I summarize the theoretical predictions for the spin–dependent nucleon polarizabilities based on chiral effective field theory approaches.

1 Introduction

Low–energy Compton scattering off the nucleon is an important probe to unravel the nonperturbative structure of QCD since the electromagnetic interactions in the initial and final state are well understood. In the long wavelength limit, only the charge of the target can be detected, which is nothing but the celebrated Thomson low energy theorem. At higher photon energies, \(50 < \omega < 100\) MeV, the internal structure of the system slowly becomes visible. These nucleon structure–dependent effects in unpolarized Compton scattering are taken into account by introducing two free parameters into the cross-section formula, commonly denoted the electric (\(\bar{\alpha}\)) and magnetic (\(\bar{\beta}\)) polarizabilities of the nucleon in analogy to the structure dependent response functions for light–matter interactions in classical electrodynamics. Over the past few decades several experiments on low energy Compton scattering off the proton have taken place, resulting in several extractions of the electromagnetic polarizabilities of the proton. At present, the commonly accepted numbers are \(\bar{\alpha}_p = (12.1 \pm 0.8 \pm 0.5) \times 10^{-4} \text{ fm}^3\), \(\bar{\beta}_p = (2.1 \pm 0.8 \pm 0.5) \times 10^{-4} \text{ fm}^3\), indicating that the proton compared to its volume of \(\sim 0.5 \text{ fm}^3\) is a rather stiff object. At present, several quite different theoretical approaches find qualitative and quantitative explanations for these two polarizabilities, but they also constitute one of the striking successes of chiral perturbation theory. The lowest order predictions stem from finite loop graphs and can be expressed in terms of well–known parameters

\[
\bar{\alpha}_p = \bar{\alpha}_n = 10\bar{\beta}_p = 10\bar{\beta}_n = \frac{5e^2g_A^2}{384\pi^2F_{\pi}^2M_\pi} = 12.4 \cdot 10^{-4} \text{ fm}^3, \quad (1)
\]

with \(g_A = 1.26\) the axial–vector coupling measured in neutron \(\beta\)–decay, \(F_\pi = 92.4\) MeV the pion decay constant, \(M_\pi = 139.57\) MeV the pion mass and \(e^2/4\pi = 1/137.036\) the fine structure constant. These numbers are in
striking agreement with the experimental values, demonstrating the importance of the pion cloud to the structure of the nucleon. In ref. 3 it was further demonstrated that the fourth order corrections to these results are modest and in particular, a novel large and negative pion loop contribution to the magnetic polarizability was found, leading to a substantial reduction of the sizeable positive $\Delta$ contribution.

Quite recently, with the advent of polarized targets and new sources with a high flux of polarized photons, the case of polarized Compton scattering off the proton $\gamma p \rightarrow \gamma p$ has come close to experimental feasibility. On the theoretical side it has been shown that one can define four spin-dependent electromagnetic response functions $\gamma_i, \ i = 1 \ldots 4$, which in analogy to $\bar{\alpha}, \bar{\beta}$ are commonly called the “spin-polarizabilities” of the proton. I remark that one can also use a more physical basis in terms of electric and magnetic dipole, quadrupole, … excitations, as discussed by Hemmert in these proceedings. 4

First studies have been published claiming that the such parameterized information on the low-energy spin structure of the proton can really be extracted from the upcoming double-polarization Compton experiments. A success of this program would clearly shed new light on our understanding of the internal dynamics of the proton and at the same time serve as a check on the theoretical explanations of the polarizabilities. The new challenge to theorists will then be to explain all six of the leading electromagnetic response functions simultaneously. At present there exist two experimental analysis that have shed some light on the magnitude of the (essentially) unknown spin-polarizabilities $\gamma_i^{(p)}$ of the proton. First, the LEGS group has reported a result for a linear combination involving three of the $\gamma_i$ (for the proton), namely

$$\gamma_1^{(p)} = \gamma_1^{(p)} + \gamma_2^{(p)} + 2\gamma_4^{(p)},$$

whose magnitude could not be explained by theoretical approaches. We note that this pioneering result was obtained from an analysis of an unpolarized Compton experiment in the backward direction, where the spin-polarizabilities come in as one contribution in a whole class of subleading order nucleon structure effects in the differential cross-section. However, the Göttingen group has performed a similar measurement at MAMI and their value for $\gamma_2^{(p)}$ comes out consistent with theoretical expectations. 5

While such indirect determinations of some linear combinations of the spin-polarizabilities are very valuable, we can only reemphasize the need for the

\[a\] I remark that as explained below, one should subtract the pion pole (anomaly) contribution from the true nucleon-structure effects. This is often not done.
upcoming polarized Compton scattering experiments to disentangle the contributions from the four spin–polarizabilities. It is therefore a challenge on the theory side to work out the complete one–loop (that is fourth order) representation of the spin–polarizabilities within the context of Heavy Baryon Chiral Perturbation Theory (HBCHPT), extending previous efforts in a significant way. These results have been reported in and a similar study can be found in. The active degrees of freedom in HBCHPT are the asymptotically observable pion and nucleon fields. The various contributions from tree and loop diagrams are organized according to power counting rules, i.e. one expands in small momenta and pion masses ($M_\pi$), collectively denoted by $p$. Previously an order $O(p^3)$ SU(2) HBCHPT calculation was performed, which showed that the leading (i.e. long–range) structure effects in the spin–polarizabilities are given by 8 different $\pi N$ loop diagrams giving rise to a $1/M_\pi^2$ behavior in the $\gamma_i$ (which is known since long, see ). Subsequently, it was shown in a third order SU(2) calculation in which the first nucleon resonance, the $\Delta(1232)$, was included as an explicit degree of freedom that two ($\gamma_2$, $\gamma_4$) of the four spin–polarizabilities receive large corrections due to $\Delta(1232)$ related effects, resulting in a big correction to the leading $1/M_\pi^2$ behavior. In that phenomenological extension of HBCHPT, one also counts the nucleon-delta mass splitting as an additional small parameter and collectively denotes all small parameters as $\epsilon$. The corresponding expansion, which also has a consistent power counting, is called the “small scale expansion” (SSE) (because it differs from a chiral expansion due to the non–vanishing of the $N\Delta$ mass splitting in the chiral limit). Another important conclusion of was that any HBCHPT calculation that wants to calculate $\gamma_2$ and $\gamma_4$ would have to be extended to $O(p^5)$ before it can incorporate the large $\Delta(1232)$ related corrections found already at $O(\epsilon^3)$ in. Recently, two $O(p^4)$ SU(2) HBCHPT calculations of polarized Compton scattering in the forward direction appeared, from which one can extract one particular linear combination of three of the four $\gamma_i$, usually called $\gamma_0$:

$$\gamma_0 = \gamma_1 - (\gamma_2 + 2\gamma_4) \cos \theta |_{\theta \to 0}. \quad (3)$$

The authors of claimed to have found a huge correction to $\gamma_0$ at $O(p^4)$ relative to the $O(p^3)$ result already found in, casting doubt on the useful-

\footnote{It is important to note that the terms of order order $\epsilon^3$ in the SSE that are proportional to the delta–nucleon mass splitting $m_\Delta - m_N$ appear at order $p^4$ in the chiral expansion. $\gamma_0$ can also be calculated from the absorption cross sections of polarized photons on polarized nucleons via the GGT sum rule as pointed out in. In the absence of such data several groups have tried to extract the required cross sections via a partial wave analysis of unpolarized absorption cross sections. Recent results of these efforts are given in table 2.}
ness/convergence of HBCHPT for spin-polarizabilities. Given that $\gamma_0$ involves the very two polarizabilities $\gamma_2, \gamma_4$, the (known) poor convergence for $\gamma_0$ found in [12, 13] should not have come as a surprise. We will come back to this point later.

2 Invariant amplitudes and polarizabilities

We first want to comment on the extraction of polarizabilities from nucleon Compton scattering amplitudes. In previous analyses [10, 11] it has always been stated that in order to obtain the spin-polarizabilities from the calculated Compton amplitudes, one only has to subtract off the nucleon tree-level (Born) graphs from the fully calculated amplitudes. The remainder in each (spin-amplitude) then started with a factor of $\omega^3$ and the associated Taylor-coefficient was related to the spin-polarizabilities. Due to the (relatively) simple structure of the spin-amplitudes at this order, this prescription gives the correct result in the $O(p^3)$ HBCHPT [13] and the $O(\varepsilon^3)$ SSE [14] calculations. However, at $O(p^4)$ (and also at $O(\varepsilon^4)$) one has to resort to a definition of the (spin-) polarizabilities that is soundly based on field theory, in order to make sure that one only picks up those contributions at $\omega^3$ that are really connected with (spin-) polarizabilities. In fact, at $O(p^4)$ ($O(\varepsilon^4)$) the prescription given in [10, 11] leads to an admixture of effects resulting from two successive, uncorrelated $\gamma NN$ interactions with a one nucleon intermediate state. In order to avoid these problems we advocate the following definition for the spin-dependent polarizabilities in (chiral) effective field theories: Given a complete set of spin-structure amplitudes for Compton scattering to a certain order in perturbation theory, one first removes all one-particle (i.e. one-nucleon or one-pion) reducible (1PR) contributions from the full spin-structure amplitudes. To be more precise, at order $O(p^4)$ one removes $F(\omega)/\omega$ terms from the amplitude, where $F(\omega)$ denotes the energy dependence of the $\gamma NN$ vertex function. This prescription has been challenged, see the contribution of McGovern to these proceedings [20]. I will come back to this later. Specifically, starting from the general form of the T-matrix for real Compton scattering assuming invariance under parity, charge conjugation and time reversal symmetry, we utilize the following six structure amplitudes $A_i(\omega, \theta)$ in the Coulomb gauge, $\epsilon_0 = \epsilon'_0 = 0$,

$$
T = A_1(\omega, \theta)\varepsilon^{*'} \cdot \hat{\epsilon} + A_2(\omega, \theta)\varepsilon^{*'} \cdot \hat{k} \hat{k}' \cdot \hat{\epsilon'} \\
+ iA_3(\omega, \theta)\hat{\sigma} \cdot (\varepsilon^{*'} \times \hat{\epsilon}) \\
+ iA_4(\omega, \theta)\hat{\sigma} \cdot (\hat{k}' \times \hat{k})\varepsilon^{*'} \cdot \hat{\epsilon}
$$
\[ + iA_5(\omega, \theta) \vec{\sigma} \cdot [(\vec{e}^{*'} \times \hat{k})\vec{e} \cdot \hat{k}' - (\vec{e} \times \hat{k}')\vec{e}^{*'} \cdot \hat{k}] \\
+ iA_6(\omega, \theta) \vec{\sigma} \cdot [(\vec{e}^{*'} \times \hat{k}')\vec{e} \cdot \hat{k} - (\vec{e} \times \hat{k})\vec{e}^{*'} \cdot \hat{k}], \]

where \( \theta \) corresponds to the c.m. scattering angle, \( \vec{e}, \hat{k}, (\vec{e}', \hat{k}') \) denote the polarization vector, direction of the incident (final) photon while \( \vec{\sigma} \) represents the (spin) polarization vector of the nucleon. Each (spin-)structure amplitude is now separated into 1PR contributions and a remainder, that contains the response of the nucleon’s excitation structure to two photons:

\[ A_i(\omega, \theta) = A_i(\omega, \theta)^{1PR} + A_i(\omega, \theta)^{exc}, \quad i = 3, \ldots, 6. \]

Taylor-expanding the spin-dependent \( A_i(\omega, \theta)^{1PR} \) for the case of a proton target in the c.m. frame into a power series in \( \omega \), the leading terms are linear in \( \omega \) and are given by the venerable low–energy theorems (LETs) of Low, Gell-Mann and Goldberger:

\[ A_3(\omega, \theta)^{1PR} = \frac{1 + 2\kappa(p)}{2M_N^2} \omega^2 e^2 + \mathcal{O}(\omega^2), \]
\[ A_4(\omega, \theta)^{1PR} = -\frac{(1 + \kappa(p))e^2}{2M_N^2} \omega + \mathcal{O}(\omega^2), \]
\[ A_5(\omega, \theta)^{1PR} = \frac{(1 + \kappa(p))^2e^2}{2M_N^2} \omega + \mathcal{O}(\omega^2), \]
\[ A_6(\omega, \theta)^{1PR} = -\frac{(1 + \kappa(p))e^2}{2M_N^2} \omega + \mathcal{O}(\omega^2). \]

While it is not advisable to really perform this Taylor-expansion for the spin-dependent \( A_i(\omega, \theta)^{1PR} \) due to the complex pole structure, one can do so without problems for the \( A_i(\omega, \theta)^{exc} \) as long as \( \omega \ll M_\pi \). For the case of a proton one then finds

\[ A_3(\omega, \theta)^{exc} = 4\pi \left[ \gamma_1^{(p)} - (\gamma_2^{(p)} + 2\gamma_4^{(p)}) \cos \theta \right] \omega^3 + \mathcal{O}(\omega^4), \]
\[ A_4(\omega, \theta)^{exc} = 4\pi\gamma_2^{(p)} \omega^3 + \mathcal{O}(\omega^4), \]
\[ A_5(\omega, \theta)^{exc} = 4\pi\gamma_4^{(p)} \omega^3 + \mathcal{O}(\omega^4), \]
\[ A_6(\omega, \theta)^{exc} = 4\pi\gamma_3^{(p)} \omega^3 + \mathcal{O}(\omega^4). \]

We therefore take Eq.(7) as starting point for the calculation of the spin-polarizabilities, which are related to the \( \omega^3 \) Taylor-coefficients of \( A_i(\omega, \theta)^{exc} \).

As noted above, both the \( \mathcal{O}(p^3) \) HBCHPT \[5\] and the \( \mathcal{O}(e^3) \) SSE \[4\] results are consistent with this definition.
3 Isoscalar polarizabilities

Utilizing Eqs. (5,7) we have calculated the first subleading correction, \( O(p^4) \), to the four isoscalar spin-polarizabilities \( \gamma_i^{(s)} \) already determined to \( O(p^3) \) in SU(2) HBCHPT. We employ here the convention

\[
\gamma_i^{(p)} = \gamma_i^{(s)} + \gamma_i^{(v)}; \quad \gamma_i^{(n)} = \gamma_i^{(s)} - \gamma_i^{(v)} .
\] (8)

Contrary to popular opinion we show, that even at subleading order all four spin-polarizabilities can be given in closed form expressions which are free of any unknown chiral counterterms! The only parameters appearing in the results are the axial-vector nucleon coupling constant, the pion decay constant \( F_\pi \), the pion mass \( M_\pi \), the mass of the nucleon \( m_N \) as well as its isoscalar, \( \kappa^{(s)} = -0.12 \), and isovector, \( \kappa^{(v)} = 3.7 \), anomalous magnetic moments. All \( O(p^4) \) corrections arise from 25 one-loop \( \pi N \) continuum diagrams, with the relevant vertices obtained from the well-known SU(2) HBCHPT \( O(p) \) and \( O(p^2) \) Lagrangians given in detail in ref [11]. To \( O(p^4) \) we find

\[
\begin{align*}
\gamma_1^{(s)} &= \frac{e^2 g_A^2}{96\pi^3 F_\pi^2 M_\pi^2} \left[ 1 - \mu M_\pi \right] , \\
\gamma_2^{(s)} &= \frac{e^2 g_A^2}{192\pi^3 F_\pi^2 M_\pi^2} \left[ 1 + \mu \left( -6 + \kappa^{(v)} \right) \pi \right] , \\
\gamma_3^{(s)} &= \frac{e^2 g_A^2}{384\pi^3 F_\pi^2 M_\pi^2} \left[ 1 - \mu \pi \right] , \\
\gamma_4^{(s)} &= -\frac{e^2 g_A^2}{384\pi^3 F_\pi^2 M_\pi^2} \left[ 1 - \mu \frac{11}{4} \pi \right] ,
\end{align*}
\] (9-12)

with \( \mu = M_\pi / m_N \simeq 1/7 \) and the the numerical values given in table 1. The leading \( 1/M_\pi^2 \) behavior of the isoscalar spin-polarizabilities is not touched by the \( O(p^4) \) correction, as expected. With the notable exception of \( \gamma_4^{(s)} \), which even changes its sign due to a large \( O(p^4) \) correction, we show that this first subleading order of \( \gamma_1^{(s)}, \gamma_2^{(s)}, \gamma_3^{(s)} \) amounts to a 25-45% correction to the leading order result. This does not quite correspond to the expected \( M_\pi / m_N \) correction of (naive) dimensional analysis, but can be considered acceptable. The large correction in \( \gamma_4^{(s)} \) should be considered accidental. It is not related to the the large \( \Delta \) effects found in the SSE calculation of [10] because these will only show up at \( O(p^5) \) in the HBCHPT framework. This can easily be understood: When the delta is not an active dof in the effective field theory, it can not modify the leading singularity of order \( 1/M_\pi^2 \) due to decoupling. In principle, the spin–dependent local operator of dimension two \( \sim c_4 \) inserted in
Table 1. Predictions for the spin-polarizabilities in HBCHPT in comparison with the dispersion analyses of refs. (Mainz1, Mainz2, BGLMN) and the $O(\epsilon^3)$ results of the small scale expansion (SSE1). All results are given in the units of $10^{-4}$ fm$^4$.

| $\gamma_{(s)}^i$ | $O(p^3)$ | $O(p^4)$ | Sum | Mainz1 | Mainz2 | BGLMN | SSE1 |
|------------------|----------|----------|------|--------|--------|--------|------|
| $\gamma_{1(s)}$ | $+4.6$   | $-2.1$   | $+2.5$ | $+5.6$ | $+5.7$ | $+4.7$ | $+4.4$ |
| $\gamma_{2(s)}$ | $+2.3$   | $-0.6$   | $+1.7$ | $-1.0$ | $-0.7$ | $-0.9$ | $-0.4$ |
| $\gamma_{3(s)}$ | $+1.1$   | $-0.5$   | $+0.6$ | $-0.6$ | $-0.5$ | $-0.2$ | $+1.0$ |
| $\gamma_{4(s)}$ | $-1.1$   | $+1.5$   | $+0.4$ | $+3.4$ | $+3.4$ | $+3.3$ | $+1.4$ |
| $\gamma_{1(v)}$ | $-1.3$   | $-1.3$   | $-1.3$ | $-0.5$ | $-1.3$ | $-1.6$ | $-1$  |
| $\gamma_{2(v)}$ | $-0.2$   | $-0.2$   | $-0.2$ | $-0.2$ | $+0.0$ | $+0.1$ | $-$   |
| $\gamma_{3(v)}$ | $+0.1$   | $+0.1$   | $+0.1$ | $-0.0$ | $+0.5$ | $+0.5$ | $-$   |
| $\gamma_{4(v)}$ | $+0.0$   | $+0.0$   | $+0.0$ | $+0.0$ | $-0.5$ | $-0.6$ | $-$   |

One loop graphs could modify the $1/M_\pi$ terms. However, such contributions cancel in the sum of all loop graphs. This is different from the spin-independent case, where a particular combination of certain $c_i$ generates an important pion loop correction at fourth order.

4 Isovector polarizabilities

We further report the first results for the four isovector spin-polarizabilities $\gamma_{(v)}^i$ obtained in the framework of chiral effective field theories. Previous calculations at $O(p^3)$ [11] and $O(\epsilon^3)$ [10] were only sensitive to the isoscalar spin-polarizabilities $\gamma_{(s)}^i$, therefore this calculation gives the first indication from a chiral effective field theory about the magnitude of the difference in the low-energy spin structure between proton and neutron. As in the case of the isoscalar spin-polarizabilities there are again no unknown counterterm contributions to this order in the $\gamma_{(v)}^i$. All $O(p^4)$ contributions arise from 16 one-loop $\pi N$ continuum diagrams with the relevant $O(p)$, $O(p^2)$ vertices again obtained from the Lagrangians given in ref. 11. To $O(p^4)$ one finds

$$\gamma_1^{(v)} = \frac{e^2 g_A^2}{96 \pi^3 F_\pi^2 M_\pi^2} \left[ 0 - \mu \frac{5\pi}{8} \right],$$

$$\gamma_2^{(v)} = \frac{e^2 g_A^2}{192 \pi^3 F_\pi^2 M_\pi^2} \left[ 0 - \mu \frac{(1 + \kappa^{(s)})\pi}{4} \right].$$

gdh2000: submitted to World Scientific on December 12, 2018
\[ \gamma_3^{(v)} = \frac{e^2 g_A^2}{384 \pi^3 F \pi^2 M^2} \left[ 0 + \mu \frac{\pi}{4} \right], \quad (15) \]

\[ \gamma_4^{(v)} = 0, \quad (16) \]

with the numerical values again given in table 1. The result of our investigation is that the size of the \(\gamma_i^{(v)}\) really tends to be an order of magnitude smaller than the one of the \(\gamma_i^{(s)}\) (with the possible exception of \(\gamma_1^{(v)}\)), supporting the scaling expectation, \(\gamma_i^{(v)} \sim (M_\pi/m_N)^{\gamma_i^{(s)}}\) from (naive) dimensional analysis. This is reminiscent of the situation in the spin-independent electromagnetic polarizabilities \(\bar{\alpha}^{(v)}\), \(\bar{\beta}^{(v)}\) which are also suppressed by one chiral power relative to their isoscalar partners \(\bar{\alpha}^{(s)}, \bar{\beta}^{(s)}\).

5 Forward and backward polarizabilities

Finally, we want to comment on the comparison between our results and existing calculations using dispersion analyses. Given our comments on the convergence of the chiral expansion for the (isoscalar) spin-polarizabilities\(\text{[1]}\), we reiterate that we do not believe our \(O(p^4)\) HBCHPT result for \(\gamma_2^{(s)}, \gamma_4^{(s)}\) to be meaningful. Their large inherent \(\Delta(1232)\) related contribution just cannot be included (via a counterterm) before \(O(p^5)\) in HBCHPT that only deals with pion and nucleon degrees of freedom. In table 1 it is therefore interesting to note that by adding ("by hand") the delta-pole contribution of \(\sim -2.5 \cdot 10^{-4} \text{fm}^4\) found in \(\text{[1]}\) to \(\gamma_2^{(s)}\) one could get quite close to the range for this spin-polarizability as suggested by the dispersion analyses\(\text{[22, 27, 6]}\). Similarly, adding \(\sim +2.5 \cdot 10^{-4} \text{fm}^4\) to \(\gamma_4^{(s)}\) as suggested by \(\text{[1]}\) also leads quite close to the range advocated by the dispersion results\(\text{[22, 27, 6]}\). However, such a procedure is of course not legitimate in an effective field theory, but it raises the hope that an extension of the \(O(\epsilon^3)\) SSE calculation of \(\text{[1]}\) that includes explicit delta degrees of freedom could lead to a much better behaved perturbative expansion for the isoscalar spin-polarizabilities. Whether this expectation holds true will be known quite soon\(\text{[1]}\). For the isovector spin-polarizabilities we have given the first predictions available from effective field theory. In general the agreement with the range advocated by the dispersion analyses is quite good.

Furthermore, in table 2 we give a comparison of our results for those linear combinations of the \(\gamma_i\) that typically are the main focus of attention in the literature. However, we re-emphasize that we do not consider our \(O(p^4)\) HBCHPT predictions for \(\gamma_0^{(s)}, \gamma_2^{(s)}\) to be meaningful, because they involve...
Table 2. Predictions for the so-called forward (backward) spin-polarizabilities $\gamma_0$ ($\gamma_\pi$). For a definition of units and references see table 1.

| $\gamma_i^{(N)}$ | $O(p^4)$ | $O(p^4)$ | Sum | Mainz1 | Mainz2 | BGLMN | SSE1 |
|------------------|----------|----------|-----|--------|--------|--------|------|
| $\gamma_0^{(s)}$ | +4.6     | -4.5     | +0.1| -0.2   | -0.4   | -1.0   | +2.0 |
| $\gamma_0^{(v)}$ | -        | -1.1     | -1.1| -0.3   | -0.4   | -0.5   | -    |
| $\gamma_\pi^{(s)}$ | +4.6     | +0.3     | +4.9| +11.4  | +11.8  | +10.4  | +6.8 |
| $\gamma_\pi^{(v)}$ | -        | -1.5     | -1.5| -0.7   | -2.4   | -2.7   | -    |

$\gamma_2^{(s)}, \gamma_4^{(s)}$. The corresponding isovector combinations, however, again seem to agree quite well with the dispersive results and so far we have no reason to suspect that they might be affected by the poor convergence behavior of some of their isoscalar counterparts. We further note that our $O(p^4)$ HBCHPT predictions for $\gamma_0^{(s,v)}$ differ from the ones given in two recent calculations. As noted above this difference solely arises from a different definition of nucleon spin-polarizabilities. If we (“by hand”) Taylor-expand our $\gamma^{NN}$ vertex functions $F(\omega)$ in powers of $\omega$ and include the resulting terms into the the $\gamma_0$ structure, we obtain the $O(p^4)$ corrections $\gamma_0^{(s)} = -6.9, \gamma_0^{(v)} = -1.6$ in units of $10^{-4}$fm$^4$, in numerical (and analytical) agreement with. This brings us to an important point: Once the first polarized Compton asymmetries have been measured, it has to be checked very carefully whether the same input data fitted to the terms we define as 1PR plus the additional free $\gamma_i$ parameters leads to the same numerical fit-results for the spin-polarizabilities as in the dispersion theoretical codes usually employed to extract polarizabilities from Compton data. Small differences for example in the treatment of the pion/nucleon pole could lead to quite large systematic errors in the determination of the $\gamma_i$. Such studies are under way. Also, the controversy about how to subtract the 1PR pieces in a non–relativistic approach like HBCHPT can be settled employing a Lorentz–invariant formulation of baryon CHPT since in that context the subtraction is unambiguous and the heavy fermion limit can be obtained easily.

Acknowledgments

These results have been obtained in collaboration with George Gellas and Thomas Hemmert, to whom I express my sincere gratitude. I also would like to thank the organizers for putting together such an interesting program.
References

1. F.J. Federspiel et al., Phys. Rev. Lett. 67 (1991) 1511; E.L. Hallin et al., Phys. Rev. C48 (1993) 1497; A. Zieger et al., Phys. Lett. B278 (1992) 34; B.E. MacGibbon et al., Phys. Rev. C52 (1995) 2097.
2. V. Bernard, N. Kaiser and Ulf-G. Meißner, Phys. Rev. Lett. 67 (1991) 1515.
3. V. Bernard, N. Kaiser, A. Schmidt and Ulf-G. Meißner, Phys. Lett. B319 (1993) 269; Z. Phys. A348 (1994) 317.
4. See e.g. S. Ragusa, Phys. Rev. D47, 3757 (1993); D49 (1994) 3157.
5. T.R. Hemmert, “Theory of the nucleon spin–polarizabilities II”, these proceedings.
6. D. Babusci et al., Phys. Rev. C58 (1998) 1013.
7. D. Drechsel, M. Gorchtein, B. Pasquini and M. Vanderhaeghen, Phys. Rev. C61 (2000) 015204.
8. J.I. Tonnison et al., Phys. Rev. Lett. 80 (1998) 4382.
9. F. Wissmann et al., Nucl. Phys. A660 (1999) 232.
10. T.R. Hemmert, B.R. Holstein, J. Kambor and G. Knöchlein, Phys. Rev. D57 (1998) 5746.
11. V. Bernard, N. Kaiser and Ulf-G. Meißen, Int. J. Mod. Phys. E4 (1995) 193.
12. X. Ji, C.-W. Kao and J. Osborne, Phys. Rev. D61 (2000) 074003.
13. K.B. Vijaya Kumar, J.A. McGovern and M.C. Birse, hep-ph/9909442.
14. G.C. Gellas, T.R. Hemmert and Ulf-G. Meißen, Phys. Rev. Lett. 85 (2000) 14.
15. K. B. Vijaya Kumar, J.A. McGovern and M.C. Birse, Phys. Lett. B479 (2000) 167.
16. V. Bernard, N. Kaiser, J. Kambor and Ulf-G. Meißen, Nucl. Phys. B388 (1992) 315.
17. T.R. Hemmert, B.R. Holstein and J. Kambor, Phys. Lett. B395 (1997) 89; J. Phys. G24 (1998) 1831.
18. M. Gell-Mann, M.L. Goldberger and W.E. Thirring, Phys. Rev. 95 (1954) 1612.
19. G.C. Gellas, T.R. Hemmert and Ulf-G. Meißen, forthcoming.
20. J.A. McGovern, M.C. Birse and K.B. Vijaya Kumar, nucl-th/0007015.
21. M. Gell-Mann and M.L. Goldberger, Phys. Rev. 96, (1954) 1433; F.E. Low, Phys. Rev. 96 (1954).
22. D. Drechsel, G. Krein and O. Hanstein, Phys. Lett. B420 (1998) 248.
23. T. Becher and H. Leutwyler, Eur. Phys. J. C9 (1999) 643.