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Joint Institute for Nuclear Research, Dubna

National Center for Particle and High Energy Physics
of the Belarussian State University, Minsk, Belarus

B. Stepanov Institute of Physics of the National Academy
of Sciences of Belarus, Minsk, Belarus

Institute of Mechanics of Metal Polymer Systems
of the National Academy of Sciences of Belarus, Gomel, Belarus

F. Skorina Gomel State University, Gomel, Belarus

P. Sukhoi Gomel State Technical University, Gomel, Belarus

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ON DEUTERON COMPTON SCATTERING IN A POTENTIAL MODEL

M.I. Levchuk[*]

B.I. Stepanov Institute of Physics, Minsk, Belarus

A.I. L'vov[†]

P.N. Lebedev Physical Institute, Moscow, Russia

Abstract

We revise our calculation [Nucl. Phys. A **674**, 449 (2000)] of deuteron Compton scattering below pion photoproduction threshold which is based on a nonrelativistic one-boson-exchange model of NN-interaction (the Bonn OBEPR potential) and related meson-exchange currents and seagulls. A bug in the computer code for evaluation of the Δ-isobar contribution to the so-called resonance amplitude was fixed. Also, effects of NN-rescattering in the intermediate state have been evaluated self-consistently, with the same NN-potential, thus avoiding approximations used previously and leading to some problems. Treating the isospin-averaged electric and magnetic polarizabilities of the nucleon, α_N and β_N, as free parameters, we fit all available data on the differential cross section of $\gamma d \to \gamma d$ and obtain

$\alpha_N + \beta_N = [16.9 \pm 1.5\ (\text{stat.} + \text{syst.})] \times 10^{-4}\ \text{fm}^3$,

$\alpha_N - \beta_N = [10.6 \pm 1.7\ (\text{stat.} + \text{syst.})] \times 10^{-4}\ \text{fm}^3$.

For the first time all the data, including those at the highest energies and backward angles, have been satisfactory described (with $\chi^2/n \simeq 1.2$).

Introduction

A possibility to measure electromagnetic polarizabilities of the neutron in the reaction of elastic Compton scattering on the deuteron [1] was considered during last years in quite a few theoretical works [2, 3, 4, 5, 6, 7, 8, 9, 10], in which

*E-mail: levchuk@dragon.bas-net.by

†E-mail: lvov@pluton.lpi.troitsk.ru

different methods of taking into account effects of the nuclear environment have been tried. Part of these analyses [3, 2, 4, 5] relied on realistic, high-quality phenomenological nucleon-nucleon potentials. In other approaches [6, 7, 8, 9, 10], methods of effective field theories, including chiral perturbation theory, have been applied.

In our opinion, among all of these investigations, our work [5] still remains the most complete and accurate (up to a bug in the computer code, see below). The aim of the present talk is to explain subtleties and advantages of the approach of Ref. [5], in which the amplitude of elastic γd-scattering is found in a physical model of NN-interaction represented by the Bonn one-boson-exchange (OBE) potential, with the corresponding realistic wave functions of the np-system both in the discrete and continuous spectrum, with the use of meson-exchange currents and seagulls (effective two-photon vertices) consistent with the given OBE NN-interaction, with inclusion of main relativistic corrections, corrections for retardation, effects of the Δ-isobar excitation, etc. Many of the effects that we consider sizeable and deliberately take into account have been neglected in works of other authors.

Generally, the dipole electric and magnetic polarizabilities of the nucleon, α and β, are structure parameters characterizing the particle's ability to get induced electric and magnetic dipole moments in external soft electromagnetic fields. These parameters enter the low-energy expansion of the nucleon Compton scattering amplitude to second order in the photon energy [11]. The proton polarizabilities have been successfully measured in dedicated experiments on low-energy γp-scattering in 1960–2001. The world-average of the latest measurements is quoted as [12]

$$\alpha_p = 12.0 \pm 0.6, \quad \beta_p = 1.9 \pm 0.5 \tag{1}$$

in the units of 10^{-4} fm^3 used hereinafter for the polarizabilities.

Similar measurements of γn-scattering can be done with neutrons weekly bound in deuterons. The deuteron is clearly very good as a neutron container owing to minimal distortions introduced by nearby protons and owing to a good theoretical control over the distortion effects.

Our analysis of the reaction of low-energy elastic Compton scattering on the deuteron, in which a coherent sum of γp- and γn-scattering amplitudes is probed, is carried out using a nonrelativistic diagrammatic approach with a nonrelativistic version OBEPR of the Bonn OBE-potential [13, 14]. This choice of the NN-interaction model is motivated by the following:

i) it is well suited for the used nonrelativistic formalism in the momentum representation;

ii) the OBEPR version provides a rather good description of the deuteron binding energy and the NN-scattering amplitude at all energies below pion threshold;

iii) the OBEPR potential suggests a simple diagrammatic picture of the strong interaction via meson exchanges from which meson-exchange currents (MEC) and seagulls consistent with the potential can be straightforwardly constructed;

iv) those MEC as well as a related technique of loop calculations have already been carefully tested [15] in applications to the reaction of deuteron photodisintegration, including polarization observables.

Actually, deuteron-structure dependent effects in low-energy γd-scattering are strongly dominated by a well-known long-ranged single-pion exchange (see below). Moreover, electromagnetic interactions related with a short-range part of the NN-interaction (the latter is parameterized by heavy-meson exchanges and form factors in the Bonn OBE potential) are mostly fixed at low energies in a unique way by gauge invariance and its consequence — Siegert theorem. For this reason, a model dependence related with unreliably known features of the short-range NN-interaction and with therefore somewhat arbitrary choice of the NN-potential is not large provided the used potential includes the long-ranged single-pion exchange and the short-range part is adjusted to give an accurate description of the binding energy and the NN-scattering amplitude. In a forthcoming work we plan to explicitly demonstrate this fact that is quite important for the problem of measuring polarizabilities of the neutron and studying the model dependence.

Having fixed the strong and electromagnetic interactions in the specific Bonn model, we take into account a full set of diagrams (up to four loops!) thus guaranteeing the gauge invariance of the resulting amplitude of $\gamma d \to \gamma d$. Beyond that additional two-body contributions due to Δ-isobar excitation and retardation in pion-exchange diagrams are added. One-photon interaction of the free nucleon is considered together with the important spin-orbital term which is the most essential relativistic correction in the considered problem. Two-photon structure of the free nucleon is accounted through its dipole electric and magnetic polarizabilities (also together with a relativistic correction) describing effects to second order in the photon energy. Important higher-order corrections are also included using results of phenomenological calculations through dispersion relations for spin, quadrupole, etc polarizabilities.

In the previous version of the calculation we met a problem with a description of the Saskatoon laboratory data [16] on γd-scattering at 94 MeV. Recently we have found that this failure is caused by a sign mistake in the computer code for the Δ-isobar contribution to the two-body electromagnetic current in crossed diagrams which resulted in a noticeable underestimation of the differential cross section at backward angles and a strong shift of extracted polarizabilities of the nucleon. After correction of this mistake we arrive at a very satisfactory agreement with all the data.

Another shortcoming of the previous calculation was in using a too poor approximation for four-loop diagrams with two meson-exchange currents and intermediate NN-rescattering. Actually, in order to facilitate four-loop computations, a simplified off-shell NN-rescattering amplitude was used, namely the amplitude found with a separable NN-potential [17] built as a truncation of another, the Paris potential. As a result of using different NN-potentials in different pieces of the whole amplitude, some mismatch appeared between the resonance and seagull parts at all energies, so that gauge invariance was not exactly maintained and the corresponding low-energy theorem was also not exactly fulfilled. For example, at zero energy and the forward scattering angle we have got the spin-averaged amplitude of deuteron Compton scattering to be -0.47 (in the units of e^2/M, M is the nucleon mass) instead of the correct value of -0.50. With the advent of faster computers, we are capable today to do four-loop calculations directly, thus avoiding the Paris-potential separable approximation and using the NN-rescattering amplitude consistently as it is in the Bonn model. In the presented now results, the mentioned mismatch is cured and the total Compton scattering amplitude obeys nicely the low-energy theorem within the accuracy of numerical calculations.

It is worth to mention that problems of that sort are completely beyond the scope of recent calculations [8, 9, 10] done in the framework of Effective Field Theories because the intermediate NN-interaction is included there only up to the leading term, which is the one-pion exchange plus a constant, and any further part of the NN-potential is counted as an effect of higher-order in the generic momentum Q or energy scale ϵ and thus is considered negligible. Generally, this is a weak point of all evaluations in EFTs which use the procedure of truncation of diagrams (up to $\mathcal{O}(\epsilon^3)$ in [10] or $\mathcal{O}(Q^4)$ in [8, 9]) according to power counting rules and thus actually use different (differently truncated) NN-interactions for finding the nuclear brackets (wave functions) and for evaluating the kernel. The truncations lead to violation of gauge invariance of the resulting amplitude and to model dependence which is numerical quite large.

In this respect our approach that takes the NN-interaction the same everywhere is clearly advantageous.

Before going to our final results, we briefly describe the main steps of the present calculation which mostly repeat those in Ref. [5].

Kinematics

Keeping in mind some future developments, we choose now to perform calculations of the γd-scattering amplitude T in the Breit frame that is most symmetric for the considered problem. In this frame, 3-momenta and total energies of the initial and final deuterons are related with the photon momenta $\mathbf{k}$, $\mathbf{k}'$ as

$$\mathbf{p}_d = -\mathbf{p}'_d = \frac{\mathbf{k}' - \mathbf{k}}{2}, \quad E_d = E'_d = \sqrt{M_d^2 - \frac{t}{4}}. \tag{2}$$

Here $M_d = 2M - E_b$ is the deuteron mass, $M = 938.9$ MeV is the (average) mass of the nucleon, $E_b = 2.2246$ MeV is the deuteron binding energy, and $t = -(\mathbf{k}' - \mathbf{k})^2$.

Comparing with experimental data, we calculate the differential cross section $d\sigma/d\Omega$ of γd-scattering in the c.m. frame at given photon energy E_γ in the lab frame and the c.m. scattering angle Θ using

$$\frac{d\sigma}{d\Omega} = \left(\frac{E_d}{4\pi W}\right)^2 \frac{1}{6} \sum_{\text{spins}} |T|^2 \tag{3}$$

where $W = \sqrt{M_d^2 + 2M_d E_\gamma}$ is the total energy of the γd-system.

Resonance amplitude

In accordance with splitting the electromagnetic interaction to pieces of order $\mathcal{O}(e)$ and $\mathcal{O}(e^2)$, the full Compton scattering amplitude $\mathcal{O}(e^2)$ consists of the two pieces too, the so-called resonance R and seagull S amplitudes. The amplitude R corresponds to the two-step process of photon absorption, leading to the excitation of low-lying (two-nucleon) states of the deuteron, followed by photon emission (plus the crossed term) as shown schematically in Fig. 1. Since the propagator $G(E) = (E - H + i0)^{-1}$ of interacting nucleons can be written in the form

$$G(E) = G_0(E) + G_0(E)\, T_{NN}(E)\, G_0(E), \tag{4}$$

where $G_0(E) = (E - H_0 + i0)^{-1}$ is the propagator of free nucleons, the amplitude R can be written respectively as the sum of terms without and with NN-rescattering in the intermediate state (Fig. 1a and 1b).

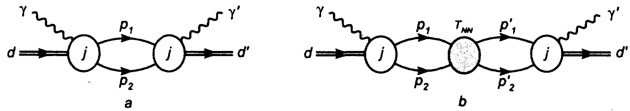

Figure 1. The resonance diagrams without (a) and with (b) NN-rescattering. Crossed diagrams are not shown.

The symbol $\boldsymbol{j}$ standing in Figs. 1a and b denotes the electromagnetic current that has to be consistent with the nuclear Hamiltonian H in order to ensure gauge invariance:

$$\left[j_0^{[1]}(\mathbf{x}), H\right] = -i\boldsymbol{\nabla}\cdot\boldsymbol{j}(\mathbf{x}) = -i\boldsymbol{\nabla}\cdot\left(\boldsymbol{j}^{[1]}(\mathbf{x}) + \boldsymbol{j}^{[2]}(\mathbf{x})\right). \qquad (5)$$

Here $\boldsymbol{j}^{[1]}(\mathbf{x})$ and $\boldsymbol{j}^{[2]}(\mathbf{x})$ are the one-body and two-body currents, respectively. It is assumed here that in absence of the energy dependence of the Hamiltonian the charge density $j_0(x)$ is not affected by meson exchanges and therefore coincides with the one-body charge density of free nucleons $i = 1, 2$:

$$j_0(\mathbf{x}) = j_0^{[1]}(\mathbf{x}) = \sum_{i=1,2} eZ_i\delta(\mathbf{x} - \mathbf{r}_i), \qquad Z_i = \frac{1+\tau_i^z}{2}. \qquad (6)$$

We take a nonrelativistic energy-independent nuclear Hamiltonian $H = H_0 + V$ as given by the Bonn OBEPR model [13, 14]. Note that there are three versions of that model which we label OBEPR(A), OBEPR(B) and simply OBEPR. We always give our predictions obtained with the OBEPR version unless other is stated explicitly. The NN-potential V is represented by the sum of one-boson exchanges and has the form

$$V(\mathbf{p}_1', \mathbf{p}_2'; \mathbf{p}_1, \mathbf{p}_2) = \sum_{\alpha=\pi,\eta,\delta,\sigma,\omega,\rho} V^\alpha(\mathbf{p}_1', \mathbf{p}_2'; \mathbf{p}_1, \mathbf{p}_2). \qquad (7)$$

Here $\mathbf{p}_i$ and $\mathbf{p}_i'$ are the initial and final momenta of the i-th nucleon subject to the constraint $\mathbf{p}_1 + \mathbf{p}_2 = \mathbf{p}_1' + \mathbf{p}_2'$, and V^α are potentials stemming from exchanges with the specified mesons $\alpha = \pi$, η, δ (which is $a_0(980)$ in the modern notation), σ, ω, and ρ. Masses of the mesons are taken as they were then known experimentally. Tuning coupling constants $g_{\alpha NN}$, the tensor-to-vector ratio $\kappa_{\alpha NN}$ in the case of the vector mesons ω and ρ, and cutoff parameters

Diagrams $3a$ and b include a block with the Δ excitation. It is well known from the study of deuteron photodisintegration that the tail of the Δ isobar manifests itself even below pion threshold giving quite a visible contribution to observables [15, 18]. As for deuteron Compton scattering, the contribution of the pure transverse current $j^{[2]\pi\Delta}$ to the differential cross section was considered for the first time by Weyrauch and Arenhövel [19] [1]. They found a noticeable effect of Δ above 80 MeV and backward angles leading to a sizable increase in $d\sigma/d\Omega$.

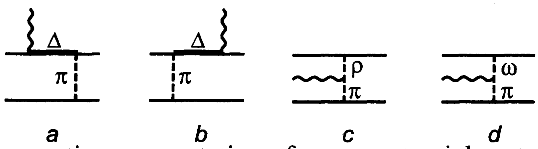

Figure 3. Diagrammatic representation of a non-potential part of the two-body current $j^{[2]}$.

In our previous work [5] we had also considered this part of the current $j^{[2]}$ but wrongly found its contribution to the differential cross section of γd-scattering very small and negative. As we discovered recently, that was a result of a sign error in the code evaluating the Δ term in the crossed diagram. After fixing the error we now obtain agreement with the results of Weyrauch and Arenhövel (see the next section). There are also diagrams similar to $3a$ and b with the intermediate ρ meson but their contribution is negligible at energies under consideration.

The considered Δ effects, though very sizable, are neglected in EFT calculations [8, 9, 10] because they are counted there as being of higher order.

Another type of non-potential two-body currents is shown in Figs. $3c$ and d. They include blocks with the $\gamma\pi\omega$ and $\gamma\pi\rho$ vertices. The currents $j^{[2]\gamma\pi\omega}$ and $j^{[2]\gamma\pi\rho}$ are pure transverse so that their form is not determined by the NN-potential and the continuity equation (5). They have to be determined from diagrams. Explicit expressions for them can be found in Ref. [20]. Diagrams $3c$ and d are usually taken into account, for example, in evaluations of the deuteron magnetic moment [20], deuteron photodisintegration [15] and deuteron electrodisintegration [21]. In Ref. [5] we disregarded these currents, but now we include them. Their contribution is not large but still leads to the 1% and 3%

[1] Since the deuteron has isospin $I = 0$ and therefore transitions $d \to N\Delta$ and $N\Delta \to d$ are forbidden, only one of two diagrams $3a$ or b contributes to the reaction amplitude. For example, only the diagram $3a$ contributes in case of the transition $\langle np|j^{[2]\pi\Delta}|\gamma d\rangle$.

increase in the differential cross section of γd-scattering at forward and backward angles at 100 MeV, respectively.

The total two-body current $\boldsymbol{j}^{[2]}$ we use now consists of contributions of diagrams in Figs. 2 and 3 and it is exactly the same as used in Ref. [15] in studies of deuteron photodisintegration. In that work a good agreement with data on the total and differential cross sections, the photon asymmetry and neutron polarization was achieved at energies below pion photoproduction threshold. This is a good evidence for the correctness of the current $\boldsymbol{j}^{[2]}$ that also supports that our present calculation of the resonance amplitude R is reliable and complete.

Evaluating effects of NN-rescattering in the resonance amplitude R (Fig. 1b), we need the NN-scattering amplitude T_{NN} fully off-shell. It is found from the Lippmann-Schwinger equation for the corresponding NN-potential.

Seagull amplitude

The seagull part of the amplitude of γd-scattering corresponds to processes in which photon absorption and photon emission happens at indistinguishable time moments as seen at the considered energy scale. Accordingly, its energy dependence is very smooth: it is either a constant or a polynomial in the photon energy and it has no absorptive part.

As in the case of the electromagnetic current, the seagull operator S obeys a relation that follows from the current conservation and gauge invariance. For the energy-independent seagull

$$\Big[j_0(\mathbf{x}), j_l(\mathbf{y})\Big] = i\,\frac{\partial S_{kl}(\mathbf{x},\mathbf{y})}{\partial x_k}. \tag{9}$$

Generally, the operator S can be split into one-body and two-body parts, $S^{[1]}$ and $S^{[2]}$. Diagrammatic representation of $S^{[1]}$ is shown in Fig. 4. Polarizabilities of the nucleon just belong to $S^{[1]}$.

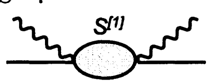

Figure 4. Diagrammatic representation of the one-body seagull $S^{[1]}$.

Before isospin averaging, the one-body seagull operator $S^{[1]}$ has the following form valid up to and including 4th order in the photon energy [5]:

$$\epsilon'^{*\mu}\epsilon^{\nu}S^{[1]}_{\mu\nu}(-k',k) = -\frac{e^2Z^2}{m}\,\boldsymbol{\epsilon}\cdot\boldsymbol{\epsilon}'^{*} + \frac{e^2Z}{4m^2}(Z+2\kappa)(\omega+\omega')\,i\boldsymbol{\sigma}\cdot\boldsymbol{\epsilon}'^{*}\times\boldsymbol{\epsilon}$$
$$+\,4\pi\omega\omega'(\alpha+\delta\alpha_0)\,\boldsymbol{\epsilon}\cdot\boldsymbol{\epsilon}'^{*} + 4\pi\omega\omega'\beta\,\mathbf{s}\cdot\mathbf{s}'^{*} + \epsilon'^{*\mu}\epsilon^{\nu}\delta S^{[1]}_{\mu\nu}(-k',k) \tag{10}$$

where

$$\delta\alpha_0 = -\frac{e^2}{4\pi}\frac{\kappa^2 + Z\kappa}{4m^3} = \begin{cases} -0.85, & \text{proton,} \\ -0.62, & \text{neutron} \end{cases} \tag{11}$$

is a relativistic correction. The first term in the r.h.s. of Eq. (10) is the well-known Thomson amplitude (a constant) and the second one (linear in the photon energy) comes from the spin-orbit interaction. Second-order terms in the photon energy are given by α and β, the nucleon dipole electric and magnetic polarizabilities which we are interested in. Higher-order terms are denoted by $\delta S^{[1]}_{\mu\nu}$. Up to and including fourth order, they are described by 8 parameters — four spin polarizabilities, two quadrupole polarizabilities and two so-called dispersion polarizabilities [22]. In the present evaluation we use the same numerical values of these parameters as in [5], found through dispersion relations.

The sum of the dipole polarizabilities is fixed by the Baldin sum rule [1] which gives in particular

$$\alpha_N + \beta_N = \int_{\omega_{\rm thr}}^{\infty} \sigma_{\rm tot}(\omega)\,\frac{d\omega}{2\pi^2\omega^2}. \tag{12}$$

Here $\sigma_{\rm tot}(\omega) = \frac{1}{2}[\sigma^p_{\rm tot}(\omega) + \sigma^n_{\rm tot}(\omega)]$ is the isospin averaged total photoabsorption cross section for the nucleon. The dispersion integral was evaluated in Ref. [5] leading to $\alpha_N + \beta_N = 14.6 \pm 0.5$. We can, however, consider the parameters α_N and β_N as unconstrained, fit them using experimental data on γd-scattering, and thus check the Baldin sum rule prediction.

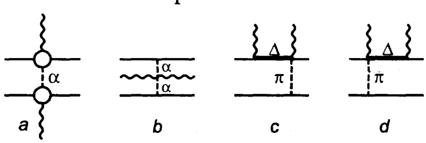

Figure 5. Diagrammatic representation of the two-body seagull $S^{[2]}$. Gray circles mean the sum of the vertices $\gamma N \to \alpha N$ as in Fig. 2 and the vertices $\gamma N \to \pi N$ as in Fig. 3*a* and *b*.

The two-body seagull $S^{[2]}$ consistent with the OBEPR potential and MEC can be obtained from the diagrammatic representation shown in Fig. 5. Explicit expressions for $S^{[2]}$ are given in Ref. [5].

Results and discussion

We discuss now some results obtained in the revised version of our calculation. First we check how the zero-energy limit of the forward spin-averaged

γd-scattering amplitude is saturated. Contributions of the diagrams $1a$ and $1b$ found with the OBEPR-potential wave functions and the OBEPR-potential NN-rescattering are numerically equal to $+0.84$ and -0.10, respectively (in the units of e^2/M). The one-body and two-body seagulls add -1 and -0.24, respectively. All together they give the correct Thomson limit of the amplitude, -0.50, without a previous 6%-mismatch [5] mentioned in Introduction.

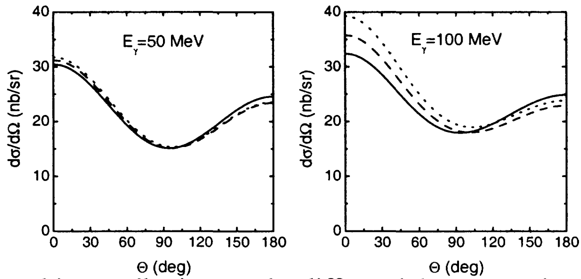

Figure 6. Spin-orbit contributions to the differential cross section. Dotted lines: all SO terms are turned off. Dashed and solid lines: SO terms from $\boldsymbol{j}^{[1]}$, Eq. (8) and $S^{[1]}$, Eq. (10) are successively added. Dipole polarizabilities are set to zero.

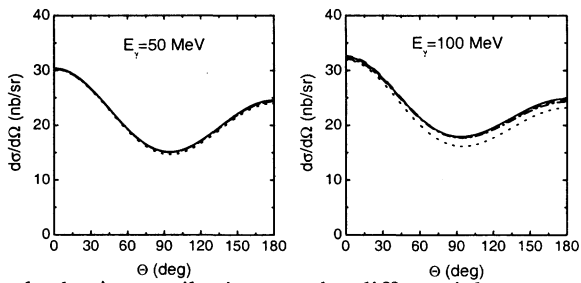

Figure 7. Two-body Δ contributions to the differential cross section. Dotted lines: all the Δ-contributions are turned off. Dashes, dash-dotted and solid lines: Δ in $j^{[2]}$, in the seagull diagram $5a$, and in the seagull diagrams $5c$-d is successively added. Dipole polarizabilities are set to zero.

An influence of the spin-orbit interaction on $d\sigma/d\Omega$ is shown in Fig. 6. As one can expect, it increases with the energy and becomes very large at 100 MeV. We have again to stress that some recent calculations [4, 7, 8, 10] neglect these important contributions or take them only through $S^{[1]}$. In this respect it is instructive to notice that SO terms in both Eq. (8) and Eq. (10) are equally important at 100 MeV, see Fig. 6.

Figure 7 shows the role of two-body Δ-contributions in the differential cross section. They all are quite negligible at 50 MeV. But this is not the case at 100 MeV, especially at backward angles. The main effect stems from the Δ-term in the resonance amplitude (see Figs. 3*a* and *b*) which leads to a noticeable increase in the differential cross section. In essence, this result was obtained long ago in Ref. [19].

Now we consider fits of available data on the differential cross section [16, 23, 24] and determinations of the isospin-averaged nucleon dipole polarizabilities α_N and β_N. Taking both α_N and β_N as free parameters, we obtain from two-parameter fits:

$$\alpha_N + \beta_N = 16.9 \pm 1.5, \quad \alpha_N - \beta_N = 10.5 \pm 1.7, \quad \chi^2/n = 33/(29-2), \tag{13}$$
$$\alpha_N + \beta_N = 17.0 \pm 1.5, \quad \alpha_N - \beta_N = 10.7 \pm 1.7, \quad \chi^2/n = 33/(29-2), \tag{14}$$
$$\alpha_N + \beta_N = 18.8 \pm 1.6, \quad \alpha_N - \beta_N = 12.6 \pm 1.8, \quad \chi^2/n = 35/(29-2) \tag{15}$$

for the OBEPR, OBEPR(A), and OBEPR(B) potentials, respectively. Here both statistical and systematic uncertainties have been combined, the later being taken into account through a rescaling of measured cross sections within their normalization uncertainties. One can see that the sum of the extracted polarizabilities is slightly above the Baldin sum rule prediction, 14.6. The difference $\alpha_N - \beta_N$ is now in good agreement with the proton's value of 10.1 ± 0.6 that means the absence of a visible isovector component in the nucleon dipole polarizabilities α and β.

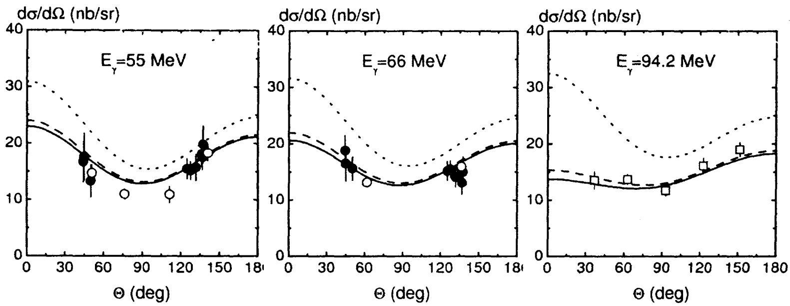

Figure 8. The differential cross section of γd-scattering at three selected energies. Dotted lines: the nucleon dipole polarizabilities are turned off. Dashed lines: $\alpha_N + \beta_N = 14.6$ and $\alpha_N - \beta_N = 9.4$. Solid lines: $\alpha_N + \beta_N = 16.9$ and $\alpha_N - \beta_N = 10.5$. Data are from [23] (○), [16] (□) and [24] (•).

The extracted values of α_N and β_N are very close for the OBEPR and

OBEPR(A) potentials but some different in the case of OBEPR(B). Note, however, that there are reasons to consider the OBEPR(B) potential less accurate than the other two because our analysis of deuteron photodisintegration [15] has already showed that the OBEPR(B)-interaction does not lead to a satisfactory description of data on that reaction. Therefore, in our opinion the results obtained with OBEPR and OBEPR(A) are physically more justified than those with OBEPR(B).

Note also that this subtle difference shows that the task of extraction of the nucleon polarizabilities with one or another NN-potential requires a very careful cross-check of the theory through analysis of observables both in γd-scattering and photodisintegration.

If $\alpha_N + \beta_N$ is fixed to be 14.6 (Baldin sum rule), we get from such a one-parametric fit:

$$\alpha_N - \beta_N = 9.4 \pm 1.6, \quad \chi^2/n = 35/(29-1), \tag{16}$$
$$\alpha_N - \beta_N = 9.5 \pm 1.6, \quad \chi^2/n = 35/(29-1), \tag{17}$$
$$\alpha_N - \beta_N = 10.3 \pm 1.7, \quad \chi^2/n = 44/(29-1) \tag{18}$$

for OBEPR, OBEPR(A), and OBEPR(B), respectively. Again both statistical and systematic uncertainties have been taken into account.

An inspection of Fig. 8 shows that two points from the the Illinois experiment [23] at 55 MeV lie well below theoretical curves. When these two points are excluded, the two-parameter fit of the reduced data set leads to a nice χ^2/n with essentially the same polarizabilities. For instance, with the OBEPR potential we get

$$\alpha_N + \beta_N = 16.9 \pm 1.5, \quad \alpha_N - \beta_N = 10.8 \pm 1.7, \quad \chi^2/n = 25/(27-2). \tag{19}$$

A conclusion is that our revised model gives for the first time a good description of all available data on γd-scattering and thus it is well suited for theoretical interpretation of new, soon expected data from MAX-lab at Lund [25].

Our final comment concerns further developments. It is desirable to perform computations using contemporary high-quality NN-potentials like CD-Bonn, Nijmegen or potentials created recently with the use of EFT, which achieve a very good $\chi^2/n \approx 1$ in description of data for the NN-system. One problem in doing that is a construction of two-body MECs and seagulls consistently with the potential used. Here the following strategy can be applied. All the considered potentials V contain explicitly the one-pion exchange V^π which gives a dominating contribution to $j^{[2]}$ and $S^{[2]}$. The rest part of the potential $\tilde{V} = V - V^{\pi}$ is so short-ranged that the related MEC and seagulls at energies and momenta characteristic for the dynamics of γd-scattering can be replaced by a Siegert-like operators which can be found in the momentum space for any potential through the minimal substitution, see e.g. [26]. The corresponding calculations are now underway.

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