Integral equation calculations for the photodisintegration process ^{4}He$(\gamma,n)^3$He

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Results obtained by solving Alt-Grassberger-Sandhas (AGS)-type integral equations for the photodisintegration of ^{4}He, employing the Malfliet-Tjon potential, are compared with the latest experimental data. Good agreement between theory and experiment is found in electric dipole approximation for the total cross section, but the differential cross sections differ at higher energies. This discrepancy is reduced, but not fully removed by taking into account the electric quadrupole contributions. In order to get some feeling for the sensitivity to the underlying potential, we also show calculations based on the Yamaguchi potential. They differ from the Malfliet-Tjon results in a way which resembles the trends known from triton photodisintegration. [S0556-2813(96)06305-4]

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I. INTRODUCTION

The photodisintegration of ^{4}He into $n+{}^3$He or $p+{}^3$H has for a long time been a rather controversial topic of few-body physics. Early data appeared consistent with the picture of a giant dipole resonance at low energies, a picture also supported by shell model [1] and resonating group [2] calculations. Generalizing the four-nucleon Alt-Grassberger-Sandhas (AGS) formalism [3], exact integral equations have been derived by Casel and Sandhas for the ^{4}He$(\gamma,n)^3$He photodisintegration amplitude which led to a completely different prediction [4,5]. The corresponding calculations, in fact, indicated a flat nonresonant behavior.

When these results were first presented, they were in flagrant disagreement with experiment (and all other theoretical attempts). Almost at the same time, however, new data were published for the total cross section of ^{4}He$(\gamma,n)^3$He which showed a similar flattening [6]. In a report on this development the authors of [4], therefore, came to the conclusion that the integral equation approach, despite some drastic approximations unavoidable in those days, represents a rather reliable tool for treating the problem [7].

Some discrepancies, however, were still evident. At low energies the levelling off of the theoretical curve appeared to be less pronounced than indicated by the data, and at higher energies the theoretical values lay above the experimental ones. This suggested the replacement of the simple separable potentials with Gaussian form factors, employed in the early calculations [4,5,7], by the semirealistic Malfliet-Tjon (MT I+III) potential [8]. The results presented at the Kalinin (now Tver) few-body workshop [9] showed, indeed, an additional flattening at low energies and the desired lowering at higher energies. In these calculations the ^{3}He and ^{4}He bound states entering the plane-wave (Born) amplitude were determined by solving the homogenous three- and four-nucleon AGS equations, while the final-state interaction was simply taken over from the previous separable-potential calculations.

In the present investigation we repeat the bound-state calculations of [9] with higher accuracy, and solve the full photodisintegration integral equation for energies below the three-body break-up threshold, employing in both cases the MT I+III potential. In contrast to [9], the final state interaction, hence, is consistently taken into account. The results obtained in this way are in good agreement with most recent low-energy data [6,10,11]. At higher energies such an agreement was achieved from the very beginning [5,12]. Our present, considerably improved calculations do not only confirm this trend, but coincide rather well with the newest data [14].

Let us add some technical remarks. The W-matrix method [15] and the energy-dependent pole expansion (EDPE) [16] were employed in order to reduce the original three- and four-body relations to (one-dimensional) effective two-body equations. In the pure nuclear case these two approximations have been demonstrated to lead to very accurate results [17–19].

There were good reasons for choosing just the MT I+III potential. Being comparatively realistic, this potential is nonetheless simple enough to be treated without too much effort. An additional advantage is that the corresponding binding energies lie comparatively close to the experimental values, a property of relevance when taking into account meson exchange currents via Siegert's theorem [20]. Moreover, for triton photodisintegration it has been shown that in the energy region below the three-body breakup threshold the Malfliet-Tjon results are rather similar to the ones obtained for the Paris potential [21]. We, therefore, expect the MT calculations to simulate the Paris-potential treatment also in the present ^{4}He case.

For triton photodisintegration a considerable potential dependence was found. The results obtained there for the Yamaguchi potential lie clearly below the MT curves [21]. To see whether this is a generally valid feature, we have performed a Yamaguchi calculation also in the present ^{4}He

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case and found a similarly lower curve, in fact too low as compared to the data.

The previous and the present calculations demonstrate the necessity of incorporating meson exchange currents via Siegert's theorem [20]. At low energies there is quite a difference between the plane-wave (Born) approximation and the full solution of the integral equation, by which the final-state interaction is taken into account. Only with both these contributions is good agreement between theory and experiment achieved.

At higher energies the plane-wave approximation appears justified. The $E2$ contributions moreover turn out to be small in the total cross section. But they are important in the differential cross section, due to the interference between the $E1$ and $E2$ amplitudes. The incorporation of meson exchange currents remains essential.

II. FORMALISM

The properly antisymmetrized amplitude for the photodisintegration of the ^{4}He bound state $|\psi_{\rm IV}\rangle$ into a three-body bound state $|\psi_{\rm III}\rangle$ and a nucleon of relative momentum $|\mathbf{q}\rangle$ is given, in plane-wave (Born) approximation, by

$$B^{\lambda}(\mathbf{q})=\langle \mathbf{q}|\langle \psi_{\rm III}|H^{\lambda}_{\rm em}|\psi_{\rm IV}\rangle. \tag{1}$$

Replacing the channel state $\langle \mathbf{q}|\langle \psi_{\rm III}|$ by the corresponding scattering state $^{(-)}\langle \mathbf{q};\psi_{\rm III}|$, we get the full amplitude

$$M^{\lambda}(\mathbf{q})=2\,^{(-)}\langle \mathbf{q};\psi_{\rm III}|H^{\lambda}_{\rm em}|\psi_{\rm IV}\rangle. \tag{2}$$

It is generally accepted that at low energies the process under consideration takes place primarily via electric dipole transition. Assuming pointlike charges, the $E1$ operator is given by

$$H^{\lambda(1)\prime}_{\rm em}=-\frac{e}{mc}\sum_{j=1}^{4}\frac{1+\tau_z^j}{2}\hat{\epsilon}_\lambda\cdot\mathbf{p}_j\,, \tag{3}$$

where $\hat{\epsilon}_\lambda$ denotes one of the two polarization directions of the incident photon. Applying Siegert's theorem [20], Eq. (3) is replaced by

$$H^{\lambda(1)}_{\rm em}=-\frac{ie}{\hbar c}(E_f-E_i)\sum_{j=1}^{4}\frac{1+\tau_z^j}{2}\hat{\epsilon}_\lambda\cdot\mathbf{x}_j\,, \tag{4}$$

with (E_f-E_i) being the difference of the final-state energy and the binding energy of $|\psi_{\rm IV}\rangle$ in the initial state. After going over to Jacobi coordinates, Eqs. (3) and (4) each become a sum of three terms, two of them acting within $|\psi_{\rm III}\rangle$, the third one depending only on the relative momentum $\mathbf{q}$ or the corresponding coordinate $\mathbf{r}$ between the nucleon and the center of mass of $|\psi_{\rm III}\rangle$. Since we employ throughout the following s wave projected potentials, and moreover neglect p-wave contributions in the three-nucleon subclusters of ^{4}He, the two internal terms vanish. In other words, what enters our calculations is only the relative-momentum part of (3)

$$\hat{H}^{\lambda(1)\prime}_{\rm em}=-\frac{e}{2mc}\left(\frac{\tau_z^1+\tau_z^2+\tau_z^3}{3}-\tau_z^4\right)\hat{\epsilon}_\lambda\cdot\mathbf{q}, \tag{5}$$

TABLE I. $4N$ binding energy for the MT I+III potential depending on the number of expansion terms in the EDPE.

Number of EDPE terms	1	2	3	4	5	6
$4N$ binding energy [MeV]	−30.1	−30.1	−30.1	−30.2	−30.2	−30.2

or the relative-coordinate part of (4),

$$\hat{H}^{\lambda(1)}_{\rm em}=-\frac{ie}{2\hbar c}(E_f-E_i)\left(\frac{\tau_z^1+\tau_z^2+\tau_z^3}{3}-\tau_z^4\right)\hat{\epsilon}_\lambda\cdot\mathbf{r}, \tag{6}$$

where $\mathbf{q}=(\mathbf{p}_1+\mathbf{p}_2+\mathbf{p}_3-3\mathbf{p}_4)/4$ and $\mathbf{r}=(3/4)[(\mathbf{x}_1+\mathbf{x}_2+\mathbf{x}_3)/3-\mathbf{x}_4]$.

As mentioned already, the dipole aproximation is no longer sufficient at higher energies. What has to be, and will be taken into account is the quadrupole operator

$$H^{\lambda(2)\prime}_{\rm em}=-\frac{ie}{2mc}\sum_{j=1}^{4}\frac{1+\tau_z^j}{2}(\hat{\epsilon}_\lambda\cdot\mathbf{p}_j\mathbf{k}_\gamma\cdot\mathbf{x}_j+\hat{\epsilon}_\lambda\cdot\mathbf{x}_j\mathbf{p}_j\cdot\mathbf{k}_\gamma), \tag{7}$$

or its Siegert form [22]

$$H^{\lambda(2)}_{\rm em}=\frac{e}{2\hbar c}(E_f-E_i)\sum_{j=1}^{4}\frac{1+\tau_z^j}{2}\hat{\epsilon}_\lambda\cdot\mathbf{x}_j\mathbf{k}_\gamma\cdot\mathbf{x}_j\,, \tag{8}$$

with $\mathbf{k}_\gamma$ being the photon momentum.

Taking over the Alt-Grassberger-Sandhas (AGS) reduction technique developed for the four-nucleon collision problem [3], one arrives at the set of integral equations [4]

$$\mathcal{M}(\mathbf{q})=\mathcal{B}(\mathbf{q})+\int d^3q'\,\mathcal{V}(\mathbf{q},\mathbf{q}')\mathcal{G}_0(\mathbf{q}')\mathcal{M}(\mathbf{q}'), \tag{9}$$

i.e., at an effective two-body matrix equation for an off-shell extension $\mathcal{M}(\mathbf{q})$ of the amplitude (2). This relation is almost identical to the set of integral equations for four-nucleon rearrangement processes [3],

$$\mathcal{T}(\mathbf{q},\mathbf{q}')=\mathcal{V}(\mathbf{q},\mathbf{q}')+\int d^3q''\,\mathcal{V}(\mathbf{q},\mathbf{q}'')\mathcal{G}_0(\mathbf{q}'')\mathcal{T}(\mathbf{q}'',\mathbf{q}'). \tag{10}$$

TABLE II. $3N$ and $4N$ binding energies for the MT I+III potential.

$3N$ binding energy [MeV]	$4N$ binding energy [MeV]	Ref.
−8.56	−29.6	[30]
−8.592	−30.36	[19]
−8.54	−29.74	[31], SIDE
−8.86	−31.02	[31], IDEA
−8.536	−30.312	[32]
−8.54	−30.29	[33]
−8.87	−31.99	[34]
−8.595	−30.2	This work

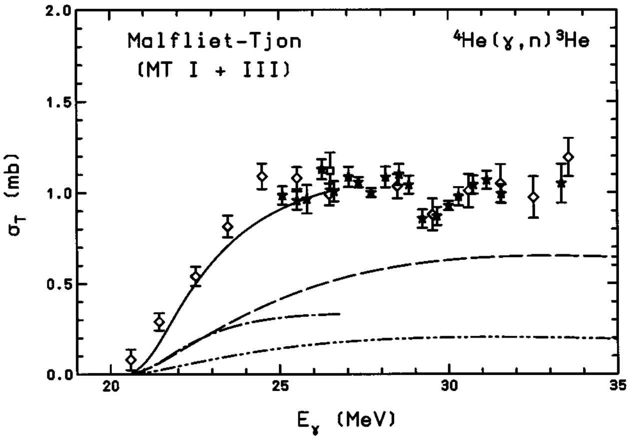

FIG. 1. Total cross section for ^{4}He$(\gamma,n)^3$He at low energies: Full solution with exchange currents (——) and without $(-\cdot-)$; plane-wave approximation with exchange currents $(--)$ and without $(-\cdot\cdot-)$; the data are from [6] ◇, [10] ⋆, and [11] □.

The kernel of both these equations is built up by the effective potential $\mathscr{V}$ and the effective Green's function $\mathscr{G}_0$, whose explicit definitions are found in [3,23]. But the potential $\mathscr{V}$ in the inhomogeneous term of Eq. (10) is replaced in Eq. (9) by an off-shell extension $\mathscr{B}(\mathbf{q})$ of the Born amplitude (1).

The construction of this amplitude requires the knowledge of the bound states $|\psi_{\rm III}\rangle$ and $|\psi_{\rm IV}\rangle$. Correspondingly one needs for the determination of $\mathscr{B}(\mathbf{q})$ the solutions of the homogenous integral equations of the effective three- and four-nucleon AGS formalism [3]. For details we refer to [5,24]. Having provided all these ingredients, in fact the main step in treating our problem, it remains to solve again a set of effective two-body integral equations, Eq. (9), but now in the continuous spectrum.

For completeness we finally mention that the disintegration cross section for an unpolarized incident photon beam is given by

$$\frac{d\sigma}{d\Omega}=\frac{\mu q}{2\pi\hbar^2 k_\gamma}\sum_{\lambda=1}^{2}|M^\lambda(\mathbf{q})|^2, \tag{11}$$

with μ being the reduced mass of the two outgoing fragments.

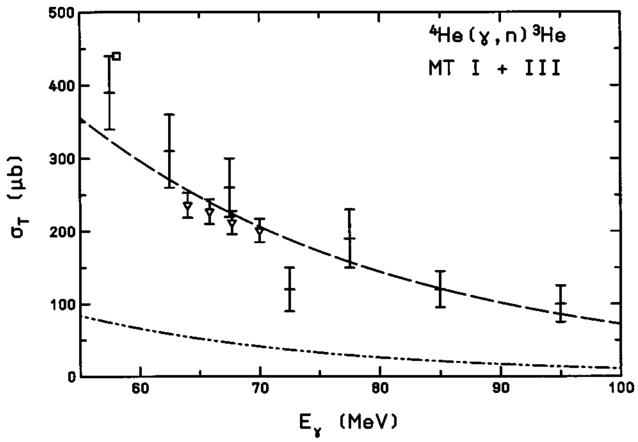

FIG. 3. Total cross section for ^{4}He$(\gamma,n)^3$He at higher energies: Plane-wave approximation with exchange currents $(--)$ and without $(-\cdot\cdot-)$; the data are from [14] ▽, [12] +, and [13] □.

III. RESULTS

For the MT I+III potential the 3N binding energy is 8.595 MeV. The 4N binding energy obtained in first-order EDPE is 30.1 MeV, a value which changes minimally to 30.2 MeV when taking into account up to six EDPE terms (see Table I). Our result moreover agrees very well with alternative calculations given in Table II. This demonstrates the high accuracy of the approximation or expansion technique underlying the following investigations. The 3N and 4N binding energies for the Yamaguchi potential are 9.97 MeV and 39.1 MeV, respectively.

Figure 1 shows our results for the ^{4}He$(\gamma,n)^3$He total cross section at low photon energies. The solid curve corresponds to the full solution of (9), the electromagnetic operator being given by the $E1$ Siegert operator (6). The dashed curve is the corresponding plane-wave (Born) result. The dashed-dotted and dashed-double-dotted curves are the analogous full and plane-wave results obtained with the electromagnetic operator (5), i.e., without inclusion of meson exchange currents. All these curves show a fairly flat nonresonant behavior. Due to the complicated cut structure of its kernel, the solution of the integral equation (9), which takes into account the final-state interaction (FSI), could only be

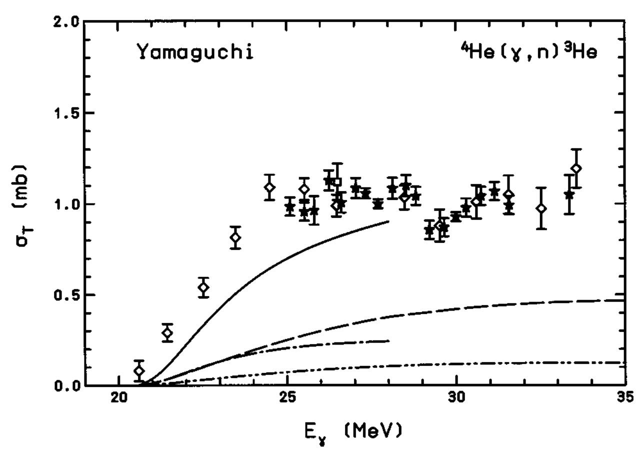

FIG. 2. Same as Fig. 1, but with the Yamaguchi potential.

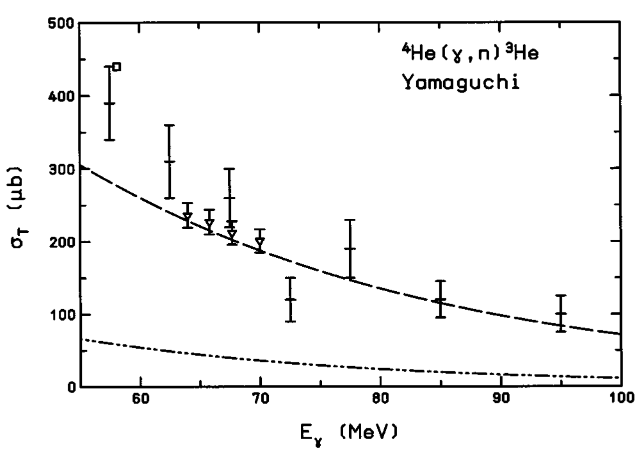

FIG. 4. Same as Fig. 3, but with the Yamaguchi potential.

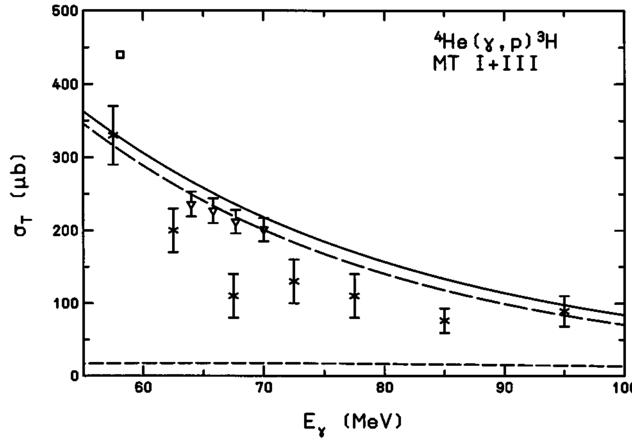

FIG. 5. Total cross section for ^{4}He$(\gamma,p)\,^3$H at higher energies: Plane-wave approximation including $E1$ and $E2$ (——). Contributions of $E1$ (— —) and $E2$ (− − −); the data are from [14] ▽, [12] ×, and [13] □.

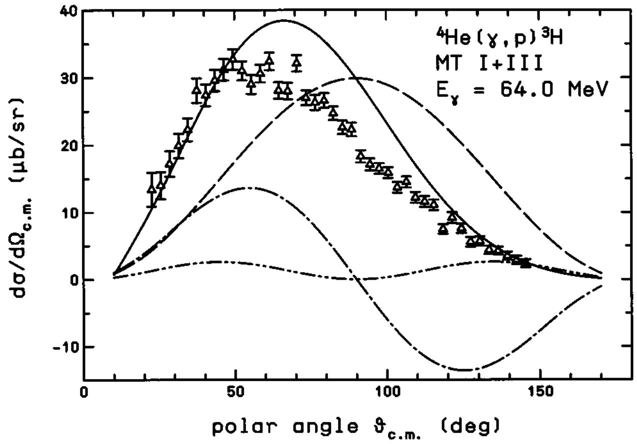

FIG. 7. Differential cross section: Plane-wave approximation including $E1$ and $E2$ (——). Contribution of $E1$ only (− −), of $E2$ only (− · · −), and of the $E1$-$E2$ interference term (− · −); the data are from [14] △.

performed below the three-fragment breakup threshold at 26.3 MeV. The agreement with the ^{4}He$(\gamma,n)\,^3$He data of Berman *et al.* [6], Ward *et al.* [10], and Asai *et al.* [11] is remarkably good for the solid curve, but poor for all other curves. This shows the necessity of fully solving Eq. (9) and including meson exchange currents.

The corresponding results obtained for the Yamaguchi potential are given in Fig. 2. They lie below the MT curves and also below the data. A similar trend was observed in ^{3}H photodisintegration calculations [21].

In Fig. 3 we present our ^{4}He$(\gamma,n)\,^3$He results at photon energies above 55 MeV, obtained by means of the plane-wave (Born) amplitude (1). To work with this approximation, i.e., to neglect the FSI, appears justified at these energies. Since in this region the Coulomb FSI is also expected to be negligible, we compare our ^{4}He$(\gamma,n)\,^3$He results not only with the old (γ,n) data by Gorbunov [12], but also with the (γ,p) data by Bernabei *et al.* [13] and the very accurate recent (γ,p) measurements by Jones *et al.* [14]. The agreement between theory and experiment is quite satisfactory and even excellent in the latter case, of course, only for the dashed curve obtained by means of the Siegert operator (4). The dashed-double-dotted curve, corresponding to the non-Siegert operator (3), is far off the data. As in the low-energy region, the meson exchange contributions, thus, are seen to be absolutely essential.

The corresponding Yamaguchi results presented in Fig. 4 lie again below the MT results, but the differences are much smaller than at low energies. In other words, the potential dependence is fairly diminished in this region, a feature also known from ^{3}H photodisintegration [21]. The agreement with the data of [14] is as good as in the MT case: the MT curve in Fig. 3 lies close to the upper bounds, the corresponding Yamaguchi curve close to the lower bounds of the error bars.

Figures 5 and 6 show our total cross sections for the process ^{4}He$(\gamma,p)\,^3$H at the same photon energies, but with the electric quadrupole operator $E2$ in Siegert form (8) being taken into account. Its contribution is seen to be comparatively small. The agreement with the data of [14] found in the pure $E1$ case, thus, persists. However, the MT curve is now slightly above the error bars, and the Yamaguchi curve sits almost exactly on the data. All this indicates trends, but is certainly not sufficient to establish any preference of one of these potentials. For the ^{4}He$(\gamma,n)\,^3$He total cross section, the $E2$ contribution is negligible due to much smaller isospin factors.

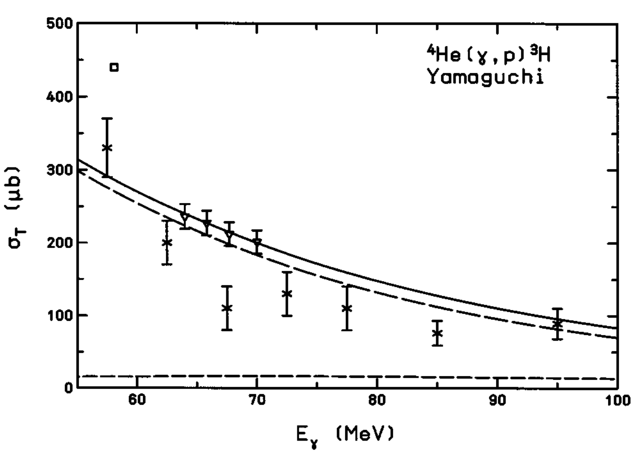

FIG. 6. Same as Fig. 5, but with the Yamaguchi potential.

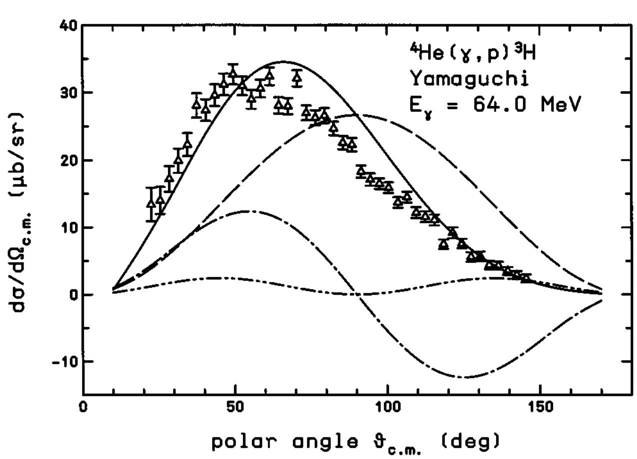

FIG. 8. Same as Fig. 7, but with the Yamaguchi potential.

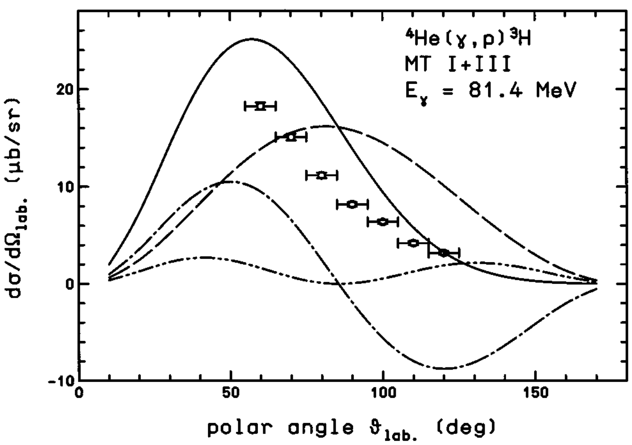

FIG. 9. Differential cross section: Plane-wave approximation including $E1$ and $E2$ (——). Contribution of $E1$ only (− −), of $E2$ only (− · · · −), and of the $E1$-$E2$ interference term (− · −); the data are from [25] ○.

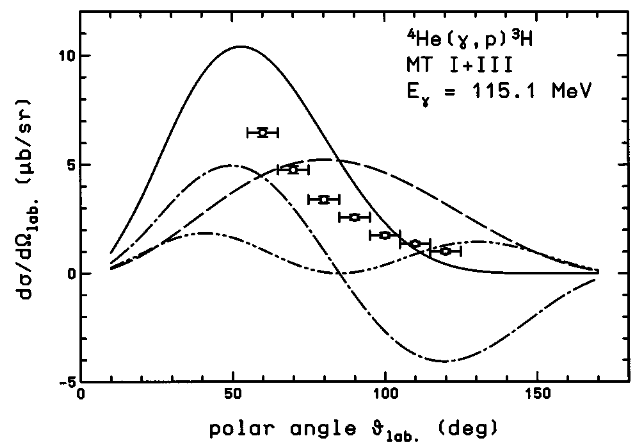

FIG. 11. Differential cross section: Plane-wave approximation including $E1$ and $E2$ (——). Contribution of $E1$ only (− −), of $E2$ only (− · · · −), and of the $E1$-$E2$ interference term (− · −); the data are from [25] ○.

In Figs. 7 and 8 our ^{4}He(γ,p)^{3}H differential cross sections at 64 MeV photon energy are compared with the data of [14]. In Figs. 9–14 our calculations at higher energies are compared with the measurements of [25]. The interference between the $E1$ and $E2$ amplitudes implies now a fairly big $E2$ contribution which is, in fact, essential for achieving better agreement between theory and experiment. As in the total cross section, this agreement is somewhat better for the Yamaguchi potential. To clarify the origin of the remaining discrepancies, higher-order electric or magnetic multipoles, higher partial waves in the effective three-nucleon potential [26,27], or the final-state interaction that may have an effect on the differential cross section even at these high energies, should be considered.

IV. CONCLUSIONS

By solving integral equations of the AGS type with the semirealistic Malfliet-Tjon potential, a remarkably good description of the ^{4}He(γ,n)^{3}He photoprocess has been achieved below the three-fragment breakup threshold. At photon energies above 55 MeV, the plane-wave approximation was assumed to be justified. The agreement between theory and experiment is good for the total cross section, but less satisfactory for the differential cross section even after incorporation of the electric quadrupole contribution. In the intermediate region, where the final-state interaction is still expected to play a noticeable role, the present state of the art does not allow us to fully solve the underlying integral equations. Up until now model treatments, therefore, had to be employed in this region. We recall that the one proposed in [28] has led to a good interpolation of the theoretical low- and high-energy results. Attempts to treat the integral equations above the threshold in a reliable manner are in progress.

Let us add some remarks. Our treatment was based on exact integral equations, simplified by two approximations, the W-matrix approximation and the EDPE. Their accuracy was demonstrated in [17–19] and in the present paper. The above calculations, thus, were performed within a practically exact framework. There are no open parameters or corrections that could have been used to adjust the theoretical results to the experimental ones.

This is true with but one trivial exception. The triton and ^{4}He binding energies, and consequently the photodisintegration thresholds, obtained for the Malfliet-Tjon, the Yamagu-

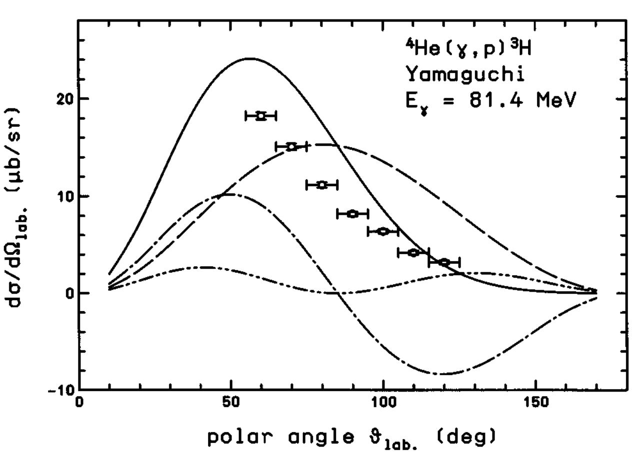

FIG. 10. Same as Fig. 9, but with the Yamaguchi potential.

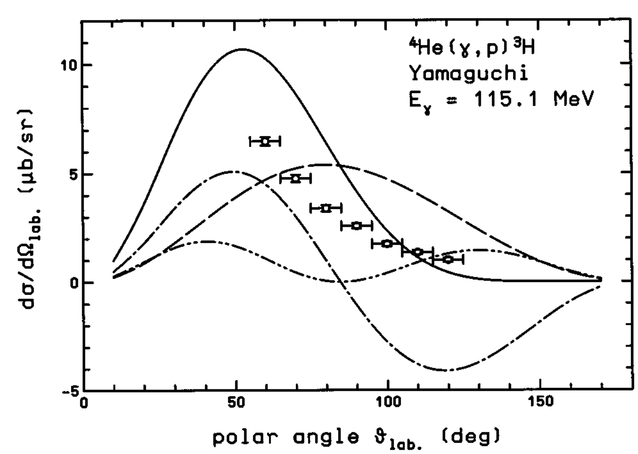

FIG. 12. Same as Fig. 11, but with the Yamaguchi potential.

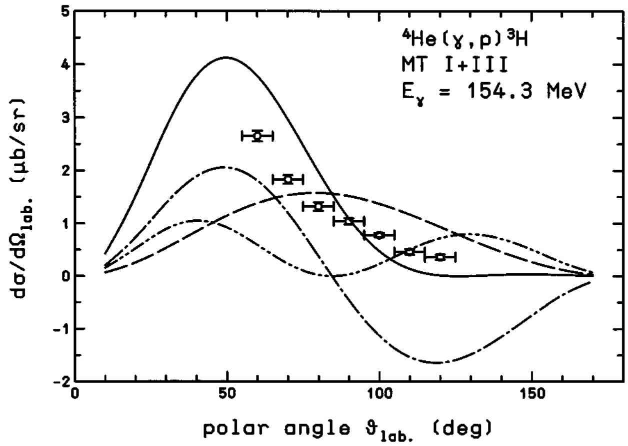

FIG. 13. Differential cross section: Plane-wave approximation including $E1$ and $E2$ (——). Contribution of $E1$ only (− −), of $E2$ only (− · · · −), and of the $E1$-$E2$ interference term (− · −); the data are from [25] ○.

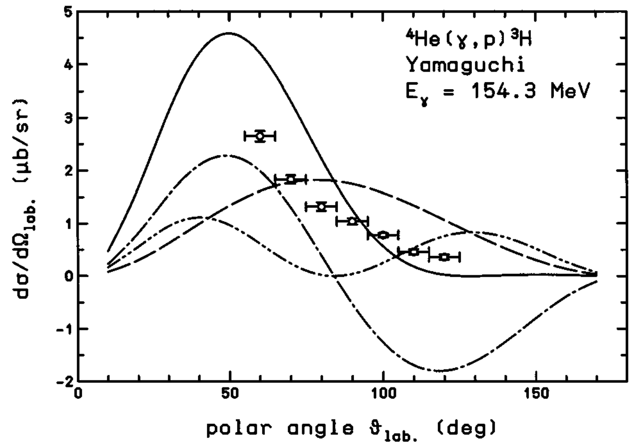

FIG. 14. Same as Fig. 13, but with the Yamaguchi potential.

chi, and all more realistic potentials, differ from the experimental ones. To achieve kinematical consistency, our curves had to be shifted correspondingly, i.e., such that the theoretical and experimental thresholds coincide.

It should be emphasized that this pure kinematical shift does not imply or require any corresponding replacement of energy or momentum variables in the photodisintegration amplitude. This concerns, in particular, the energies E_f and E_i in the Siegert operators (4) and (8), which occur there as eigenvalues of the total nuclear Hamiltonian. The theoretical energy values, corresponding to the respective nuclear wave functions, consequently had to be, and were chosen in the present and previous applications of our formalism. Just for the Malfliet-Tjon (MT I+III) potential the differences between the theoretical and experimental energy values are comparatively small, so that in this case the above issue is of minor relevance. This has to be contrasted with the resonating group method treatment of [2]. There, a Gaussian potential has been chosen with energy eigenvalues far from experiment. The authors of [2], therefore, were led to replace these values by the experimental ones, reproducing in this way the apparent giant resonance. Later on, by going back to the unrealistic theoretical energies, the by now accepted flat behavior was also reproduced within the same approach [29]. Further comments on these ambiguities are found in [28].

In conclusion, the treatment of the photodisintegration process ^{4}He$(\gamma,n)\,^3$He within an exact theory is essential for achieving satisfactory agreement with the data in an unambiguous way. The potential dependence found for our total cross sections offers the possibility of testing nuclear forces when going over to more realistic interactions, although one may expect the magnitude of the low-energy photodisintegration cross section to be related to the ^{4}He binding energy, analogous to the correlation of the size of the photodisintegration peak and the binding energy known from the triton case [21]. Using realistic potentials or including the Coulomb force in the exact four-nucleon integral equation theory is, of course, a nontrivial task. At higher photon energies, where differences between our results and the differential cross section data persist, further investigations are required.

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