NUCLEON COMPTON SCATTERING IN CHIRAL EFFECTIVE FIELD THEORIES

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Chiral effective field theories have a long history studying the process of Compton scattering on the nucleon. In this contribution I want to focus on the new developments that have occurred since the last Chiral Dynamics conference in Mainz 1997. Moreover, in view of the limited time, I will focus exclusively on the spin-dependent sector, where most of the recent work has been done.

1 Spin Polarizabilities in Real Compton Scattering

As introduced in the previous talk by Barry Holstein, the leading spin-structure dependent response of a nucleon in the presence of external electromagnetic fields can be parameterized via 4 dipole spin-polarizabilities $\gamma_i$, defined in complete analogy to the more familiar spin-independent electromagnetic (dipole) polarizabilities $\bar{\alpha}_E, \bar{\beta}_M$.

$$
\begin{align*}
\gamma_{E1} : & (E1 \rightarrow E1) \quad \gamma_{M1} : (M1 \rightarrow M1) \\
\gamma_{E2} : & (M1 \rightarrow E2) \quad \gamma_{M2} : (E1 \rightarrow M2).
\end{align*}
$$

The physics behind these spin-polarizabilities thus involves excitation of the spin 1/2 nucleon target via an electric/magnetic dipole transition and a successive de-excitation back into a spin 1/2 nucleon final state via an electric/magnetic dipole or quadrupole transition. As discussed by Holstein, none of these (dipole) spin-polarizabilities has been measured directly up to now—results of future double-polarized Compton scattering experiments on the proton $\vec{\gamma} \vec{p} \rightarrow \vec{\gamma}' \vec{p}'$ at MAMI, BNL-LEGS and TUNL, which would suppress the dominant spin-independent physics by measuring Compton asymmetries, are eagerly awaited. In chiral perturbation theory the dipole spin-polarizabilities have now been calculated to next-to-leading order (NLO), the results and a comparison with recent dispersion theoretical calculations are

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*a Some information is known about the particular linear combination $\gamma_e = -\gamma_{M2} - \gamma_{E1} + \gamma_{E2} + \gamma_{M1}$ from a dispersion analysis of unpolarized Compton scattering experiments.

For recent efforts to extract the so-called “forward spin polarizability” $\gamma_0 = -\gamma_{M2} - \gamma_{E1} + \gamma_{E2} + \gamma_{M1}$ from measurements of double-polarized pion photoproduction in the resonance region at MAMI/ELSA and the resulting model dependence see ref. [7].
given in Table 1. The NLO heavy baryon chpt predictions for the essentially unknown dipole spin polarizabilities of the nucleon can be obtained in closed form expressions without any free fit parameters. In general the qualitative agreement between the NLO chiral predictions and the dispersive analyses is quite good, with the marked exception of the isoscalar spin polarizability $\gamma_{M2}$, which consistently shows a different sign between the 2 approaches. One also notes that the isovector spin polarizabilities (which only start at $O(p^4)$ in the chiral expansion) tend to be much smaller than their corresponding isoscalar counterparts, leading to the non-trivial prediction that the low energy spin structure of proton and neutron is quite similar. In column A I have extracted the $O(p^4)$ dipole spin polarizabilities from the full Compton matrix element by only subtracting the Born terms as advocated in refs. 9, 11, whereas in column B I have subtracted all one-particle reducible graphs from the full Compton matrix element as advocated in ref. 10. The difference between these two approaches lies in one particular one-loop diagram, which can be seen to primarily affect the large magnetic spin-polarizability $\gamma_{M1}$, both in the isoscalar and in the isovector channel. Note that in a $O(p^4)$ heavy baryon chpt calculation of the isoscalar $\gamma_{M1}$ spin polarizability the large contribution from an intermediate $\Delta(1232)$ state is still missing, as this effect can only be included at $O(p^5)$ in a full 2-loop calculation. Adding this contribution “by hand” as indicated with the asterisk in Table 1 shows that procedure B then also leads to a consistent picture for $\gamma_{M1}$, whereas procedure A requires an unnaturally large higher order cancellation of yet unknown origin (if the values of the dispersion analyses are to be taken as a serious benchmark). We note that procedure B has been criticised in the past few months. In my opinion the only serious criticism towards procedure B lies in the fact that the one-to-one correspondence between a dispersion theoretical ansatz, which only relies on analyticity, crossing-symmetry etc., and the diagrammatic microscopic chpt calculation gets lost, because dispersion theory cannot distinguish between one-particle reducible vs. one-particle irreducible contributions, this clearly being a microscopic concept having nothing to do with the global properties of dispersion theory. Further theoretical investigations are clearly needed to identify the physics behind the cancellation mechanism needed for procedure A. Let me close this discussion by restating the trivial fact that both procedure A and procedure B lead to the absolutely identical Compton matrix element—the differences are only pertaining to a different separation into a polarizability dependent part and a remainder.

Finally I want to note that a brief summary on recent work regarding the extension of the nucleon’s spin-dependent response to multipolarities beyond the dipole truncation can be found in ref. 13.
Table 1. Predictions for the isoscalar and isovector dipole spin-polarizabilities of the nucleon found via 2 different prescriptions to $\mathcal{O}(\alpha^2)$ (i.e. NLO) in heavy baryon chpt, to $\mathcal{O}(\alpha)$ (i.e. LO) in the small scale expansion (SSELO) and in dispersion analyses (DKH,DGPV,BGLM). All results are given in the units $10^{-4}$ fm$^4$. ($\ast$: $+2.5 \times 10^{-4}$ fm$^4$ from $\Delta(1232)$ pole still missing).

| $\gamma^{(s)}_{1}$ | chpt$^4_{NLO}$ | chpt$^B_{NLO}$ | DKH | DGPV | BGLM | SSELO |
|-------------------|----------------|----------------|------|-------|-------|-------|
| $\gamma^{(v)}_{1}$ | $-2.8$         | $-3.0$         | $-5.0$ | $-5.2$ | $-4.5$ | $-5.2$ |
| $\gamma^{(v)}_{M1}$ | $+2.8^{*}$    | $+0.4^{*}$    | $+3.4$ | $+3.4$ | $+3.3$ | $+1.4$ |
| $\gamma^{(s)}_{E1}$ | $+2.0$         | $+2.0$         | $+2.4$ | $+2.7$ | $+2.4$ | $+1.0$ |
| $\gamma^{(s)}_{M1}$ | $+0.3$         | $+0.6$         | $-0.6$ | $-0.5$ | $-0.2$ | $+1.0$ |
| $\gamma^{(v)}_{E2}$ | $+1.4$         | $+1.2$         | $+0.5$ | $+0.8$ | $+1.1$ | -     |
| $\gamma^{(v)}_{M1}$ | $+0.5$         | $+0.0$         | $+0.0$ | $-0.5$ | $-0.6$ | -     |
| $\gamma^{(s)}_{E2}$ | $-0.2$         | $-0.2$         | $-0.2$ | $-0.5$ | $-0.5$ | -     |
| $\gamma^{(s)}_{M2}$ | $-0.1$         | $+0.1$         | $-0.0$ | $+0.5$ | $+0.5$ | -     |

2 Generalized Spin Polarizabilities

As discussed in Hyde-Wright’s plenary talk [4], the pioneering Virtual Compton Scattering (VCS) experiment on the proton $e\,p \to e'\,p'\,\gamma$ at MAMI is now analyzed [5]. From the viewpoint of the low energy structure of the nucleon the only difference to a real Compton scattering experiment is the fact that the incoming photon is virtual, $Q^2 \neq 0$. (Of course one also has to add the Bethe-Heitler contribution coherently to obtain measurable quantities as discussed by Hyde-Wright, but here we want to focus on the “proper” VCS contribution alone.) The Mainz experiment was performed in a special kinematic regime where one restricts oneself to small energies $\omega'$ of the outgoing real photon at fixed three-momentum transfer $|\vec{q}| = 600$ MeV stemming from the virtual incoming photon. From a theorist point of view this means that in the experiment one wants to be sure, that the de-excitation of the target back to a spin 1/2 proton via the real photon in the final state can be described with an electromagnetic dipole transition. This is the kinematic condition that Guichon, Liu and Thomas used for their ground-breaking definition of generalized polarizabilities (GPs) [6] and up to now this is the only theoretical framework we have in order to analyze/discuss low energy VCS experiments.

It was shown that in this “Guichon limit” (i.e. the dipole truncation for the final state radiation) there is a total of 6 generalized (i.e. momentum-transfer dependent) polarizabilities $P_{(Y1,X1)}(\vec{q})$, where $Y1$ ($X1$) denotes...
the multipolarity of the final state (initial state) radiation. In the long wavelength limit 2 of these GPs can be identified with the familiar (spin-independent) electric/magnetic polarizabilities:

\[
\bar{\alpha}_E = -\frac{e^2}{4\pi}\sqrt{\frac{3}{2}} \lim_{\vec{q} \to 0} P_{(E1,E1)}(\vec{q}),
\]

\[
\bar{\beta}_M = -\frac{e^2}{4\pi}\sqrt{\frac{3}{8}} \lim_{\vec{q} \to 0} P_{(M1,M1)}(\vec{q}).
\] (2)

(Based on this identification one can also define the momentum-transfer dependent generalized dipole polarizabilities \(\bar{\alpha}_E(\vec{q}), \bar{\beta}_M(\vec{q})\).) Here we want to focus on the remaining 4 GPs of the “Guichon set” 

\[P_{(C1,M2)}(\vec{q}), P_{(M1,C2)}(\vec{q}), P_{(M1,C0)}(\vec{q}), \hat{P}_{(C1,(C1,E1))}(\vec{q})\]

which can be shown to be spin-dependent polarizabilities. In the long-wavelength limit one can establish a connection to two of the dipole spin-polarizabilities of polarized real Compton scattering introduced in the previous section:

\[
\gamma_{M2} = -\frac{e^2}{4\pi}\frac{3}{\sqrt{2}} \lim_{\vec{q} \to 0} P_{(C1,M2)}(\vec{q})
\]

\[
\gamma_{E2} = -\frac{e^2}{4\pi}\frac{3\sqrt{3}}{2\sqrt{2}} \lim_{\vec{q} \to 0} P_{(M1,C2)}(\vec{q})
\] (3)

The remaining 2 generalized spin polarizabilities \(P_{(M1,C0)}(\vec{q}), \hat{P}_{(C1,(C1,E1))}(\vec{q})\) involve longitudinal multipole excitations \(C0, C1\) due to the incoming virtual photon and therefore do not have an analogue in real Compton scattering—they correspond to new low energy nucleon structure terms which can only be accessed via VCS!

The “Guichon set” of 6 generalized polarizabilities thus has a very intuitive explanation in terms of multipole excitation/de-excitation, similar to the multipole basis of the real Compton polarizabilities \(\gamma_i\) introduced in the previous section. Given a theoretical prediction for these 6 GPs, one can calculate the experimentally accessible response functions \(P_{LL}(\vec{q}), P_{LT}(\vec{q}), P_{TT}(\vec{q})\) and judge how well the theoretical understanding of the microscopic dynamics behind the (generalized) polarizabilities matches with the real world. For the Mainz experiment the results are given in Table 2. One can see that the theoretical values of the response functions—which use the SU(2) \(O(p^3)\) heavy

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Table 2. Experimental values of the response functions measured at MAMI at $|\vec{q}| = 600$ MeV compared with the leading order (i.e. $\mathcal{O}(p^3)$) heavy baryon chpt predictions. For comparison the corresponding values from real Compton scattering which can be obtained in the long-wavelength limit $|\vec{q}| \to 0$ are given as well.

| Response | Expt. [GeV$^{-2}$] | Chpt [GeV$^{-2}$] |
|----------|------------------|------------------|
| $[P_{LL} - \frac{1}{2}P_{TT}]|_{|q|=0}$ | $81.0 \pm 5.4 \pm 3.3$ | $83.5$ |
| $[P_{LT}]|_{|q|=0}$ | $-7.0 \mp 2.7 \mp 1.7$ | $-4.2$ |
| $[P_{LL} - \frac{1}{2}P_{TT}]|_{|q|=600\text{MeV}}$ | $23.7 \pm 2.2 \pm 0.6 \pm 4.3$ | $26.0$ |
| $[P_{LT}]|_{|q|=600\text{MeV}}$ | $-5.0 \pm 0.8 \pm 1.1 \pm 1.4$ | $-5.3$ |

baryon chpt calculation of the GPs from 1997 predating the experiment—are in a good agreement with the experimental numbers, even though the momentum-transfer in the Mainz experiment is rather large for such a leading order calculation! Keeping in mind that the Mainz experiment only determines one point in the $\vec{q}$-evolution of the response functions and that certainly further tests (especially at smaller values of momentum-transfer) are needed before one can establish a definite picture of the momentum-dependence of the underlying GPs, I want to discuss the physics behind the $\vec{q}$-evolution of the generalized polarizabilities from the point of view of chiral effective field theory.

For a $\mathcal{O}(p^3)$ Heavy Baryon calculation of the “proper” VCS matrix-element—which provides the leading order result for the GPs—one has to evaluate the same 9 one-loop $\pi N$-diagrams shown in Figure 1 as for leading order real Compton scattering with the initial photon now being virtual. To this (leading) order no new diagrams/counterterms appear beyond the ones also present in real Compton scattering. It is precisely the simple physics contained in the diagrams of Figure 1 evaluated in ref. 17 which completely determines the $\vec{q}$-dependence of the GPs and results in the theoretical numbers displayed in Table 2 when put into the formulae of the response functions. The remarkably simple picture behind the physics of the GPs that emerges suggests that it is the pion-cloud of the nucleon—which can be so easily excited because the chiral symmetry of the underlying QCD-lagrangian is spontaneously broken at low energies—that governs the $\vec{q}$-evolution at small momentum transfer, leading to a typical scaling behavior of

$$P_{(Y_{1,X})}(\vec{q}) \sim \frac{\vec{q}^2}{m_\pi^2}$$  \hspace{1cm} (4)$$

The relevant mass scale thus is the light mass of the (quasi-) Goldstone bo-
son which in the low momentum regime makes the effects of the pion-cloud dominant over other $\vec{q}$-dependent effects like the excitation/de-excitation of nucleon-resonances via transition form factors with a typical mass scale of $M_B \sim 1$ GeV. It is obvious that such a regime should exist for the VCS process due to the (naive) scaling arguments given above, however, the positive surprise seems to be that this regime where the chiral physics dominates over the usual baryon (and vector-meson) resonance physics seems to extend over a larger range of $|\vec{q}|$ than expected. By the time of the next Chiral Dynamics workshop we should have a better understanding of how far in $|\vec{q}|$ the leading order chiral dynamics given in Figure 1 provides a sufficient description and where the baryon/vector-meson resonance physics starts taking over. On the experimental side this fall the second VCS experiment on the proton will start taking data at Bates 18 at much lower $\vec{q}$, which should provide a strong constraint on any theoretical description of the GPs.

The final development in VCS which I want to discuss brings me back to the spin-sector. Although the pioneering VCS experiment at MAMI was an unpolarized experiment, it was noted 19 that the chpt predictions for the spin-dependent generalized polarizabilities entering the response functions give quite large contributions, even changing the sign in the case of $P_{LT}(|\vec{q}| = 600\text{MeV})$ for the Mainz kinematics! This situation is quite in contrast to the situation in unpolarized real Compton scattering, where the spin-polarizabilities tend to be small effects completely masked by spin-independent physics at low energies. On the experimental side this observation has led to a new proposal 20 to perform a second VCS experiment at MAMI,
which would use polarized electrons to determine a double-polarization asymmetry via measuring the average polarization of the recoiling final state proton. This experiment would give access to a different set of response functions possibly allowing for enough constraints to separate spin-dependent and spin-independent GPs. Note that such an experiment also holds the prospect of determining the essentially unknown real Compton spin polarizabilities $\gamma_{M2}, \gamma_{E2}$ via measuring $P_{(C1,M2)}(\vec{q})$, $P_{(M1,C2)}(\vec{q})$ at small values of momentum-transfer and then extrapolating $|\vec{q}| \to 0$.

On the theoretical side this prominence of the spin-effects certainly needs to be further investigated. In fact, one might worry that this large sensitivity to spin-effects gives an indication for the breakdown of the $O(p^3)$ heavy baryon calculation—especially if one remembers the tremendous sensitivity of some of the real Compton dipole spin-polarizabilities $\gamma_i$ on effects connected with $\Delta(1232)$ intermediate states. There it was argued that only a $O(p^5)$ (i.e. 2-loop) heavy baryon calculation (which no group seems to have on the “things-to-do” list until the next Chiral Dynamics workshop) posses the right operator structure with which one can hope to arrive at a decent description of these important spin-structure quantities. However, nature seems to be very kind to us in the case of VCS—as long as we stay in the “Guichon limit”. First, I want to show you the result of a recent calculation which gives the chiral $O(p^3)$ (i.e. leading order) momentum dependence of the 4 generalized spin-polarizabilities in closed form expressions

\begin{align}
P_{(C1,M2)}^{(3)}(\vec{q}) &= -\frac{g_A^2}{24\sqrt{2} \pi^2 F_\pi^2 |\vec{q}|^2} \left[ 1 - g\left(\frac{|\vec{q}|}{2m_\pi}\right)\right] \\
P_{(M1,C2)}^{(3)}(\vec{q}) &= -\frac{g_A^2}{12\sqrt{6} \pi^2 F_\pi^2 |\vec{q}|^2} \left[ 1 - g\left(\frac{|\vec{q}|}{2m_\pi}\right)\right] \\
P_{(M1,C0)}^{(3)}(\vec{q}) &= \frac{g_A^2}{12\sqrt{3} \pi^2 F_\pi^2} \left[ 2 - \left( 2 + \frac{3|\vec{q}|^2}{4m_\pi^2}\right) g\left(\frac{|\vec{q}|}{2m_\pi}\right)\right] \\
\hat{P}_{(C1,(C1,E1))}^{(3)}(\vec{q}) &= \frac{g_A^2}{24\sqrt{6} \pi^2 F_\pi^2 |\vec{q}|^2} \left[ 3 - \left( 3 + \frac{|\vec{q}|^2}{m_\pi^2}\right) g\left(\frac{|\vec{q}|}{2m_\pi}\right)\right], \quad (5)
\end{align}

with the functional dependence given by

$$g(x) = \frac{\sinh^{-1}(x)}{x\sqrt{1 + x^2}}.$$ 

One can clearly see that the scale of the $\vec{q}$-variation in the generalized spin polarizabilities is given by the pion mass $m_\pi$, as promised. The only other parameters entering the expressions in Eq.(5) are the axial vector coupling constant $g_A$ measured in neutron beta-decay and the pion-decay constant $F_\pi$. 
Figure 2. $O(\epsilon^3)$ SSE results for the four generalized spin polarizabilities, compared to the $O(p^3)$ heavy baryon chpt results of Eq.(5) in gray shading.

The leading order chiral prediction is therefore completely fixed. What about possible large corrections due to \(\Delta(1232)\)? Looking at the multipole content of the generalized spin-polarizabilities in the Guichon limit Eq.(5), one realizes that \(\Delta(1232)\) pole-contributions can only enter via an interference between the large \(\gamma N \Delta M_1\) coupling and the very small \(\gamma N \Delta E_2, C_2\) quadrupole couplings, triggering some optimism that a theoretical description of generalized spin polarizabilities in the Guichon limit can be reasonable without explicit Delta degrees of freedom. We note again that once (in future calculations) one goes beyond the Guichon kinematical limit and allows for additional independent generalized spin polarizabilities which for example involve two large \(M_1\) transitions to a \(\Delta(1232)\) in the intermediate state, then one will encounter the same situation in spin-dependent VCS as in polarized real Compton scattering— in heavy baryon chpt there would then be no way around a full $O(p^5)$ (i.e. 2-loop) calculation.

Having argued that for the Guichon set of generalized spin-polarizabilities we do not expect large pole contributions from \(\Delta(1232)\), what about possible large corrections due to \(\Delta \pi\) intermediate continuum states? A recent calculation [3] also shows that these contributions are quite small, completely analogous to the situation of the dipole spin polarizabilities \(\gamma_i\) in real Compton scattering. Figure 3 shows the results of a leading order (i.e. $O(\epsilon^3)$) calculation in the small scale expansion (SSE) which contains explicit nucleon, pion and delta degrees of freedom. As one can clearly see, the heavy baryon calculation lies quite close to their SSE counterpart, with the largest discrepancy
showing up in $P_{\text{MLCD0}}(q)$. The SSE calculation involves a lot more diagrams but only 2 additional parameters—$\Delta = M_\Delta - M_N = 292\text{MeV}$ the location of the Delta resonance pole and the strong $\pi N\Delta$ coupling constant $g_{\pi N\Delta}$, determined consistently within SSE from the width of the Delta resonance. No other (normalization) adjustment was made. Figure 2 shows the absolute predictions of the 2 quite distinct theoretical frameworks, yielding surprisingly similar results. Unfortunately the SSE results cannot be written into a nice analytic form as in Eq.(5) but need a numerical evaluation of the integrals, details can be found in ref.[21]. Certainly further investigations into the generalized spin polarizabilities are needed, but the new calculations[21] suggest, that chiral effective theories can also provide meaningful interpretation for the planned polarized VCS experiments focusing on the spin-dependent response of the nucleon in the presence of external electromagnetic source-terms. For recent progress utilizing a dispersion theoretical analysis to predict the $q^2$-dependence of some of the generalized spin-polarizabilities starting from pion photo-/electroproduction data, I refer to the talk by M. Vanderhaeghen in the working group[22].

3 Summary

I have reported about new developments since the last Chiral Dynamics workshop in Mainz in the field of nucleon Compton scattering using chiral effective field theories. The spontaneous breaking of chiral symmetry in QCD at low energies leads to a prominence of $\pi N$ intermediate states which often dominate the leading structure dependent response (i.e. the polarizabilities) of the nucleon when probed via external electromagnetic fields. Real and Virtual Compton scattering on the nucleon thus provide an excellent laboratory to uncover these signatures of chiral symmetry breaking amidst the usually dominant/overwhelming baryon-resonance physics. A microscopic understanding of polarizabilities (and their momentum dependence) in terms of a few simple Feynman diagrams connected to chiral effective field theories can be given in many cases, leading to a physical intuition/understanding of the numbers extracted in experiment. Several new experiments will have taken data by the time of the next Chiral Dynamics meeting improving the data base considerably. The main challenge for the theoretical efforts is to go beyond the leading order results in a controlled and effective approximation.

I would like to thank the organizers for giving me the opportunity to present this overview talk. Many interesting topics had to be skipped (for example the active field of Compton scattering on the Deuteron[23]) due to lack of time. See you in Jülich in 2003!
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