Generalized spin polarizabilities of the nucleon in Heavy Baryon Chiral Perturbation Theory at order $\mathcal{O}(p^4)$

Chung Wen Kao\(^1\) and Marc Vanderhaeghen\(^2\)

\(^1\)Department of Physics and Astronomy, University of Manchester, Manchester, M13 9PL UK
\(^2\)Institut für Kernphysik, Johannes Gutenberg-Universität, D-55099 Mainz, Germany

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We present the first heavy baryon chiral perturbation theory calculation at order $\mathcal{O}(p^4)$ for the spin-dependent amplitudes for virtual Compton scattering off the nucleon, and extract the $\mathcal{O}(p^3)$ results for the generalized spin polarizabilities of the nucleon.

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Over the past years, the virtual Compton scattering (VCS) process on the nucleon, accessed through the $ep \rightarrow ep\gamma$ reaction, has become a powerful new tool to study the internal structure of the nucleon both at low and high energies (see \cite{1} for a review). Being an electromagnetic process, it provides us with a well understood probe, as the electron interacts with the nucleon, to a very good precision, through the exchange of one (virtual) photon. Furthermore, in contrast to the real Compton scattering (RCS) process, VCS allows us to vary the virtuality of the initial photon, which plays the role of the resolution with which one “sees” the constituents within the nucleon. At very high virtuality and energy, the photon interacts with a single quark in the nucleon, and one accesses Compton scattering at the quark level. The VCS process on the nucleon tells us then how this quark is embedded in the nucleon, which is parametrized through so-called generalized parton distributions. At low virtuality and energy, the outgoing real photon plays the role of an applied quasi-static electromagnetic field, and the VCS process measures the response of the nucleon to this applied field, which can be parametrized in nucleon structure quantities, termed generalized polarizabilities (GPs) \cite{2}. In this case, the virtuality of the initial photon can be dialed so as to map out the spatial distribution of the electric polarization, the magnetization, or the spin densities of the nucleon \cite{3}. In particular, this allows us to study, at different length scales, the interplay of the role of the pion cloud and quark core contributions to the nucleon GPs. Because the chiral dynamics plays a dominant role in this regime, at low virtuality, Heavy Baryon Chiral Perturbation Theory (HBChPT) provides a natural way to study the GPs of the nucleon.

On the experimental side, the measurement of the VCS process became only possible in recent years, with the advent of high performance electron accelerator facilities, as these experiments require accurate measurements of small cross sections. At low photon energy and virtuality, first unpolarized VCS observables have been measured at the MAMI accelerator \cite{4} at a virtuality $Q^2 = 0.33$ GeV$^2$, and recently at JLab \cite{5} at higher virtualities, $1 < Q^2 < 2$ GeV$^2$, as well as at MIT-Bates \cite{6}. All those experiments measure two combinations of GPs.

The nucleon structure functions extracted from the first unpolarized VCS experiment \cite{4}, at a moderately large virtuality of $Q^2 = 0.33$ GeV$^2$, are in surprisingly good agreement with the $\mathcal{O}(p^4)$ HBChPT predictions for the GPs, performed in Refs. \cite{7,8}. In comparison with other model approaches, the good agreement between the unpolarized VCS response functions of \cite{9} and the $\mathcal{O}(p^3)$ HBChPT result, is for an important part due to the large spin GPs of the nucleon at this order, which also contribute to the unpolarized observables. This has triggered a further experimental program \cite{10} to directly access and separate the spin GPs of the nucleon by measuring VCS double polarization observables \cite{11}. To sort out the importance of the spin GPs of the nucleon, we calculate in this letter the spin-dependent VCS amplitudes in HBChPT at order $\mathcal{O}(p^4)$. At this order, no unknown low-energy constants enter the calculations so that one can extract the $\mathcal{O}(p^4)$ results for the spin GPs of the nucleon, without free parameters.

We begin by specifying our notation for the VCS process:

$$\gamma^* (\epsilon_1, q) + N(p_1) \rightarrow \gamma (\epsilon_2, q') + N(p_2). \quad (1)$$

In calculating the VCS amplitudes, we work in the c.m. frame and choose the Coulomb gauge. In the VCS process, the initial spacelike photon is characterized by its four momentum $q = (\omega, \hat{q})$, virtuality $Q^2 \equiv -q^2$, and polarization vector $\epsilon_1 = (0, \epsilon_1)$. The outgoing real photon has four momentum $q' = (\omega' = |q'|^2, \hat{q}')$ and polarization vector $\epsilon_2 = (0, \epsilon_2)$. We define $\hat{q} \equiv |\hat{q}|$, and denote $\theta$ as the scattering angle between virtual and real photons, i.e., $\cos \theta = \hat{q} \cdot \hat{q}'$. The polarization vector $\epsilon_1$ of the virtual photon can be decomposed into a longitudinal component $\epsilon_{1L} = (\epsilon_1 \cdot \hat{q}) \hat{q}$ and a transverse component $\epsilon_{1T}$. For further use, we introduce the virtual photon energy in the limit $\omega' = 0$:

$$\omega_0 \equiv \omega(\omega' = 0, \hat{q}) = M_N - \sqrt{M_N^2 + \hat{q}^2},$$

$$= -\hat{q}^2/(2M_N) + \mathcal{O}(1/M_N^2). \quad (2)$$

with $M_N$ the nucleon mass, and where the last line in Eq. (2) indicates the heavy baryon expansion.
The VCS amplitude $\mathcal{M}_{VCS}$ can be expressed in terms of twelve structure functions as [3]:

$$
\mathcal{M}_{VCS} = i e^2 \left\{ (\vec{c}_2 \cdot \tilde{c}_{1T}) A_1 + (\vec{c}_2' \cdot q)(\tilde{c}_{1T} \cdot \tilde{q}') A_2 
+ i \vec{q}' \cdot (\vec{c}_2' \times \tilde{c}_{1T}) A_3 + i \vec{q}' \cdot (\tilde{q}' \times \tilde{q}')(\vec{c}_2' \cdot \tilde{c}_{1T}) A_4 
+ i \vec{q}' \cdot (\vec{c}_2' \times \tilde{q}') (\tilde{c}_{1T} \cdot \tilde{q}') A_5
+ i \vec{q}' \cdot (\vec{c}_2' \times q)(\tilde{c}_{1T} \cdot \tilde{q}') A_6 
- i \vec{q}' \cdot (\tilde{c}_{1T} \times q)(\vec{c}_2' \cdot \tilde{q}') A_7 
- i \vec{q}' \cdot (\tilde{c}_{1T} \times \tilde{q}')(\vec{c}_2' \cdot q) A_8 
+ (\vec{c}_{1L} \cdot \tilde{q}')(\vec{c}_2' \cdot \tilde{q}') A_9 + i \vec{q}' \cdot (\tilde{q}' \times \tilde{q}')(\vec{c}_2' \cdot \tilde{q}') A_{10}
+ i \vec{q}' \cdot (\vec{c}_2' \times \tilde{q}') A_{11} + i \vec{q}' \cdot (\vec{c}_2' \times q) A_{12} \right\}
$$

To extract the GPs, we first calculate the complete fourth order VCS amplitudes $A_4$ in HBChPT, and subsequently separate the amplitudes in a Born part $A_4^{Born}$, and a non-Born part $\bar{A}_4$ as:

$$
A_4(\omega', \tilde{q}, \theta) = A_4^{Born}(\omega', \tilde{q}, \theta) + \bar{A}_4(\omega', \tilde{q}, \theta),
$$

In the Born process, the virtual photon is absorbed on a nucleon and the intermediate state remains a nucleon. We calculate them, following the definition of Ref. [2], from direct and crossed Born diagrams with the electromagnetic vertex given by:

$$
\Gamma^\mu(q^2) = F_1(q^2)\gamma^\mu + F_2(q^2)i\sigma^{\mu\nu}\frac{q_\nu}{2M_N},
$$

where $F_1$ ($F_2$) are the nucleon Dirac (Pauli) form factors respectively. The non-Born terms contain the information on the internal structure of the nucleon. In the framework of HBChPT, the non-Born amplitudes start at order $O(p^4)$, where $p$ ($p = m_N, \tilde{q}, \omega'$) is a momentum scale of the problem, which is small compared with the nucleon mass $M_N$, such that a systematic expansion can be made in powers of $p/M_N$. In particular, we are interested in this work in the nucleon response to an applied quasi-static electromagnetic dipole field. This is accessed by performing a low energy expansion (LEX), in the outgoing photon energy $\omega'$, of the non-Born VCS amplitudes, and by selecting the term linear in $\omega'$, as:

$$
\left[ \frac{\partial \bar{A}_4}{\partial \omega'}(\omega', \tilde{q}, \theta) \right]_{\omega'=0} = S_4(\tilde{q}) + \mathcal{P}_4(\tilde{q}) \cos \theta.
$$

In Eq. (5), the nucleus structure quantities $S_4$ and $\mathcal{P}_4$ can be expressed in terms of the GPs of the nucleon, as introduced in [3], which are functions of $\tilde{q}$ and which are denoted by $D^\rho L_{\rho L} S(\tilde{q})$. In this notation, $\rho$ ($\rho'$) refers to the electric ($E$), magnetic ($M$) or longitudinal ($L$) nature of the initial (final) photon, $L$ ($L'$) represents the angular momentum of the initial (final) photon, and $S$ differentiates between the spin-flip ($S = 1$) and non-spin-flip ($S = 0$) character of the transition at the nucleon side. Restricting oneself to a dipole transition for the final photon (i.e. $L' = 1$), angular momentum and parity conservation lead to 3 scalar and 7 spin GPs [3]. The 7 spin GPs, which are the subject of investigation in this letter, are obtained as [11]:

$$
\begin{align*}
\mathcal{P}^{(M;L)1} &= \frac{-2\sqrt{2}}{3\sqrt{3}} \frac{1}{\tilde{q}\omega_0} \sqrt{\frac{M_N}{E_N}} S_{10}, \\
\mathcal{P}^{(L;L)1} &= \frac{-2}{3\sqrt{3}} \frac{1}{\tilde{q}\omega_0} \sqrt{\frac{M_N}{E_N}} S_{11}, \\
\mathcal{P}^{(M;L)0} &= \frac{2}{\sqrt{3}} \frac{\tilde{q}}{\omega_0} \sqrt{\frac{M_N}{E_N}} \left[ S_{12} - \frac{2}{3} S_{10} \right], \\
\mathcal{P}^{(M;M)1} &= \frac{2}{3\sqrt{3}} \frac{1}{\sqrt{E_N}} \sqrt{\frac{M_N}{E_N}} S_8, \\
\mathcal{P}^{(M;L)1} &= \frac{-4}{3\sqrt{10}q^3} \sqrt{\frac{M_N}{E_N}} \left[ S_4 + S_7 - S_{10} \right], \\
\mathcal{P}^{(L;L)1} &= \frac{-2\sqrt{2}}{3\sqrt{3}g} \sqrt{\frac{M_N}{E_N}} \left[ S_3 + \frac{1}{2} S_8 - S_{11} \right],
\end{align*}
$$

where $E_N = \sqrt{M_N^2 + q^2}$, and where the GPs denoted by $\mathcal{P}$ correspond with mixed electric and longitudinal multipoles, as introduced in [3]. We will call a GP as $O(p^\nu)$ in HBChPT if it is extracted from the $O(p^\nu)$ VCS amplitudes using Eq. (4). As the linear term in the LEX of the 9 spin-dependent VCS amplitudes of Eq. (3) depends upon only 7 independent GPs, there is a redundancy among the coefficients of Eq. (6). In particular, one has the following constraints between the $S_i$ and $\mathcal{P}_i$:

$$
\begin{align*}
S_3 &= -S_7, \\
S_{11} &= -S_{10}, \\
S_4 &= -S_5, \\
S_6 &= S_4 = S_5 = S_6 = S_7 = S_8 = S_9 = S_{10} = S_{12} = 0.
\end{align*}
$$

which provide consistency checks for the HBChPT results of the VCS amplitudes at a given order, and have been verified up to fourth order here. Furthermore, it has been shown [11] that nucleon crossing combined with charge conjugation symmetry of the VCS amplitudes provides 3 additional constraints among the 7 spin GPs:

$$
\begin{align*}
S_4 &= 0, \\
S_3 &= \frac{\tilde{q}}{\omega_0} S_7, \\
S_{10} - S_{12} &= \frac{\tilde{q}}{\omega_0} S_{11},
\end{align*}
$$

leaving 4 independent spin GPs. By expanding $\omega_0$ as in [3], it is obvious that the relations (6) connect quantities of different order in the heavy baryon expansion. For the first non-trivial check of these relations in HBChPT, one has to calculate the VCS amplitudes to $O(p^4)$. This is performed for the first time for the spin-dependent VCS amplitudes in this letter, and we verified that the VCS amplitudes in HBChPT at $O(p^4)$ satisfy the relations (6).

Furthermore, in the limit $\tilde{q} \to 0$, the GPs are related with the Ragusa spin polarizabilities [12] of RCS, which have been calculated before in HBChPT at order $O(p^2)$ [13, 14, 15]. We verified that our $O(p^4)$ results reduce at the real point to those of Refs. [13, 14].

By calculating the VCS amplitudes in HBChPT at
\( \mathcal{O}(p^3) \), one extracts from Eq. (3) the following expressions for the \( \mathcal{O}(p^3) \) GPs (see Ref. 8):

\[
\begin{align*}
\mathcal{P}^{(M1, L2)}_1(q) &= \mathcal{P}^{(M1, L0)}_1(q) = \hat{\mathcal{P}}^{(M1, L2)}_1(q) = 0, \\
\mathcal{P}^{(M1, M1)}_1(q) &= \mathcal{P}^{(L1, L1)}_1(q) = 0, \\
\mathcal{P}^{(L1, M2)}_1(q) &= \frac{-g_4^2}{24 \sqrt{2} F_2^2 M^2 q^2} \left[ 1 - \frac{4}{w \sqrt{w^2 + 4}} \sinh^{-1} \frac{w}{2} \right], \\
\hat{\mathcal{P}}^{(L1, L1)}_1(q) &= \frac{g_4^2}{24 \sqrt{5} \pi F_2^2 q^2} \left[ 3 - \frac{4w^2 + 12}{w \sqrt{w^2 + 4}} \sinh^{-1} \frac{w}{2} \right].
\end{align*}
\]

where \( w \equiv q / m_\pi \), with \( m_\pi \) the pion mass. Furthermore, throughout this paper we use the values: \( g_A = 1.267 \), \( F_\pi = 0.9924 \text{ GeV} \), and \( m_\pi = 0.14 \text{ GeV} \).

Our calculation of the VCS amplitudes in HBChPT at \( \mathcal{O}(p^3) \), yields the \( \mathcal{O}(p^3) \) GPs, from Eq. (3), as :

\[
\begin{align*}
\mathcal{P}^{(M1, L2)}_1(q) &= \frac{-g_4^2}{12 \sqrt{6} \pi F_2^2 q^2} \left[ 1 - \frac{4}{w \sqrt{w^2 + 4}} \sinh^{-1} \frac{w}{2} \right], \\
\mathcal{P}^{(M1, L0)}_1(q) &= \frac{g_4^2}{12 \sqrt{3} \pi F_2^2 q^2} \left[ 2 - \frac{3w^2 + 8}{w \sqrt{w^2 + 4}} \sinh^{-1} \frac{w}{2} \right], \\
\mathcal{P}^{(M1, M1)}_1(q) &= \frac{g_4^2}{24 \sqrt{2} \pi F_2^2 q^2} \left[ 1 - \frac{w^2 + 4}{w \sqrt{w^2 + 4}} \sinh^{-1} \frac{w}{2} \right], \\
\hat{\mathcal{P}}^{(M1, L2)}_1(q) &= \frac{-g_4^2}{24 \sqrt{10} \pi F_2^2 q^2} \left[ 3 - \frac{2w^2 + 12}{w \sqrt{w^2 + 4}} \sinh^{-1} \frac{w}{2} \right], \\
\mathcal{P}^{(L1, L1)}_1(q) &= 0, \\
\mathcal{P}^{(L1, M2)}_1(q) &= \frac{g_4^2}{96 \sqrt{2} \pi F_2^2 q^2} \frac{q}{MN} \times \left[ \frac{1}{2w} + \frac{2w^2 + 4}{w (w^2 + 4)} + \frac{5}{4} \cdot \frac{3}{w} \tan^{-1} \frac{w}{2} \right] + \gamma (\frac{1}{2w} + \frac{1}{w^2} \tan^{-1} \frac{w}{2}) \right], \\
\hat{\mathcal{P}}^{(L1, L1)}_1(q) &= \frac{-g_4^2}{96 \sqrt{6} \pi F_2^2 q^2} \frac{q}{MN} \times \left[ \frac{11}{2w} + \frac{2w^2 + 4}{w (w^2 + 4)} + \frac{25}{4} + \frac{9}{w} \tan^{-1} \frac{w}{2} \right] + \gamma (\frac{3}{2w} + \frac{5}{4} + \frac{3}{w} \tan^{-1} \frac{w}{2}).
\end{align*}
\]

We can get more predictions by use of the crossing relations (8), which hold in general in a relativistic quantum field theory, and which we verified here in HBChPT at \( \mathcal{O}(p^4) \). By plugging in the fourth order amplitudes in (8), we can extract fifth order predictions from \( S_3^{(4)} = -2M_N/q \sigma_7^{(5)} \), and \( S_1^{(4)} = -S_2^{(4)} = -2M_N/q \sigma_1^{(5)} \).

In this way, these relations allow us to extract two \( \mathcal{O}(p^5) \) spin GPs:

\[
\begin{align*}
p^{(M1, M1)}_1(q) &= \frac{-g_4^2 q}{192 \sqrt{2} \pi F_2^2 M q^2} \left[ \frac{3}{w} + \frac{5}{2} + \frac{6}{w^2} \tan^{-1} \frac{w}{2} \right] + \gamma (\frac{1}{w} - \frac{1}{2} + \frac{2}{w^2} \tan^{-1} \frac{w}{2}) \right], \\
\hat{p}^{(L1, L1)}_1(q) &= \frac{g_4^2 q}{48 \pi^2 F_2^2 M} \left[ 1 - \frac{2w^2 + 4}{w \sqrt{w^2 + 4}} \sinh^{-1} \frac{w}{2} \right].
\end{align*}
\]

The direct verification of these two predictions would imply a two-loop calculation, which has not been done so far. Furthermore, it is interesting to mention that the \( \mathcal{O}(p^3) \) calculation (8) was able to obtain some of these \( \mathcal{O}(p^3) \) and one \( \mathcal{O}(p^3) \) results, assuming the crossing relations (8), in particular the \( \mathcal{O}(p^4) \) results for \( p^{(M1, M1)}_1 \) and \( \hat{p}^{(M1, L2)}_1 \), as well as the \( \mathcal{O}(p^5) \) result for \( p^{(L1, L1)}_1 \).

Also, the \( \mathcal{O}(p^4) \) results for \( p^{(M1, L2)}_1 \) and \( p^{(M1, L0)}_1 \) were obtained in (8), by performing the LEX of the amplitudes \( A_1 \) to second order in \( \omega' \), and, by isolating two terms in \( \omega^2 \) which depend on those spin GPs. Besides confirming these previous results, our genuinely new predictions at \( \mathcal{O}(p^4) \) are for \( p^{(L1, M2)}_1 \) and \( \hat{p}^{(L1, L1)}_1 \), and at \( \mathcal{O}(p^5) \) for \( p^{(M1, M1)}_1 \). It is also worth noting that the relations (8) allow to classify the GPs in two groups :

\[
\begin{align*}
G1 &= \{ p^{(M1, L2)}_1, p^{(M1, L0)}_1, p^{(M1, M1)}_1, \hat{p}^{(M1, L2)}_1 \}, \\
G2 &= \{ p^{(L1, L1)}_1, p^{(L1, M2)}_1, \hat{p}^{(L1, L1)}_1 \}.
\end{align*}
\]

It is interesting to observe that their analytical forms alternate from one order to the next. At \( \mathcal{O}(p^3) \) one has :

\[
\begin{align*}
G1 : 1/\pi \tan^{-1}[q/2m_\pi], \\
G2 : 1/\pi^2 \cdot \sinh^{-1}[q/2m_\pi].
\end{align*}
\]

At \( \mathcal{O}(p^4) \), one has :

\[
\begin{align*}
G1 : 1/\pi^2 \cdot \sinh^{-1}[q/2m_\pi], \\
G2 : 1/\pi \cdot \tan^{-1}[q/2m_\pi].
\end{align*}
\]

At \( \mathcal{O}(p^5) \) one again has the same analytical structure as at \( \mathcal{O}(p^3) \), etc. The alternating analytical forms teach us how the crossing relations (8), connecting the GPs between two different orders in HBChPT, work.

![FIG. 1: Results for the spin GPs: LO HBChPT results (dotted curves); NLO HBChPT results (dashed curves); dispersive evaluations of Ref. 10 (solid curves).](image-url)
In the following, we denote the leading order non-vanishing results for the GPs in HBChPT as LO HBChPT, and refer to the results including the first chiral corrections as NLO HBChPT. In Fig. 4 we show the 4 independent spin GPs, and compare the LO HBChPT results of [3, 4], with the NLO results from the present work. Furthermore, we also show a comparison with a dispersion relation (DR) evaluation [16, 17]. The GPs $P^{(M_1, M_0)}$, $P^{(M_1, M_2)}$, and $P^{(L_1, L_1)}$ are expected to be promising observables to study the effects of pionic effects for these GPs in HBChPT as LO calculations in Fig. 4, indeed show large contribution of pionic effects for these GPs. For $P^{(L_1, L_1)}$, this result is confirmed by our $O(p^4)$ calculation. For the GP $P^{(M_1, M_1)}$, one sees that the NLO result yields a large reduction compared to the LO one, and calls the convergence of the HBChPT result for this observable into question. Such a reduction has also been noticed in the linear $\sigma$-model calculation [3], which takes account of part of the higher order terms of a consistent chiral expansion, resulting in general in smaller values for $P^{(L_1, L_1)}$ and $P^{(M_1, M_1)}$ compared with the LO calculations in HBChPT. For the GP $P^{(L_1, M_2)}$, we also notice that the NLO HBChPT result yields a relatively large correction. It is of course no surprise that the NLO correction are very large since the corresponding calculations at the real photon point have already shown it [3, 4]. However it is worth noting that the NLO results show a smoother $Q^2$ dependence, and tend to lie closer to the phenomenological DR estimates.

The unpolarized VCS observables at low energy can be expressed in terms of 3 structure functions, one of which, denoted by $P_{TT}$, involves only spin GPs [3]:

$$P_{TT} = -3G_M \frac{q^2}{\omega_0} \left( P^{(M_1, M_1)} - \sqrt{2}\omega_0 P^{(L_1, M_2)} \right),$$

with $G_M = F_1 + F_2$. We show the result for $P_{TT}$ in HBChPT in Fig. 2, and see that it receives a very sizeable change at NLO. The measurement of $P_{TT}$ is planned in the near future [3] and will provide an interesting check on the spin densities of the nucleon.

In this work, we provided the first HBChPT calculation at $O(p^4)$ for the spin-dependent VCS amplitudes. We extracted analytical formulas for the $O(p^4)$ spin GPs of the nucleon, which involve no free parameter at this order. Our results show sizeable chiral corrections for the spin GPs, which are important for the interpretation of existing and forthcoming experiments.

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