SPIN EFFECTS IN HARD EXCLUSIVE REACTIONS

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Abstract

The present status of applications of perturbative QCD to large momentum transfer exclusive reactions is discussed. It is argued that in the region of momentum transfer accessible to present day experiments soft contributions dominate the exclusive processes.

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1. Introduction

This talk is not meant as a review but rather represents my opinion about the current status of the theoretical understanding of hard exclusive reactions. Although we are mainly interested in spin physics at this conference I am compelled to discuss unpolarized observables first since their understanding is prerequisite for a reliable treatment of the in general smaller spin effects.

There is general agreement that perturbative QCD in the framework of the hard-scattering approach (HSA) \cite{1} is the correct description of form factors and perhaps other exclusive reactions at asymptotically large momentum transfer. In that approach a form factor or a scattering amplitude is expressed by a convolution of distribution amplitudes (DA) with hard scattering amplitudes to be calculated in collinear approximation within perturbative QCD. The universal, process independent DAs, which represent hadronic wave functions integrated over transverse momenta ($k_\perp$), are controlled by long distance physics in contrast to the hard scattering amplitudes which are governed by short distance physics. Explicitly, a helicity amplitude for a process $AB \rightarrow CD$ reads

$$M_{CDAB}(s, t) = \int \prod_{i=A,B,C,D} [dx_i] \phi^*_C(x_C, \mu_F)\phi^*_D(x_D, \mu_F)T_H(x_i, s, t, \mu) \times \phi_A(x_A, \mu_F)\phi_B(x_B, \mu_F) \quad (1.1)$$

where helicity labels are omitted for convenience. $[dx_i]$ is short for

$$[dx_i] = \prod_{j=1}^{n_i} dx_{ij} \delta(1 - \sum x_{ij}) \quad (1.2)$$

$x_{ij}$ is the momentum fraction the constituent $j$ of hadron $i$ carries. $\mu_F$ is the scale at which short and long distance physics factorizes and $\mu$ is the renormalization scale.

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ization scale. If one of the external particles is point-like, e. g. a photon, the corresponding DA is to be replaced by \( \delta(1-x_{1}) \).

The HSA possesses two characteristic properties, the quark counting rules and the helicity sum rule. The first property says that the fixed angle cross-section behaves at large Mandelstam \( s \) as

\[
\frac{d\sigma}{dt} = f(\theta) s^{2-n} \quad \text{(modulo log} \ s) \quad (1.3)
\]

where \( n \) is the minimum number of external particles in the hard scattering amplitude. The counting rules also apply to form factors: a baryon form factor behaves as \( 1/Q^4 \), a meson form factor as \( 1/Q^2 \) at large momentum transfer \( Q \).

These counting rules are in surprisingly good agreement with experimental data. Even at momentum transfers as low as 2 GeV the data seem to respect the counting rules.

The second characteristic property of the HSA is the conservation of hadronic helicity. For a two-body process the helicity sum rule reads

\[
\lambda_A + \lambda_B = \lambda_C + \lambda_D. \quad (1.4)
\]

It appears as a consequence of utilizing the collinear approximation and of dealing with (almost) massless quarks which conserve their helicities when interacting with gluons. The collinear approximation implies that the relative orbital angular momentum between the constituents has a zero component in the direction of the parent hadron. Hence the helicities of the constituents sum up to the helicity of their parent hadron. As will be discussed below hadronic helicity is not conserved; the ratio of hadronic helicity flip matrix elements to non-flip ones is about 0.2.

Many hard exclusive reactions have been analysed within the framework of the HSA: Electromagnetic form factors of mesons and baryons, Compton scattering off nucleons, photoproduction of mesons, two-photon annihilations into pairs of mesons or baryons, decays of heavy mesons etc. No clear picture has emerged as yet; there are successes and failures. It however seems that results of the order of the experimental values are only obtained if, at least for the proton and the pion, DAs are used which are strongly concentrated in the end-point regions. Chernyak and Zhitnitsky (CZ) \[2\] claimed that such DAs find a certain justification in QCD sum rules by means of which a few moments of the DAs have been calculated.

The CZ moments are subject of considerable controversy: Other QCD sum rule studies as well as lattice gauge theory provide other values for the moments. The asymptotic forms of the DAs (\( \sim x_1 x_2 \) for the pion, \( \sim x_1 x_2 x_3 \) for the proton), into which any DA evolves for \( Q \to \infty \), lead to results which are typically orders of magnitudes too small as compared with data. Consider as an example the magnetic form factor of the proton. For the DAs of the CZ type one obtains \( Q^4 G_M \simeq 1 \text{GeV}^4 \) in agreement with experiment, whereas a vanishing result is found for the asymptotic DA.

Purely hadronic reactions, as for instance elastic proton-proton scattering, have not yet been studied in the framework of the HSA. The reason for that fact is, on
the one hand, the huge number of Feynman graphs contributing to such reactions and, on the other hand, the occurrence of multiple scatterings (pinch singularities [3]), i.e. the possibility that pairs of constituents scatter independently in contrast to the HSA in which all constituents collide within a small region of space-time. A general framework for treating multiple scattering contributions has been developed by Botts and Sterman [4].

2. The modified perturbative approach

The applicability of the HSA at experimentally accessible momentum transfers, typically a few GeV, was questioned [5]. It was asserted that in the few GeV region the HSA accumulates large contributions from the soft end-point regions, rendering the perturbation calculation inconsistent. This is in particular the case for the end-point concentrated DAs. Another theoretical defect is caused by the collinear approximation: The neglect of the transverse momentum dependence of the hard scattering amplitude is a poor approximation in the end-point regions. The magnitude of the errors in the final results for, say, the pion’s or the nucleon’s form factor, generated by the collinear approximation depends on the shape of the DAs. Obviously, the CZ-like DAs entail large errors in contrast to the asymptotic DA and similar forms for which the errors are sufficiently small.

Recently a modification of the HSA was suggested [4, 6] in which the transverse momentum dependence of the hard scattering amplitude is retained and Sudakov corrections are taken into account. Let me discuss the characteristics of that approach on the example of the \( \pi \gamma \) transition form factor which is written as [7]

\[
F_{\pi\gamma}(Q^2) = \int dx \frac{d^2b}{4\pi} \hat{\Psi}_0(x, -b, \mu_F) \hat{T}_H(x, b, Q) \exp[-S(x, b, Q)],
\]

(2.1)

where \( b \) is the quark-antiquark separation in the transverse configuration space. \( \hat{T}_H \) denotes the Fourier transform of the momentum space hard scattering amplitude. Neglecting masses, it reads at lowest order (QED)

\[
\hat{T}_H(x, b, Q) = \frac{2}{\sqrt{3\pi}} K_0(\sqrt{(1-x)Qb}) + \mathcal{O}(\alpha_S),
\]

(2.2)

where \( K_0 \) is the modified Bessel function of order zero. The Sudakov exponent \( S \) in (2.1) comprising those gluonic radiative corrections not taken into account by the evolution of the wave function, is given by

\[
S(x, b, Q) = s(x, b, Q) + s(1 - x, b, Q) - \frac{4}{\beta_0} \ln \left( \frac{\mu}{\Lambda_{QCD}} \right) \ln \left( \frac{1}{b\Lambda_{QCD}} \right)
\]

(2.3)

where a Sudakov function \( s \) appears for each quark line entering the hard scattering amplitude. The last term in (2.3) arises from the application of the renormalization group equation (\( \beta_0 = 11 - 2/3 n_f \)). A value of 200 MeV for \( \Lambda_{QCD} \) is...
used throughout and the renormalization scale \( \mu \) is taken to be the largest mass scale appearing in \( \hat{T}_H \), i.e., \( \mu = \max(\sqrt{1-xQ}, 1/b) \). For small \( b \) there is no suppression from the Sudakov factor; as \( b \) increases the Sudakov factor decreases, reaching zero at \( b = 1/\Lambda_{QCD} \). For even larger \( b \) the Sudakov factor is set to zero. The Sudakov function \( s \) is explicitly given in [4, 6]. \( b \) plays the role of an infrared cut-off; it sets up the interface between non-perturbative soft gluon contributions – still contained in the hadronic wave function \( \hat{\Psi}_0 \) – and perturbative soft gluon contributions accounted for by the Sudakov factor. Hence, the factorization scale \( \mu_F \) is taken to be \( 1/b \).

The quantity \( \hat{\Psi}_0 \) appearing in (2.1) represents the soft part of the transverse configuration space pion wave function, i.e., the full wave function with the perturbative tail removed from it. The wave function is parameterized as [8]

\[
\hat{\Psi}_0(x, b, \mu_F) = \frac{f_\pi}{2\sqrt{6}} \varphi(\mu_F) \hat{\Sigma}(\sqrt{x(1-x)} b).
\] (2.4)

It is subject to the auxiliary conditions

\[
\hat{\Sigma}(0) = 4\pi, \quad \int_0^1 dx \varphi(x, \mu_F) = 1.
\] (2.5)

\( f_\pi (= 130.7 \text{MeV}) \) is the usual \( \pi \)-decay constant. The wave function does not factorize in \( x \) and \( b \), but in accord with the basic properties of the HSA [9, 10] the \( b \)-dependence rather appears in the combination \( \sqrt{x(1-x)} b \). The transverse part of the wave function is assumed to be a simple Gaussian (see [8] for a discussion of this ansatz)

\[
\hat{\Sigma} = 4\pi \exp \left[-x(1-x) \frac{b^2}{4a^2}\right].
\] (2.6)

The transverse size parameter \( a \) is fixed for a given DA from the \( \pi^0 \to \gamma\gamma \) decay providing the relation [10]

\[
\int d\mathbf{x} d^2b \hat{\Psi}_0(x, b, \mu_0) = \sqrt{6}/f_\pi.
\] (2.7)

There is no other parameter to adjust. Utilizing two frequently used DAs, namely the asymptotic one

\[
\varphi_{AS}(x) = 6x(1-x)
\] (2.8)

and a form proposed by Chernyak and Zhitnitsky [2]

\[
\varphi_{CZ}(x) = 30x(1-x)(1-2x)^2,
\] (2.9)

Jakob et al. [7] evaluated the \( \pi\gamma \) form factor. The results are compared to CELLO [11] and CLEO [12] data in Fig. 1. The predictions obtained with the asymptotic wave function are in perfect agreement with the data whereas the CZ wave
function, i.e. the CZ DA multiplied by the Gaussian (2.6), leads to results in dramatic conflict with the data. In the perturbative approach the wave function is an universal, i.e. process-independent object. Hence, in analyses of other hard exclusive reactions the AS wave function (or eventually slightly modified versions of it) should be applied. The use of the CZ wave function, on the other hand, appears to be inconsistent in the light of the observations made in [7].

Applications of the modified perturbative approach to the π’s and the nucleon’s electromagnetic form factors [8, 13] reveals that the perturbative contributions, although self-consistently evaluated, are too small. In Fig. 2 the results for the magnetic form factor of the proton are shown [13]. The CZ-like DAs [14] lead at best, namely when the wave functions are normalized to unity, to perturbative results amounting to about 30% of the experimental values for \( Q^2 \geq 10 \text{GeV}^2 \).

3. Soft contributions

The above mentioned results on form factors clearly indicate the dominance of soft physics, i.e. of higher twist contributions, in the experimentally accessible region of momentum transfer. At this point I have to recall the following fact: The elastic form factors also receive contributions from the overlap of the initial and final state wave functions \( \Psi_0 \). In the case of the pion’s electromagnetic form factor, the overlap contribution reads in the transverse configuration space

\[
F_{\pi}^{\text{soft}}(Q^2) = \frac{1}{4\pi} \int dx d^2b \exp \left[i(1-x)b\cdot q_\perp\right] |\Psi_0(x, b)|^2,
\]

where \( Q^2 = q_\perp^2 \). As shown in [7, 8] the AS wave function provides an overlap contribution of the right magnitude to fill in the gap between the perturbative contribution and the and the experimental data. As required by the consistency of the entire picture, \( F_{\pi}^{\text{soft}} \) decreases faster with increasing \( Q \) than the perturbative contribution. The broad flat maximum of the overlap contribution in the few GeV region simulates the \( Q \)-dependence predicted by the dimensional counting rules. For the CZ wave function the overlap contribution exceeds the data significantly; the maximum value of \( Q^2 F_{\pi} \) amounts to about 2.1 GeV\(^2\). This result is to be considered as a serious failure of the CZ wave function. Soft contributions of the overlap type have also been discussed in [5, 16, 17] and observations have been made similar to those in [7, 8]. In the case of the \( \pi\gamma \) transition form factor the overlap contribution is expected to be very small due to a helicity mismatch.

In Ref. [18] the overlap contribution to the nucleon’s magnetic form factor is evaluated where again a Gaussian \( b \) dependence analogue to (2.6) is employed. A rather small overlap contribution is obtained with the asymptotic wave function

\[^2\text{Note that formally the perturbative contribution to elastic form factors represents the overlap of the large transverse momentum tails of the wave functions.}\]
\( \phi_{AS} = 120 x_1 x_2 x_3 \) for the proton and a strict zero for the neutron. The CZ-like wave functions given in [14], on the other hand, lead to very large overlap contributions exceeding the data by huge factors. Again, as in the case of the pion, the strongly end-point concentrated wave functions while providing large but theoretically inconsistent leading twist contributions, have to be rejected. In order to find a proper wave function for the nucleon the following strategy is adopted: Starting point is the expansion of the nucleon’s DA over the eigenfunctions of the evolution equation which are appropriate linear combinations of the Appell polynomials truncated at the same order as the CZ-like DAs [14]:

\[
\phi(x, \mu_F) = 120 x_1 x_2 x_3 [1 + \sum_{n=1}^{5} B_n(\mu_F) \tilde{\phi}_n(x)]. 
\tag{3.2}
\]

The expansion coefficients \( B_n \) and \( f_N \), playing the rôle of the wave function at the origin of the configuration space, as well as the transverse size parameter are fitted to the experimental data of the proton’s magnetic form factor, the decay width for \( \Psi \to p\bar{p} \) and the inclusive valence quark distribution functions. Thereby the form factor is evaluated from the overlap contribution whereas the \( \Psi \) decay is calculated within the modified perturbative approach. For good reasons the dominance of the perturbative contribution is to be expected for the \( \Psi \) decay. Indeed fair agreement with the experimental value for the decay width is found. The fitted parameters of that new nucleon wave function are compiled in Table 1 where, for comparison, also the parameters for the asymptotic and the COZ [19] wave functions are shown. The transverse size parameter has a value of 0.75 GeV\(^{-1}\) for the fitted wave function. In Fig. 3 the overlap contributions are compared to the data for the proton and the neutron magnetic form factors. Obviously the magnitudes of the overlap contributions are correct in the \( Q^2 \) range from about 8 to 20 GeV\(^2\). For smaller values of \( Q^2 \) similar contributions from higher Fock states are expected to be important (for the fitted wave function the probability of the valence Fock state amounts to 0.205). For \( Q^2 \) larger than about 20 GeV\(^2\) the fit is slightly below the data. Little modifications of the \( b \)-dependence (2.6) may cure that defect [3].

Table 1: \( f_N \) and the expansion coefficients \( B_i \), \( i = 1 - 5 \) for the AS, the COZ [19] and the fitted wave functions [18] at a factorization scale \( \mu_F \) of 1 GeV.

| DA    | \( f_N \) [GeV\(^2\)] | \( B_1 \)     | \( B_2 \)     | \( B_3 \)     | \( B_4 \)     | \( B_5 \)     |
|-------|-----------------|----------------|----------------|----------------|----------------|----------------|
| asympt. | 5.0000 \cdot 10^{-3} | 0.0000         | 0.0000         | 0.0000         | 0.0000         | 0.0000         |
| COZ    | 5.0000 \cdot 10^{-3} | 3.6750         | 1.4840         | 2.8980         | -6.6150        | 1.0260         |
| Fit    | 7.0671 \cdot 10^{-3} | -0.1081        | 0.5167         | -0.2879        | -3.8884        | -0.0371        |
4. Spin effects

As mentioned in the introduction a characteristic property of the HSA is the conservation of hadronic helicity. Yet the helicity sum rule (1.4) is violated at moderately large momentum transfer. For instance, the single spin asymmetry in proton-proton elastic scattering amounts to about 20% at a momentum transfer of $6\text{ GeV}^2$ \[21\]. Another example is the Pauli form factor of the proton which is measured to be very large \[22\]. Its $Q^2$ dependence is seen to be compatible with a higher twist contribution ($\sim 1/Q^6$), see Fig. 4. Charmonium decays into proton and antiproton also show evidence for violations of the helicity sum rule. Consider the spin zero states, the $\eta_c$ and the $\chi_0$. In the rest frame of the mesons the decay products, proton and antiproton, must have the same helicity in obvious contradiction to the helicity sum rule. Hence, at leading twist, the widths for these two decays are zero. For the other three reactions, $\Psi, \chi_1, \chi_2 \rightarrow p\bar{p}$, the leading twist analysis provides non-zero decay widths which, at least for the COZ DA \[13\] are in apparent agreement with the data leaving aside the theoretical difficulties with the strong end-point contributions \[3\] \[19\]. The leading twist pattern does, however, not match with the experimentally observed pattern of branching ratios ($BR(\eta_c \rightarrow p\bar{p}) = (1.2 \pm 0.4) \cdot 10^{-3}; BR(\chi_0 \rightarrow p\bar{p}) \leq 0.9 \cdot 10^{-3}$ \[23\]).

I would like to mention that the E605 collaboration is going to measure, hopefully precisely enough, the decay $\chi_0 \rightarrow p\bar{p}$ at FERMILAB. It will be very interesting to see whether the measured width is close to the present day upper bound \[23\] which would indicate a strong spin effect, or much smaller as helicity conservation demands.

How can we understand these large violations of hadronic helicity conservation in perturbation theory? The simplest idea to generate helicity flips is to take into account quark masses. According to Efremov and Teryaev \[25\] the relevant mass should be of the order of the hadron mass since helicity flips are of twist three type. Model builders take this result as justification for the use of constituent quark masses. Still this mechanism does not lead to sizeable spin effects. For example, only values of the order of $10^{-6}$ are obtained for the $\eta_c, \chi_0 \rightarrow p\bar{p}$ branching ratios \[21\].

Another and perhaps more appealing idea is to take into account non-zero orbital angular momentum in the wave functions. In this case the hadron spin fails to equal the sum of the quark spins; a basic presumption in the derivation of the helicity sum rule is not satisfied. For example the pion wave function

\[3\]

The $\alpha_S$ values used in these analyses are typically too small as compared with the expectations for a characteristic mass scale of the order of the charm quark mass and a typical value of 200 MeV for $\Lambda_{QCD}$. Since the decay widths are proportional to $\alpha_S^6$ a large factor of uncertainty is therefore hidden in these calculations. In the modified perturbative analysis \[18\] of the $\Psi$ decay the average value of $\alpha_S$ has the reasonable value of 0.41.
(including its spin part) may be written as

\[ \hat{\Psi}_0(x, b) \not{p} \gamma_5 + \hat{\Psi}_1(x, b)[\not{p}, \not{b}] \gamma_5. \] (4.1)

\( b \) represents one unit of orbital angular momentum.\(^4\) In combination with a \( b \) dependent hard scattering amplitude such a wave function can provide violations of helicity conservation within the modified perturbative approach. Recently, Gousset, Pire and Ralston \[27\] applied this idea to elastic scattering. They argue that the multiple scattering mechanism \[3, 4\] provides the \( b \) dependent hard scattering amplitude since the transverse distance between the various elementary scattering planes provide a characteristic direction. Hence, helicity non-conservation is obtained in the large \( Q^2 \) limit without flipping a quark helicity. However, the multiple scattering mechanism only provides violations of the helicity sum rule by two units ("two flip rule"); single spin asymmetries like that one in \( pp \) elastic scattering, remain unexplained. So far Gousset et al. applied their mechanism only to \( \pi \pi \rightarrow \pi \pi \) and \( \pi \pi \rightarrow \rho \rho \). As a consequence of the two flip rule the amplitude for \( \pi \pi \rightarrow \rho_0 \rho_+ \) is zero while that of \( \pi \pi \rightarrow \rho_+ \rho_- \) is non-zero. There is an objection against the mechanism proposed by Gousset et al: The multiple scattering contribution only amounts to a small fraction of the \( pp \rightarrow pp \) cross section \[28, 29\]. Hence, any particular property of it is perhaps not relevant to what we may see in data.

Last I want to discuss briefly the contribution of the Wuppertal group to that field. For baryons one may think of quark-quark correlations in the wave functions which also constitute higher twist effects. In a series of papers \[24, 30, 31, 32\] we have advocated that correlations of this type can effectively be described as quasi-elementary diquarks. That diquark model is a variant of the HSA: The valence Fock state of baryons consists of quark and diquark with a corresponding DA and the hard scattering amplitude is the amplitude for a subprocess involving quarks and diquarks. In order to take care of the composite nature of the diquarks and in order to guarantee that asymptotically the standard HSA emerges phenomenological vertex functions are introduced, their parameterizations bear resemblance to meson form factors. A systematic study of all exclusive photon-nucleon reactions has been performed \[30\]: Form factors in the space-like and time-like regions, real and virtual Compton scattering off protons, two-photon annihilations into proton-antiproton, photoproduction of mesons. A fair description of all the large momentum transfer data for these reactions has been achieved, utilizing in all cases the same baryon DAs as well as the same values for the few parameters specifying the diquarks. Due to the occurrence of vector diquarks the model provides non-zero helicity flip amplitudes and consequently violations of the helicity

\(^4\)Dorokhov \[17\] has constructed such a pion wave function from the helicity and flavour changing instanton force.
sum rule at finite $Q^2$. Therefore, also the process $\eta_c \to p\bar{p}$ can be investigated \cite{31}. The prediction for the Pauli form factor of the proton is shown in Fig. 4. Its fair agreement with the data as well as the reasonable value obtained for the $\eta_c \to p\bar{p}$ branching ratio ($0.38 \pm 0.15$) indicates the correct size of the spin effects generated by diquarks.

As a last example of our results I want to mention the electron asymmetry in the reaction $ep \to ep\gamma$:

$$A_L = \frac{\sigma(+) - \sigma(-)}{\sigma(+) + \sigma(+)}$$  \hspace{1cm} (4.2)

where $\pm$ indicates the helicity of the incoming electron. $A_L$ measures the imaginary part of the longitudinal ($\lambda_{\gamma^*} = 0$) -- transverse ($\lambda_{\gamma^*} = \pm 1$) interference. The longitudinal amplitudes for virtual Compton scattering $\gamma^*p \to \gamma p$ turn out to be small in the diquark model (hence $A_{V\gamma}^L$ is small). However, according to the model, $A_L$ is large in the region of strong Bethe-Heitler contamination (see Fig. 5). In that region, $A_L$ measures the relative phase (being of perturbative origin from on-shell going internal gluons, quarks and diquarks \cite{33}) between the Bethe-Heitler amplitudes and the virtual Compton ones. The magnitude of the effect shown in Fig. 5 is sensitive to details of the model and, therefore, should not be taken literally. Despite of this our result may be taken as an example of what may happen. The measurement of $A_L$, e. g. at CEBAF, will elucidate strikingly the underlying dynamics of the virtual Compton process.

5. Summary

There is compelling evidence for the smallness of the perturbative contributions to the pion’s and nucleon’s electromagnetic form factors. With a few exceptions the perturbative contributions to other hard exclusive reactions will also be too small in the experimentally accessible range of momentum transfer. Among these exceptions are the $\pi\gamma$ transition form factor and the decay $\Psi \to p\bar{p}$. For both these reactions fair agreement between data and perturbation theory is found. In general however, soft, higher twist contributions seem to dominate. For the pion’s and nucleon’s electromagnetic form factors the overlap contribution evaluated from plausible wave functions, seem to have the right magnitude to account for the data.

Spin effects while experimentally large in many cases, are difficult to explain in a theoretically satisfactory way. There are many attempts to be found in the literature but, so far, only the diquark model provides quantitative predictions which are in fair agreement with data. Despite of this unsettled and unsatisfactory situation it is important to persist in the attempt to understand the dynamical foundation for exclusive processes in QCD.
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Figure 1: The $\pi\gamma$ transition form factor vs. $Q^2$. The solid (dashed) line represents the prediction obtained with the modified HSA using the asymptotic (CZ) wave function. The dotted line represents the results of the standard HSA (for the asymptotic wave function). Data are taken from [11, 12].
Figure 2: The proton’s magnetic form factor vs. $Q^2$. Data are taken from [15] (filled(open) circles: $G_M (F_1)$). The strip of theoretical results is obtained from the DAs given in [14]. The wave functions are normalized to unity. The plot is taken from [13].

Figure 3: The overlap contributions to the proton’s and neutron’s magnetic form factors [18]. Data are taken from [15, 20].
Figure 4: The Pauli form factor of the proton scaled by $Q^6$. Data are taken from [22]. The solid line represents the result obtained from the diquark model [24].

Figure 5: The electron asymmetry in $ep \rightarrow ep\gamma$ as predicted by the diquark model [32]. $\phi$ denotes the angle between the hadronic and the leptonic scattering planes; $\theta$ is the scattering angle of the photon in the $p\gamma^*$ cms.