GENERALIZED PARTON DISTRIBUTIONS
FROM NUCLEON FORM FACTORS

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Abstract
Results from a recent analysis of the zero-skewness generalized parton distributions (GPDs) for valence quarks are reviewed. The analysis bases on a physically motivated parameterization of the GPDs with a few free parameters adjusted to the nucleon form factor data. The Fourier transforms of the GPDs representing quark densities in the impact parameter plane, as well as momentum sums of the GPDs are also discussed. The 1-x momentum in particular form the soft physics input to Compton scattering off protons within the handbag approach. The Compton cross section evaluated from this information is found to be in good agreement with experiment.

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1 Introduction

In recent years we have learned how to deal with hard exclusive reactions within QCD. In analogy to hard inclusive reactions the process amplitudes factorize in partonic subprocess amplitudes, calculable in perturbation theory, and in GPDs which parameterize soft hadronic matrix elements. In some cases rigorous proofs of factorization exist. For other processes factorization is shown to hold in certain limits, under certain assumptions or is just a hypothesis. The GPDs which are defined by Fourier transforms of biquark interaction matrix elements of quark field operators \[ F \], describe the emission and reabsorption of partons by the proton. For equal helicities of the emitted and reabsorbed parton the structure of the nucleon is described by four GPDs, termed \( H \), \( \bar{F} \), \( E \) and \( \bar{E} \), for each quark flavor and the gluons. The GPDs are functions of three variables, the momentum transfer from the initial to the final proton, \( t \), a momentum fraction \( x \) and the skewness, \( \xi \). The latter variable is related to the difference of the individual momentum fractions the emitted and reabsorbed partons carry. The GPDs are subject to evolution and, hence, depend on the factorization scale \( \Lambda \), too. They satisfy the reduction formulas

\[ H_q(x; \xi = 0; t = 0) = q(x); \quad \bar{F}_q(x; \xi = 0; t = 0) = q(x); \quad E_q(x; \xi = 0; t = 0) = q(x); \quad \bar{E}_q(x; \xi = 0; t = 0) = q(x); \quad (1) \]

i.e. in the forward limit \( H \) and \( \bar{F} \) reduce to the usual unpolarized and polarized parton distributions (PDFs), respectively. Another property of the GPDs is the polynomiality which comes about as a consequence of Lorentz covariance

\[ \int_{-1}^{1} \frac{1}{x^m} H_q(x; t; \xi) dx \sim \sum_i h_{m,i}^q (t; \xi) x^i; \quad (2) \]
where \( m = 2 \) denotes the largest integer smaller than or equal to \( m = 2 \). Eq. (2) holds analogously for the other GPDs and, for \( m = 1 \), implies sum rules for the form factors of the nucleon, e.g.

\[
F_1^q(t) \quad H_{A_0}^q(t) = \frac{Z_1}{1} \int dx H_0^q(x; \ t; t_1); 
\]

(3)

Reinterpreting as usual a parton with a negative momentum fraction \( x \) as an antiparton with positive \( x \) \( (H_0^q(x) = H_q^q(x)) \), one becomes aware that only the difference of the contributions from quarks and antiquarks of given flavor contribute to the sum rules. Introducing the combination

\[
H_0^q(x; \ t; t_1) = H_q^q(x; \ t; t_1) - H_q^q(x; \ t; t_1); 
\]

(4)

for positive \( x \) which, in the forward limit, reduces to the usual valence quark density \( q^v(x) = q(x) - q(x) \), one finds for the Dirac form factor the representation

\[
F_1^p(t) = e_u(q) \frac{Z_1}{1} \int dx H_u^q(x; \ t; t_1) + e_d(q) \frac{Z_1}{1} \int dx H_d^q(x; \ t; t_1); 
\]

(5)

Here, \( e_q \) is the charge of the quark \( q \) in units of the positron charge. Possible contributions from other flavors, \( s \) or \( c \) or \( c \), are neglected in the sum rule (6). A representation analogous to (5) holds for the Pauli form factor with \( H \) being replaced by \( E \).

The isovector axial-vector form factor satisfies the sum rule \( (\tilde{A}_0^q(x) = \tilde{A}_0^q(x)) \)

\[
F_A(t) = \frac{Z_2}{1} \int dx \tilde{A}_0^u(x; \ t; t_1) + \tilde{A}_0^d(x; \ t; t_1) + 2 \tilde{A}_0^u(x; \ t; t_1) + \tilde{A}_0^d(x; \ t; t_1); 
\]

(6)

At least for small \( t \) the magnitude of the second integral in Eq. (6) reflects the size of the flavor non-singlet combination \( u(x) - d(x) \) of forward densities. This difference is poorly known, and at present there is no experimental evidence that it might be large [2]. For instance, in the analysis of the polarized PDFs performed by Blumlein and Bottcher Ref. [3] it is zero. In a perhaps crude approximation the second term in (6) can be neglected. These simplifications to the sum rules (5) and (6) do not imply that the nucleon is assumed to consist solely of valence quarks. Sea quarks are there but the virtual photon that probes the quark content of the nucleon, sees only the differences between quark and antiquark distributions which are likely small.

2 A determination of the zero-skewness GPDs

Since the GPDs cannot be determined from QCD with a sufficient degree of accuracy at present we have to either rely on models or to extract them from experiment as it has been done for the PDFs, see for instance Refs. [3, 4]. A first attempt to extract the GPDs phenomenologically has been carried through in Ref. [5] (see a similar analysis Ref. [6]). The idea is to exploit the sum rules (5) and (6) and to determine the GPDs from the data (7) on the nucleon form factors, \( F_1 \) and \( F_2 \) for proton and neutron as well as \( F_A \). Since the sum rules represent only the first moments of the GPDs this task is an ill-posed problem in a strict mathematical sense. In many applications the moments are needed to deduce the integrand, i.e. the GPDs, unambiguously from an integral. However, from a phenomenological point of view one expects the GPDs...
to be smooth functions and a small number of moments will likely suffice to \( x \) the GPDs. The extreme and at present the only feasible point of view is that the lowest moment of a GPD alone is sufficient [5]. Indeed, using recent results on PDFs [3, 4] and form factor data [7] in combination with physically motivated parametrizations of the GPDs, one can carry through this analysis. Needless to say that this method while phenomenologically successful as will be discussed below, does not lead to unique results. Other parametrizations which may imply different physics, cannot be excluded at the present stage of the art.

The sum rules (5), (6) are valid at all but guessing a plausible parametrization of a function of three variables is a nearly hopeless task. Choosing the special value \( q = 0 \) for which the emitted and reabsorbed partons carry the same momentum fractions, has many advantages. One exclusively works in the so-called DGLAP region where \( x \). In this region parton ideas apply and the Fourier transform of a GPD with respect to the momentum transfer \( (\frac{2}{\sqrt{t}}) \) has a density interpretation in the impact parameter plane [5]. Moreover, as shown in Refs. [3, 12], wide-angle exclusive processes are controlled by generalized form factors that represent \( l=x \) momenta of zero-skewness GPDs. The parametrization that is exploited in Ref. [5] combines the usual PDFs with an exponential dependence (the arguments \( x = 0 \) and are omitted in the following)

\[
H_3^0(x; t) = q(x) \exp[ f_q(x)];
\]

where the profile function reads

\[
f_q(x) = [0 \log(l=x) + B_q(l \times)^{x+1} + A_q x (l \times)]^2
\]

This ansatz is motivated by the expected Regge behavior at low \( t \) and low \( x \) [11] (\( 0 \) is the Regge slope for which the value 0.9 GeV \(^2\) is imposed). For large \( t \) and large \( x \), on the other hand, one expects a behavior like \( f_q \) \( 1 \times \) from overlaps of light-cone wavefunctions [3, 12]. The ansatz (7), (8) interpolates between the two limits smoothly \( 0 \) and allows for a stronger suppression of \( f_q \) in the limit \( x \rightarrow 1 \). It matches the following criteria for a good parametrization: simplicity, consistency with theoretical and phenomenological constraints, stability with respect to variations of PDFs and stability under evolution (scale dependence of the GPDs can be absorbed into parametters).

Using the CTEQ PDFs [3], one obtains a reasonable ansatz (9) of the ansatz (7), (8) to the data on the Dirac form factor (from \( t = 0 \) up to \( 30 \) GeV \(^2\)) with the parameters

\[
B_u = B_d = (0.59 \quad 0.03) \text{ GeV}^2;
A_u = (1.22 \quad 0.020) \text{ GeV}^2; \quad A_d = (2.59 \quad 0.29) \text{ GeV}^2;
\]

quoted for the case \( n = 2 \) and at a scale of \( \mu = 2 \) GeV. In Fig. 11 the results for \( H_3^0 \) are shown at three values of \( t \). While at small \( t \) the behavior of the GPD still reflects that of the parton densities it exhibits a pronounced maximum at larger values of \( t \). The maximum becomes more pronounced with increasing \( t \) and its position moves towards \( x = 1 \). In other words only a limited range of \( x \) contributes to the form factor substantially. The quality of the \( t \) is very similar in both the cases, \( n = 1 \) and 2; the results for the GPDs agree well with each other. Substantial differences between the two

\[1\] The parameter \( B_q \) is not needed if \( 0 \) is fixed. A \( t \) to the data of about the same quality and with practically the same results for the GPDs is obtained for \( 0' \) \( 1s \).
results only occur for very low and very large values of \( x \), i.e., in the regions which are nearly insensitive to the present form factor data. It is the physical interpretation of the results which favours the \( t \) with \( n = 2 \). Indeed the average distance between the struck quark and the cluster of spectators becomes unphysical large for \( x = 1 \) in the case \( n = 1 \); it grows like \( (1 - x)^3 \) while, for \( n = 2 \), it tends towards a constant value of about 0.5 fm [5].

The analysis of the axial and Pauli form factors, with parametrizations analogous to Eqs. (7), (9), provides the GPDs \( A \) and \( E \). In general they behave similar to \( H \). Note-worthy differences are the opposite signs of \( A^u \), \( E^u \), and \( A^d \), \( E^d \) and the approximately same magnitude of \( E^u \) and \( E^d \) at least for not too large values of \( t \). For \( H^3 \) and \( H^5 \), on the other hand, the d-quark contributions are substantially smaller in magnitude than the u-quark ones, see Fig. 1. Since there is no data available for the pseudoscalar form factor of the nucleon the GPD \( E \) cannot be determined this way.

3 Moments of the GPDs

Having the GPDs at disposal one can evaluate various moments, some of them are displayed in Fig. 2. Comparison with recent results from lattice QCD [13] reveals remarkable agreement of their \( t \) dependencies given the uncertainties of the GPD analysis [5] and of the lattice calculations. It is to be stressed that the results presented in Refs. [13] are provided for a scenario in which the pion has a mass of 600 to 800 MeV. The extrapolation to the chiral limit is therefore problematic and may entail failures of the normalization of the moments. An interesting property of the moments is that the u and d quark contributions decrease with different rates at large \( t \). Because at large \( t \) the dominant contribution to the form factors comes from a narrow region of large \( x \), see Fig. 2, one can use the large \( x \) approximations

\[
q \quad (1 - x)^3; \quad f_q \quad A_q (1 - x)^3; \quad (10)
\]
Figure 2: Left: The first three moments of $H_u$ scaled by $t^2$. The error bands denote the parametric uncertainties. Right: The lowest moment of $H_d$ scaled by $t^2$. The shaded band represents the results obtained in Ref. [5]. Data are taken from [7,15].

and evaluate the sum rule (5) in the saddle point approximation. This leads to

$$h_{1,0}^q = \frac{i}{j} (\frac{1}{b})^{2-n} ; \quad 1 \quad x_s = \frac{n}{q} A_q \frac{j}{1-n} \quad (11)$$

where $x_s$ is the position of the saddle point. For $n = 2$ the saddle point lies within the region where the GPD is large $\mathcal{G}$, and using the the CTEQ values for $q (u ' 3 \pi \text{ and } d ' 5 \mathcal{Q})$ one obtains a drop of the form factor $h_{1,0}^q$ slightly faster than $t^2$ while the d-quark form factor falls as $t^3$. Strengthened by the charge factor the u-quark contribution dominates the proton's D irac form factor for $t$ larger than about 5 GeV$^2$, the d-quark contribution amounts to less than 10%. This is what can be seen in Figs. 2. The power behavior bears resemblance to the Drell-Yan relation (14). In fact the common underlying dynamics is the Feynman mechanism $\mathcal{G}$. The Drell-Yan relation, however, an asymptotic result $(x \gg 1, t \gg 1)$ which bases on the assumption of valence Fock state dominance, i.e. on the absence of sea quarks. The different powers for $u$ and $d$ quarks signal that the asymptotic region where the dimensional counting rules apply, has not yet been reached.

One may object to the different powers that they are merely a consequence of the chosen parameterization (7), (5). However, this is likely not the case. The moments $h_{1,0}^q$ can directly be extracted from the form factor data (7). The experimental results for the d-quark moment are shown in Fig. 2. The sharp drop of it for increasing $t$ is clearly visible. The small deviations between data and the moment obtained in [5] are caused by the very precise but still preliminary CLAS data (15) for $G_\pi^5$. These data have been

$^2$ For $n = 1$ this is only the case for $t$ larger than 30 GeV$^2$. The onset of the corresponding power behaviour of the form factors occurs for larger $t$.

$^3$ This implies the ratio $F_1^p/F_1^n = e_u = e_d$ at large $t$.

$^4$ The Feynman mechanism applies in the soft region where $1 - x = P - t$ and the virtualities of the active partons are $\frac{P - t}{t}$ (is a typical hadronic scale).
utilized in the extraction of the moments $h_{i,\rho}^3$ but not in the analysis in Ref. [15]. Worst measured of the four form factors is the electric one of the neutron. The recent data [16] on it are shown in Fig. 3. Above 1.5 GeV$^2$ no data is available yet and extrapolations to larger $t$ are to be used. The solid line has been used in Ref. [15] and in the determination of $h_{i,\rho}^3$ from data. If the electric form factor is smaller than that extrapolation (e.g. the dashed line) $h_{i,\rho}^3$ falls even faster. Thus, except $G_E^n$ is much larger than expected, the rapid decrease of $h_{i,\rho}^3$ seems to be an experimental fact. Future JLab data on $G_E^n$ will settle this question.

A combination of the second moments of $H$ and $E$ at $t = 0$ is Ji's sum rule [17] which allows for an evaluation of the valence quark contribution to the orbital angular momentum the quarks inside the proton carry

$$h_{L,i}^v = \frac{1}{2} \int_0^1 dx \ x E^q_v(x; t = 0) + x q_v(x) : q_v(x) :$$  \hbox{(12)}

The analysis of the GPDs leads to

$$h_{L,i}^v = (0.24 \quad 0.27); \quad h_{d,i}^v = 0.15 \quad 0.19; \quad (13)$$

for the valence quark contributions to the orbital angular momentum at a scale of $= 2$ GeV.

While the parton distributions only provide information on the longitudinal distribution of quarks inside the nucleon, GPDs also give access to the transverse position distributions of partons within the proton. Thus, the Fourier transform of $H$

$$q_v(x; b) = \frac{Z}{(2 \pi)^2} e^{ib \cdot \mathbf{x}} H^q_v(x; t = -2);$$  \hbox{(14)}

gives the probability of finding a valence quark with longitudinal momentum fraction $x$ and impact parameter $b = (b^X; b^Y)$ as seen in a frame in which the proton moves rapidly
in the Z direction. Together with the analogous Fourier transform of \( E_q^Y(x;t) \) one can form the combination (\( m_p \) being the mass of the proton)

\[
q^X(x;b) = q_v(x;b) \frac{b^x}{m_p} \frac{\Theta}{b^2} e^Y_q(x;b);
\]

which gives the probability to find an unpolarized valence quark with momentum fraction \( x \) and impact parameter \( b \) in a proton that moves rapidly along the Z direction and is polarized along the X direction [3]. As shown in Ref. [4], for small \( x \) one observes a very broad distribution while at large \( x \) it becomes more focused on the center of m om entum defined by \( \int x_b = 0 \) (\( x = 1 \)). In a proton that is polarized in the X direction the symmetry around the Z axis is lost and the center of the density is shifted in the Y direction away from the center of m om entum, downward for d quarks and upward for u ones. In other words, a polarization of the proton induces a flavor segregation in the direction orthogonal to the direction of the polarization and the proton m om entum, see Fig. 1. This effect may be responsible for certain asymmetries as, for instance, that one observed in p+p!X.

4 Wide-angle scattering

As mentioned previously the analysis of the GPDs gives insight in the transverse distribution of quarks inside the proton. However, there is more in it. With \( Q^2 = 0 \) GPDs at hand one can predict hard wide-angle exclusive reactions like Compton scattering off protons or meson photo- and electroproduction. For these reactions one can work in a so-called symmetric frame where the skewness is zero. It has been argued [9,10] that, for large Mandelstam variables \( (s; t; u) \), the amplitudes for these reactions factorize into a hard partonic subprocess, e.g. Compton scattering off quarks, and in form factors representing \( l=x-m \) on of zero-skewness GPDs (see Fig. 4). For Compton scattering these form factors read

\[
\begin{align*}
R_V(t) &= \int_{-1}^{+1} \frac{dx}{x} e^2_q \int_0^Z e^{1-x} dx H^q_v(x;t) ; \\
R_A(t) &= \int_{-1}^{+1} \frac{dx}{x} e^2_q \int_0^Z e^{1-x} dx A^q_v(x;t) ; \\
R_T(t) &= \int_{-1}^{+1} \frac{dx}{x} e^2_q \int_0^Z e^{1-x} dx E^q_v(x;t) ;
\end{align*}
\]

In these expressions contributions from sea quarks have been neglected \(^5\). An analogous pseudoscalar form factor related to the GPD \( \overline{E} \), decouples in the symmetric frame. Numerical results for the Compton form factors are shown in Fig. 4. To leading-order of pQCD the Compton cross section reads [5,13]

\[
\frac{d}{dt} = \frac{d^2}{dt} \left( \frac{1}{2} R^Y_V(t) + \frac{t}{4m^2} R^Y_A(t) + R^Y_T(t) \right)
\]

\(^5\) An estimate of the sea quark contributions may be obtained by using the ansatz (7) with the same profile function for the sea quarks as for the valence ones but replacing the valence quark density with the CTEQ [11] antiquark ones. The so estimated contributions are indeed small.
where $d^\omega dt$ is the Klein-Nishina cross section for Compton scattering of massless, point-like spin-1/2 particles of charge unity. Inserting the Compton form factors (16) into Eqs. (17), one can predict the Compton cross section in the wide-angle region. The results for sample values of $s$ are shown in Fig. 4 and compared to recent measurements from JLab [19]. The inner bands of the predictions for $d^\omega dt$ reflect the parametric errors of the form factors. The outer bands indicate the size of target mass corrections, see Ref. [20]. In order to comply with the kinematical requirement for handbag factorization, at least in a minimal fashion, predictions are only shown for $t$ and $u$ larger than about 2.5 GeV$^2$. Fair agreement between theory and experiment is to be seen. Next-to-leading order QCD corrections to the subprocess have been calculated in Ref. [19]. They are not displayed in (17) but taken into account in the numerical results.

The handbag approach also applies to wide-angle photon- and electroproduction of pseudoscalar and vector mesons [21]. The amplitudes again factorize into a parton-level subprocess, $q \rightarrow M q$, and form factors which represent $1=x$ moments of GPDs. Their flavor composition reflects the valence quark content of the produced meson. Therefore, the form factors which occur in photo- and electroproduction of pions and mesons, can also be evaluated from the GPDs given in Ref. [5].

One may also consider the time-like process $\gamma p \rightarrow \gamma p$ within the handbag approach [22]. Similar form factors as in the space-like region occur but they are now functions of $s$ and represent moment of the pp distribution amplitudes, time-like versions of GPDs. With sufficient data on the time-like electron magnetic form factors at disposal one may attempt a determination of the time-like GPDs.
5 Outlook

Results from a rst analysis of GPDs at zero skewness have been reviewed. This analysis, has been performed in analogy to those of the usual parton distributions, rests upon a physically motivated param e terization of the GPDs set ted to the available nucleon form factor data. The analysis provides results on the valence-quark GPDs \( h, \tilde{h} \) and \( E \). The distribution of the quarks in the impact parame ter plane transverse to the direction of the nucleon’s mom entum can be evaluated from them. Polarizing the nucleon induces a anor segmentation in the direction orthogonal to the those of the nucleon’s mom entum and of its polarization. The average orbital angular mom entum of the valence quarks can be estim ated from the obtained GPDs, too. Within the handbag approach the soft physics input to hard wide-angle exclusive reactions is encoded in speci rstorm factors which represent 1=x mom ents of zero-skewness GPDs. Using the GPDs deter mined in Ref. [5], one can evaluate these form factors and, for instance, predict the wide-angle Compton cross section. The results are found to be in good agreement with experiment.

An analysis as that one performed in Ref. [5] can only be considered as a rst attempt towards a deter mination of the GPDs. It needs in proven ents in various aspects. High quality data on the form factors at larger t are required in order to stabilize the param e teriza tion. JLab will provide such data in the near future. As already mentioned CLAS [15] will come up with data on \( G_n^m \) up to about 5 GeV\(^2\). \( G_n^p \) will be measured up to about 3.3 GeV\(^2\) next year and \( G_E^p \) up to 9 GeV\(^2\) in 2007. The upgraded JLab will allow measurements of the nucleon form factors up to about 13 GeV\(^2\). Data on the axial form factor are needed in order to improve our knowledge of \( \tilde{h} \).

Up to now only the sum rules (5), (6) have been utilized in the GPD analysis. As mentioned above, in this situation param e terizations of the GPDs are required with the consequence of non-unique results. An alternative ansatz [5] is for instance

\[
H_v^x(x; t) = \sum_{i=1}^{h} f_i(x) l t f_i(x) = p
\]

Although reasonable results to the form factors are obtained with (18) for \( p > 2.5 \), it is physically less appealing than the param e terization (7): the combination of Regge behavior at small x and t with the dynamics of the Feynman mechanism is lost. The resulting GPDs have a broader shape and \( H(x = 0; t) \) remains finite. Thus, small x also contribute to the large-t form factors. Higher mom ents are needed in order to lessen the dependence on the chosen param e terization for the GPDs. Such mom ents can be provided by lattice gauge theories. The present lattice results [13] are however calculated in scenarios where the pion is heavy (typically 600 800 MeV) and the extrapolation to the chiral limit is uncertain. Obviously such results are inappropriate for use in a GPD analysis. In a few years the quality of the mom ents from lattice gauge theories may su ce. It is also tempting to use the data on wide-angle Compton scattering [19] which imply information on 1=x mom ents of the GPDs. However, even at the highest measured energy, \( s = 11 \text{ GeV}^2 \), the Compton cross section may still be cont aminated by power corrections rendering its use in a GPD analysis dubious. For data at, say, \( s = 20 \text{ GeV}^2 \) the situation may be different. Instead of a param e terization one may use the maximum entropy method [23] for the extraction of the GPDs from the sum rules. This method has not yet been utilized for this purpose.

Up to now only the zero-skewness GPDs have been deter mined. Although the sum
rules (5), (6) are valid at all values of the skewness it seems a hopeless task to extract functions of three variables from them. Additional information is demanded in the case of \( \xi \neq 0 \) and will be provided by deeply virtual exclusive scattering (DVES). The skewness is related to Bjorken-\( x \) in DVES and, hence, non-zero. At large \( Q^2 \) but small the amplitudes for DVES are given by a convolution of GPDs and appropriate subprocess amplitudes. The convolution possesses poles at \( x = x_0 \). For Compton scattering the pole terms can directly be isolated in the interference region with the Bethe-Heitler process and will lead to a nearly model-independent determination of \( H(\xi, t) \). For \( x \neq x_0 \) parameterizations of the GPDs will most likely be needed. At present a number of experiments on DVES are running (HERA, COMPASS, HERMES, JLab). In Fig. 5 kinematical regions are shown as shaded areas in which information on GPDs is currently accumulated. A lot of work is ahead of us before we can say that we have a fair knowledge on the GPDs.

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