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REACTION MECHANISMS IN TWO-BODY
PHOTODISINTEGRATION AND
ELECTRODISINTEGRATION OF ^4He

DAPNIA

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Reaction mechanisms in two-body photodisintegration and electrodisintegration of ${}^4\text{He}$.

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Abstract

One and two body mechanisms cannot alone reproduce the experimental values of the ${}^4\text{He}(\gamma, p)\text{T}$ reaction cross section in the Δ region. Meson double scattering is shown to account for the major part of these discrepancies. All together these mechanisms account for the near equality of the cross section of the ${}^4\text{He}(\gamma, p)\text{T}$ and ${}^4\text{He}(\gamma, n){}^3\text{He}$ channels. Consequences on the ${}^4\text{He}(e, e'p)\text{T}$ and the ${}^4\text{He}(e, e'n){}^3\text{He}$ channels are investigated, with a special emphasis on the determination of the ${}^4\text{He}$ wave function and the study of the behavior of bound nucleon form-factors. The sensitivity of spin observables, to various interaction effects, is investigated.

1 Introduction.

The cross section of the ${}^4\text{He}(\gamma, p)T$ reaction has been measured for a long time [1]. However, its behavior in the Δ energy range is not satisfactorily understood, although some efforts have been devoted to extend, above the pion production threshold, low energy theories based on correlated shell model four body wave function [2], or to estimate the contribution of the Δ in the intermediate state [3]. Taking full advantage of the progresses which have been made in our knowledge of the four body ground state wave function [4] and in our understanding of the dominant mechanisms in photo and electrodisintegration of ${}^3\text{He}$ [5,6], we are now in a good position to go beyond, and to provide a unified description of photo- and electro-disintegration of ${}^4\text{He}$.

Fig.1 shows the excitation functions of the ${}^4\text{He}(\gamma, p)T$ reaction at $\theta_p^{c.m.} = 60^\circ$, 90° and 120° . The contribution of the one-body mechanisms exhibits the characteristic node of the S-wave $\langle {}^4\text{He}|T \rangle$ overlap integral, but is strongly suppressed due to the large momentum transfer. The contribution of two-body meson exchange mechanisms leads to a good agreement with the data below the pion production threshold, but fails to reproduce them in the Δ energy range, by a factor two at 90° and six at 60° . The meson double scattering mechanism accounts for most of the discrepancies between the theory and the experiment. As in the ${}^3\text{He}(\gamma, p)D$ channel [6], the momentum transfer is more likely to be shared by three nucleons rather than two, and the on-shell nature of the photoproduced pion leads to a triangular singularity which enhances the corresponding contribution. Moreover, contrary to the ${}^3\text{He}(\gamma, p)D$ channel no isospin selection rules prevent the Δ to be formed at the pion photoproduction vertex : its strong contribution is however quenched by the extra form factor due to the presence of a fourth nucleon.

On the contrary, the ${}^4\text{He}(e, e'p)T$ reaction has been studied [7,8,9] in a kinematical range where the contribution of two and three body meson exchange mechanisms is small, but where the contribution of final state interactions is significant. Here special care of the orthogonality between initial and final states has to be taken [10]. Although the major trends of the reaction mechanisms are well understood [8] the size of the corresponding corrections to quasi free scattering prevented us to extract accurately the ${}^4\text{He}$ momentum distribution.

The aim of this article is precisely to deal with all these interaction mechanisms, to calibrate them in the relevant kinematical domains and to look for kinematics and observables which allow for a reliable determination of the ${}^4\text{He}$ wave function, or the form-factors of bound nucleons. Section 2 summarizes the expressions of the matrix elements of the dominant mechanisms. Since the model has already been compared to experiments performed at Amsterdam [8] or Saclay [9], section 3 puts a special emphasis on the comparison with other data and on future developments. Appendices deal with more technical points : relationships between cross-sections,

spin observables and matrix elements; elementary currents; etc. . . .

2 The Model.

The method is the straightforward extension, to four nucleons, of the diagrammatic method which has been widely used in the analysis of photo and electrodisintegration of three-nucleon systems [11]. The general relations between the reaction amplitude and the cross-section or the various spin observables are gathered in Appendix A. The amplitude is expanded in terms of a few relevant diagrams which are computed in the momentum space. The energy and the momentum are conserved at each vertex, and the kinematics as well as the phase space are relativistic. The non relativistic reduction (including all the terms of order $1/m^3$) is only done for the elementary operators. The ${}^4\text{He}$ and T wave functions are the variational solutions of the Urbana group [4] for the Argonne V14 potential [12]. Let me recall the expressions [5,6] of each amplitude, with a special emphasis on the further assumptions needed in the four-nucleon case. In the laboratory frame, the amplitudes of each diagram in Figs.2-4 are the following :

2.1 The one body mechanisms.

The one-nucleon exchange amplitude (Fig.2) takes the form

$$T^{(1)} = \sum_{m_p} \langle m_1 | j_N | m_p \rangle \langle \frac{1}{2} m_p \frac{1}{2} m_T | 00 \rangle \phi_0(\vec{p}_T) / \sqrt{4\pi} \quad (1)$$

where $(\vec{p}, m_p)$, $(\vec{p}_1, m_1)$ and $(\vec{p}_T, m_T)$ are respectively the momenta and magnetic quantum numbers of the exchanged nucleon, the emitted proton and the spectator Triton ($\vec{p} = -\vec{p}_T$ in the Lab. frame). The dipole fit to the Sachs form factors is used in the space and time components $j_N = (\vec{j}_N, j_N^0)$ of the nucleon current which are expanded up to and including terms of order $1/m^3$: for free nucleon it is gauge invariant at this order (see Appendix B). The $\langle {}^4\text{He} | T \rangle$ overlap integral ϕ_0 has been computed numerically starting from the variational wave functions [4] of the three and four nucleon systems.

The three nucleon exchange amplitude (Fig.2) comes from the antisymmetry of the four nucleons in the final state and takes the same form as the one nucleon exchange amplitude, under the assumption that a Triton is exchanged : the magnetic quantum numbers and the momenta of the proton and the Triton have to be interchanged and the Triton electromagnetic form factors [13] have to be used. This graph dominates [14] in a restricted part of the phase space (when the proton is emitted at very backward angles), but it is not important in most of the kinematics covered by this study.

2.2 The two body mechanisms.

The two nucleon meson exchange amplitude and the nucleon-nucleon rescattering amplitude (Fig.3) are computed in the same way and can be combined in the same compact expression. For an active np pair in T=0, J=1 state :

$$\begin{aligned}
T^{(2)} = & \sum_{\Lambda, M_J, m_n} \int \frac{d^3 \vec{\xi}}{(2\pi)^3} \tilde{T}(\vec{k}, -\vec{\xi} 1 M_J \rightarrow \vec{p}_1 m_1, \vec{n} m_n) \\
& \left\{ \frac{\chi_0(\vec{\xi})}{\sqrt{4\pi}} \frac{\psi_0(\vec{\xi} - 2\vec{p}_T/3)}{\sqrt{4\pi}} \langle \frac{1}{2} m_n 1 \Lambda | \frac{1}{2} m_T \rangle \langle 1 M_J 1 \Lambda | 00 \rangle \right. \\
& + \sum_{m_l, m_j} \frac{\chi_0(\vec{\xi})}{\sqrt{4\pi}} \psi_2(\vec{\xi} - 2\vec{p}_T/3) Y_2^{m_l}(\vec{\xi} - 2\vec{p}_T/3) \\
& \langle \frac{3}{2} m_j 1 \Lambda | \frac{1}{2} m_T \rangle \langle 2 m_l \frac{1}{2} m_n | \frac{3}{2} m_j \rangle \langle 1 M_J 1 \Lambda | 00 \rangle \\
& + \sum_{m_l, m_s} \chi_2(\vec{\xi}) Y_2^{m_l}(\vec{\xi}) \frac{\psi_0(\vec{\xi} - 2\vec{p}_T/3)}{\sqrt{4\pi}} \\
& \left. \langle \frac{1}{2} m_n 1 \Lambda | \frac{1}{2} m_T \rangle \langle 2 m_l 2 m_s | 00 \rangle \langle 1 M_J 1 \Lambda | 2 m_s \rangle \right\} \quad (2)
\end{aligned}$$

where $(-\vec{\xi}, M_J)$ are the momentum and the magnetic quantum number of the active pair, where $(\vec{\xi}, \Lambda)$ and $(\vec{n}, m_n)$ are those of the spectator nucleon pair and the neutron which recombine into the final Triton, where $(\vec{p}_1, m_1)$ and $(\vec{p}_T, m_T)$ are those of the outgoing proton and Triton, and where $(\vec{k}, \omega)$ is the four-momentum of the incoming photon. The antisymmetrized Triton wave function [4] is projected on a cluster basis, where two nucleons are coupled to angular momentum L' , spin J' and isospin T' , the third nucleon moving with angular momentum l' . It is a good approximation to factorise its dominant S and D parts as $\Phi_{l'L'}^{(3)}(p, q) = \psi_{l'}(p) v_{L'}(q)$. An integration over the relative momentum of the two spectator nucleons is implicit in eq.2 : it is performed by computing the overlap between the antisymmetrized ${}^4\text{He}$ wave function [4] and the relative wave function $v_{L'}(q)$ of two nucleons in a T=0 state in the Triton. The dominant S and D parts of this overlap integral is also factored as $\Phi_{l'L'}^{(4)}(\xi, q) = \chi_{l'}(\xi) U_{L'}(q)$, where $U_{L'}$ and $\chi_{l'}$ stand respectively for the relative wave function of the two active nucleons and the wave function of their center of mass. All these overlap integrals and wave functions are computed numerically; their norm contains the relevant isospin coupling coefficient. The three-fold integral over the momentum of the center of mass of the two spectator nucleons is performed numerically according to the fifteen points Gauss-Kronrod rule [15]; their mass is assumed to be twice the nucleon mass. The three terms in eq.2 correspond to the transitions between the S and D waves of ${}^4\text{He}$ and Triton; the D to D state transition is neglected.

The antisymmetrized two-body matrix element $\tilde{T}(\vec{k}, -\vec{\xi} 1 M_J \rightarrow \vec{p}_1 m_1, \vec{n} m_n)$ de-

scribes the absorption of the photon by a neutron-proton pair in the $T=0$ state. It is directly deduced from the two-body part of the matrix element of the $D(e,e'p)n$ reaction (see [5] and references therein), and it is made of two parts (see Appendix C for details). The meson exchange amplitude involves π and ρ exchange, and takes into account the creation of the Δ , besides non resonant meson production terms; the three fold integral is performed numerically. The nucleon rescattering amplitude takes into account the full half-off shell expansion, up to and including D-waves, of the nucleon-nucleon scattering matrix. Due to antisymmetry, charge exchange nucleon-nucleon rescattering is automatically included.

When the active pair is in a $T=1$ isospin state (neutron-proton or two identical nucleons), the two nucleon amplitude exhibits a similar form : only S-waves are relevant and trivial changes have to be done in the spin couplings.

2.3 The three body mechanisms.

The meson double scattering amplitude (Fig.4) relates the ${}^4\text{He}$ four-body break-up and the subsequent Triton recombination matrix elements. The pair, made of the spectator nucleon and the nucleon which recoil at the pion photoproduction vertex, is projected on the corresponding state (of given spin and isospin) of a nucleon pair in the Triton. Let's assume, for the moment, that it is a $T=0$ neutron-proton pair :

$$T^{(3)} = \sum_{\Lambda, m_p} \int \frac{d^3 \vec{\xi}}{(2\pi)^3} T(\gamma^4 He \rightarrow (np)_0 np) \left\{ \frac{\psi_0(\vec{\xi} - 2\vec{p}_T/3)}{\sqrt{4\pi}} \langle \frac{1}{2} m_n \frac{1}{2} m_T | \frac{1}{2} m_T \rangle \right. \\ \left. \psi_2(\vec{\xi} - 2\vec{p}_T/3) \sum_{m_j, m_l} Y_2^{m_l}(\vec{\xi} - 2\vec{p}_T/3) \langle \frac{3}{2} m_j \frac{1}{2} m_T | \frac{1}{2} m_T \rangle \langle 2 m_l \frac{1}{2} m_n | \frac{3}{2} m_j \rangle \right\} \quad (3)$$

where the two parts correspond respectively to the S and D parts of the motion of center of mass of the np pair with respect to the third nucleon in the final Triton. The notations are the same as in eq.2, and the integral is performed numerically according to the Gauss-Kronrod rule [15].

Since pion absorption by a $T=1$ pair is strongly suppressed [16], the dominant graph in the ${}^4\text{He}$ four-body photodisintegration is the photoproduction, on a $T=0$ np pair, of a neutral pion which is absorbed by the second $T=0$ np pair of ${}^4\text{He}$. The corresponding matrix element takes the following form :

$$T(\gamma^4 He \rightarrow (np)_0 np) = \int \frac{d^3 \vec{\nu}}{(2\pi)^3} \frac{\chi_0(\vec{\nu})}{\sqrt{4\pi}} \frac{1}{q_\pi^2 - m_\pi^2 + i\epsilon} \sum_{M_J} \langle 1 M_J \frac{1}{2} -M_J | 00 \rangle \\ T_{\gamma(np)_0 \rightarrow \pi^0(np)_0}(\vec{k}, -\vec{\nu} M_J \rightarrow \vec{q}_\pi, \vec{\xi} \Lambda) T_{\pi^0(np)_0 \rightarrow np}(\vec{q}_\pi, \vec{\nu} - M_J \rightarrow \vec{n} m_n, \vec{p}_1 m_1) \quad (4)$$

where $(-\vec{\nu}, M_J)$ and $(\vec{\nu}, -M_J)$ are respectively the momenta and the quantum numbers of the pair on which the pion is created and the pair which reabsorbs it, and

where the notations are the same as in eq.2. Only relative motion of this two pairs in a S wave is retained : this was found to be a good assumption in a recent study [17] of related three body mechanisms in the ${}^3\text{He}(\gamma, 2p)$ reaction.

The antisymmetrized matrix element of the pion absorption process $T_{\pi^0(np)_0 \rightarrow np}$ is directly related to the matrix element of the reaction $\pi^+ d \rightarrow pp$, by a trivial transformation in the isospin space and by replacing the components of the deuteron wave function by the corresponding components U_0 and U_2 of the relative wave function of an np pair in ${}^3\text{He}$ (see Appendix C). Both π and ρ are exchanged and the three-fold integration, over the relative momentum in the pair, is done numerically. The matrix element of the Δ term can be found in the appendix of Ref.[18]. All the other πN partial waves, up to D wave, are added and parameterized by the experimental phase shifts [19]. Below the pion production threshold, only the real part of the Δ propagator is retained and the S-wave πN amplitudes are approximated by the corresponding scattering lengths.

The pion production amplitude takes the form :

$$T_{\gamma(np)_0 \rightarrow \pi^0(np)_0} = -F \left(\frac{\vec{k} - \vec{q}_\pi}{2} \right) \sum_{\lambda_p \lambda'_p \lambda_n} \left\langle \frac{1}{2} \lambda_p \frac{1}{2} \lambda_n \mid 1 M_J \right\rangle \left\langle \frac{1}{2} \lambda'_p \frac{1}{2} \lambda_n \mid 1 \Lambda \right\rangle \\ \left\{ T_{\gamma p \rightarrow \pi^0 p}(\vec{k}, \vec{p} \lambda_p \rightarrow \vec{q}_\pi, \vec{p}' \lambda'_p) + T_{\gamma n \rightarrow \pi^0 n}(\vec{k}, \vec{p} \lambda_p \rightarrow \vec{q}_\pi, \vec{p}' \lambda'_p) \right\} \quad (5)$$

where $(\vec{p}, \lambda_p)$ and $(\vec{p}', \lambda'_p)$ are respectively the momenta and magnetic quantum numbers of the target and the recoiling nucleons. The pion photoproduction elementary amplitudes [20] have been factored out the integral which reduces to the transition form factor between two pairs in a T=0 state :

$$F \left(\frac{\vec{k} - \vec{q}_\pi}{2} \right) = F \left(\frac{\vec{\xi} - \vec{\nu}}{2} \right) = \int \frac{d^3 \vec{n}}{(2\pi)^3} \frac{U_0(\vec{n} - \vec{\nu}/2)}{\sqrt{4\pi}} \frac{v_0(\vec{n} - \vec{\xi}/2)}{\sqrt{4\pi}} \quad (6)$$

where $(\vec{n}, \lambda_n)$ are the momentum and the magnetic quantum number of the spectator neutron in the pair. This assumption reduces the twelvefold integral, which is implicit in eq.3, to a nine fold integral.

To further reduce it to a six fold integral which is numerically tractable, the pion photoproduction and absorption amplitudes are factored out the integral (and computed for a target pair at rest, *i.e.* $\vec{\nu} = 0$), in eq.4, which becomes :

$$T(\gamma {}^4\text{He} \rightarrow (np)_0 np) = \sum_{M_J} \langle 1 M_J 1 -M_J \mid 00 \rangle \\ T_{\gamma(np)_0 \rightarrow \pi^0(np)_0}(\vec{k}, -\vec{0} M_J \rightarrow \vec{q}_\pi, \vec{\xi} \Lambda) T_{\pi^0(np)_0 \rightarrow np}(\vec{q}_\pi, \vec{0} -M_J \rightarrow \vec{n} m_n, \vec{p}_1 m_1) \\ \int \frac{d^3 \vec{\nu}}{(2\pi)^3} \frac{\chi_0(\vec{\nu})}{\sqrt{4\pi}} \frac{1}{q_\pi^2 - m_\pi^2 + i\epsilon} \quad (7)$$

Expanding the relative wave function χ_0 as a sum of poles, the remaining integral is computed in a compact analytical form according to Ref.[21]. Above the pion threshold, its logarithmic singularity, associated with the on-shell propagation of the pion, enhances the contribution of the three body mechanism. It selects the low momentum components of the wave function, makes more likely the contribution of its S part which is retained in eq.7, and justifies the factorization procedure. It turns out that all these approximations, which have been made to make tractable the evaluation of the meson double scattering amplitude, have been shown to lead to a fair understanding of related three body processes in the ${}^3\text{He}(\gamma, 2p)n$ reaction [17].

The two pion photoproduction elementary amplitudes add coherently : the plus sign in eq.5 comes from the $T=0$ nature of both the initial and the recoiling np pairs. Contrary to the ${}^3\text{He}(\gamma, p)D$ reaction [6], the Δ creation term is not forbidden. However, the transition form factor F damps the matrix element.

If a $T=1$ np pair recoils at the pion photoproduction vertex, a minus sign appears in eq.5 : it comes from the $T=0$ to $T=1$ nature of the transition between the two pairs. The contributions of the Δ creation terms, as well as those of the ω exchange terms, cancel. The corresponding meson double scattering contribute very little and is discarded in the present study.

When a neutron pair recoils at the pion photoproduction vertex, the meson double scattering matrix element exhibits a form very similar to eqs.3-7. Besides trivial changes in spin and isospin couplings and under the same assumptions, only a π^+ can be exchanged and only S wave is relevant in eq.3.

2.4 The ${}^4\text{He}(e, e'n){}^3\text{He}$ channel.

The extension of the model to this channel is straightforward. Except for trivial changes in the isospin coupling coefficients, the matrix elements of the two body and of the three body mechanisms are the same as in the ${}^4\text{He}(e, e'p)T$ channel. The only major changes concern the charge and the magnetic moment of the active nucleon in the one body matrix elements. The yields of these two channels are expected to be almost the same, in kinematics where one body mechanisms are suppressed.

3 Comparison with Experiments and Perspectives.

3.1 Photodisintegration.

Fig.5 shows the excitation functions of the ${}^4\text{He}(\gamma, n){}^3\text{He}$ at $\theta_p^{c.m.} = 60^\circ, 90^\circ$ and 120° . As in the case of the ${}^4\text{He}(\gamma, p)\text{T}$ reaction (Fig.1), the contribution of two and three body mechanisms dominates over the contribution of one body mechanisms which exhibit the characteristic node of the S-wave $\langle {}^4\text{He} | T \rangle$ overlap integral, except at $\theta_p^{c.m.} = 120^\circ$ where the contribution of the ${}^3\text{He}$, or triton, exchange graph begins to play a role.

At very low energies, one-body mechanisms dominate the cross-section of the ${}^4\text{He}(\gamma, p)\text{T}$ channel, while two-body exchange currents are necessary to get a fair understanding of the experiment : nucleon rescattering mechanisms play a minor role. This is not the case of the ${}^4\text{He}(\gamma, n){}^3\text{He}$ channel. As expected one body mechanisms are strongly suppressed and interfere destructively with two-body meson exchange currents : the cross-section is dominated by nucleon rescattering mechanisms. However, we shall see later that the way we have estimated this nucleon rescattering contribution may not be reliable at energies as low as 50 MeV, and this result should be taken as an indication of the main trends of the model in this energy range.

In the Δ energy range, the model leads to a fair agreement with the experimental data of which it reproduces the main trends. As expected, it predicts roughly the same strength in the two channels. However the model systematically underestimates the data at $\theta_p^{c.m.} = 60^\circ$ and overestimates them at $\theta_p^{c.m.} = 120^\circ$, while it reproduces them fairly well at $\theta_p^{c.m.} = 90^\circ$. This trend is also apparent in Fig.6 which shows the angular distribution of the cross section of the ${}^4\text{He}(\gamma, p)\text{T}$ reaction [1,22] at $E_\gamma = 246$ MeV. In that energy range, the Δ formation terms dominate and it is very likely that the approximations which have been made to reduce the numerical integration in the three body amplitudes to a reasonable time (around one hour of Cray cpu for an angular distribution at eleven angles) is a major source of this disease. A more exact treatment of the Fermi motion, in the first loop of the three body amplitudes (it is correctly taken into account in the various two nucleon amplitudes), would spread the Δ contribution and lead to smoother excitation functions.

At lower energies, below and around the pion threshold, the agreement between theoretical and experimental cross sections [23] is very good, as can be seen in Fig.7 : even at such a low energy the contribution of many body mechanisms is dominant. Here the Born terms in the pion photoproduction amplitude play a significant role.

Spin observables are very scarce. Fig.8 compares the model to the proton polar-

ization [22] : the main trends of the data are reproduced and three-body mechanisms are necessary to reproduce them.

To summarize, multinucleon mechanisms dominate the two body photodisintegration of ${}^4\text{He}$ at high energy, where they can be calibrated : their contribution to the electrodisintegration can be safely estimated.

3.2 Electrodisintegration.

A systematic investigation of the ${}^4\text{He}(e,e'p)\text{T}$ reaction in various kinematics has been performed at Amsterdam [8] : variations of the cross section against the three momentum of the virtual photon in the range 300 to 500 MeV/c, against the relative energy between the outgoing proton and triton in the range 30 to 110 MeV, and against the triton recoil momentum up to 350 MeV/c. The ${}^4\text{He}(e,e'n){}^3\text{He}$ reaction has also been investigated [24] but to a lower extent. The model accounts for these data within twenty percent over several decades and I refer the reader to refs.[8,24] for a detailed comparison. The main results are that, in this kinematical domain, two and three body meson exchange mechanisms contribute little and that the main correction to the quasi free mechanism comes from nucleon final state interactions. It decreases when the three momentum of the virtual photon increases, according to eq.2 which involves a transition form factor besides the nucleon electromagnetic form factors which are the only source of variation with the photon momentum in the quasi free matrix element (eq.1).

A more stringent test of the model has been performed at Saclay where the transverse and the longitudinal cross sections have been accurately determined in collinear kinematics (see Appendix A, and section 6 of Ref.[25]). The results [9,27] are summarized in Fig.9 which shows the variation of the corresponding spectral functions against the three momentum of the virtual photon, for a given value of the recoiling triton momentum. While nucleon rescattering induces the main effect on both spectral functions, meson exchange currents contribute mainly to the transverse one. The model reproduces fairly well the transverse spectral function and accounts for half of the quenching of the longitudinal spectral function. It is worthy to point out that very similar results have been obtained in ref. [10] using a different formalism. The remaining discrepancy is a good measure of the uncertainty of the theoretical predictions, rather than an evidence for possible exotic effects.

The degree to which gauge invariance is violated by the calculation is also a measure the inaccuracy of the model. As usual in nuclear physics the conservation of the electromagnetic current has been used to eliminate its longitudinal component : its time component is used to compute the longitudinal response function. With such a choice the contribution of exchange current to the longitudinal response function is negligible (of the order of 1% in Fig.9). Equivalently, the longitudinal

response function can also be computed with the longitudinal component of the current, after the elimination of its time component : the result should be the same. This is the case for the $D(e,e'p)n$ reaction [5] : gauge invariance is satisfied within a few percent at the two nucleon level (see Appendix C). This is not the case in Fig.9, where the choice of the longitudinal component of the current leads to a better agreement with the experiment. The reason for this departure from gauge invariance lies in the various approximations which have been made. Among them, the most important concerns the treatment of the continuum as a truncated multiple scattering series : the corresponding wave function is not strictly orthogonal to the four body ground state wave function. Such a lack of orthogonality is more severe when the longitudinal response function is computed from the time component of the current : the charge operator connects states with the same quantum numbers and, at low momentum transfer, spurious contributions appear if they are not strictly orthogonal. The way to cure this disease would be to solve the four nucleon problem in the continuum with the same degree of accuracy as in the ground state, but this seems out of reach. A more sensible way would be to suppress its consequences by performing experiments at momentum of the virtual photon large enough to strongly reduce rescattering contributions.

It should also be pointed out that the truncated multiple scattering series is a good approximation of the continuum at high energy only (let say relative energies between the proton and the triton above 40 MeV). This condition has been achieved in the kinematics studied in Ref.[8]. In the kinematics of Fig.9 the relative energy between the outgoing proton and triton is close to 100 Mev, except for the point at $k=380$ MeV/c where it is only 30 MeV : the violation of gauge invariance is the largest here.

The ^4He momentum distribution has been also investigated at momentum higher than 300 MeV/c at Saclay [28]. As can be seen in Fig.10, the effects of multinucleon mechanisms are huge in this momentum range where the $\langle ^4\text{He}|T \rangle$ overlap integral exhibits a zero. Nucleon final state interaction shift this zero toward a lower momentum, while two and three body meson exchange mechanisms fill in the corresponding dip and bring the theory in close agreement with the data for the highest values of the recoiling triton momentum. However a discrepancy still remains at lower momenta. It may come partly from the uncertainties of the model, which were just discussed, but also from possible inadequacies of the variational method [4] to predict the high momentum tail of the momentum distribution. Repeating the present analysis with more recent four body wave functions, obtained with the Green function Monte-Carlo method [26], would certainly be interesting.

It has already been noticed that both the Amsterdam [8] and the Saclay [27] data prefer, at low values of the triton momentum, the four body wave function computed with the Urbana potential rather than with the Argonne potential, with which all the preceding results were obtained. Fig.11 summarizes the cross section of the

${}^4\text{He}(e,e'p)T$ reaction which was measured at Amsterdam [7] and Saclay [28] in order to determine the ${}^4\text{He}$ momentum distribution. It is compared to the prediction of the model when the $\langle {}^4\text{He}|T\rangle$ overlap integral is computed with the Urbana potential : the agreement is fairly good below 300 MeV/c and above 500 MeV/c, but a problem still subsists around 400 MeV/c.

To summarize, although the model reproduces the main trends of the available data, the size of the contribution of multinucleon mechanisms prevents a reliable determination of the high momentum part of the ${}^4\text{He}$ wave function, in the kinematics so far achieved. A significant increase of the momentum of the virtual photon will strongly reduce the contribution of these multinucleon mechanisms and will offer us a way to overcome this drawback. To go further, and to get rid of the contribution of meson exchange currents, the longitudinal cross section must be singled out. This appears clearly by comparing Fig.12, which shows the transverse and the longitudinal cross sections in the kinematics already achieved at Saclay ($|\vec{k}| = 280$ MeV/c), Fig.13 which corresponds to a kinematics soon achievable at Amsterdam or Mainz ($|\vec{k}| = 426$ MeV/c), and Fig.14 which corresponds to a kinematics typical of CE-BAF where the photon momentum ranges from $|\vec{k}| = 1.15$ to 2.5 GeV/c. However, even for such high values of the photon momentum, nucleon rescattering effects always make very difficult a precise determination of the $\langle {}^4\text{He}|T\rangle$ overlap integral in the vicinity of its zero. On the other hand, the interference cross sections are also sensitive to the various reaction mechanisms : for instance, two and three body meson exchange currents strongly affect the Transverse-Transverse interference cross section but contribute weakly to the Transverse-Longitudinal cross section. While their determination would certainly be a way to refine our understanding of the $(e,e'p)$ reaction, it requires out of the plane experiments. The determination of spin observables provides us with a way to avoid the related experimental problems.

3.3 Spin observables.

Spin observables depend upon various interferences between the Transverse and the Longitudinal components of the nuclear current (see Appendix A). They can be classified into two categories. The first vanish in Plane Wave Impulse Approximation (PWIA) : their determination is a good way to study and calibrate many body and rescattering mechanisms which contribute beyond quasi free scattering. The second do not vanish in PWIA : they are directly related to the nucleon electromagnetic form factors and to the nuclear wave function. However, it turns out that, in the particular case of the ${}^4\text{He}(\vec{e}, e'\vec{p})T$ reaction, the ground state wave function factorizes in the various cross sections and does not contribute to the spin observables in PWIA (see Ref.[29]). Contrary to the ${}^3\text{He}(\vec{e}, e'\vec{p})D$ reaction, where the S and D parts of the $\langle {}^3\text{He}|D\rangle$ overlap integral interfere [30], spin observables are of no use to determine the pure S wave $\langle {}^4\text{He}|T\rangle$ overlap integral. They could be rather used to determine

the behavior of the form factors of a bound nucleon, provided that many body and rescattering contributions are suppressed. This is the case in Fig.15, which shows the variation of the three proton polarizations which do not vanish when the proton is emitted along the direction of the virtual photon, in the same kinematics as in Fig.14 where the large value of the photon momentum strongly reduces the effects of many body mechanisms. In such collinear kinematics, the sideways (P'_X) and the longitudinal (P'_Z) transferred polarizations depends on the ratio between the electric and the magnetic form factors of the struck proton : this is the good place to study their behavior in the nuclear medium. The normal polarization (P'_Y), which is measured with unpolarized electrons, vanishes in PWIA and provides us with a way to calibrate the various corrections.

The contributions of many body and rescattering mechanisms are more important at a lower value of the photon three-momentum. Fig.16 shows the variations of the proton polarizations which do not vanish in coplanar geometry ($\Phi = 180^\circ$). The kinematics is the same as in Fig.13 and is typical of the Amsterdam or Mainz facilities.

For the sake of completeness, Fig.17 shows the electron asymmetry A_e and the proton transferred polarization P'_Y which vanish in coplanar kinematics, since they are proportional to $\sin \Phi$. Both depend only on a transverse-longitudinal interference cross-section, but A_e vanishes in PWIA while P'_Y does not. They exhibit a different sensitivity to the various interaction effects. The azimuthal angle is chosen as $\Phi = 90^\circ$: for symmetry reason $P'_Y = P'_X$ in collinear kinematics ($\Theta_p = 0$).

Figs. 18 and 19 show the sensitivity, to charge exchange nucleon rescattering, of selected spin observables in the reactions ${}^4\text{He}(e, e'p)T$ and ${}^4\text{He}(e, e'n){}^3\text{He}$ respectively, in kinematics already achieved at Amsterdam when determining unpolarized cross-sections. As expected the induced normal polarization P'_Y is more sensitive, and the effects are larger in the $(e, e'n)$ channel.

Finally, Fig. 20 shows the sensitivity to the small neutron electric form-factor of the various spin observables in the ${}^4\text{He}(\vec{e}, e'\vec{n}){}^3\text{He}$ reaction. Here multinucleon mechanisms completely change the simple and elegant picture of the PWIA treatment : when the target nucleon is at rest, the sign of the outgoing neutron polarization P'_X is reversed!. This is a good illustration of the fact that the medium modifications of the bound nucleon form factors and the various interaction effects must be treated on the same footing : any attempt to disentangle them might be irrelevant.

To summarize, at low momentum transfer, spin observables provide us with further constraints on the reaction mechanisms. At high momentum transfer, the contribution of many body and rescattering mechanisms are strongly suppressed and spin observables provide us with a way to study the behavior of the nucleon form factors in the nuclear medium.

4 Conclusion.

We have reached a fair understanding of the dominant mechanisms in the two-body photodisintegration and electrodisintegration of ${}^4\text{He}$. Although the diagrammatic method, which was used, relies on a truncated multiple scattering series and does not treat the final state with the same accuracy as the ground state, the degree to which gauge invariance is fulfilled provides us with a way to estimate the theoretical uncertainties. This method allows us to study the relative contributions of the various reaction mechanisms. It offers us a powerful way to select kinematics where each of them dominates and can be calibrated.

The photodisintegration channels are dominated by two and three body mechanisms. The Δ formation terms dominate and the nature prefers to share the momentum transfer among three nucleons rather to transfer it to a single nucleon. The model leads to a fair agreement with the existing set of data up to $E_\gamma = 500$ MeV and reproduces the near equality of the cross-section of the (γ, p) and (γ, n) channels.

The electrodisintegration channels are dominated by one body mechanisms, when the recoil momentum is not too high ($P_R \leq 200$ MeV/c). However, nucleon rescattering mechanisms are not negligible, while the contribution of two and three body meson exchange mechanisms are less important than in photodisintegration. Although the contribution of all these interaction mechanisms always prevents an accurate determination of the $\langle {}^4\text{He}|T \rangle$ overlap integral in the vicinity of its zero ($P_R \simeq 470$ MeV/c), it is strongly suppressed when the three momentum of the virtual photon increases above 1 GeV/c. This is the kinematical domain of CEBAF, where the separation between the Transverse and the Longitudinal parts of the cross section should be performed in order to further suppress the contribution of two and three body exchange currents.

In the particular case of the ${}^4\text{He}(\vec{e}, e'\vec{p})T$ reaction, spin observables are not at all sensitive (in PWIA) to the $\langle {}^4\text{He}|T \rangle$ overlap integral. At low momentum transfer (photon momentum of the order of 400 MeV/c), they provide us with further constraints on reaction mechanisms. At momentum higher than 1 GeV/c, the contribution of interaction effects is very small and spin observables are only sensitive to the nucleon form factors : they provide us with a powerful way to study their behavior in the nuclear medium. However, at such high momentum it could be likely that medium modifications are not large and relativistic effects should be taken into account in a more systematic way.

Recent progresses in the quality and the reliability of polarized electron beams and of polarimeters open up a new field. The determination of the relevant spin observables should be systematically undertaken at Amsterdam and Mainz, and completed at CEBAF.

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APPENDICES

A Cross sections and spin observables

The general expression of the cross section, of the $(\vec{e}, e' \vec{N})$ exclusive reaction induced by a polarized electron, when the polarization of the outgoing nucleon is determined, can be cast in the form [25,30,31] :

$$\frac{d\sigma(h, \vec{S})}{d\Omega_e dE_e d\Omega_N} = \Gamma_v \frac{d\sigma^0}{d\Omega_N} \frac{1}{2} \left[1 + \vec{S} \cdot \vec{P}^0 + h (A_e + \vec{S} \cdot \vec{P}') \right] \quad (\text{A.1})$$

where h is the helicity of the electron, $\vec{S}$ the direction along which the polarization of the nucleon is determined, σ^0 the unpolarized cross section, $\vec{P}^0$ the nucleon polarization when the beam is unpolarized, A_e the electron asymmetry when the nucleon polarization is not determined and $\vec{P}'$ the spin transfer polarization to the nucleon.

The flux of the virtual photon is :

$$\Gamma_v = \frac{\alpha}{2\pi^2} \frac{E'}{E} \frac{|\vec{k}|}{1-\epsilon} \frac{1}{-q^2} \quad (\text{A.2})$$

and the degree of polarization of the virtual photon is :

$$\epsilon = \left[1 - 2 \frac{\vec{k}^2}{q^2} \tan^2 \frac{\theta}{2} \right]^{-1} \quad (\text{A.3})$$

The fine structure constant is α . The energy of the incident electron is E while the energy and the polar angle of the scattered electron are respectively E' and θ . The four-momentum of the virtual photon is $q = (\omega, \vec{k})$, with $q^2 = \omega^2 - \vec{k}^2$.

The quantization axis lies along the direction of the momentum of the virtual photon. When the X component lies in the nucleon emission plane, the various cross-sections and polarizations take the following forms.

The unpolarized cross-section is made of four parts :

$$\frac{d\sigma^0}{d\Omega_N} = \frac{d}{d\Omega_N} \left[\sigma_T + \epsilon \sigma_L + \epsilon \cos 2\phi \sigma_{TT} - \sqrt{\frac{-q^2 \epsilon (\epsilon + 1)}{2\omega^2}} \cos \phi \sigma_{TL} \right] \quad (\text{A.4})$$

where ϕ is the angle between the electron scattering plane and the nucleon emission plane. It vanishes when the nucleon is emitted on the same side as the scattered electron, with respect to the direction of the virtual photon. Each cross-section is related to the relevant components of the unpolarized hadronic tensor

$$W_{\mu\nu} = \sum_{m_1, m_T} \langle m_1, m_T | J_\nu^* | 0 \rangle \langle m_1, m_T | J_\mu | 0 \rangle \quad (\text{A.5})$$

where the hadronic current $J = (J^0, \vec{J})$ is the coherent sum of the matrix elements $T^{(i)}$, which are given in the main body of this paper

$$\langle m_1, m_T | J | 0 \rangle = \sum_i T^i(\vec{k}, 0 \rightarrow \vec{p}_1 m_1, \vec{p}_T m_T) \quad (\text{A.6})$$

Namely :

$$\begin{aligned} \sigma_T &= K (W_{XX} + W_{YY}) \\ \sigma_L &= 2K \frac{-q^2}{\omega^2} W_{ZZ} \\ \sigma_{TT} &= K (W_{XX} - W_{YY}) \\ \sigma_{TL} &= 2K (W_{ZX} + W_{XZ}) \end{aligned} \quad (\text{A.7})$$

where K is the relativistic phase space factor. In the c.m. frame of the emitted proton-triton system (of which W is the invariant mass) it takes the simple form :

$$K_{c.m.} = \frac{1}{(2\pi)^2} \frac{m M_T M_{He}}{4W^2} \left[\frac{|\vec{p}_1|}{|\vec{k}|} \right]_{c.m.} \quad (\text{A.8})$$

where m , M_T and M_{He} are respectively the masses of the nucleon, the emitted Triton and the target 4He . In the Lab. system it takes the form :

$$K_{Lab.} = K_{c.m.} \frac{d\Omega_{c.m.}}{d\Omega_{Lab.}} \quad (\text{A.9})$$

Throughout this paper, all the matrix elements of the hadronic tensor have been computed in the Lab. frame. The electrodisintegration cross sections are given and plotted in the Lab. frame, while the (real or virtual) photon absorption cross sections are given are plotted in the c.m. frame.

For symmetry reasons, σ_{TT} and σ_{TL} vanish in collinear kinematics. At the photon point ($q^2 = 0$), σ_T reduces to the photon absorption cross-section, while σ_L vanishes.

The induced proton polarization, which arises when the electron is not polarized, has the following normal (P_Y^0 , measured along the direction Y normal to the nucleon emission plane), sideways (P_X^0 , measured along the direction X normal to the direction of the incoming virtual photon in the nucleon emission plane) and

longitudinal (P_Z^0 , parallel to the virtual photon direction) cartesian components in the Lab. frame :

$$\begin{aligned}
P_Y^0 \frac{d\sigma^0}{d\Omega_N} &= \frac{d}{d\Omega_N} [\sigma_T(Y) + \epsilon \sigma_L(Y) + \epsilon \cos 2\phi \sigma_{TT}(Y) \\
&\quad - \sqrt{\frac{-q^2 \epsilon (\epsilon + 1)}{2\omega^2}} \cos \phi \sigma_{TL}(Y)] \\
P_{X,Z}^0 \frac{d\sigma^0}{d\Omega_N} &= \frac{d}{d\Omega_N} \left[\epsilon \sin 2\phi \sigma_{TT}(X, Z) - \sqrt{\frac{-q^2 \epsilon (\epsilon + 1)}{2\omega^2}} \sin \phi \sigma_{TL}(X, Z) \right]
\end{aligned} \tag{A.10}$$

The various cross sections are expressed in terms of the relevant components of the polarized hadronic tensor :

$$\vec{W}_{\mu\nu} = \sum_{m_1, m_1', m_T} \langle m_1, m_T | J_\nu^* | 0 \rangle \vec{\sigma}_{m_1, m_1'} \langle m_1', m_T | J_\mu | 0 \rangle \tag{A.11}$$

where $\vec{\sigma} = (\sigma(X), \sigma(Y), \sigma(Z))$ are the cartesian components of the Pauli matrices. Namely :

$$\begin{aligned}
\sigma_T(Y) &= K (W_{XX}(Y) + W_{YY}(Y)) \\
\sigma_L(Y) &= 2K \frac{-q^2}{\omega^2} W_{ZZ}(Y) \\
\sigma_{TT}(X) &= K (W_{XY}(X) + W_{YX}(X)) \\
\sigma_{TT}(Y) &= K (W_{XX}(Y) - W_{YY}(Y)) \\
\sigma_{TT}(Z) &= K (W_{XY}(Z) + W_{YX}(Z)) \\
\sigma_{TL}(X) &= 2K (W_{ZY}(X) + W_{YZ}(X)) \\
\sigma_{TL}(Y) &= 2K (W_{ZX}(Y) + W_{XZ}(Y)) \\
\sigma_{TL}(Z) &= 2K (W_{ZY}(Z) + W_{YZ}(Z))
\end{aligned} \tag{A.12}$$

Similarly, the transferred polarization takes the form :

$$\begin{aligned}
P_Y' \frac{d\sigma^0}{d\Omega_N} &= \frac{d}{d\Omega_N} \left[-\sqrt{\frac{-q^2 \epsilon (1 - \epsilon)}{2\omega^2}} \sin \phi \sigma'_{TL}(Y) \right] \\
P_{X,Z}' \frac{d\sigma^0}{d\Omega_N} &= \frac{d}{d\Omega_N} \left[\sqrt{(1 - \epsilon^2)} \sigma'_{TT}(X, Z) - \sqrt{\frac{-q^2 \epsilon (1 - \epsilon)}{2\omega^2}} \cos \phi \sigma'_{TL}(X, Z) \right]
\end{aligned} \tag{A.13}$$

with

$$\begin{aligned}
\sigma'_{TT}(X) &= iK (W_{YX}(X) - W_{XY}(X)) \\
\sigma'_{TT}(Z) &= iK (W_{YX}(Z) - W_{XY}(Z)) \\
\sigma'_{TL}(X) &= 2iK (W_{YZ}(X) - W_{ZY}(X)) \\
\sigma'_{TL}(Y) &= -2iK (W_{XZ}(Y) - W_{ZX}(Y)) \\
\sigma'_{TL}(Z) &= 2iK (W_{YZ}(Z) - W_{ZY}(Z))
\end{aligned} \tag{A.14}$$

Finally, the beam asymmetry takes the form :

$$A_e \frac{d\sigma^0}{d\Omega_N} = \frac{d}{d\Omega_N} \left[-\sqrt{\frac{-q^2\epsilon(1-\epsilon)}{2\omega^2}} \sin \phi \sigma_{TL}^e \right] \quad (\text{A.15})$$

with

$$\sigma_{TL}^e = iK [W_{XZ} - W_{ZX}] \quad (\text{A.16})$$

Note that σ_{TL} and σ_{TL}^e are respectively the real and imaginary parts of the same quantity, and that σ_{TL}^e vanishes in Plane Wave Impulse Approximation.

Current conservation ($\omega J^0 = \vec{k} \cdot \vec{J}$) has been used to eliminate the time component in the previous expressions : this is the usual choice in particle physics. If, instead, the third component J_Z of the current had been eliminated, it would have been replaced by $\omega J^0 / |\vec{k}|$ in those equations : this choice is usually done in nuclear physics, and it has been used along this paper, unless otherwise stated. Since the current is conserved, both prescriptions must lead to the same result.

All these expressions of the ejected nucleon polarization are valid in the Lab. frame, when the Z axis is parallel to the virtual photon direction and the X axis lies in the nucleon emission plan. In an actual experiment, the Z axis of the polarimeter is usually aligned along the direction of the emitted nucleon : a trivial rotation against the Y axis has been performed in order to express, in this new frame, the polarizations given in this paper.

If, in addition, the X axis of the polarimeter does not rotate out of the plane in the same way as the nucleon emission plane, and for instance stays parallel to the electron scattering plane, additional dependence on $\sin \phi$ and $\cos \phi$ appear in the expressions of the polarizations. The relevant formulae depend on the way the experiment is performed and are not given here. This point has to be kept in mind especially when dealing with the finite angular acceptance of the polarimeter

The determination of the angular distributions of the unpolarized cross sections and the proton polarizations allows the various structure functions, which they are made of, to be disentangled. A few relevant kinematics are particularly interesting.

In coplanar geometry (when the ejectile lies in the electron scattering plane) :

- $P_Y' = P_Z^0 = P_X^0 = 0$;
- $P_Y^0 = 0$ in Plane Wave Impulse Approximation;
- $P_Y^0 \neq 0$ if Final State Interactions or Meson Exchange Currents are taken into account;
- $P_X' \neq 0$ and $P_Z' \neq 0$.

In collinear kinematics, when the ejectile is emitted along the direction of the virtual photon, P'_X depend only on the transverse-longitudinal interference $\sigma'_{TL}(X)$, whereas P'_Z depend only on the transverse-transverse interference $\sigma'_{TT}(Z)$. The measurement of P'_X in collinear kinematics is therefore a good way to enhance the small contribution of the Coulomb part of the cross section, using the transverse part as an amplifier. This idea is at the basis of the determination of the small electric form factor of the neutron and of the small Coulomb form factor of the $N \rightarrow \Delta$ transition or the $N \rightarrow$ Roper resonance transition.

B The nucleon current

At order $1/m^3$ the space component of the nucleon current $J = (J^0, \vec{J})$ takes the form :

$$\begin{aligned} \vec{J} = & -i \frac{\vec{P}}{2m} \left[F_1(q^2) - F_2(q^2) \frac{\vec{k}^2}{8m^2} \right] + i \frac{\omega \vec{k}}{8m^2} G_M(q^2) \\ & + \frac{\vec{\sigma} \times \vec{k}}{2m} \left[G_M(q^2) + F_2(q^2) \frac{\vec{P}^2}{8m^2} \right] - \frac{\omega}{8m^2} \vec{\sigma} \times \vec{P} [2F_2(q^2) + F_1(q^2)] \\ & - F_2(q^2) \frac{\vec{\sigma} \cdot \vec{P} \vec{P} \times \vec{k}}{16m^3} + O\left(\frac{1}{m^4}\right) \end{aligned} \quad (\text{B.1})$$

where terms in $\omega/(8m^2)$ are in fact of order $1/(m^3)$, since ω is of order $1/m$.

The time component takes the form :

$$\begin{aligned} J_0 = & -i F_1(q^2) \left[1 + \frac{\vec{p}_1^2 + \vec{p}^2}{4m^2} \right] + i \frac{\vec{k}^2}{8m^2} [F_1(q^2) + 2F_2(q^2)] \\ & + i \frac{\vec{\sigma} \cdot \vec{P} \times \vec{k}}{8m^2} [F_1(q^2) + 2F_2(q^2)] + O\left(\frac{1}{m^4}\right) \end{aligned} \quad (\text{B.2})$$

These expressions have to be used when the norm of the Dirac spinors is $\bar{u}u = 1$. The nucleon mass is m . The momenta $\vec{p}$ of the initial and $\vec{p}_1$ of the final nucleons are related as follows :

$$\begin{aligned} \vec{p}_1 &= \vec{p} + \vec{k} \\ \vec{P} &= \vec{p} + \vec{p}_1 \\ &= 2\vec{p}_1 - \vec{k} \end{aligned} \quad (\text{B.3})$$

The Dirac form factors F_1 and F_2 are related to the Sachs form factors G_E and G_M in the following way :

$$\begin{aligned} G_M(q^2) &= F_1(q^2) + F_2(q^2) \\ G_E(q^2) &= F_1(q^2) + \frac{q^2}{4m^2} F_2(q^2) \end{aligned} \quad (\text{B.4})$$

where $F_1^n(0) = 0$, $F_1^p(0) = e$, $F_2^n(0) = e\kappa_n$ and $F_2^p(0) = e\kappa_p$, being e , κ_p and κ_n the charge and the anomalous moments of the proton and the neutron. G_E and G_M are parametrized by the dipole fit.

For free nucleon, this current is gauge invariant up to and including order $1/m^3$. This is not anymore the case when nucleons are off-shell. The inclusion of Meson Exchange Currents is necessary to restore gauge invariance.

C The two nucleon matrix elements

The expressions of the two-body matrix element of $\gamma(np)_0 \rightarrow np$ has been derived in [32]. The extension to virtual photons has been performed in [33], where the treatment of Meson Exchange Currents has been improved. A more complete treatment of Final State Interactions has been achieved in [5], where the model has been calibrated against the latest results of the experimental studies of the $D(e, e'p)n$ and ${}^3\text{He}(e, e'p)D$ channels. The expression of the Δ contribution to the $\pi^+(np)_0 \rightarrow pp$ amplitude has been derived in [18]. This appendix summarizes these works and gathers the expressions of the two body matrix elements needed in the anlysis of the ${}^4\text{He}(e, e'p)T$ channel.

C.1 Meson Exchange Currents

For a np pair of isospin $T = 0$, moving with a momentum $-\vec{\xi}$ and a magnetic quantum number M_J , the pion exchange antisymmetrized matrix element takes the form :

$$\begin{aligned} \tilde{T}_\pi(\vec{k}, -\vec{\xi} 1M_J \rightarrow \vec{p}_1 m_1, \vec{n} m_n) &= T_\pi(\vec{k}, -\vec{\xi} 1M_J \rightarrow \vec{p}_1 m_1, \vec{n} m_n) \\ &\quad - T_\pi(\vec{k}, -\vec{\xi} 1M_J \rightarrow \vec{n} m_n, \vec{p}_1 m_1) \end{aligned} \quad (\text{C.1})$$

where[32] :

$$\begin{aligned} T_\pi(\vec{k}, -\vec{\xi} 1M_J \rightarrow \vec{p}_1 m_1, \vec{n} m_n) &= \frac{g_0}{2m} \sum_{\lambda_p \lambda_n} \int \frac{d^3 \vec{n}'}{(2\pi)^3} \left(m_n \mid \vec{\sigma} \cdot \left[\vec{q} - \frac{\vec{n}' + \vec{n}}{2m} q^0 \right] \mid \lambda_n \right) \\ [F_\pi(q_\pi^2)]^2 \frac{-T(\gamma p \rightarrow p\pi^0) - \sqrt{2}T(\gamma n \rightarrow p\pi^-)}{q_\pi^2 - m_\pi^2 + i\epsilon} &\left\{ \frac{1}{4\pi} U_0(\mid \vec{n}' + \frac{\vec{\xi}}{2} \mid) \left(\frac{1}{2} \lambda_p \frac{1}{2} \lambda_n \mid 1M_J \right) \right. \\ &\quad \left. + U_2(\mid \vec{n}' + \frac{\vec{\xi}}{2} \mid) \sum_{m_l m_s} \left(\frac{1}{2} \lambda_p \frac{1}{2} \lambda_n \mid 1m_s \right) (2m_l 1m_s \mid 1M_J) Y_2^{m_l}(\widehat{\vec{n}' + \frac{\vec{\xi}}{2}}) \right\} \end{aligned} \quad (\text{C.2})$$

The momenta and the magnetic quantum numbers of the active proton and neutron (in the target pair) are respectively $(\vec{p}', \lambda_p)$ and $(\vec{n}', \lambda_n)$. The squared four momentum and the mass of the exchanged pion are respectively $q_\pi^2 = (q^0)^2 - \vec{q}^2$ and m_π . Energy and momentum are conserved at each vertex, for instance : $-\vec{\xi} = \vec{p}' + \vec{n}'$, $\vec{n} = \vec{n}' + \vec{q}$ and $\vec{k} + \vec{p}' = \vec{p}_1 + \vec{q}$. The S and D parts of the relative wave function of the two active nucleons are respectively U_0 and U_2 , and are obtained by projecting the four-body wave function [4] on the relevant partition. The three-fold integral is performed numerically.

The spatial part of the current corresponding to the combination of the elementary nucleon amplitude is derived from [20,33] :

$$\begin{aligned} \tilde{T}_{\gamma N \rightarrow N\pi}(\vec{p}'\lambda_p \rightarrow \vec{p}_1 m_1) &= -T(\gamma p \rightarrow p\pi^0) - \sqrt{2}T(\gamma n \rightarrow p\pi^-) \\ &= -\frac{4}{3} \frac{G_1}{2M} G_3 G_\Delta(q^2) \frac{(m_1 | \vec{S} \cdot [\vec{q} - \vec{R}_M^2] \vec{S}^\dagger \times [\vec{k} - \vec{p}' \frac{M-m}{m}] | \lambda_p)}{R^0 - E_R + \frac{1}{2}i\Gamma} \\ &\quad -i \frac{g_0}{m} \left(m_1 | \vec{\sigma} G_A(q^2) + \frac{\vec{\sigma} \cdot (\vec{p}_1 - \vec{p}') \vec{q}}{(\vec{q} - \vec{k})^2 + m_\pi^2} G_\pi(q^2) | \lambda_p \right) \\ &\quad + \frac{1}{m_\pi} \frac{g_{\omega_1} g_{\omega_2}}{(\vec{q} - \vec{k})^2 + m_\omega^2} \vec{q} \times \vec{k} \delta_{\lambda_p m_1} G_\omega(q^2) \end{aligned} \quad (C.3)$$

The contributions of the Δ formation term and of each Born (contact, pion photoelectric and ω exchange) term can be easily identified. The various coupling constants, the mass M and the width Γ of the Δ and the various electromagnetic form factors are chosen as in [20]).

The small Coulomb part has been neglected, except in Fig. 9 where it has been determined according the expressions of the elementary currents given in [33].

It should be noted that the nucleon Born terms have been removed : they are already contained in the wave function of the initial and the final state (see [32]). Also the pion photoelectric term has been divided by two, in order to avoid double counting when antisymmetrizing the amplitude (see [33]).

The ρ meson exchange amplitude is easily obtained by making the substitution

$$[F_\pi(q_\pi^2)]^2 \frac{\vec{S} \cdot \vec{q} \vec{\sigma} \cdot \vec{q}}{q_\pi^2 - m_\pi^2} \rightarrow \frac{G_\rho^2}{G_\pi^2} [F_\rho(q_\rho^2)]^2 \frac{\vec{q}^2 \vec{S} \cdot \vec{\sigma} - \vec{S} \cdot \vec{q} \vec{\sigma} \cdot \vec{q}}{q_\rho^2 - m_\rho^2} \quad (C.4)$$

in the Δ part of Eq. C.3, and neglecting the small contribution [32] of the Born terms. The ratio of the relevant combinations of π -baryon and ρ -baryon coupling constants is chosen as :

$$\frac{G_\rho^2}{G_\pi^2} = 1.8 \quad (C.5)$$

In order to take into account the off-shell behaviour of the exchanged meson, a

monopole form factor is used at each π -baryon vertex :

$$\begin{aligned} F_\pi(q^2) &= \frac{\Lambda_\pi^2 - m_\pi^2}{\Lambda_\pi^2 - q^2} \\ \Lambda_\pi &= 1200 \text{ MeV} \end{aligned} \quad (\text{C.6})$$

and a dipole form factor at each ρ -baryon vertex :

$$\begin{aligned} F_\rho(q^2) &= \left[\frac{\Lambda_\rho^2 - m_\rho^2}{\Lambda_\rho^2 - q^2} \right]^2 \\ \Lambda_\rho &= 2m \end{aligned} \quad (\text{C.7})$$

For an active NN pair with isospin $T = 1$, the expressions of the meson exchange current amplitudes are easily deduced. The relative wave functions are to be replaced by the relevant projection, of the four-body wave function, on the corresponding partition. It turns out that the S wave component dominates (this is the only component which survives in a pair at rest), and the contributions of the other components are neglected. Due to different isospin couplings the combination of elementary amplitudes is different. For an initial pn pair, the Δ parts cancel and only Born terms survive. On the contrary, for an initial pp pair only neutral mesons can be exchanged and only the Δ and ω exchange parts survive. However, spin-isospin selection rules prevent a $p\Delta^+$ pair, in a relative S wave, to decay into a pp pair. The photodisintegration of a pp pair is therefore strongly suppressed [11,34] and contribute little in the two-body disintegration of ${}^4\text{He}$.

C.2 Nucleon rescatterings

For an active pn pair with isospin $T = 0$, the antisymmetrized nucleon rescattering amplitude takes the following form, according to the isospin T' of the final pn pair :

$$\begin{aligned} \tilde{T}_{pn}(\vec{k}, -\xi | 1M_J \rightarrow \vec{p}_1 m_1, \vec{n} m_n) &= m \sum_{\lambda_p \lambda_n} \int \frac{d^3 \vec{n}'_{c.m.}}{(2\pi)^3} \frac{1}{\vec{p}_{c.m.}^2 - \vec{n}'_{c.m.}^2 + i\epsilon} \\ T_{NN}^{T'=0,1}(-\vec{n}'_{c.m.}, \lambda_p \vec{n}'_{c.m.}, \lambda_n \rightarrow \vec{p}_1 m_1 \vec{n} m_n) &\sum_{m_l m_s} (\lambda_p | J_p(q^2) + (-)^{T'} J_n(q^2) | m_s - \lambda_n) \\ &\times \left\{ \frac{1}{4\pi} U_0(|\vec{n}'_{c.m.} + \frac{\vec{k}}{2}|) (\frac{1}{2} \lambda_n \frac{1}{2} (m_s - \lambda_n) | 1m_s) \delta_{M_J m_s} \delta_{m_l 0} \right. \\ &\left. + U_2(|\vec{n}'_{c.m.} + \frac{\vec{k}}{2}|) (\frac{1}{2} \lambda_n \frac{1}{2} (m_s - \lambda_n) | 1m_s) (2m_l 1m_s | 1M_J) Y_2^{m_l}(\widehat{|\vec{n}'_{c.m.} + \frac{\vec{k}}{2}|}) \right\} \end{aligned} \quad (\text{C.8})$$

In the center of mass (c.m.) of the scattering pair, the off-shell and on-shell nucleon momenta are respectively $\vec{n}'_{c.m.}$ and $\vec{p}_{c.m.}$, with

$$|\vec{p}_{c.m.}| = \frac{W^2 - 4m^2}{2m} \quad (C.9)$$

where W is the invariant mass of the pair. The integral runs over the momentum of the spectator on-shell nucleon in the loop.

The space and time components of the combination $J_{pn} = J_p + (-)^{T'} J_n$ of the elementary nucleon current are the following :

$$\begin{aligned} \vec{J}_{pn} &= i \frac{\vec{n}'_{c.m.}}{m} \left[\frac{1 + \tau_3}{2} \right] G_E^p(q^2) + \frac{\vec{\sigma} \times \vec{k}}{2m} [G_M^p(q^2) + (-)^{T'} G_M^n(q^2)] \\ J_{pn}^0 &= -i G_E^p(q^2) \end{aligned} \quad (C.10)$$

The half-off shell nucleon-nucleon scattering amplitude is expanded in terms of partial waves :

$$\begin{aligned} T_{NN}^{T'=0,1}(-\vec{n}'_{c.m.} \lambda_p \vec{n}'_{c.m.} \lambda_n \rightarrow \vec{p}_1 m_1 \vec{n} m_n) &= -\frac{8\pi^2 W}{|\vec{p}_{c.m.}| m^2} \sum_{LL'SJ} \sum_{\mu\mu'\sigma\sigma'} \left(\frac{1}{2} \lambda_p \frac{1}{2} \lambda_n \mid S\sigma' \right) \\ &\quad (L'\mu' S\sigma' \mid IM_I) \left(\frac{1}{2} m_1 \frac{1}{2} m_n \mid S\sigma \right) (L\mu S\sigma \mid IM_I) Y_L^\mu(\vec{n}) Y_{L'}^{\mu'*}(\vec{n}') (LSI \mid f_{off}^{0,1} \mid L'SI) \end{aligned} \quad (C.11)$$

In each partial wave, the half-off shell and the on-shell scattering amplitudes are related as follows :

$$f_{off}(\vec{n}'^2, \vec{p}^2) = f_{on}(\vec{p}^2, \vec{p}^2) F(\vec{n}'^2, \vec{p}^2) \quad (C.12)$$

where the form factor $F(\vec{n}'^2, \vec{p}^2)$ is nothing but the ratio of the numerical off- and on-shell solutions of the Lippman-Schwinger equation. As explained in [34] it is parametrized, at each energy, as a series of monopoles :

$$F(\vec{n}'^2, \vec{p}^2) = \sum_i \frac{c_i(\vec{p}^2)}{\vec{n}'^2 + \beta_i^2(\vec{p}^2)} \quad (C.13)$$

The coefficients c_i and β_i depend on the on-shell energy.

The wave functions of the active np pair are also parametrized as a sum of monopoles :

$$\begin{aligned} U_0(p) &= 4\pi N \sum_i \frac{a_i}{p^2 + \alpha_i^2} \\ \sum_i a_i &= 0 \\ U_2(p) &= 4\pi N \sum_i \frac{b_i}{p^2 + \gamma_i^2} \\ \sum_i b_i &= \sum_i b_i \gamma_i^2 = \sum_i \frac{b_i}{\gamma_i^2} = 0 \end{aligned} \quad (C.14)$$

Under those assumptions, all the integrals can be performed analytically with the method outlined in [32].

For an active pn pair with isospin $T = 1$, only the S wave component is to be considered and the isospin couplings change accordingly. For instance, the combination of the elementary electromagnetic currents becomes :

$$J_{pn}(q^2) = J_p(q^2) - (-)^{T'} J_n(q^2) \quad (\text{C.15})$$

For a pp active pair it becomes :

$$J_{pp'}(q^2) = J_p(q^2) + J_{p'}(q^2) \quad (\text{C.16})$$

It should be noted that in [33] only the transitions between the S and D part of the wave function of the active pair and the S part of the final state, together with the transition between the S and P part, have been considered. On the contrary, all the transitions between the S and D components of the bound state and the S , P and D components of the final state have been considered in [5] and in the present work.

C.3 Gauge invariance

All the above matrix elements have been calibrated against the latest experimental values of the cross-section of the $D(e, e'p)n$ reaction [5], and selected values of the cross-section of the $D(\gamma, p)n$ reaction [33]. As it is demonstrated in Fig. 21 the model satisfies gauge invariance in this channel. One body mechanisms lead to a very different result whether the time component (J_0) or the third component of the current (J_z) is used. The agreement is restored when two-body mechanisms are taken into account. The reason is that the wave functions of the bound state and of the scattering state are solutions of the same Lippman-Schwinger equation for the Paris potential, and that Meson Exchange Currents are related in a consistent way to this potential. A schematic formal demonstration can be found in an appendix of Ref. [33].

C.4 Pion absorption

For the absorption of a π^0 by a $T = 0$ pn pair at rest, the antisymmetrized matrix element takes the form :

$$\begin{aligned} \tilde{T}_{\pi^0(np)_0}(\vec{q}_\pi, \vec{\nu} - M_J \rightarrow \vec{n}m_n, \vec{p}_1m_1) &= T_{\pi^0(np)_0}(\vec{q}_\pi, \vec{\nu} - M_J \rightarrow \vec{n}m_n, \vec{p}_1m_1) \\ &\quad - T_{\pi^0(np)_0}(\vec{q}_\pi, \vec{\nu} - M_J \rightarrow \vec{p}_1m_1, \vec{n}m_n) \quad (\text{C.17}) \end{aligned}$$

where[18] :

$$\begin{aligned}
T_{\pi^0(np)_0}(\vec{q}_\pi, \vec{v} - M_J \rightarrow \vec{p}_1 m_1, \vec{n} m_n) &= \frac{g_0}{2m} \sum_{\lambda_p \lambda_n} \int \frac{d^3 \vec{n}'}{(2\pi)^3} \left(m_n \mid \vec{\sigma} \cdot \left[\vec{q} - \frac{\vec{n}' + \vec{n}}{2m} q^0 \right] \mid \lambda_n \right) \\
[F_\pi(q_\pi^2)]^2 \frac{-T(\pi^0 p \rightarrow p \pi^0) - \sqrt{2} T(\pi^0 n \rightarrow p \pi^-)}{q_\pi^2 - m_\pi^2 + i\epsilon} &\left\{ \frac{1}{4\pi} U_0 \left(\mid \vec{n}' - \frac{\vec{v}}{2} \mid \right) \left(\frac{1}{2} \lambda_p \frac{1}{2} \lambda_n \mid 1 - M_J \right) \right. \\
&\left. + U_2 \left(\mid \vec{n}' - \frac{\vec{v}}{2} \mid \right) \sum_{m_l m_s} \left(\frac{1}{2} \lambda_p \frac{1}{2} \lambda_n \mid 1 m_s \right) (2 m_l 1 m_s \mid 1 - M_J) Y_2^{m_l}(\widehat{\vec{n}' - \frac{\vec{v}}{2}}) \right\}
\end{aligned} \tag{C.18}$$

The combination of elementary πN elastic scattering amplitude is splitted into the dominant Δ formation amplitude, which is treated analytically, and the other partial waves, which are parametrized in terms of the corresponding phase shifts, as follows :

$$\begin{aligned}
\tilde{T}_{\pi N \rightarrow N \pi}(\vec{q}_\pi, \vec{p}' \lambda_p \rightarrow \vec{p}_1 m_1, \vec{q}) &= -T(\pi^0 p \rightarrow p \pi^0) - \sqrt{2} T(\pi^0 n \rightarrow p \pi^-) \\
&= -\frac{8}{3} G_3^2 \frac{(m_1 \mid \vec{S} \cdot [\vec{q} - \vec{R}_M^0] \vec{S}^\dagger \cdot [\vec{q}_\pi - \vec{R}_M^0] \mid \lambda_p)}{R^0 - E_R + \frac{1}{2} i \Gamma} \\
&\quad + \text{other partial waves}
\end{aligned} \tag{C.19}$$

The notations are the same as in the matrix element which describes the absorption of a photon by a nucleon pair. The coupling constants and the off-shell form factors are the same, and the ρ meson exchange mechanisms are treated in the same way.

References

- [1] S.E. Kiergan, A.O. Hanson and L.J. Koester, Phys. Rev. C8 (1973) 431;
P.E. Argan *et al.*, Nucl. Phys. A237 (1975) 447;
J. Arends *et al.*, Nucl. Phys. A322 (1979) 253;
R.A. Schumacher *et al.*, Phys. Rev. C33 (1986) 50;
R.T. Jones *et al.*, Phys. Rev. C43 (1991) 2052;
H. Schmieden *et al.*, Phys. Lett. B314 (1993) 284.
- [2] M. Gari and H. Hebach, Phys. Rep. 72 (1981) 1.
- [3] C.Y. Cheung and B.D. Keister, Phys. Rev. C33 (1986) 776.
- [4] R. Schiavilla, V. Pandharipande and R. Wiringa, Nucl. Phys. A449 (1986) 219.
- [5] J.M. Laget, Phys. Lett. 199B (1987) 493.
- [6] J.M. Laget, Phys. Rev. C38 (1988) 2993.
- [7] J. van den Brand *et al.*, Phys.Rev.Lett. 60 (1988) 2006.
- [8] J. van den Brand *et al.*, Phys.Rev.Lett. 66 (1991) 409; Nucl. Phys. A534 (1991) 637.
- [9] A. Magnon *et al.* Phys. Lett. B222 (1989) 352.
- [10] R. Schiavilla, Phys. Rev. Lett. 65 (1990) 835.
- [11] J.M. Laget, in "New Vistas in Electronuclear Physics". E.L. Tomusiak, H.S. Caplan and E.T. Dressler Eds., (Plenum Press, 1986), p.361.
- [12] R.B. Wiringa, R.A. Smith and T.L. Ainsworth, Phys. Rev. C29 (1984) 1207.
- [13] F-P. Schuster *et al.*, Phys. Rev. Lett. 55 (1985) 2261;
A. Amroun *et al* Phys. Rev. Lett. 69 (1992) 253.
- [14] J. van den Brand, Thesis (University of Amsterdam, 1988, unpublished), and to be published.
- [15] R. Piessens *et al.*, Quadpack (Springer, Berlin, 1983).
- [16] P. Weyer, Phys. Rep. 195 (1990) 297.
- [17] G. Audit *et al.*, Phys. Rev. C44 (1991) R575.
- [18] J.M. Laget, J.F. Lecolley and F. Lefevbre, Nucl. Phys. A370 (1981) 479.
- [19] J. Hamilton, in "High Energy Physics". E.H.S. Burhop, Ed. (Academic Press, NY, 1967) p.193.

- [20] I. Blomquist and J.M. Laget, Nucl. Phys. A280 (1977) 405;
J.M. Laget, Nucl. Phys. A481 (1998) 765.
- [21] J.M. Laget, J. Phys. G: Nucl. Phys. 14 (1988) 1445.
- [22] J. Thekkumthala *et al.*, Nucl. Phys. A455 (1986) 687.
- [23] J.P. Didelez and H. Langevin-Joliot, Il Nuovo Cimento, LXVII A (1970) 388.
- [24] M.A. Daman, Thesis (University of Amsterdam, 1991, unpublished), and to be published.
- [25] J.M. Laget, in “Modern Topics in Electron Scattering”. B. Frois and I. Sick Eds. (World Scientific, 1991), p.290.
- [26] J. Carlson, Nucl. Phys. A508 (1990) 141c.
- [27] J-E. Ducret *et al.*, Nucl. Phys. A556 (1993) 373.
- [28] J.M. Legoff *et al.*, in “4th Workshop on Perspectives in Nuclear Physics at Intermediate Energies”. S. Boffi *et al.*, eds. (World Scientific, 1989) p.376.
- [29] A. Picklesimer and J.W. Van Orden, Phys. Rev. C40 (1989) 290.
- [30] J.M. Laget, Phys. Lett. B276 (1992) 398.
- [31] J.M. Laget, Phys. Lett. B273 (1991) 765.
- [32] J.M. Laget, Nucl. Phys. A312 (1978) 265.
- [33] J.M. Laget, Can. J. Phys. 62 (1984) 1046.
- [34] J.M. Laget, Phys. Rev. C35 (1987) 832.
- [35] J.-E. Ducret *et al.*, Nucl. Phys. A533 (1993) 697c.

Figure Captions.

Fig.1 : The excitation functions of the ${}^4\text{He}(\gamma, p)\text{T}$ reaction at $\theta_p^{c.m.} = 60^\circ, 90^\circ$ and 120° are plotted against the energy E_γ of the incoming photon. One-body contributions (including Triton exchange) : dashed lines. Two-body contributions : dash-dotted lines. Three-body contributions : full lines. Final State Interactions are small on this scale. The symbols for experiments [1] are as follows : Urbana (Δ and $\blacksquare$), MIT ($\circ$), Mainz ($\blacktriangledown$), Saclay ($\bullet$).

Fig.2 : The one body mechanisms. The Triton exchange graph on the right comes from the antisymmetry among the nucleons.

Fig.3 : The two body mechanisms. Top part : π and ρ Meson Exchange Currents. Bottom part : nucleon-nucleon rescattering Final State Interactions. The diagrams at the right come from the antisymmetry of the two active nucleons.

Fig.4 : The three body mechanisms. The diagrams in the bottom part come from the antisymmetry of the two nucleons in the pair which absorbs the intermediate pion.

Fig.5 : The excitation functions of the ${}^4\text{He}(\gamma, n){}^3\text{He}$ reaction at $\theta_p^{c.m.} = 60^\circ, 90^\circ$ and 120° . The meaning of the curves, and of the symbols for experiments, is the same as in Fig.1. In addition Final State Interactions are included in the dotted lines.

Fig.6 : The angular distribution of the reaction ${}^4\text{He}(\gamma, p)\text{T}$ at $E_\gamma = 246$ MeV. The meaning of the curve, and of the symbols of photodisintegration data, is the same as in Fig. 1. In addition the experimental cross-section deduced from radiative capture [22] is also plotted ($\square$).

Fig.7 : The angular distribution of the reaction ${}^4\text{He}(\gamma, p)\text{T}$ at $E_\gamma = 131$ MeV. The meaning of the curves is the same as in Fig. 1. The experimental cross-sections have been deduced from radiative capture [23].

Fig.8 : Proton polarization [22] in the ${}^4\text{He}(\gamma, p)\text{T}$ reaction at $E_\gamma = 246$ MeV. The meaning of the curve is the same as in Fig. 1. Note that P_Y^0 vanishes in PWIA.

Fig.9 : The Coulomb (S_L) and Transverse (S_T) spectral functions, as extracted in the PWIA analysis [27] of the ${}^4\text{He}(e, e'p)\text{T}$ reaction in collinear kinematics, are plotted against the three momentum k of the virtual photon. The relative kinetic energy T_L between the outgoing proton and Triton is also plotted on abscissa. Open circles correspond to a recoil momentum $p_m = 90\text{MeV}/c$, open squares to $p_m = 30\text{MeV}/c$ and open triangles to $p_m = 190\text{MeV}/c$. Dotted curves correspond to the theoretical value $S_T = S_L = \rho(90\text{MeV}/c)$ of the variational wave function [Sc86] for the Argonne potential. Full circles (and full curves) correspond to the theoretical calculation which includes final state interaction and meson-exchange currents. The small contribution of meson-exchange currents to the Coulomb spectral function is

depicted as the dash-dotted curve. The dashed curve is obtained when the Coulomb spectral function is computed with the third component J_Z of the nuclear current.

Fig.10 : The ${}^4\text{He}(e,e'p)\text{T}$ reaction cross-section for the high recoil momentum (P_R) Saclay kinematics [28] : $E_- = 560$ MeV/c, $\omega = 200$ MeV and $\theta_e = 25^\circ$. All the theoretical curves have been obtained with the Argonne potential. Plane-wave (including Triton exchange) : dotted curve. Nucleon-nucleon rescatterings included : dashed curve. Two-body Meson-Exchange Currents included : dash-dotted curve. Three body mechanisms included : full curve.

Fig.11 : The cross-section of the ${}^4\text{He}(e,e'p)\text{T}$ reaction, in the two Amsterdam [7] kinematics (full circles and boxes) and the Saclay [28] kinematics (full triangle), are plotted against the momentum P_R of the recoiling Triton. The theoretical curves have been obtained with the Urbana potential. Their meaning is the same as in Fig. 10.

Fig.12 : The Transverse and the Longitudinal (virtual photon absorption) cross-sections in the same kinematics as in Fig. 10. The meaning of the curves is the same.

Fig.13 : The Transverse, Longitudinal, Transverse-Transverse and Transverse-Longitudinal interference components of the cross-section (for the absorption of a virtual photon) in a kinematics typical of the present generation of high duty factor electron accelerators. The meaning of the curve is the same as in Fig. 10.

Fig.14 : The Transverse and Longitudinal components of the (virtual photon absorption) cross-section in a collinear kinematics typical of CEBAF. Plane-Wave : dashed curves. Full calculations : full curves. The momentum of the recoiling Triton is P_R .

Fig.15 : Proton polarization components in the same collinear kinematics as in Fig 14. The meaning of the curves is the same. Note that the X axis lies on the same side as the scattered electron with respect to the virtual photon direction ($\Phi = 0$).

Fig.16 : Proton polarization components in the same coplanar kinematics as in Fig. 13. The meaning of the curves is the same. Note that the X axis lies on the side opposite to the scattered electron with respect to the virtual photon direction ($\Phi = 180^\circ$).

Fig.17 : Out of the plane proton polarization components in the same kinematics as in Fig. 13. The meaning of the curves is the same. Note that the Y axis is normal to the proton emission plane.

Fig.18 : The effect of charge exchange neutron-proton rescattering, $np \rightarrow pn$, on the components of the proton polarization in the ${}^4\text{He}(\vec{e}, e'\vec{p})\text{T}$ reaction, in a kinematics already achieved at Amsterdam with unpolarized electrons [7]. The dashed curves correspond to the Plane Wave approximation (including Triton exchange).

The long-dashed curves include only elastic $pn \rightarrow pn$ rescattering, beyond Plane Wave approximation. The dash-dotted curves correspond to the complete FSI calculation. The full curves include also Meson Exchange Currents. The momentum P_R of the recoiling Triton is plotted on abscissa.

Fig.19 : The effect of the charge exchange neutron-proton rescattering, $pn \rightarrow np$, on the components of the neutron polarization in the ${}^4\text{He}(\vec{e}, e'\vec{n}){}^3\text{He}$ reaction, in a kinematics already achieved at Amsterdam with unpolarized electrons [24]. The dashed curves are the result of Plane Wave Impulse Approximation (without ${}^3\text{He}$ exchange). The dotted curves correspond to the Plane Wave approximation with both the nucleon and the ${}^3\text{He}$ exchange mechanisms. The double-dot-dashed curves include only elastic $np \rightarrow np$ rescattering, beyond Plane Wave approximation. The dot-dashed curves correspond to the complete FSI calculation. The full curves include also Meson Exchange Currents. The momentum P_R of the recoiling ${}^3\text{He}$ is plotted in abscissa.

Fig.20 : The sensitivity to G_E^n of the sideways neutron polarization P'_X in the ${}^4\text{He}(\vec{e}, e'\vec{n}){}^3\text{He}$ reaction. Dot-dashed curve : Plane Wave approximation, when $G_E^n = 0$. Dashed curve : Plane Wave approximation, when $G_E^n \neq 0$. Complete calculation including nucleon rescattering and meson exchange current contributions. The momentum P_R of the recoiling ${}^3\text{He}$ is plotted in abscissa.

Fig.21 : The Coulomb component of the (virtual photon absorption) cross-section of $D(e, e'p)n$ reaction is computed with the time component J_0 or the third component J_Z of the nuclear current. Plane Wave approximation : dotted and dash-dotted curves. Full calculation : dashed and full curves. This is one of the kinematics already achieved at Saclay [35].

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