THE GDH SUM RULE FOR THE DEUTERON

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As first topic, the GDH sum rule is discussed in the context of a more general class of sum rules associated with the various contributions to the total photoabsorption cross section for target and beam polarization. Then I address the question of whether the GDH sum rule for the neutron can be determined from the one for the deuteron. It appears that this will not be possible in a simple manner. The spin response of the deuteron is calculated including contributions from the photodisintegration channel and from coherent and incoherent single pion production as well, and the GDH integral is evaluated up to a photon energy of 550 MeV. The photodisintegration channel converges fast enough and gives a large negative contribution, essentially from the $^1S_0$ state near threshold and its absolute size is about the same than the sum of proton and neutron GDH values. It is only partially cancelled by the single pion production contribution. But the incoherent channel has not reached convergence at 550 MeV.

1 Introduction

The Gerasimov-Drell-Hearn (GDH) sum rule connects the anomalous magnetic moment of a particle with the energy weighted integral - henceforth denoted by $I^{GDH}$ - from threshold up to infinity over the difference of the total cross sections for the absorption of circularly polarized photons on a target with spin parallel ($\sigma^P(k)$) and antiparallel ($\sigma^A(k)$) to the spin of the photon. In detail it reads for a particle of mass $M_t$, charge $eQ$, anomalous magnetic moment $\kappa$ and spin $I$

$$I^{GDH} = 4\pi^2\kappa^2 \frac{e^2}{M_t^2} = \int_0^\infty \frac{dk}{k} \left( \sigma^P(k) - \sigma^A(k) \right), \quad (1)$$

where the anomalous magnetic moment is defined by the total magnetic moment operator of the particle $\vec{M} = (Q + \kappa)\frac{\vec{S}}{M_t}$, where $\vec{S}$ denotes the spin operator of the target.

This sum rule gives a very interesting relation between a magnetic ground state property of a particle and an integral property of its whole excitation spectrum. It shows that the existence of a nonvanishing anomalous magnetic moment is directly tied to an internal dynamic structure of the particle. Fur-
thermore, it tells us, because the lhs of (1) is positive, that the integrated, 
energy-weighted total absorption of a circularly polarized photon on a particle 
with its spin parallel to the photon spin is larger than the one on a target 
with its spin antiparallel, if the anomalous magnetic moment does not vanish.

The GDH sum rule has first been derived by Gerasimov and, shortly 
afterwards, independently by Drell and Hearn and also, less well known, by 
Hosada and Yamamoto. The last authors have used current algebra relations 
while the others based the derivation on the low energy theorem for the 
Compton scattering amplitude of a particle and the assumption of an unsub-
tracted dispersion relation for the difference of the elastic forward scattering 
amplitudes for circularly polarized photons and a completely polarized target 
with spin parallel and antiparallel to the photon spin.

First, I will briefly discuss a general class of photoabsorption sum rules of 
which the GDH sum rule is a special case. Then I will address the question 
whether the GDH sum rule of the neutron can be determined from the GDH 
sum rule of the deuteron in the absence of free neutron targets. Subsequently, 
I will present results on an evaluation of the GDH sum rule for the deuteron 
by explicit integration of the GDH integral up to a photon energy of 550 MeV 
including the photodisintegration channel as well as coherent and incoherent 
single pion photoproduction channels. I will close with some conclusions and 
an outlook.

2 A general class of photoabsorption sum rules

The GDH sum rule belongs to a larger class of photoabsorption sum rules 
related to the various contributions to the total photoabsorption cross section 
for the general case of beam and target polarization.

\[
\sigma_{\text{tot}}(k, \rho^\gamma, \rho^t) = \frac{1}{2} \sum_{J=0}^{2I} P_J^t \left[ (1 + (-)^J) \sigma_{11}^J(k) \right. \\
+ (1 - (-)^J) P_J^\gamma \sigma_{11}^J(k) P_J(\cos \theta_t) \\
+ (1 + (-)^J) P_J^\gamma \sigma_{-11}^J(k) d_{20}^J(\theta_t) \cos(2\phi_t) \big],
\]

where \(P_J^\gamma\) and \(P_J^t\) denote the degree of linear and circular photon polarization, 
respectively. Furthermore, the target polarization parameters \(P_J^t\) with respect 

to an orientation direction, characterized by the angles \(\theta_t\) and \(\phi_t\), are defined 
by the target polarization density matrix

\[
\rho_{M M'}^t = \frac{(-)^{I-M'}}{I} \sum_{J, m} \hat{J} \left( \begin{array}{cc} I & J \\ M' & -M \\ \end{array} \right) P_J^t e^{im\phi_t} d_{m0}^J(\theta_t).
\]
The separate contributions \( \sigma^J_\lambda \) are related to the forward Compton scattering amplitude via the optical theorem

\[
\sigma^J_\lambda(k) = \frac{4\pi}{k} 3m T^J_\lambda(k), \tag{4}
\]

where \( T^J_\lambda \) is defined by

\[
T^J_{\lambda':\lambda M}(k) = (-)^{I-M} \sum_{J=0}^{2I} \mathcal{J} \left( \begin{array}{ccc} I & J & I \\ -M' & -\lambda' & M \end{array} \right) T^J_{\lambda':\lambda M}(k). \tag{5}
\]

It can be expressed in terms of generalized polarizabilities

\[
T^J_{\lambda':\lambda M}(k) = \mathcal{J} \sum_{L',L} (-)^{L'+L} \left( \begin{array}{ccc} L & L' & J \\ \lambda & -\lambda' & -\lambda \end{array} \right) P^J_{L'\lambda'\lambda}(k), \tag{6}
\]

with

\[
P^J_{L'\lambda'\lambda}(k) = \sum_{\nu',\nu=0,1} \lambda^{\nu'} \lambda^{\nu} \mathcal{P}_J(M^{\nu'}L', M^{\nu}L; k), \tag{7}
\]

where \( M^0 = E \) (electric) and \( M^1 = M \) (magnetic) multipole. The \( T^J_\lambda \) are also related to the expansion of the scattering amplitude in terms of a complete set of operators \( \tau^J[I] \) with \( J = 0, 1, \ldots, 2I \) in the ground state spin space with reduced matrix elements \( \langle I||\tau^J[I]|I\rangle = \mathcal{J} \)

\[
T^J_{\lambda':\lambda M}(k) = \sum_{J=0}^{2I} (-)^{I+\lambda} \langle IM'\rangle \langle \tau^J[J] \rangle_{\lambda',\lambda}(IM) \Omega^{[J]}_{\lambda',\lambda}(k). \tag{8}
\]

Comparison with (5) leads to the simple relation

\[
T^J_{\lambda':\lambda M}(k) = \Omega^{[J]}_{\lambda',\lambda}(k). \tag{9}
\]

Specifically one has

\[
\sigma^{11}_J = \mathcal{J} \sum_M (-)^{I-M} \left( \begin{array}{ccc} I & J & I \\ -M & 0 & M \end{array} \right) \sigma_{1M}, \tag{10}
\]

where \( \sigma_{1M} \) denotes the total cross section for the absorption of a photon with helicity \( \lambda = 1 \) by a target with definite spin projection \( M \) on the photon momentum. Corresponding expressions hold for \( \sigma^{-11}_J \) with respect to the absorption of linearly polarized photons. In detail one finds for \( J = 0, 1, 2 \)

\[
\sigma^{11}_0 = \frac{1}{I^2} \sum_M \sigma_{1M}, \tag{11}
\]
\[ \sigma_{11}^1 = \frac{\sqrt{3}}{I^2 \sqrt{I(I+1)}} \sum_M M \sigma_{1M}, \]  
\[ \sigma_{21}^1 = \frac{\sqrt{5}}{I^2 \sqrt{I(I+1)}} \sum_M (3M^2 - I(I+1)) \sqrt{(2I-1)(2I+3)} \sigma_{1M}, \]  
and for the spin asymmetry \( (\sigma_{P/A} = \sigma_{1,\pm 1}) \)
\[ \sigma^p - \sigma^A = i \sum_j \hat{J} (1 - (-)^J) \left( \begin{array}{cc} I & J \\ -I & 0 \end{array} \right) \sigma_{j1}^{11} \]
\[ = \frac{2\sqrt{3}}{\sqrt{I+1}} \sigma_{11}^{11} + \ldots \]  
(14)

Now, crossing symmetry implies
\[ (T^j_{\lambda \lambda}(-k))^* = (-)^J T^j_{\lambda \lambda}(k). \]  
(15)

For \( J = \) even one takes a once-subtracted dispersion relation for \( T^j_{\lambda \lambda}(k) \)
\[ \Re \left( T^j_{\lambda \lambda}(k) - T^j_{\lambda \lambda}(0) \right) = \frac{2k^2}{\pi} P \int_0^\infty \frac{dk'}{k'} \frac{3m T^j_{\lambda \lambda}(k')}{k'^2 - k^2} \]
\[ = \frac{k^2}{2\pi^2} P \int_0^\infty dk' \frac{\sigma^j_{\lambda \lambda}(k')}{k'^2 - k^2}, \]  
(16)

while for \( J = \) odd an unsubtracted dispersion relation applies
\[ \Re T^j_{\lambda \lambda}(k) = \frac{2k}{\pi} P \int_0^\infty dk' \frac{3m T^j_{\lambda \lambda}(k')}{k'^2 - k^2} \]
\[ = \frac{k}{2\pi^2} P \int_0^\infty dk' \frac{\sigma^j_{\lambda \lambda}(k')}{k'^2 - k^2}. \]  
(17)

A power series expansion according to (13)
\[ \Re T^j_{\lambda \lambda}(k) = \left\{ \begin{array}{ll} \sum_{\nu=0}^\infty t^j_{\lambda \lambda, J k^\nu} & \text{for } J \text{ even,} \\
\sum_{\nu=0}^\infty t^j_{\lambda \lambda, J k^{\nu+1}} & \text{for } J \text{ odd,} \end{array} \right. \]  
(18)
yields a class of sum rules
\[ t^j_{\lambda \lambda, J} = \left\{ \begin{array}{ll} \frac{1}{2\pi^2} \int_0^\infty dk' \frac{\sigma^j_{\lambda \lambda}(k')}{k'^2} & \text{for } J \text{ even and } \nu = 1, 2, \ldots, \\
\frac{1}{2\pi^2} \int_0^\infty dk' \frac{\sigma^j_{\lambda \lambda}(k')}{k'^2} & \text{for } J \text{ odd and } \nu = 0, 1, \ldots, \end{array} \right. \]  
(19)
one of which is the GDH, namely for $J = \text{odd}$ and $\nu = 0$. Because from the low-energy expansion of the Compton amplitude

$$T_{\lambda M, \lambda M}(k) = -e^2 \frac{Q^2}{M_t} + \lambda \kappa^2 \frac{e^2}{M_t^2} (S_z)_{1M} k + \mathcal{O}(k^2),$$

one finds specifically

$$T_{\lambda M}^J(k) = \begin{cases} -\delta_{\lambda M} \delta_{j0} \frac{e^2}{M_t} + \mathcal{O}(k^2) & \text{for } J \text{ even}, \\ k \left( \delta_{\lambda M} \delta_{j1} \lambda \kappa^2 \frac{e^2}{M_t^2} \frac{\sqrt{T(T+1)}}{\sqrt{3}} + \mathcal{O}(k^2) \right) & \text{for } J \text{ odd}. \end{cases}$$

The latter yields the GDH sum rule in the form

$$4 \pi^2 \frac{\kappa^2 e^2}{M_t^2} I = 2 \frac{\sqrt{3T}}{\sqrt{T+1}} \int_0^\infty \frac{dk'}{k'} \sigma_{11}^{j1}(k'),$$

from which (1) follows because of (14) and the fact that the higher order terms $\sigma_{11}^{j1}$ for $J > 1$ do not contribute to the left hand side.

3 Is it possible to get the GDH sum rule for the neutron from the one of the deuteron?

With respect to the GDH sum rule for the neutron, it has been suggested to measure in the absence of neutron targets its spin asymmetry using a polarized deuteron target. It would rest on two assumptions:

(i) A vector polarized deuteron constitutes effectively a polarized neutron target.

(ii) The contribution of the meson production to the spin asymmetry of the deuteron is dominated by the quasifree process, and one can neglect binding and final state interaction effects arising from the presence of the spectator nucleon, so that the deuteron spin asymmetry is an incoherent sum of proton and neutron contributions.

However, I would like to point out a few “caveats” which make it very unlikely that one can determine the spin asymmetry of the neutron in this way:

(i) First of all, the neutron is not completely polarized in a completely vector polarized deuteron target, i.e., $P(n) = 1 - 1.5 p_D$, and its polarization degree is slightly model dependent due to the appearance of the deuteron $D$-wave probability $p_D$ which is not observable.

(ii) A model calculation of the impulse approximation (IA) for the incoherent pion photoproduction on the deuteron by R. Schmidt et al [12] shows
already for the unpolarized cross section that the complete IA is not the incoherent sum of proton and neutron contributions. In addition, final state interaction and other two-body effects will very likely add further complications thus spoiling this simple idea.

(iii) Since polarization observables show in general a stronger sensitivity to small dynamical and coherence effects, the spin asymmetry might be even more sensitive to the mentioned disturbing effects.

(iv) The contribution from coherent \(\pi^0\) production on the deuteron is non-negligible which certainly is not an incoherent sum of \(\pi^0\) production on proton and neutron. Thus all these factors will prohibit a simple determination of the neutron spin asymmetry by subtracting from the meson production part of the deuteron’s spin asymmetry the proton one. That does not mean, that one does not learn anything about neutron properties. On the contrary, pion photoproduction on the deuteron will provide a very important test for the understanding of the production process on the neutron. But this can be achieved only in the context of a reliable theoretical model which takes into account all important two-body effects.

4 The GDH sum rule for the deuteron

In the case of the deuteron, one finds a very interesting cancellation of large contributions. The deuteron has isospin zero, excluding most of the contribution of the large nucleon anomalous magnetic moments to its magnetic moment, and thus one finds a very small anomalous magnetic moment, namely \(\kappa_d = -1.43\) resulting in a GDH prediction of \(I_d^{GDH} = 0.65\,\mu_b\), which is more than two orders of magnitude smaller than the nucleon values. Considering the possible absorption processes, one notes first that the incoherent pion production on the deuteron is dominated by the quasifree production on the nucleons bound in the deuteron. This gives a rough estimate for its contribution to the GDH value, namely the sum of the proton and neutron GDH values of \(438\,\mu_b\).

Another contribution arises from the coherent \(\pi^0\) production channel. On the other hand, for the additional photodisintegration channel which is the only photoabsorption process below the pion production threshold, one finds at very low energies near threshold a sizeable negative contribution which arises from the \(M1\)-transition to the virtual \(^1S_0\) state, because this state can only be reached if the spins of photon and deuteron are antiparallel, and is forbidden for the parallel situation as has been pointed out, for example in Ref.6.

We have evaluated explicitly the finite GDH sum rule for the deuteron by integrating up to a photon energy of 550 MeV. Three contributions have
Table 1. Various contributions of the photodisintegration channel to the GDH integral for the deuteron integrated up to 550 MeV in µb.

|          | N   | N+MEC | N+MEC+IC | N+MEC+IC+RC |
|----------|-----|-------|----------|-------------|
|          | -619| -689  | -665     | -413        |

been included: (i) the photodisintegration channel \( \gamma d \rightarrow np \), (ii) the coherent pion production \( \gamma d \rightarrow \pi^0 d \), and (iii) the incoherent pion production \( \gamma d \rightarrow \pi NN \). The upper integration limit of 550 MeV has been chosen because on the one hand one finds sufficient convergence for the photodisintegration channel, while on the other hand only single pion photoproduction has been considered, thus limiting the applicability of the present theoretical treatment to energies not too far above the two pion production threshold as long as significant contributions from multipion production cannot be expected. I will now discuss the three contributions separately and refer for details to Ref. [7].

The photodisintegration channel is evaluated within the nonrelativistic framework as is described in detail in Ref. [8] but with inclusion of the most important relativistic contributions. Explicitly, all electric and magnetic multipoles up to the order \( L = 4 \) are considered which means inclusion of the final state interaction in all partial waves up to \( j = 5 \). For the calculation of the initial deuteron and the final n-p scattering wave functions we use the realistic Bonn potential (r-space version) [9]. In the current operator we distinguish the one-body currents with Siegert operators (N), explicit meson exchange contributions (MEC) beyond the Siegert operators, essentially from \( \pi^- \) and \( \rho^- \)-exchange, contributions from isobar configurations of the wave functions (IC), calculated in the impulse approximation [10], and leading order relativistic contributions (RC) of which the spin-orbit current is by far the most dominant part.

The results are summarized in Fig. 1 where the spin asymmetry and the GDH integral is shown. The GDH values are listed in Tab. 1. One readily notes the huge negative contribution from the \( ^1S_0 \)-state at low energies (see the upper left panel of Fig. 1). Here, the effects from MEC are relatively strong, resulting in an enhancement of the negative value by about 15 percent. Isobar effects are significant in the region of the \( \Delta \)-resonance, as expected. They give a positive contribution, but considerably smaller in absolute size than MEC. The largest positive contribution stems from RC in the energy region up to about 100 MeV (see the upper right panel of Fig. 1) reducing the GDH value in absolute size by more than 30 percent. This strong influence from relativistic.
Figure 1. Contribution of the photodisintegration channel to the GDH sum rule for the deuteron. Two upper and lower left panels: difference of the cross sections in various energy regions; lower right panel: $I_{\gamma d\rightarrow np}^{GDH}$ as function of the upper integration energy. Dashed curves: N, dash-dot: N+MEC, dotted: N+MEC+IC, and full curves N+MEC+IC+RC.

effects is not surprising in view of the fact, that the correct form of the term linear in the photon momentum of the low energy expansion of (20) is only obtained if leading order relativistic contributions are included. The total sum rule value from the photodisintegration channel then is $I_{\gamma d\rightarrow np}^{GDH}(550\text{ MeV}) = -413\,\mu\text{b}$. Its absolute value almost equals within less than ten percent the sum of the free proton and neutron values. This may not be accidental since the large value is directly linked to the nucleon anomalous magnetic moment as is demonstrated by the fact that one finds indeed a very small but positive value $I_{\gamma d\rightarrow np}^{GDH}(550\text{ MeV}) = 7.3\,\mu\text{b}$ if the nucleon anomalous magnetic moment is switched off in the e.m. one-body current operator.
The theoretical model used to calculate the contribution of the coherent pion production channel is described in detail in Ref. [11]. The reaction is clearly dominated by the magnetic dipole excitation of the $\Delta$ resonance from which one obtains a strong positive $I_{\gamma d \rightarrow d\pi^0}^{GDH}$ contribution because the $\Delta$-excitation is favoured if photon and nucleon spins are parallel. The model takes into account pion rescattering by solving a system of coupled equations for the $N\Delta$, $NN\pi$ and $NN$ channels. The inclusion of the rescattering effects is important and leads in general to a significant reduction of the unpolarized cross section in reasonable agreement with the differential cross section data available in the $\Delta$ region. Fig. 2 shows the result of our calculation. One sees the strong positive contribution from the $\Delta$-excitation giving a value $I_{\gamma d \rightarrow d\pi^0}^{GDH}(550 \text{ MeV}) = 63 \mu b$. The comparison with the unpolarized cross section, also plotted in Fig. 2, demonstrates the dominance of $\sigma^P$ over $\sigma^A$. Furthermore, one notes quite satisfactory convergence.

The calculation of the $\gamma d \rightarrow \pi NN$ contributions to the GDH integral is based on the spectator nucleon approach discussed in Ref. [12]. In this framework, the reaction proceeds via pion production on one nucleon while the other nucleon acts merely as a spectator. Thus, the $\gamma d \rightarrow \pi NN$ operator is given as the sum of the elementary $\gamma N \rightarrow \pi N$ operators of the two nucleons. For this elementary operator, we have taken the standard pseudovector Born terms and the contribution of the $\Delta$ resonance, and a satisfactory description of pion photoproduction on the nucleon is achieved in the $\Delta$-resonance region [12]. Although the spectator model does not include any final state interaction except for the resonant $M_{1+}^{3/2}$ multipole, it gives quite a good description.
Figure 3. Contribution of the incoherent π production to the GDH sum rule for the deuteron and the nucleon. Upper part: difference of the cross sections; lower part: \( I_{\gamma d \rightarrow NN\pi}^{GDH} \) as function of the upper integration energy. Full curves for the deuteron, dotted curves for the nucleon. In the case of \( \pi^0 \) production, the dotted curve shows the summed proton and neutron contributions.

of available data on the total cross section demonstrating the dominance of the quasifree production process, for which the spectator model should work quite well. The results are collected in Fig. 3. The upper part shows the individual contributions from the different charge states of the pion and their total sum to the cross section difference for pion photoproduction on both the deuteron and for comparison on the nucleon. One notes qualitatively a similar behaviour although the maxima and minima are smaller and also slightly shifted towards higher energies for the deuteron. In the lower part of Fig. 3 the corresponding GDH integrals are shown. A large positive contribution comes from \( \pi^0 \)-production whereas the charged pions give a negative but - in absolute size - smaller contribution to the GDH value. Up to an energy of 550 MeV one finds for the total contribution of the incoherent pion production channels a value \( I_{\gamma d \rightarrow NN\pi}^{GDH}(550 \text{ MeV}) = 167 \, \mu \text{b} \) which is remarkably close to the sum of the neutron and proton values for the given elementary model \( I_{n}^{GDH}(550 \text{ MeV}) + I_{p}^{GDH}(550 \text{ MeV}) = 163 \, \mu \text{b} \). However, as is evident from Fig. 3 convergence is certainly not reached at this energy. Thus it is not clear whether this result is accidental, and in addition, one has to see what the effect of neglected rescattering and other two-body contributions will be. Furthermore, the elementary pion production operator used in this work had been constructed primarily to give a realistic description of the \( \Delta \) resonance
Table 2. Contributions of the different absorption channels to the GDH integral for the deuteron integrated up to 550 MeV in µb.

| Channel                  | Contribution (µb) |
|--------------------------|-------------------|
| $\gamma d \rightarrow np$ | $-413$            |
| $\gamma d \rightarrow d\pi^0$ | $63$              |
| $\gamma d \rightarrow np\pi^0$ | $288$            |
| $\gamma d \rightarrow np\pi^+$ | $-35$            |
| $\gamma d \rightarrow pp\pi^-$ | $-86$            |
| Total                    | $-183$            |

In fact, it underestimates the GDH integral up to 550 MeV by about a factor two compared to a corresponding evaluation based on a multipole analysis of experimental pion photoproduction data.

The contributions from all three channels and their sum are listed in Tab. 2. A very interesting and important result is the large negative contribution from the photodisintegration channel near and not too far above the break-up threshold with surprisingly large relativistic effects below 100 MeV. Hopefully, this low energy feature of the spin asymmetry can be checked experimentally in the near future. For the total GDH value from explicit integration up to 550 MeV, we find a negative value $I_{GDH}^d(550 \text{ MeV}) = -183 \mu b$. However, as I have mentioned above, some uncertainty lies in the contribution of the incoherent pion production channel because of shortcomings of the model of the elementary production amplitude above the $\Delta$ resonance. If one uses instead of the model value $I_{GDH}^d(550 \text{ MeV}) = 167 \mu b$ the sum of the GDH values of neutron and proton by integrating the cross section difference obtained from a multipole analysis of experimental data (fit SM95 from the VPI-analysis), giving $I_{GDH}^n(550 \text{ MeV}) + I_{GDH}^p(550 \text{ MeV}) = 331 \mu b$, one finds for the deuteron $I_{GDH}^d(550 \text{ MeV}) = -19 \mu b$, which I consider a more realistic estimate. Since this value is still negative, a positive contribution of about the same size should come from contributions at higher energies in order to fulfill the small GDH sum rule for the deuteron, provided that the sum rule is valid. These contributions should come from the incoherent single pion production above 550 MeV, because for this channel convergence had not been reached in contrast to the other two channels, and in addition, from multipion production.

5 Conclusions and Outlook

In order to summarize let me draw a few important conclusions: The spin asymmetry of the deuteron is a very interesting observable of its own value because of a strong anticorrelation of photodisintegration and pion production. It is also very sensitive to relativistic effects at quite low energies which have never been tested in detail in this observable.
It is very doubtful, if not impossible, that one can extract in a simple manner the neutron spin asymmetry from the spin asymmetry of the deuteron. However, the spin asymmetry of the deuteron will provide more detailed tests of our understanding of pion photoproduction in a nucleus, and in particular on the neutron.

Future theoretical and experimental studies should be devoted to the following topics:

(a) Theory: Improvement of the elementary pion photoproduction amplitude above the two pion threshold and inclusion of multiple pion production. For photodisintegration the inclusion of a retarded $\Delta N$ interaction, of higher nucleon resonances, and of relativistic contributions is needed. Furthermore, other sum rules, alluded to in Sect. 2, should be studied.

(b) Experiment: A careful measurement of the spin asymmetry for the proton over a larger energy range. With respect to the deuteron, a measurement of the spin asymmetry is needed, separately for the photodisintegration channel, in particular close to the break-up threshold, as well as for the meson production channel.

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