Nuclear Compton scattering in the $\Delta$-resonance region
with polarized photons

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Abstract

Nuclear Compton scattering in the $\Delta$-resonance region is reconsidered
within the framework of the $\Delta$-hole model. The different role of the
resonant and non-resonant contributions to the transition amplitudes
is discussed and their effect is investigated by comparing the results
of calculation with recent data also taken with polarized photons.

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1 Introduction

Photon scattering on nuclei is a pure electromagnetic process providing a very
useful tool for investigating the nuclear structure and dynamics. The photon
probe has a well-known interaction and a mean free path much longer than
the target dimensions, thus exploring the entire target volume. The scattered
photon emerges in general through a two-step mechanism involving the whole
internal dynamics of the target nucleus through virtual excitations in the in-
termediate state. In the elastic scattering case, i.e. Compton scattering,
coherence among the different transition amplitudes is demanded to recover
the ground state in the final state. On the one hand, this implies the possi-
bility of using closure and calculating ground-state expectation values with a
minimum uncertainty since ground-state wave functions are best known. On
the other hand, the role of the transition operator and the relevant degrees
of freedom is emphasized.

Depending on the photon energy, different nuclear degrees of freedom
come into play. At low energy, low-energy theorems can be derived from
basic principles, such as Lorentz invariance and gauge invariance of the elec-
 tromagnetic interaction \[1,2\]. For spin-1/2 particles the elastic photon scatter-
ing amplitude is determined up to first order in the photon energy by global
properties such as its charge, mass and magnetic moment. In general, up to
second order in the photon energy the nuclear response can be described in
terms of the static electric and magnetic polarizabilities \[3,4,5\]. A correct
treatment of the c.m. contribution is also important \[6\] and must be included
consistently \[7\].

By expanding the elastic photon scattering amplitude into multipole fiel-
descending as well as the scattered photon, generalized polarizabilities
can be defined and an extension of the low-energy theorems has been de-
 rived \[8\]. A consistent study is then possible up to and above the pion
production threshold \[9\] to test the stressed importance of meson-exchange
effects \[10\].

Cross sections for elastic photon scattering are very small and the energy
resolution in the detection of the scattered photon must be sufficient to ex-
clude inelastic scattering. Thus, elastic-scattering experiments are difficult
to carry out and existing data are limited and only refer to zero-spin nu-
 clei. Below pion threshold, combined analyses of absorption and scattering
data \[11,12,13\] were successful in describing the scattering process in terms
of bound-nucleon properties and in extracting information about the nuclear
polarizabilities. Further references together with a detailed discussion of the
theoretical frame and experimental methods can be found in refs. \[14\] and
\[15\], respectively.

At higher energies, where the \(\Delta\) resonance becomes important, early mea-
urements \[16\] on \(^{12}\text{C}\) and \(^{208}\text{Pb}\) have been followed by an extensive inves-
tigation on \(^{4}\text{He}\) from pion threshold to the \(\Delta\) region with bremsstrahlung
beams \[17,18,19,20\]. In this energy domain Compton scattering provides
information about the intermediate formation and propagation of the \(\Delta\) iso-
bar in the nuclear medium, which is complementary to that obtained from
pion scattering and pion photoproduction. Different \(\Delta\)-hole models have
been proposed \[21,22,23\] to account for a unified description of such pro-
cesses. In particular, a consistent picture can be obtained between data from pion scattering \[24, 25, 26\] and pion photoproduction \[27, 28, 29\]. The Compton scattering data have been compared with results obtained within the same \(\Delta\)-hole model \[23\]. The data at 320 MeV \[18\] near the peak of the \(\Delta\) resonance compare well with the predictions of the \(\Delta\)-hole model. At lower energies, however, the model fails to reproduce the magnitude of the cross section at forward angles and the strong back-angle rise. At backward angles inelastic contributions are important \[30, 31\] as in the case of pion photoproduction \[32\]. However, as repeatedly noted, also non-resonant background contributions to coherent scattering are relatively large outside the \(\Delta\)-resonance region and may not have been correctly included in the model. On the other hand, the impulse approximation with \(\gamma N\) amplitudes from relativistic dispersion relations, supplemented by Siegert-like arguments to include the main E1-part of the meson-exchange contributions \[33\], seems to give a reasonable description of the \(^4\)He data.

Recently, elastic and inelastic scattering of monochromatic photons from \(^{12}\)C was investigated in the energy range across the \(\Delta\) resonance \[34\]. The observed elastic cross section at a scattering angle of 40° was compared with a recent \(\Delta\)-hole calculation \[35\] which differs from \[23\] by a changed pion coupling and by inclusion of the \(\rho\)-meson into the final-state interaction. As was also found for \(^4\)He the \(\Delta\)-hole model gives a good agreement in the region of the resonance, but it fails to reproduce the data below the \(\Delta\) peak where they are dominated by a non-resonant background.

With polarized photons two structure functions contribute to the cross section \[36\]. The first structure function, \(W_T\), is the same quantity determined by scattering of unpolarized photons and is the incoherent sum of photon-helicity flip and non-flip contributions. The second structure function, \(W_{TT}\), contains interference contributions from helicity flip and non-flip amplitudes. Thus the resonant and non-resonant amplitudes enter quite differently in \(W_T\) and \(W_{TT}\). Their separated determination is made possible by a combined measurement of the differential cross sections for photons linearly polarized in the reaction plane and perpendicular to it. Such a separation has been achieved in a recent experiment on \(^4\)He \[37\] covering the same energy range previously explored with unpolarized bremsstrahlung beams \[17, 18, 19, 20\]. This adds further information to the precise data recently obtained at Mainz with tagged photons \[38\].

The enriched quality of the data opens the possibility of a better un-
standing of the reaction mechanism. In this paper the problem is reconsidered within the frame of the \( \Delta \)-hole model. In order to reduce the computational effort, the local version of the model developed in refs. [39, 40, 41] is adopted to describe the resonant \( \Delta \) formation and propagation. This simplified version retains essential ingredients of the original \( \Delta \)-hole model and provides a useful framework to describe photon absorption [42, 43], pion-nucleus scattering [41], coherent \( \pi^0 \) photo- and electro-production (refs. [44] and [45], respectively) and Compton scattering [46].

In sect. 2, the general formalism for photon scattering is reviewed and the model used to describe nuclear Compton scattering in the \( \Delta \)-resonance region is described in sect. 3. Results for \( ^{12}\text{C} \) and \( ^4\text{He} \) are presented and compared with the experimental data in sect. 4. The conclusions are drawn in the final section.

2 General formalism

The scattering of a photon with momentum \( \vec{k} \) and polarization \( \lambda = \pm 1 \) into a photon with momentum \( \vec{k}' \) and polarization \( \lambda' \) is described by the scattering amplitude \( T^{\lambda \lambda'}(\vec{k}', \vec{k}) \). During the scattering process the nuclear target undergoes a transition from the initial state \( |I_i M_i \rangle \) to the final state \( |I_f M_f \rangle \).

According to refs. [8, 14], a convenient decomposition of the scattering amplitude is provided in terms of an expansion into generalized polarizabilities \( P^J(M^{\nu'} L', M^{\nu} L; k', k) \). They correspond to an expansion of the incoming and scattered photon into multipole fields of order \( L \) and \( L' \), respectively, with \( M^0 = E \) (\( \nu = 0 \), electric) and \( M^1 = M \) (\( \nu = 1 \), magnetic), while the total angular momentum transferred to the target nucleus is \( J \), which is constrained by the conditions \( |L - L'| \leq J \leq L + L' \) and \( |I_i - I_f| \leq J \leq I_i + I_f \). In general, these polarizabilities are defined by

\[
P_J(M^{\nu'} L', M^{\nu} L; k', k) =
\]

\[
= (-)^{L+L'-I_f} \hat{L}^2 \sum_{M_i M_f M M'} \frac{1}{4} \sum_{\lambda \lambda'} \lambda^{\nu'} \lambda^{\nu}
\times (-)^{M_i} \left( \begin{array}{c}
I_f & J & I_i \\
-M_f & m & M_i
\end{array} \right) \left( \begin{array}{c}
L & L' & J \\
M & M' & -m
\end{array} \right)
\]
\[ \times \frac{1}{(8\pi^2)^2} \int dR \int dR' \mathcal{D}^{L'\ast}_{M'\lambda}(R') T^{\lambda\Lambda}_{M_iM_f}(\vec{k}', \vec{k}) \mathcal{D}^{L\ast}_{M\Lambda}(R), \]  

where \( \hat{L}^2 = 2L + 1 \), \( R \) and \( R' \) denote rotations of the quantization axis into the direction of \( \vec{k} \) and \( \vec{k}' \), respectively.

In the case of elastic photon scattering, i.e. Compton scattering, and for a zero-spin target nucleus (\( I_i = I_f = 0 \)), only the scalar electric and magnetic polarizabilities, \( P_0(EL, EL; k, k) \) and \( P_0(ML, ML; k, k) \), respectively, survive in eq. (1).

In the c.m. frame of reference the differential cross section for scattering of polarized photons on a nucleus reads [36]

\[
\frac{d\sigma}{d\Omega} = \frac{k'}{k} \left( \frac{M_T}{4\pi E_T} \right)^2 \frac{1}{2J_i + 1} \sum_{\lambda\lambda'} \sum_{M_iM_f} (T^{\lambda\Lambda}_{M_iM_f}) \rho_{\lambda\bar{\lambda}}(T^{\lambda\bar{\lambda}}_{M_iM_f})^*, \]

where \( M_T \) is the target mass and \( E_T \) is the total c.m. energy.

The density matrix \( \rho_{\lambda\bar{\lambda}} \) describing the polarization of the incident photon is given by

\[
\rho_{\lambda\bar{\lambda}} = \frac{1}{2} \left( \begin{array}{cc} 1 + \mathcal{P} & -\mathcal{L}e^{-2i\phi} \\ -\mathcal{L}e^{2i\phi} & 1 - \mathcal{P} \end{array} \right),
\]

where \( \mathcal{L} (\mathcal{P}) \) is the relative intensity of linearly (circularly) polarized photons. Assuming the \( xz \)-plane as the photon scattering plane with the quantization axis along \( \vec{k} \), \( \phi \) is the angle between the direction of the linear polarization and the \( x \)-axis.

Using standard Racah algebra, the cross section (2) becomes

\[
\frac{d\sigma}{d\Omega} = \frac{k'}{k} \left( \frac{M_T}{4\pi E_T} \right)^2 \frac{1}{2J_i + 1} \sum_{\lambda\lambda'} \sum_{M_iM_f} \sum_{LL'} \sum_{JK} \sum_{J'K'} (-)^{J+K} j^2 \hat{K}^2 \mathcal{D}^{K}_{M\lambda M_f \lambda'} \rho_{\lambda\bar{\lambda}} \rho_{\lambda'\bar{\lambda}'} P_J(M'\bar{L}', M\bar{L}; k', k) \rho_{\lambda\bar{\lambda}} P_J(M'\bar{L}', M\bar{L}; k', k),
\]

where \( \hat{L}^2 = 2L + 1 \), \( R \) and \( R' \) denote rotations of the quantization axis into the direction of \( \vec{k} \) and \( \vec{k}' \), respectively.

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\]
Performing the summations over $\lambda$, $\lambda'$ and $\bar{\lambda}$ and making explicit the $\phi$-dependence of the cross section, one has

\[
\frac{d\sigma}{d\Omega} = \frac{k'}{k} \left( \frac{M_T}{4\pi E_T} \right)^2 \frac{1}{2J_i + 1} \sum_{\nu\nu'} \sum_{LL'} \sum_{LL'} \sum_{JK} (-)^{J + L + \bar{L}} j^2 \hat{K}^2 \\
\times \left[ 1 + (-)^{\nu + \bar{\nu} + L + \bar{L} + K} \right] \left( \begin{array}{ccc}
\bar{L}' & \bar{L}' & K \\
-1 & 1 & 0
\end{array} \right) \left( \begin{array}{ccc}
L' & L & J \\
\bar{L} & \bar{L}' & K
\end{array} \right) \\
\times \left[ P_J(M^{\nu'} L', M^{\bar{\nu}} \bar{L}; k', k) \right]^* P_J(M^{\nu} L', M^{\bar{\nu}} \bar{L}; k, k) \\
\times \left[ \left( \begin{array}{ccc}
L & \bar{L} & K \\
1 & -1 & 0
\end{array} \right) D^K_{00} + (-)^{1 + \phi} \left( \begin{array}{ccc}
L & \bar{L} & K \\
1 & 1 & -2
\end{array} \right) D^K_{20} \cos 2\phi \right]. (5)
\]

Eq. (5) can also be rewritten as

\[
\frac{d\sigma}{d\Omega} = \frac{k'}{k} \left( \frac{M_T}{4\pi E_T} \right)^2 \left( W_T + W_{TT} \cos 2\phi \right), (6)
\]

where the two structure functions are defined in terms of helicity flip ($T_{M_f M_i}^{1-1}$) and non-flip ($T_{M_f M_i}^{11}$) amplitudes by

\[
W_T = \frac{1}{2J_i + 1} \sum_{M_f M_i} \left[ \left| T_{M_f M_i}^{11} \right|^2 + \left| T_{M_f M_i}^{1-1} \right|^2 \right], \quad (7)
\]

\[
W_{TT} = -\frac{1}{2J_i + 1} \sum_{M_f M_i} 2\text{Re} \left[ T_{M_f M_i}^{11} T_{M_f M_i}^{1-1*} \right]. \quad (8)
\]

In terms of the generalized polarizabilities (11), for a zero-spin target nucleus the Compton scattering amplitude becomes

\[
T_{M_f M_i}^{\lambda \lambda} = \sum_L (-)^{L + \bar{L}} L^{-1} D_{\lambda \lambda}^L \left[ P_0(EL, EL, k, k) + \lambda \lambda' P_0(ML, ML, k, k) \right], \quad (9)
\]

so that the two structure functions are simply given by

\[
W_T = \sum_{LL} (-)^{L + \bar{L}} L^{-1} \hat{L}^{-1}
\]

6
\[ \times \left\{ D_{1,1}^L(R') D_{-1,1}^L(R') \left[ P_0(EL, EL, k, k) + P_0(ML, ML, k, k) \right] \right. \]
\[ \times \left[ P_0^*(E\bar{L}, E\bar{L}, k, k) + P_0^*(M\bar{L}, M\bar{L}, k, k) \right] \} \]
\[ + D_{1,-1}^L(R') D_{-1,1}^L(R') \left[ P_0(EL, EL, k, k) - P_0(ML, ML, k, k) \right] \]
\[ \times \left[ P_0^*(E\bar{L}, E\bar{L}, k, k) - P_0^*(M\bar{L}, M\bar{L}, k, k) \right] \}, \quad (10) \]

\[ W_{TT} = -\sum_{LL} (\pm^L \pm^L \pm^L \pm^L) 2 \text{Re} \left\{ D_{1,1}^L(R') D_{-1,1}^L(R') \right. \]
\[ \times \left[ P_0(EL, EL, k, k) + P_0(ML, ML, k, k) \right] \]
\[ \times \left[ P_0^*(E\bar{L}, E\bar{L}, k, k) - P_0^*(M\bar{L}, M\bar{L}, k, k) \right] \}. \quad (11) \]

With linearly polarized photons (\( \mathcal{L} = 1 \)) one can separate these two structure functions by also measuring the photon asymmetry

\[ A = \frac{\left( \frac{d\sigma}{d\Omega} \right)_\parallel - \left( \frac{d\sigma}{d\Omega} \right)_\perp}{\left( \frac{d\sigma}{d\Omega} \right)_\parallel + \left( \frac{d\sigma}{d\Omega} \right)_\perp} = \frac{W_{TT}}{W_T}, \quad (12) \]

where

\[ \left( \frac{d\sigma}{d\Omega} \right)_\parallel = k' k \left( \frac{M_T}{4\pi E_T} \right)^2 (W_T + W_{TT}), \quad (13) \]
\[ \left( \frac{d\sigma}{d\Omega} \right)_\perp = k' k \left( \frac{M_T}{4\pi E_T} \right)^2 (W_T - W_{TT}), \quad (14) \]

are the cross sections for scattering of photons with polarization along the \( x \)-axis (\( \phi = 0 \)) and the \( y \)-axis (\( \phi = \pi/2 \)), respectively.
3 The model

The electromagnetic interaction Hamiltonian describing photon scattering is second order in the electromagnetic potential and the scattering amplitude is the sum of two terms, the one from second-order contributions of the current-density operator and the other from the first-order contribution of the genuine two-photon amplitude. In the $\Delta$-resonance region the first amplitude contains a resonant term with intermediate excitation of a $\Delta$-hole state which dominates the cross section. All of the other terms in the total scattering amplitude may be considered as a non-resonant background. Accordingly, the transition operator $T^{\lambda\lambda'}$ is split into the sum of a resonant ($R^{\lambda\lambda'}$) and a background ($B^{\lambda\lambda'}$) part:

$$T^{\lambda\lambda'} = R^{\lambda\lambda'} + B^{\lambda\lambda'}.$$  \hspace{1cm} (15)

In this section a model is described to define $T^{\lambda\lambda'}$ and explicit expressions for the generalized polarizabilities and the scattering amplitude are given for the case of a zero-spin nucleus.

3.1 The transition operator

In the original $\Delta$-hole model applied to elastic pion scattering [24, 25, 26] and coherent $\pi^0$ photoproduction [27, 28, 29], the $\Delta$ dynamics in the nuclear medium is dictated by the $\Delta$-hole Green’s function

$$G_{\Delta h}(E) = \left[ E - E_R(E) + \frac{1}{2}i\Gamma(E) - H_{\Delta h} \right]^{-1},$$  \hspace{1cm} (16)

where both the natural free-$\Delta$ width $\Gamma(E)$ and the resonance energy $E_R$ are modified by the effective $\Delta$-nucleus interaction. Such an interaction is modeled by the Hamiltonian $H_{\Delta h}$ which incorporates $\Delta$-propagation, binding and Pauli blocking effects as well as a contribution $W_\pi$ describing intermediate pion propagation in the presence of the nuclear ground state and corresponding to pion multiple scattering. A complex term, the so-called spreading potential $V_{sp}$, is also included in $H_{\Delta h}$ to account for coupling to multi-hole intermediate channels.

The pion absorption process described by the spreading potential is quite complicated and a large part of it is expected to be due to the coupling of the $\Delta$ to the 2p-2h continuum configurations. It has been modeled either using a
phenomenological potential with parameters fitted to the pion scattering data \cite{24, 25, 26} or performing a microscopic calculation of the process \cite{21, 22, 47}.

To reduce the computational effort, a local-density approximation to the medium-modified $\Delta$ propagator \cite{16} has been used in ref. \cite{27} to analyze the pion-nucleus data. The same approximation has been used in the approach of refs. \cite{39, 40} where the $\Delta$-hole Green’s function is written

$$G_{\Delta h}(\rho(\vec{r}), s) = \left[ \sqrt{s - M_{\Delta} + \frac{1}{2} i \tilde{\Gamma}(\rho(\vec{r}), s) - \Sigma_{\Delta}(\rho(\vec{r}), s)} \right]^{-1}.$$ \hspace{1cm} (17)

In eq. (17) $M_{\Delta} = 1238$ MeV is the $\Delta$ mass, while $\tilde{\Gamma}$ and $\Sigma_{\Delta}$ are the Pauli-blocked $\Delta$ width and self-energy, respectively. $\Sigma_{\Delta}$ is in general a non-Hermitean, non-local and energy-dependent operator. However, when evaluated in the local-density approximation, both $\tilde{\Gamma}$ and $\Sigma_{\Delta}$ become single-particle operators depending on the nucleon density $\rho(\vec{r})$ and the Mandelstam variable $s$ for the photon-nucleon system

$$s = M^2 + 2\omega \left( M + \frac{3}{5} \frac{k_F^2}{2M} \right),$$ \hspace{1cm} (18)

where $M$ is the nucleon mass, $\omega$ the photon energy and $k_F$ the (local) Fermi momentum [$k_F^2 = \frac{2}{3} \pi^2 \rho(\vec{r})$]. In eq. (18) an average over the nucleon momentum has been performed within the Fermi gas model. For zero-spin nuclei a spherical nucleon density $\rho(\vec{r})$ will be used.

A many-body expansion in terms of pf and $\Delta h$ excitations and the spin-isospin induced interaction has been proposed in ref. \cite{39} to evaluate $\Sigma_{\Delta}$ in nuclear matter accounting for quasi-elastic corrections, two-body and three-body absorption. We adopt the same analytical expression and the corresponding numerical parametrization presented in ref. \cite{39} which are supported by the microscopic evaluation of the $\Delta$-spreading potential in finite nuclei performed in ref. \cite{47}. Moreover, we include in $\Sigma_{\Delta}$ the $W_\pi$ contribution according to the indications of ref. \cite{23}.

As a consequence of the local-density approximation the part $R_{\Delta}^{\lambda \lambda}$ of the transition operator (13) due to the direct excitation of the $\Delta$ resonance becomes

$$R_{\Delta}^{\lambda \lambda} = \sum_{i=1}^{A} F_{\gamma N \Delta}(\vec{k}', i) G_{\Delta h}(\rho(\vec{r}_i), s) F_{\gamma N \Delta}(\vec{k}, i),$$ \hspace{1cm} (19)

where

$$F_{\gamma N \Delta}(\vec{k}, i) = \frac{f_\pi}{m_\pi} \vec{e}_\lambda \cdot \vec{k} \times \vec{S} \; T_3 \; e^{i\vec{k} \cdot \vec{r}_i}.$$ \hspace{1cm} (20)
is the effective $\gamma N \Delta$ vertex with $f_\gamma = 0.122$ \[4\] and $m_\pi$ the pion mass.

We include a second contribution to the resonant part $R_{\chi}^{\chi}$ coming from the crossed $\Delta$-hole excitation. It is given by

$$R_{\chi \chi}^{\chi} = \sum_{i=1}^{A} F^R_{\chi \chi} G_{\Delta h}(\rho(\vec{r}_i), s') F^R_{\chi \chi} G_{\Delta h}(\vec{k}', i),$$  \hspace{1cm} (21)

where

$$F^R_{\chi \chi}(\vec{k}', i) = \frac{f_\gamma}{m_\pi} \epsilon^*_{\lambda'} \cdot \vec{k}' \times S T_3 \epsilon'^{-i \vec{k}' \cdot \vec{r}_i},$$  \hspace{1cm} (22)

$$s' = M^2 - 2\omega \left( M + \frac{3}{5} \frac{k_F^2}{2M} \right).$$  \hspace{1cm} (23)

Thus

$$R_{\chi \chi}^{\chi} = R_{\chi \chi}^{\chi} + R_{\chi \chi}^{\chi}.$$  \hspace{1cm} (24)

Non-resonant background contributions are due to s-wave pion production and absorption on a single nucleon \[23\] and to a two-photon contact interaction described by a seagull diagram arising from the quadratic term in the vector potential of the non-relativistic interaction Hamiltonian \[18, 48, 9\].

Assuming the Kroll-Ruderman form for pion production as in ref. \[23\], the corresponding transition operator is

$$B_{\chi \chi}^{\chi} = \sum_{i=1}^{A} B(E) \epsilon^*_{\lambda'} \cdot \vec{\sigma}(i) \epsilon_{\lambda} \cdot \vec{\sigma}(i) e^{i(\vec{k} - \vec{k}') \cdot \vec{r}_i},$$  \hspace{1cm} (25)

where $B(E)$ has been derived by a dispersion integral at the energy $E = \sqrt{s}$. Equivalently, $E = \omega_q + \epsilon_q = \sqrt{m_\pi^2 + q^2} + \sqrt{M^2 + q^2}$, with $q$ being the c.m. on-shell momentum in the $\pi N$ channel. One has

$$B(E) = 4\pi \alpha \left( \frac{f}{m_\pi} \right)^2 \int \frac{dq}{(2\pi)^3} \frac{M}{\omega_q \epsilon_q} \left| \frac{v(q^2)}{v(q^2)} \right|^2 \frac{1}{E - (\omega_q' + \epsilon_q') + i\epsilon},$$  \hspace{1cm} (26)

$$\text{Im} B(E) = -2\alpha \left( \frac{f}{m_\pi} \right)^2 \left( \frac{M}{E} \right) q,$$  \hspace{1cm} (27)

with $f^2/4\pi = 0.08$ and $\alpha = e^2/4\pi = 1/137$. The form factor $v(q^2)$ governing the off-shell behaviour is parametrized as in ref. \[23\], with $v(q^2) = (1 + \beta^2/q^2)^{-1}$ and $\beta = 300$ MeV. The two-photon operator includes a term
describing the Thomson scattering by individual nucleons and a term due to two-body exchange currents \[48\]. Both contributions are required to fulfill gauge invariance. However, in ref. \[9\] the exchange contribution was shown to be rather small. Therefore it will be neglected here and the two-photon operator is simply given by

\[
B^{\lambda \lambda'}_{TP} = \frac{1}{M} \bar{e}_\lambda \cdot \bar{e}_{\lambda'} \sum_{i=1}^{A} e_i^2 e^{i(\vec{k}-\vec{k}') \cdot \vec{r}_i},
\]

where \(e_i = \frac{1}{2}[1 + \tau_3(i)]e\) is the nucleon electric charge.

Second order contributions from the one-body current with a nucleon in the intermediate state are rather small in the energy range considered here and will also be neglected. Therefore the total background part of the transition operator is

\[
B^{\lambda \lambda'} = B^{\lambda \lambda'}_{KR} + B^{\lambda \lambda'}_{TP}.
\]

Within the same local-density approximation to the \(\Delta\) propagator, eq. (17), the original \(\Delta\)-hole model of ref. \[23\] is recovered with the following modifications in the transition operator.

\(i\) In the effective \(\gamma N\Delta\) vertex of eq. (20) the coupling \(f_\gamma/m_\pi\) is replaced by \(g_{\gamma N\Delta}/M_\Delta\) and an additional factor \(V\) is introduced, which includes a background contribution describing the pion photoproduction followed by the resonant pion rescattering (Fig. 1). We have

\[
V = 1 - a_B(E)\Sigma_{\pi N}^{\Delta}, \quad a_B(E) = \frac{\bar{g}_{\gamma N\Delta} \sin \phi(E)}{g_{\gamma N\Delta} \Gamma(E)/2},
\]

where \(\Sigma_{\pi N}^{\Delta}\) is the medium-corrected self-energy corresponding to intermediate coupling to the \(\pi N\) channel, \(\bar{g}_{\gamma N\Delta} = 1.03\), \(g_{\gamma N\Delta} = 1.02\) and \(\phi(E)\) is parametrized according to ref. \[23\].

\(ii\) The crossed \(\Delta\)-hole excitation of eq. (21) is replaced by a background photopion production in the resonant channel

\[
R_{\lambda}^{\lambda'} = \sum_{i=1}^{A} F^\dagger_{\gamma N\Delta}(\vec{k}', i) A_B(E) F_{\gamma N\Delta}(\vec{k}, i),
\]

where

\[
A_B(E) = a_B^2(E)\Sigma_{\pi N}^{\Delta}.
\]
{iii} The non-resonant background is simply given by the transition operator
\( B_{K'K}^{\lambda'\lambda} \) of eq. (25) with \( B(E) \rightarrow B'(E) \), which is obtained by introducing under
the integral of eq. (26) a phenomenological energy-dependent factor \( h^2(E) \)
fitted to the total photo-absorption cross section.

### 3.2 Generalized polarizabilities for zero-spin nuclei

The transition operator for photon scattering in the model described above
turns out to be a single-particle operator. When evaluating the corresponding
transition amplitude all matrix elements can be expressed in terms of the
nuclear-matter density. Therefore the nuclear density enter twice, first in the
local-density approximation to the \( \Delta \)-hole propagator and second as a result
of an effective impulse approximation.

Confining the discussion to zero-spin nuclei, only the scalar electric and
magnetic polarizabilities are different from zero. The resonant contribution
to the polarizabilities is given by

\[
P_R^0(EL, EL; k, k) = \left( \frac{f_\gamma}{m_\pi} \right)^2 \frac{2}{9} \tilde{k} \tilde{k}' \ (-)^{L+1} \hat{L} \\
\times \int d\vec{r} \rho(r)(2L + 1) j^2_L(kr) \ [G_{\Delta h}(\rho(r), s) + G_{\Delta h}(\rho(r), s')] ,
\]

(33)

\[
P_R^0(ML, ML; k, k) = \left( \frac{f_\gamma}{m_\pi} \right)^2 \frac{2}{9} \tilde{k} \tilde{k}' \ (-)^{L+1} \hat{L} \\
\times \int d\vec{r} \rho(r) \left[ L j^2_{L+1}(kr) + (L + 1) j^2_L(kr) \right] \\
\times [G_{\Delta h}(\rho(r), s) + G_{\Delta h}(\rho(r), s')] ,
\]

(34)

where \( \tilde{k} \) is the photon energy in the photon-nucleon system \[14\].

The other background contributions are

\[
P_B^0(EL, EL; k, k) = \left[ \frac{1}{2} B(E) + \frac{e^2}{4M} \right] (-)^{L+1} \hat{L} \\
\times \int d\vec{r} \rho(r) \left[ L j^2_{L+1}(kr) + (L + 1) j^2_L(kr) \right] ,
\]

(35)
\[ P_0^B(ML, ML; k, k) = \left[ \frac{1}{2} B(E) + \frac{e^2}{4M} \right] (-)^{L+1} \hat{L} \times \int d\vec{r} \rho(r) (2L + 1) j_{L+1}^2 (kr). \]  

(36)

The corresponding expressions in the local-density approximation to the original \( \Delta \)-hole model of ref. [23] are obtained with the modifications i)–iii) explained in sect. 3.1.

### 3.3 Compton scattering amplitude

The Compton scattering amplitude for zero-spin nuclei in the c.m. frame becomes

\[
T_{\lambda' \lambda} = \frac{1 + \lambda \lambda' \cos \theta}{2} \left\{ F(q^2) \left[ B(E) + \frac{e^2}{2M} \right] + \frac{4}{9} \left( \frac{f_\gamma}{m_\pi} \right)^2 \int d\vec{r} e^{i(\vec{k} - \vec{k}')} \rho(r) \left[ G_{\Delta h}(\rho(r), s) + G_{\Delta h}(\rho(r), s') \right] \right\}, 
\]

(37)

where

\[ F(q^2) = \int d\vec{r} e^{i(\vec{k} - \vec{k}')} \rho(r) \]

is the nuclear form factor with \( \vec{q} = \vec{k} - \vec{k}' \).

In the local-density approximation to the original \( \Delta \)-hole model of ref. [23] discussed in sect. 3.1 the Compton scattering amplitude (37) becomes

\[
T_{\lambda' \lambda} = \frac{1 + \lambda \lambda' \cos \theta}{2} \left\{ F(q^2) B'(E) + \frac{4}{9} \int d\vec{r} e^{i(\vec{k} - \vec{k}')} \rho(r) \right\} \times \left[ \left( \frac{g_\gamma N_\Delta}{M_\Delta} V \right)^2 G_{\Delta h}(\rho(r), s) + \left( \frac{g_\gamma N_\Delta}{M_\Delta} A_B \right)^2 \right]. 
\]

(39)
4 Results

The model is applied to study Compton scattering from zero-spin nuclei such as $^4$He and $^{12}$C. The nuclear-matter density for such nuclei can be identified with the charge density derived from the charge form factor measured with high accuracy in elastic electron scattering.

Recently, new data on $^{12}$C have been accumulated at a scattering angle $\theta = 40^\circ$ in the photon energy range between 200 and 500 MeV [34]. At the highest energy the corresponding momentum transferred to the recoiling target nucleus reaches the value of the first minimum in the charge form factor of $^{12}$C ($q \approx 1.7$ fm$^{-1}$). This kinematical situation makes it possible to test the validity of the local-density approximation to the $\Delta$-hole model.

In fig. 2 the data are shown together with different model calculations. The dotted curve is the result of a full $\Delta$-hole calculation [35], where the residual $\Delta$-hole interaction includes $\pi$ and $\rho$ exchanges as well as short-range correlations simulated by the Landau-Migdal parameter. Including the residual interaction, a good description of the data is obtained at photon energies higher than 250 MeV. The discrepancy between theory and data at lower energies has been ascribed to the importance of background effects [35].

The solid curve is obtained with the scattering amplitude (39) corresponding to the original $\Delta$-hole model [23] treated in the local-density approximation. The adopted nuclear density is derived within a projected-Hartree-Fock approach to the description of the ground state properties of $^{12}$C and accurately accounts for the charge form factor all over the explored range of momenta [49]. The difference between the solid and dotted curves is a combined effect of the local-density approximation and the medium modifications applied to the effective $\gamma N\Delta$ vertex and to the background contributions according to ref. [23] and not included in the calculation of ref. [35].

The dashed curve corresponds to the scattering amplitude (37) with the same nuclear density used for the solid curve. The different background and the contribution of the crossed $\Delta$-hole excitation reduce the size of the peak and shift its position to higher energies.

According to the results of fig. 2 the uncertainty introduced by the local-density approximation in the case of Compton scattering is of the same magnitude as that obtained in coherent pion production [14]. The main difference between the model calculations stems from the different ingredients taken into account in the transition mechanism. In particular, comparison of the
different results with data in fig. 2 confirms the important role of the background at the lower energies. A detailed study of it can be done with the aid of the recent data on $^4$He obtained at Mainz [38]. In the explored photon energy range between 200 and 500 MeV at fixed scattering angle ($\theta_{cm} = 40^\circ$) the momentum transferred to the recoiling target nucleus corresponds to the low-$q$ part of the charge form factor where a gaussian-shaped $\rho(r)$ is a good approximation to the nuclear density. Furthermore, the relatively small variation of $F(q^2)$ in this range gives a better insight into the role of the other theoretical ingredients.

In figs. 3 and 4 the results are shown as obtained in the local-density approximation to the $\Delta$-hole of ref. [23] and with the scattering amplitude (37), respectively. As expected, the background contribution is smooth in both cases as a consequence of the small variation of $F(q^2)$ with $q$ and the weak energy-dependence of the non-resonant amplitudes $B(E)$ and $B'(E)$. However, a much larger background is produced by including the two-photon operator in the scattering amplitude (37) as required by gauge invariance. Its effect is hardly taken into account by the modification $B(E) \rightarrow B'(E)$ introduced in ref. [23] whose limitations were already stressed by the authors themselves. On the other hand, a quite different resonant contribution is provided by the two models. Here the medium effects are substantial in defining the correct position of the peak energy. While these effects were accurately taken into account in ref. [23], they are only in part considered in the calculation with the scattering amplitude (37). In addition, the crossed $\Delta$-hole excitation shifts to higher energies the peak location, also anomalously increasing the resonant-background interference. In any case, up to about 350 MeV both models account for the data satisfactorily, thus giving confidence that a local-density approximation is well suited to by-pass the technical difficulties connected with the full $\Delta$-hole calculation.

For the first time Compton scattering data with polarized photons have become available [37]. The experiment was performed on $^4$He by the LEGS collaboration at Brookhaven in the energy range between 200 and 310 MeV at five different incoming energies. The data for three of them are shown in fig. 5 as a function of the c.m. scattering angle for a photon beam fully polarized in the reaction plane ($\phi = 0^\circ$, $(d\sigma/d\Omega)_\parallel$ in eq. (13)) and perpendicular to the reaction plane ($\phi = 90^\circ$, $(d\sigma/d\Omega)_\perp$ in eq. (14)). Quite similar results are obtained in the local-density approximation to the $\Delta$-hole of ref. [23] (solid curve) and with the scattering amplitude (37) (dashed
curve). The observed behaviour at forward angles is well reproduced, while discrepancies are present in the backward scattering region. This is a well known problem already stressed in ref. [20] and sometimes related to an unsatisfactory treatment of background contributions [19].

With polarized photons it is possible to separate the two structure functions $W_T$ and $W_{TT}$ which are differently related to the resonant and background contributions. This separation is illustrated in fig. 6, where the unpolarized ($W_0$) and polarized ($W_1$) cross sections for elastic Compton scattering off $^4$He are reported. They are given in terms of the two structure functions $W_T$ and $W_{TT}$ as follows:

$$W_0 = \frac{k'}{k} \left( \frac{M_T}{4\pi E_T} \right)^2 W_T, \quad W_1 = \frac{k'}{k} \left( \frac{M_T}{4\pi E_T} \right)^2 W_{TT}. \quad (40)$$

In turn, taking benefit of the local-density approximation the structure of $W_T$ and $W_{TT}$ is quite simple:

$$W_T = \frac{1}{2} (1 + \cos^2 \theta_{cm}) \left[ |r(E)|^2 + |b(E)|^2 \right] + 2 \cos \theta_{cm} Re \left[ r(E)b^*(E) \right], \quad (41)$$

$$W_{TT} = \frac{1}{2} \sin^2 \theta_{cm} \left[ |r(E)|^2 - |b(E)|^2 \right], \quad (42)$$

where $r(E)$ and $b(E)$ represent the model resonant and background amplitudes, respectively. Therefore, according to eqs. (13) and (14), at $\theta_{cm} = 90^\circ$, $(d\sigma/d\Omega)_\parallel \propto |r(E)|^2$ and $(d\sigma/d\Omega)_\perp \propto |b(E)|^2$, so that the two cross sections are uniquely determined by either the resonant or background contributions.

The magnitude of these two terms appears in agreement with data in fig. 5. On the other hand, the angular behaviour of $(d\sigma/d\Omega)_\parallel$ is determined by the background through terms in $\cos^2 \theta_{cm}$ (pure background) and $\cos \theta_{cm}$ (background-resonant interference). At the resonance energy a fair agreement is obtained, while the deficiencies at backward angles for energies below the $\Delta$-resonance region show that some mechanism, whose effects increase with $q$, is lacking there. This is confirmed by looking at $W_0$ in fig. 6, where the forward-backward asymmetry is entirely due to the interference term proportional to $\cos \theta_{cm}$. In the case of $(d\sigma/d\Omega)_\perp$ the role of background and resonant contributions is interchanged. In the forward emisphere the correct behaviour of $(d\sigma/d\Omega)_\perp$ is mainly determined by the form factor $F(q^2)$, but again at backward angles the yield is too low. The role of $F(q^2)$ is better
seen in $W_1$ in fig. 6. The $\sin^2 \theta_{cm}$ dependence of $W_1$ is modulated by $F(q^2)$ with the result that the peak in the angular distribution of $W_1$ is shifted at angles lower than $90^\circ$. In addition, the magnitude of $W_1$ is larger when the background contribution is smaller.

The behaviour of the two structure functions $W_T$ and $W_{TT}$ also determines the photon asymmetry. The peak in the angular distribution at the different photon energies shown in fig. 7 always occurs at angles larger than those of the data as a result of the too low response $W_T$ at backward angles.

\section{Conclusions}

The high quality of recent data on elastic Compton scattering off zero-spin nuclei such as $^4$He and $^{12}$C in the $\Delta$-resonance region and the possibility of disentangling different nuclear responses with polarized photons have made possible a detailed investigation of the interplay between resonant and background contributions in the interpretation of data within the $\Delta$-hole model. A local-density approximation has been shown to give reasonable results also for such a light nucleus as $^4$He, thus giving confidence to be of help when dealing with heavier targets. Under resonance conditions a rather satisfactory agreement between theory and data is gained. Below the $\Delta$-resonance energy discrepancies persist at backward angles which can be ascribed to some lacking mechanism active in the structure function $W_T$.

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Figure captions

Fig. 1. Effective $\gamma N\Delta$ vertex in the background production.

Fig. 2. Differential cross sections for elastic Compton scattering off $^{12}\text{C}$ as a function of the incoming photon energy $\omega$ at a fixed photon angle $\theta = 40^\circ$. The data are taken from ref. [34]. The dotted curve represents the results of a full $\Delta$-hole calculation [35]. The solid curve is the local-density approximation to the $\Delta$-hole of ref. [23]. The dashed curve is obtained with the scattering amplitude (37).

Fig. 3. Differential cross sections for elastic Compton scattering off $^4\text{He}$ as a function of the incoming-photon energy $\omega$ at a fixed c.m. photon angle $\theta_{\text{cm}} = 40^\circ$. The data are taken from ref. [38]. The solid curve is the local-density approximation to the $\Delta$-hole of ref. [23]. The dashed, dot-dashed and dotted lines give the separate contributions of the resonant, background and interference parts, respectively.

Fig. 4. The same as in fig. 3 but for the model described by the scattering amplitude (37).

Fig. 5. Differential cross sections for elastic Compton scattering off $^4\text{He}$ as a function of the c.m. scattering angle for a photon beam fully polarized in the reaction plane ($\phi = 0^\circ$) and perpendicular to the reaction plane ($\phi = 90^\circ$) at the indicated values of the laboratory photon energy. The data are taken from [37]. The solid curve is the local-density approximation to the $\Delta$-hole of ref. [23]. The dashed curve is obtained with the model described by the scattering amplitude (37).

Fig. 6. The unpolarized ($W_0$) and polarized ($W_1$) cross sections for elastic Compton scattering off $^4\text{He}$ as a function of the c.m. scattering angle at the indicated values of the laboratory photon energy. The data are taken from [37]. Solid and dashed curves as in fig. 5.

Fig. 7. The photon asymmetry for elastic Compton scattering off $^4\text{He}$ as a function of the c.m. scattering angle at the indicated values of the laboratory photon energy. The data are taken from [37]. Solid and dashed curves as in fig. 5.
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