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GDH2000 Convenor's Report: Spin Polarizabilities

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Abstract

The subject of low energy polarized Compton scattering form the proton, which is characterized by phenomenological spin-polarizabilities, is introduced and connection is made to new theoretical and experimental developments which were reported to this meeting.

1 Introduction

It is a pleasure to report on what was a very stimulating session involving low energy aspects of the GDH sum rule, which can be characterized in terms of so-called "spin-polarizabilities." Earlier at this meeting we heard a quote from Jürgen Ahrens, the essence of which was that details of the GDH integrand were of more interest than the final number itself. I concur with his statement and in fact would add a corollary which asserts that the decomposition into intrinsic amplitudes is even more important than study of the cross section. I think that this will become clear in the discussion below.

One aspect of this low energy analysis was stressed by G. Krein[1], who pointed out that the component of the GDH integrand studied experimentally at MAMI and reported at this workshop involves the energy range 200 MeV$\leq \omega \leq$ 800 MeV and omits the region $\sim m_\pi \leq \omega \leq$ 200 MeV. Thus in order to perform the GDH integration between the inelastic threshold and 200 MeV requires some sort of reliable theoretical input such as provided by the MAID or SAID analysis of low energy pion photoproduction. In this regard, Krein emphasized that until recently the charge pion E_{0+} multipole, which dominates the low energy cross section, was underpredicted by SAID by about 15%, compared both to experiment and to the Kroll-Ruderman stricture[2]. To show the importance of having the correct behavior in this near threshold region, he noted that use of the SAID E_{0+} multipole leads to a 20μb ($\sim$ 10%) change in the GDH sum rule value and a 0.75×10^{-4} fm^4 ($\sim$ 100%) shift in the value of the forward spin polarizability, to be discussed below.

2 Real Compton Scattering

It is thus very important to have a proper multipole analysis of the low energy Compton amplitude in order to perform a correct GDH analysis. In the process one can also learn a good deal about nucleon structure. In order to see how this comes about, consider first *very* low energy Compton scattering—say $\omega <<$ 20 MeV—wherein the photon wavelength is much longer than the size of the nucleon. In this case, one is unable to resolve the structure of the target and is sensitive only to its overall charge–e–and mass–m. The interaction is described by the simple lowest order Hamiltonian

$$H = \frac{(\vec{p} - e\vec{A})^2}{2m} + e\phi \tag{1}$$

and the resultant Compton scattering amplitude has the canonical Thomson form

$$\text{Amp} = -\frac{e^2}{m}\hat{\epsilon}' \cdot \hat{\epsilon} \tag{2}$$

At higher energy—shorter wavelength—the internal structure becomes visible and one can describe the interaction in terms of an effective Hamiltonian having certain elementary properties—

i) quadratic in $\vec{A}$;

ii) gauge invariant;

iii) rotational scalar;

iv) P,T even, etc.

To the next order then the resultant form of the interaction is unique and must have the from

$$H_{eff} = -\frac{1}{2}4\pi\alpha_E\vec{E}^2 - \frac{1}{2}4\pi\beta_M\vec{H}^2 \tag{3}$$

where α_E, β_M are phenomenological constants having the dimensions of volume. The physical meaning of these constants can be seen from the definitions

$$\vec{p} = -\frac{\delta H_{eff}}{\delta\vec{E}} = 4\pi\alpha_E\vec{E}; \quad \vec{\mu} = -\frac{\delta H_{eff}}{\delta\vec{H}} = 4\pi\beta_M\vec{H} \tag{4}$$

where $\vec{p}$, $\vec{\mu}$ are the electric, magnetic dipole moments generated under the influence of external electric, magnetizing fields $\vec{E}$, $\vec{H}$. We recognize α_E, β_M as being the the electric, magnetic polarizabilities, which have obvious and intutive classical meanings[3]

i) in the presence of an external electric field the positive, negative components of the charge distribution move in opposite directions, resulting in an induced electric dipole moment;

ii) in the presence of an external magnetizing field intrinsic magnetic moments tend to align, creating a paramagnetic effect, while orbital moments generate a diamagnetic component, in accord with Lenz law.

Using the above effective Hamiltonians the Compton scattering amplitude becomes

$$\mathrm{Amp}^{(2)} = \hat{\epsilon}\cdot\hat{\epsilon}'\left(\frac{-e^2}{M} + \omega\omega'\,4\pi\alpha_E^p\right) + \hat{\epsilon}\times\vec{k}\cdot\hat{\epsilon}'\times\vec{k}'\,4\pi\beta_M^p + \mathcal{O}(\omega^4)\,. \quad (5)$$

and the resultant differential scattering cross section is

$$\begin{aligned}\frac{d\sigma}{d\Omega} &= \left(\frac{\alpha}{M}\right)^2\left(\frac{\omega'}{\omega}\right)^2\Bigg[\frac{1}{2}(1+\cos^2\theta) \\ &- \frac{M\omega\omega'}{\alpha}\left(\frac{1}{2}(\alpha_E^p+\beta_M^p)(1+\cos\theta)^2 + \frac{1}{2}(\alpha_E^p-\beta_M^p)(1-\cos\theta)^2\right) + \ldots\Bigg],\end{aligned} \quad (6)$$

where $\alpha = e^2/4\pi$ is the fine structure constant. It is clear then that α_E, β_M can be extracted via careful measurement of the differential cross section and previous experiements at SAL and MAMI have yielded the values[4][1]

$$\alpha_E^p = (12.1\pm0.8\pm0.5)\times10^{-4}\ \mathrm{fm}^3;\quad \beta_M^p = (2.1\mp0.8\mp0.5)\times10^{-4}\ \mathrm{fm}^3. \quad (7)$$

At this meeting Wissmann announced new values obtained from precise $p(\gamma,\gamma)p$ measurements using the TAPS and LARA spectrometers[5]

$$\alpha_E^p = (12.24\pm0.24\pm0.54)\times10^{-4}\ \mathrm{fm}^3;\quad \beta_M^p = (1.57\mp0.24\mp0.54)\times10^{-4}\ \mathrm{fm}^3. \quad (8)$$

These measured numbers can be compared to the corresponding quantities calculated in heavy baryon chiral perturbation theory (HBχpt) at $\mathcal{O}(p^3)$[6]

$$\alpha_E^p = 10K_p = 12.7\times10^{-4}\ \mathrm{fm}^3,\quad \beta_M^p = K_p = 1.3\times10^{-4}\ \mathrm{fm}^3 \quad (9)$$

where $K_p = \alpha g_A^2/192\pi F_\pi^2 m_\pi$. Here $g_A \simeq 1.266$ is the axial coupling constant in neutron beta decay and $F_\pi \simeq 92.4$ MeV is the pion decay constant. Of

[1]In order to put these numbers in perspective note that for a hydrogen atom one finds $\alpha_E(H) \sim \mathrm{Volume}(H)$ while for the proton Eq. 7 gives $\alpha_E(p) \sim 10^{-3}\mathrm{Volume}(p)$, so that the proton is a much more strongly bound system.

course, one must include higher order terms in order to properly judge the convergence behavior of the series, and such a calculation at $\mathcal{O}(p^4)$ has been performed by Bernard, Kaiser, Schmidt and Meißner (BKSM)[7]. At this order counterterms are required, which were estimated by BKSM by treating higher resonances—including $\Delta(1232)$—as very heavy with respect to the nucleon, yielding

$$\alpha_E^p = (10.5 \pm 2.0) \times 10^{-4} \text{ fm}^3; \quad \beta_M^p = (3.5 \pm 3.6) \times 10^{-4} \text{ fm}^3 \tag{10}$$

where the uncertainty is associated with the counterterm contribution from the $\Delta(1232)$ and from K, η loop effects. Agreement remains good between theory and experiment, supporting the view that the pion cloud gives a good description of such quantities.

The above results are well known and our task today is to extend this discussion to indlude spin degrees of freedom. In this case the general Compton amplitude can written in the general form

$$\begin{aligned} T &= A_1(\omega,z)\vec{\epsilon}' \cdot \vec{\epsilon} + A_2(\omega,z)\vec{\epsilon}' \cdot \hat{k}\, \vec{\epsilon} \cdot \hat{k}' \\ &+ iA_3(\omega,z)\vec{\sigma} \cdot (\vec{\epsilon}' \times \vec{\epsilon}) + iA_4(\omega,z)\vec{\sigma} \cdot (\hat{k}' \times \hat{k})\vec{\epsilon}' \cdot \vec{\epsilon} \\ &+ iA_5(\omega,z)\vec{\sigma} \cdot [(\vec{\epsilon}' \times \hat{k})\vec{\epsilon} \cdot \hat{k}' - (\vec{\epsilon} \times \hat{k}')\vec{\epsilon}' \cdot \hat{k}] \\ &+ iA_6(\omega,z)\vec{\sigma} \cdot [(\vec{\epsilon}' \times \hat{k}')\hat{\epsilon} \cdot \hat{k}' - (\vec{\epsilon} \times \hat{k})\vec{\epsilon}' \cdot \hat{k}], \end{aligned} \tag{11}$$

and each amplitude can be expanded in terms of a lowest order Born contribution plus a higher order and structure-dependent polarizability term. In the case of the spin-dependent amplitudes $A_{3,4,5,6}$ such structure effects arise at $\mathcal{O}(\omega^3)$ and can be characterized in terms of an effective Hamiltonian involving four "spin-polarizabilities"

$$H_{eff}^{(3)} = -\frac{1}{2}4\pi(\gamma_{E1}^p \vec{\sigma}\cdot\vec{E}\times\dot{\vec{E}} + \gamma_{M1}^p \vec{\sigma}\cdot\vec{H}\times\dot{\vec{H}} - 2\gamma_{E2}^p E_{ij}\sigma_i H_j + 2\gamma_{M2}^p H_{ij}\sigma_i E_j) \tag{12}$$

where

$$E_{ij} = \frac{1}{2}(\nabla_i E_j + \nabla_j E_i), \quad H_{ij} = \frac{1}{2}(\nabla_i H_j + \nabla_j H_i) \tag{13}$$

denote electric and magnetizing field gradients. While these quantities are mathematically well-defined via Eq. 12, I am unable to provide a good physical picture. The parameters γ_{E1}, γ_{M1} are similar to the classical Faraday rotation, wherein the linear polarization of the photon passing longitudinally through a magnetized medium exhibits a rotation due to the difference in index of refraction for photons with circular polarization parallel and antiparallel to the direction of magnetization. However, I don't know how to go much farther than this and will offer a bottle of fine Mainz Kupfenberg Sekt to anyone who is able to provide me such a classical picture.

On the theoretical side the chiral predictions for the spin-polarizabilities at $\mathcal{O}(p^3)$ are given by[8],[9]

$$\gamma^p_{E1} = -\frac{10K_p}{\pi m_\pi}, \quad \gamma^p_{M1} = -\frac{2K_p}{\pi m_\pi}, \quad \gamma^p_{E2} = \frac{2K_p}{\pi m_\pi}, \quad \gamma^p_{M2} = \frac{2K_p}{\pi m_\pi} \tag{14}$$

and extensions to $\mathcal{O}(p^4)$, necessary to assess the convergence, were presented to this workshop by Hemmert[10] and by McGovern[11]. While the numerical results of the two calculations are in agreement, the interpretation in terms of polarizabilities is presently in dispute. The problem arises from HBχpt diagrams wherein there is a nucleon pole on one side of which exists a pion loop renormalizing an electromagnetic vertex. In ordinary relativistic perturbation theory such diagrams are clearly one particle reducible and would be discarded as having nothing to do with polarizabilities. However, this is not so clear in the heavy baryon chiral calculation and this is where the problem lies at present. Ulf Meissner has promised a resolution in n-weeks but has not yet given a definitive value (or even a bound!) for the number n. In the meantime I shall quote the Jülich calculations

$$\begin{aligned}
\gamma^p_{E1} &= -\frac{10K_p}{\pi m_\pi}(1-\frac{29}{20}\frac{\pi m_\pi}{M}) \\
\gamma^p_{M1} &= -\frac{2K_p}{\pi m_\pi}(1-\frac{11}{4}\frac{\pi m_\pi}{M}) \\
\gamma^p_{E2} &= \frac{2K_p}{\pi m_\pi}(1-\frac{2\kappa_n+3}{4}\frac{\pi m_\pi}{M}) \\
\gamma^p_{M2} &= \frac{2K_p}{\pi m_\pi}(1-\frac{3}{4}\frac{\pi m_\pi}{M})
\end{aligned} \tag{15}$$

On the experimental side, there exist as yet no direct polarized Compton scattering measurements. However, a global analysis of unpolarized Compton data by the LEGS group has yielded the value[12][2]

$$\gamma_\pi = -\gamma_{E1} - \gamma_{M2} + \gamma_{E2} + \gamma_{M1} = (15.7 \pm 2.3 \pm 2.8 \pm 2.4) \times 10^{-4}\,\mathrm{fm}^4 \tag{16}$$

[2]Note here that we have subtracted the pion pole contribution.

in disagreement with the theoretical prediction

$$\gamma_\pi = \frac{2K_p}{\pi m_\pi}(4-(\kappa_n+\frac{9}{2})\frac{\pi m_\pi}{M}) = 3.3\times 10^{-4}\,\mathrm{fm}^4 \tag{17}$$

from Eq. 15. However, Wissman has announced a new value from the TAPS data

$$\gamma_\pi = (7.4\pm 2.3)\times 10^{-4}\,\mathrm{fm}^4 \tag{18}$$

which is in better agreement with theory. The other quantity about which much has been written is the forward spin polarizability γ_0, which is given by the first moment of the DGH sum rule

$$\gamma_0 = \gamma_{E1}+\gamma_{M2}+\gamma_{E2}+\gamma_{M1} = \int_{\omega_0}^\infty \frac{d\omega}{\omega^3}(\sigma_{\frac{3}{2}}(\omega)-\sigma_{\frac{1}{2}}(\omega)) \tag{19}$$

Drechsel at this meeting has quoted perhaps the best current value of the sum rule, based upon the MAID analysis,

$$\gamma_0 = -0.80\times 10^{-4}\,\mathrm{fm}^4 \tag{20}$$

which is in reasonable agreement with previous determinations.

While at present we do not have direct experimental values for the four spin-polarizabilities, Barbara Pasquini described a way by which they can be obtained using a dispersive analysis of the Compton process[13]. One assumes that the Compton amplitudes A_i can be represented in terms of once subtracted dispersion relations at fixed t

$$A_i(\nu,t) = A_i^{\mathrm{Born}}(\nu,t)+(A_i(0,t)-A_i^{\mathrm{Born}}(0,t))+\frac{2\nu^2}{\pi}P\int_{\nu_{thr}}^\infty \frac{\mathrm{Im}A_i(\nu',t)}{\nu'(\nu'^2-\nu^2)}. \tag{21}$$

Here $\mathrm{Im}A_i(\nu',t)$ is evaluated using empirical photoproduction data while the subtraction constant $A_i(0,t)-A_i^{\mathrm{Born}}(0,t)$ is represented via use of t-channel dispersion relations

$$A_i(0,t)-A_i^{\mathrm{Born}}(0,t) = a_i + a_i^{t-pole}+\frac{t}{\pi}\left(\int_{4m_\pi^2}^\infty - \int_{-\infty}^{-4Mm_\pi-2m_\pi^2} dt'\frac{\mathrm{Im}_t A_i(0,t')}{t'(t'-t)}\right) \tag{22}$$

with $\mathrm{Im}_t A_i$ evaluated using the contribution from the $\pi\pi$ intermediate state. In principle then there remain six unknown subtraction constants a_i to be

polarizability	HBχpt	Dispersive Evaluation
γ^p_{E1}	-1.8	-4.4
γ^p_{M1}	2.9	2.9
γ^p_{E2}	1.8	2.2
γ^p_{M2}	0.7	0.0

Table 1: Calculated and "experimental" values for spin polarizabilities obtained via dispersion relations. All are in units of 10^{-4} fm^4.

determined empirically. However, in view of the limitations posed by the the data, Drechsel et al. note that four of these quantities can be reasonably assumed to obey unsubtracted forward dispersion relations, while the remaining two—$\alpha_E - \beta_M$ and γ_π—can be treated as parameters and fitted from the data. Once this is done the other spin polarizabilities may be extracted using sum rules, as done above in the case of the forward spin polarizability. The results of this process are compared in Table 1 with predictions of chiral perturbation theory and one finds generally satisfactory agreement.

It has been noted by Babusci et al.[14] and by Holstein et al.[15] that one can extend this analysis to include terms of $\mathcal{O}(\omega^4)$ in the Compton amplitude by introducing higher order polarizabilities via

$$H^{(4)}_{eff} = -\frac{1}{2}4\pi\alpha^p_{E\nu}\dot{\vec{E}}^2 - \frac{1}{2}4\pi\beta^p_{M\nu}\dot{\vec{H}}^2 - \frac{1}{12}4\pi\alpha^p_{E2}E^2_{ij} - \frac{1}{12}4\pi\beta^p_{M2}H^2_{ij} \tag{23}$$

Likewise Holstein et al. have extended this to $\mathcal{O}(\omega^5)$ by defining higher order spin-polarizabilities—

$$\begin{aligned} H^{(5)}_{eff} &= -\frac{1}{2}4\pi\Big[\gamma^p_{E1\nu}\vec{\sigma}\cdot\dot{\vec{E}}\times\ddot{\vec{E}} + \gamma^p_{M1\nu}\vec{\sigma}\cdot\dot{\vec{H}}\times\ddot{\vec{H}} - 2\gamma^p_{E2\nu}\sigma_i\dot{E}_{ij}\dot{H}_j + 2\gamma^p_{M2\nu}\sigma_i\dot{H}_{ij}\dot{E}_j \\ &+ 4\gamma^p_{ET}\epsilon_{ijk}\sigma_iE_{j\ell}\dot{E}_{k\ell} + 4\gamma^p_{MT}\epsilon_{ijk}\sigma_iH_{j\ell}\dot{H}_{k\ell} - 6\gamma^p_{E3}\sigma_iE_{ijk}H_{jk} + 6\gamma^p_{M3}\sigma_iH_{ijk}E_{jk}\Big] \end{aligned} \tag{24}$$

where

$$\begin{aligned} (E,H)_{ijk} &= \frac{1}{3}(\nabla_i\nabla_j(E,H)_k + \nabla_i\nabla_k(E,H)_j + \nabla_j\nabla_k(E,H)_i) \\ &- \frac{1}{15}(\delta_{ij}\nabla^2(E,H)_k + \delta_{jk}\nabla^2(E,H)_i + \delta_{ik}\nabla^2(E,H)_j) \end{aligned} \tag{25}$$

polarizability	HBχpt	Dispersive value
$\alpha^p_{E\nu}$	2.4	-3.8
$\beta^p_{M\nu}$	7.5	9.3
α^p_{E2}	22.1	29.3
β^p_{M2}	-9.5	-24.3
$\gamma^p_{E1\nu}$	-2.4	-3.4
$\gamma^p_{M1\nu}$	1.8	2.2
$\gamma^p_{E2\nu}$	1.6	1.3
$\gamma^p_{M2\nu}$	-0.1	-0.6

Table 2: Calculated and "experimental" values for higher order polarizabilities obtained via dispersion relations. Spin independent and spin dependent polarizabilities are in units of 10^{-4} fm^5 and 10^{-4} fm^6 respectively

are the (spherical) tensor gradients of the electric and magnetizing fields. Each of these new higher order polarizabilities can be extracted via sum rules from the Mainz dispersive analysis and results are compared with chiral predictions in Table 2. Agreement is obviously quite satisfactory, except for $\alpha_{E\nu}$. Despite the success of this program it would be highly desirable to measure such quantities directly and Wissmann has suggested a $\vec{p}(\vec{\gamma},\gamma)p$ program by which it might be possible to achieve this.

3 Virtual Compton Scattering

Nicole d'Hose discussed the virtual Compton scattering (VCS) process by which one can measure "generalized" (q-dependent) polarizabilities[16]. In order to understand the meaning of such quantities, recall that in ordinary electron scattering measurement of the q-dependent charge form factor allows access, via Fourier transform, to the nucleon charge *density.* In an analogous fashion measurement of a generalized polarizability such as $\alpha_E(q)$ permits one to determine the polarization density of the nucleon. On the experimental side this is an extremently challenging process because the generalized polarizabilities can be determined only after (large) Bethe-Heitler and Born

	$P_{LL} - P_{TT}/\epsilon$	P_{LT}
expt.	$23.7 \pm 2.2 \pm 0.6 \pm 4.3$	$-5.0 \pm 0.8 \pm 1.1 \pm 1.4$
HBχpt[18]	26.0	-5.3
LσM[19]	11.5	0.0
ELM[20]	5.9	-1.9
NRQM[21]	11.1	-3.5

Table 3: Measured and calculated values for generalized polarizabilities. All are in units of GeV^{-2}.

diagram contributions have been subtracted.[3] One then seeks a systematic deviation growing with ω' of the measured cross section from that predicted with only Bethe-Heitler plus Born input in order to extract the desired signal. This has been achieved in a recent MAMI experiment[17] and what results is information on the two combinations

$$\begin{aligned} P_{LL} - \frac{1}{\epsilon}P_{TT} &= a_0\alpha_E(q) - c_1\gamma_{M2}(q) + c_2 M^{M1-M1}(q) \\ P_{LT} &= b_0\beta_M(q) + c_3 M^{C0-M1}(q) - c_4\gamma_{E2}(q) \end{aligned} \tag{26}$$

where here the multipoles M^{M1-M1}, M^{C0-M1} have no RCS analogs. The extracted numbers are given in Table 3 together with values calculated in various models as well as in HBχpt. It is remarkable that once again agreement with the simple chiral calculation is outstanding, despite the fact that the measurement took place at q=0.6 GeV, where one should question the validity of the chiral approach.

On the theoretical side, Marc Vanderhaeghen reported on an extension of the Mainz dispersive RCS analysis to the case of VCS[22]. This extension is not as straightforward as it might appear, since replacement of a real photon by a virtual one requires now *twelve* invariant amplitudes which are functions of three variables, which may be taken as ω, θ, Q^2. Analysis of the asymptotic dependence is correspondingly much more complex. The calculation is still in progress but Vanderhaeghen presented preliminary results for four of the generalized polarizabilities compard to the chiral predictions.

[3] Radiative corrections are also substantial here.

Agreement is satisfactory though not outstanding, but Hemmert argued that $\mathcal{O}(p^4)$ corrections will improve matters.

4 Conclusion

The subject of polarizabilities—both spin-dependent as well as spin-independent—in both RCS and VCS is undergoing a major emphasis due to the presence of high quality data from machines such as MAMI, Bates, SAL, CEBAF, etc. and we have heard lots of exciting results at this meeting. I believe that at GDH2002 we can anticipate equally exciting developments. For RCS:

i) theory—definitive $\mathcal{O}(p^4)$ calculation of all polarizabilities in HBχpt will be completed;

ii) experiment—analysis of the MAMI results will be completed and preliminary $\vec{p}(\vec{\gamma},\gamma)p$ data will be presented.

For VCS:

i) theory—existing HBχpt calculations will be extended to $\mathcal{O}(p^4)$;

ii) experiment—analysis of the MAMI experiment will be complete and results from JLab and Bates experiments will be presented. Preliminary $\vec{p}(\vec{\gamma}^*,\gamma)p$ data will be available.

These are exciting times for low energy Compton scattering and I look forward to our meeting two years hence in Italy.

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