko, V.K. Magas, Proceedings of the NATO Advanced Research Workshop DIFFRACTION 2002, Alushta, Ukraine, August 31 - September 6, 2002 - in press.

[10] M. Köbberling et al., Nucl. Phys. **82** (1974) 201.

[11] S. Stein et al., Phys. Rev. **12** (1975) 1884.

[12] F.W. Brasse et al., Nucl. Phys. **B110** (1976) 413.

[13] A. Bodek et al., Phys. Rev. **D20** (1979) 1471.

[14] L.J. Whitlow et al., Phys. Lett. **B250** (1990) 193.

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

[16] P. Stoler, Phys. Rev. **D44** (1991) 73; Phys. Rev. Lett. **66** (1991) 1003.

[17] C.E. Carlson, N.C. Mukhopadhyaya, Phys. Rev. Lett. **74** (1995) 1288; Phys. Rev. **D58**(1998) 094029; Eur. Phus. J **C**, in press.

[18] P. Freund, Phys. Rev. Lett. **20**, (1968) 235; H. Harari, Phys. Rev. Lett. **20**, (1968) 1395.

[19] G. Domokos, S. Kövesi-Domokos, E. Schoenberg, Phys. Rev. **3** (1971) 1184.

[20] V.V. Davydovsky and B.V. Struminsky, hep-ph/0205130; Proceedings of the NATO Advanced Research Workshop DIFFRACTION 2002, Alushta, Ukraine, August 31 - September 6, 2002 - in press.

[21] http://www.jlab.org/resdata.

[22] L. Csernai, L. Jenkovszky, K. Kontros, A. Lengyel, V. Magas and F. Paccanoni, Eur. Phys. J. C **24** (2002) 205.