Induced Nucleon Polarization and Meson-Exchange Currents in \((e, e'p)\) Reactions

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Nucleon recoil polarization observables in \((e, e'p)\) reactions are investigated using a semi-relativistic distorted-wave model which includes one- and two-body currents with relativistic corrections. Results for the induced polarization asymmetry are shown for closed-shell nuclei and a comparison with available experimental data for \(^{12}\)C is provided. A careful analysis of meson exchange currents shows that they may affect significantly the induced polarization for high missing momentum.

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I. INTRODUCTION

Measurements of the polarization of the ejected proton in \((e, e'p)\) reactions provide valuable information on the nucleus complementary to that extracted from unpolarized experiments. In fact a new set of 8 spin-dependent response functions that present different sensitivities to the various ingredients of the reaction mechanism enter in the general analysis. A richer source of information on nucleon properties inside the nucleus is thus embedded into the spin-dependent nuclear responses.

In a previous work we have developed a model aiming to provide a systematic investigation of spin-dependent observables in \((e, e'p)\) reactions. Relying on the distorted wave impulse approximation (DWIA), our approach includes in addition two-body meson exchange currents (MEC) and relativistic corrections based on the semi-relativistic form of the electromagnetic currents derived in the last years. A richer source of information on nucleon properties inside the nucleus is thus embedded into the spin-dependent nuclear responses.

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The focus of this paper is the analysis of the properties displayed by the polarization observables induced by unpolarized electrons, i.e., the induced polarization asymmetry \(P\) which, contrary to the transferred asymmetries, is zero in PWIA. In fact, since the target nucleus is unpolarized the electron can hit with equal probabilities all spin orientations along their orbits. In absence of FSI, these nucleons leave the nucleus as plane waves with the same amplitude, hence giving no net induced polarization. The situation clearly differs when FSI are considered in the description of the process; first, because of the relation between the spin direction and the nucleon location in the orbit which implies different FSI strengths for different spin orientations due to the central part of the optical potential (mainly the imaginary absorptive term), and second, because of the explicit spin dependence of the spin-orbit interaction in the optical potential.

One of the goals of this work is to evaluate the impact of the two-body MEC over the induced polarization components and hence to analyze the validity of the impulse approximation (IA), i.e., one-body currents only. We are guided by previous studies where the role of MEC on asymmetry observables has been found to be in general small for low missing momentum \(p\). This result is in part due to the occurrence of an effective cancellation of MEC effects between the numerators and denominators involved in the polarization ratios. The same applies to FSI effects.

The full set of polarized response functions was computed and analyzed for intermediate to high values of the momentum transfer, \(q\), at the quasielastic peak. Their dependence on the model of final state interactions (FSI) was studied and the effects of MEC were evaluated. The emphasis was placed on the proton polarization induced by polarized electrons, i.e., on the transferred polarization asymmetries \(P_{l,s}\), which only contribute when the initial electron beam polarization is measured. These transferred asymmetries survive in the plane wave impulse approximation (PWIA) limit and may provide ideal tools for studying the electromagnetic nucleon form factors in the nuclear medium.

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The structure of the paper is as follows: in Sec. II we shortly present our distorted wave model. In Sec. III we discuss the results obtained for the induced polarization asymmetry and compare with available data. Finally, the conclusions are drawn in Sec. IV.

II. DWIA MODEL FOR \((e, e'p)\)

We refer to our previous works and references therein for details on the model. Here we just set up some general definitions of interest for the reader and for the discussion that follows. In the \((e, e'p)\) process sketched in Fig. 1, we consider an electron with four-momentum \(K_e = (\varepsilon_e, k_e)\) that scatters off a nucleus transferring a four-momentum \(Q^p = (\omega, q)\). The electron scattering angle is \(\theta_e\). A proton with momentum \(p'\) and exit solid angle \(\Omega\) is detected in coincidence with the outgoing electron. The proton spin polarization along
the ejected nucleon momentum, $p_l$, in the barycentric system ($\vec{l}, \vec{t}, \vec{n}$), where the normal vector $\vec{n}$ is defined by $\vec{q} \times \vec{p'}$. An arbitrary, unitary vector $\vec{s}$ is also measured. We assume that the residual nucleus is left in a discrete state, and neglect the recoil. The cross section for this process can be written in the Born approximation and extreme relativistic limit for the electron as

$$\Sigma = \frac{d\sigma}{d\epsilon'_e d\Omega_e d\Omega_v} = K\sigma_M \left( v_L R_L + v_T R_T + v_{TL} R_{TL} + v_{TT} R_{TT} \right)$$

where $\sigma_M$ is the Mott cross section, $K$ is the kinematic factor $m_N p'/(2\pi\hbar)^3$ (being $m_N$ the nucleon mass) and the $v_\alpha$ coefficients, $\alpha = L, T, TL, TT$ are the usual ones arising from the leptonic tensor $\mathcal{S}$. Finally, the exclusive response functions $R^\alpha$ are linear combinations of the Hadronic tensor and hence they contain all the pertinent information on the nuclear reaction mechanism.

In the present paper we make use of the semi-relativistic distorted wave model developed in Refs. [6, 12], whose basic ingredients are the following: i) the final proton state is described by a solution of the Schrödinger equation with a non-relativistic optical potential, but assuming the relativistic energy-momentum relation, thus we effectively solve a Klein-Gordon kind equation. ii) Semi-relativistic (SR) operators are used for the one-body (OB) electromagnetic current and two-body MEC. These have been obtained by expanding the corresponding relativistic operators to first order in the missing momentum over the nucleon mass $p/m_N$ (being $\vec{p} = \vec{p'} - \vec{q}$), while the exact dependence on $(\omega, \vec{q})$ is maintained [12, 13]. We consider the one-pion exchange diagrams of seagull (S or contact), pion-in-flight (P or pionic) and $\Delta$-isobar kinds.

The induced polarization asymmetry $\mathbf{P}$, which is the focus of this paper, is defined by

$$\Sigma = \frac{1}{2} \Sigma_{unpol} (1 + \mathbf{P} \cdot \mathbf{s}) .$$

The vector $\mathbf{P} = (P_l, P_t, P_n)$ is usually set in the barycentric coordinate system, referred to the reaction plane, as longitudinal ($\vec{l}$), transverse or sideways ($\vec{t}$), and normal ($\vec{n}$) directions defined in Fig. 1.

These induced polarization components can be written in the form

$$P_n = \frac{2}{W_0} \left( v_L W_n^L + v_T W_n^T + v_{TL} \cos \phi W_n^{TL} + v_{TT} \cos 2\phi W_n^{TT} \right) ,$$

$$P_l = \frac{2}{W_0} \left( v_T \sin \phi W_t^{TL} + v_{TT} \sin 2\phi W_t^{TT} \right) ,$$

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where $\phi = \phi'$ is the azimuthal angle of $\mathbf{p}$, and we have defined the function

$$W_0 \equiv v_L W_0^L + v_T W_0^T + v_{TL} \cos \phi W_0^{TL} + v_{TT} \cos 2\phi W_0^{TT} ,$$

FIG. 1: Coordinate system used in the $(e,e'p)$ reaction. The $x$-$z$ coordinates span the scattering plane, with the $z$-axis pointing along the momentum transfer $\vec{q}$. The proton polarization is described in the barycentric system ($\vec{l}, \vec{t}, \vec{n}$), where the normal vector $\vec{n}$ is defined by $\vec{q} \times \vec{p'}$. An arbitrary, unitary vector $\vec{s}$ is also measured. We assume that the residual nucleus is left in a discrete state, and neglect the recoil. The cross section for this process can be written in the Born approximation and extreme relativistic limit for the electron as

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FIG. 2: Proton induced polarization in the $n$-direction for knock-out from the $p$ shells of $^{16}$O. The kinematics are $q = 460$ MeV/c, $\omega = 100$ MeV, $\theta_e = 30^\circ$. The dependence on FSI is analyzed by comparison of two-optical potentials.
and the unpolarized \( W_n^0 \) and polarized \( W_n^c \) reduced response functions have been introduced. The role of the various ingredients of our model (FSI and MEC) over the separate response functions was analyzed in [9, 18]. In the following we show results for the induced polarization components for selected kinematical conditions.

III. RESULTS FOR THE INDUCED POLARIZATION

Since the induced polarization is zero in absence of FSI, this observable is expected \textit{a priori} to be specially sensitive to details of the optical potential used to describe the final proton state. Results of FSI model dependences are presented in Figs. 2 and 3 for proton knock-out from the \( p_{1/2} \) and \( p_{3/2} \) shells in \(^{16}\)O. Quasi-perpendicular kinematics is considered [21, 22], with \( q = 460 \) MeV/c and \( \omega = 100 \) MeV, corresponding closely to the quasi-elastic peak. The electron scattering angle is \( \theta_e = 30^\circ \). In Fig. 2 we show results for the normal component \( P_n \) for three values of the proton azimuthal angle \( \phi = 0, 90^\circ \) and \( 180^\circ \) (see definition in Fig. 1). Two optical potentials widely used in the literature to describe these reactions are considered: the Schwandt [23] potential with dashed lines and the Comfort & Karp [24] one with solid lines. Note that the Schwandt potential has been extrapolated here to \( A = 16 \), since it was fitted for heavier nuclei. These two potentials differ in their spin-orbit dependence. Whereas the Comfort-Karp potential includes a purely real term, Schwandt’s has also an imaginary part; moreover the real part of the Comfort potential is more attractive near the nuclear surface. Concerning the dependence on the real and imaginary parts of the central potential, Schwandt’s is more attractive and has less absorption. However, in their gross features the two potentials do not present remarkable discrepancies. This explains why at low missing momentum they provide similar predictions for \( P_n \), starting to differ for higher \( p \)-values \( p > 200 \) MeV/c.

The corresponding \( P_1 \) and \( P_3 \) polarization components are shown in Fig. 3. Note from Eqs. [21, 22] that these components are zero for in-plane emission (\( \phi = 0, 180^\circ \)), hence we only present results for out-of-plane kinematics, \( \phi = 90^\circ \). The biggest differences between both potentials show up in \( P_3 \) (upper panels) for low and high values of the missing momentum, while \( P_1 \) (lower panels) exhibits less dependence to details of the potential. The reason why \( P_1 \) is more sensitive to the details is not clear. However both polarizations are crucially dependent on the interaction, since they are strictly zero in PWIA. This is clearly illustrated in Fig. 3 where we show with dotted lines the results obtained with the Comfort-Karp potential but neglecting its spin-orbit dependence. The drastic change produced in the two polarizations shows that both observables depend equally on the global form of the interaction. However, at the kinematics selected, \( P_1 \) presents a slightly stronger sensitivity to the “fine” details of the potential.

In the following we analyze MEC effects restricting ourselves to the use of the Comfort-Karp potential. Discussion of results for the Schwandt potential follows.
are considered. As in the previous case of (top panels) may even change its global sign when MEC currents upon \( \Delta \) currents. Quantitatively the importance of the three contribution is opposite to those provided by the pionic and \( \Delta \)-currents, results in Figs. 4 and 5 for high missing momen-
tum region, the results in Figs. 6 and 7 constitute an indication of what kind of effects can be expected from MEC in this region.

As in the previous kinematics, MEC contributions are negligible for low missing momentum, \( p \lesssim 200 \text{ MeV/c} \). This strong MEC suppression is in part due to the behavior of the pion-nucleon form factor at high \( |Q^2| \). For higher \( p \)-values, \( p > 300 \text{ MeV/c} \), MEC start to be important giving a significant contribution for \( P_n \) at \( \phi = 180^\circ \) and \( P_l \) at \( \phi = 90^\circ \). Note the difference with the previous kinematics where the largest MEC effects were exhibited by \( P_n \) at \( \phi = 0 \) and \( P_l (\phi = 90^\circ) \). Particularly noteworthy is also the clear dominance of the \( \Delta \) current over the \( P \) and \( S \) terms. These MEC effects are similar to

In Figs. 4–7 we present the impact of MEC upon the induced polarization components for two values of the momentum transfer. Fig. 4 displays the \( P_n \) polarization for intermediate \( q = 460 \text{ MeV/c} \) and \( \omega = 100 \text{ MeV} \). In addition to the OB calculation (dotted lines), in each panel of the figure we show three more curves corresponding to the additional contribution of the several MEC: seagull (OB+S), seagull plus pionic (OB+S+P) and total MEC (OB+S+P+\( \Delta \)). As shown, for low missing momentum MEC contributions are in general negligible, and tend to increase as \( p \) goes higher. In particular, for \( \phi = 0 \) (top panels) MEC are shown to modify significantly the results of \( P_n \) for \( p > 200 \text{ MeV/c} \). This effect being larger for the \( p_{1/2} \) shell. For the other \( \phi \)-values selected, MEC are smaller for all missing momenta. This outcome, which is specific of the kinematics selected, can be ascribed to a cancellation of the two-body effects in \( P_n \). Note that the contribution of the interference \( TL \) response for \( \phi = 180^\circ \) is just opposite to that one occurring at \( \phi = 0 \), while for \( \phi = 90^\circ \) the \( TL \) response does not enter, see Eq. 3.

Larger MEC effects for high missing momentum, \( p > 200 \text{ MeV/c} \), are shown in Fig. 5 in the case of the longitudinal and transverse induced polarizations. Note that \( P_l \) (top panels) may even change its global sign when MEC are considered. As in the previous case of \( P_n \), the role of MEC for \( p < 200 \text{ MeV/c} \) is again negligible. From these results, we may conclude that the description of the induced polarization observables within the IA, i.e., with OB current operators only, is expected to be quite acceptable in this low-\( p \) regime.

Concerning the separate role played by the \( S \), \( P \) and \( \Delta \)-currents, results in Figs. 4 and 5 for high missing momentum, \( p > 200 \text{ MeV/c} \), show that the seagull contribution is opposite to those provided by the pionic and \( \Delta \) currents. Quantitatively the importance of the three currents upon \( P_n \) for \( \phi = 0 \) is similar. Note also that the

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the ones found over the $A_{TL}$ asymmetry for the same kinematics \cite{12}. In Fig. 8 we show results for the normal polarization of proton knock-out from the two shells in $^{12}$C, including only OB electromagnetic operators. We have chosen three sets of kinematics following closely those of Ref. \cite{1}, $q \simeq 760$ MeV/c and $\omega \simeq 290$ MeV. These values correspond nearly to the quasielastic peak. Since several sets of values have been used in the literature when comparing with the corresponding experimental data, in Fig. 8 we show three curves for three slightly diverse ($q, \omega$) sets. The results for the $p_{3/2}$ shell (upper panel) illustrates that extreme caution is needed before final conclusions can be drawn. In fact, the experimental point at $p \simeq 40$ MeV/c is extremely sensitive to the missing momentum region allowed by the kinematics. For exact quasi-elastic conditions, this region begins at $p = 0$, corresponding to forward emission of a nucleon with momentum $p' = q$. A slight change of $(q, \omega)$ can shift the allowed region by more than 25 MeV/c. The large error bars in the first data point are then reminiscent of the large instability of $P_n$ under tiny kinematical variations. The remaining data points located in the region of larger missing momentum, and the case of the $s_{1/2}$ shell, where a great stability exists, are of more physical interest for the analysis of two-body currents.

Regarding MEC effects upon $P_n$ in $^{12}$C, we show in Fig. 9 the comparison between our calculations including the OB and the several pieces of the two-body current. In the region of low missing momentum, $p < 200$ MeV/c, where experimental data are located, MEC contributions are negligible for both shells. As in the case of $^{16}$O, here MEC lead to significant effects, particularly due to $\Delta$ which gives the main contribution, in the high missing momentum region $p > 300$ MeV/c. Results in Fig. 9 also show that the SR distorted wave model calculations agree nicely with data. Moreover, the discrepancy with results obtained within the RDWIA framework \cite{21,22}, which is better suited to describe these high $(q, \omega)$ data, is small in the low missing momentum region, and begin to be important for $p \sim 200$ MeV/c. For larger missing momenta, $p > 300$ MeV/c, dynamical relativity, not included in our calculations, plays an important role and differences between RDWIA and SR-DWIA calculations increase. Note however, that MEC make also a very significant effect in this high-$p$ region.

To finish we compare our results with previous calculations of MEC effects over the induced polarization asymmetries, namely with the models developed by the Gent \cite{15} and Pavia \cite{24} groups. In the Gent calculation \cite{15} MEC contributions upon $P_n$ for $^{12}$C were also found to be very small for low missing momenta, while they increase importantly, particularly due to $\Delta$, for high $p$. Discrepancies with our results emerge because of the different models used to describe FSI; the Gent group makes use of a real potential without absorption. In the case of the Pavia calculation \cite{24}, the induced polarization was evaluated for low missing momentum, $p < 200$ MeV/c, and FSI were computed by means of complex phenomenological optical potentials. Their $P_n$ results for $^{16}$O with the Giannini & Rico optical potential \cite{31}, and...
FIG. 9: MEC effects over the normal induced polarization for $^{12}$C. The kinematics correspond to Fig. 8 with $(q, \omega) = (760 \text{ MeV/c}, 290 \text{ MeV})$. The experimental data are from Ref. [1].

IV. CONCLUSIONS

In this work we have analyzed the induced polarization asymmetry of knocked-out protons in exclusive $(e,e'p)$ reactions to discrete residual nucleus states. We have used a semi-relativistic distorted wave model including relativistic corrections to the one- and two-body currents as well as relativistic kinematics. We have applied the model to proton knock-out from the outer shells of $^{16}$O and have compared with the experimental data available for $^{12}$C.

Regarding FSI, the $P_n$ polarization is little dependent on the details of the optical potential for low missing momentum, while its dependence increases for high $p$. The longitudinal $P_l$ polarization is more sensitive to details of the potential both for low and high values of $p$, while the sideways $P_t$ polarization does only show tiny FSI uncertainties for $p > 300 \text{ MeV/c}$.

Concerning MEC, they are negligible for low $p < 200 \text{ MeV/c}$, but they can importantly change the induced polarization components for high $p$. In this regime the largest MEC contributions are found for intermediate values of the momentum transfer. The effects upon $P_n$ are in accordance with the Gent calculation, even if differences emerge due to the different FSI models used. The comparison with the older Pavia calculation, which gives rise to some peculiar differences between the two $p$-shells in $^{16}$O, is more troublesome.

The present calculation justifies the validity of the impulse approximation for $p < 200 \text{ MeV/c}$, while it emphasizes the fact that, besides dynamical relativity, other effects beyond the IA as MEC are also expected to contribute sizeably for high missing momentum. New experimental data for these observables in this regime would be welcomed to explore this physics.

Acknowledgments

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[1] R.J. Woo et al., Phys. Rev. Lett. 80, 456 (1998).
[2] S. Frullani and J. Mougey, Adv. Nucl. Phys. 14, 1 (1984).
[3] S. Boffi, C. Giusti, and F.D. Pacati, Phys. Rep. 226, 1 (1993).
[4] J. J. Kelly, Adv. Nucl. Phys. 23, 75 (1996).
[5] S. Boffi, C. Giusti, F. D. Pacati, and M. Radici, Electromagnetic response of atomic nuclei, Oxford University Press (1996).
[6] A. Picklesimer and J.W. Van Orden, Phys. Rev. C 35, 266 (1987).
[7] A. Picklesimer and J.W. Van Orden, Phys. Rev. C 40, 290 (1989).
[8] A.S. Raskin and T.W. Donnelly, Ann. Phys. (N.Y.) 191, 78 (1989).
[9] F. Kazemi Tabatabaei, J.E. Amaro, and J.A. Caballero, Phys. Rev. C 68, 034611 (2003).
[10] J.E. Amaro, J.A. Caballero, T.W. Donnelly, E. Moya de Guerra, A.M. Lallena, and J.M. Udas, Nucl. Phys. A 602, 263 (1996).
[11] J.E. Amaro, M.B. Barbaro, J.A. Caballero, T.W. Donnelly, and A. Molinari, Nucl. Phys. A 643, 349 (1998).
[12] J.E. Amaro, M.B. Barbaro, J.A. Caballero, T.W. Don-
nelly, and A. Molinari, Phys. Rep. 368, 317 (2002).
[13] J.E. Amaro, M.B. Barbaro, J.A. Caballero, T.W. Donnelly, and A. Molinari, Nucl. Phys. A723, 181 (2003).
[14] S. Malov et al., Phys. Rev. C 62, 057302 (2000).
[15] J. Ryckebusch, D. Debruyne, W. Van Nespén, and S. Janssen, Phys. Rev. C 60, 034604 (1999).
[16] J.J. Kelly, Phys. Rev. C 59, 3256 (1999); C 60, 044609 (1999).
[17] M.C. Martínez, J.R. Vignote, J.A. Caballero, T.W. Donnelly, E. Moya de Guerra, and J.M. Udías, Phys. Rev. C 69, 034604 (2004).
[18] J.E. Amaro, M.B. Barbaro, J.A. Caballero, and F. Kazemi Tabatabaei, Phys. Rev. C 68, 014604 (2003).
[19] M. Mazziotta, J.E. Amaro, and F. Arias de Saavedra, Phys. Rev. C 65, 034602 (2002).
[20] J.M. Udías and J.R. Vignote, Phys. Rev. C 62, 034302 (2000).
[21] L. Chinitz et al., Phys. Rev. Lett. 67 (1991) 568.
[22] G.M. Spaltro et al., Phys. Rev. C48, 2385 (1993).
[23] P. Schwandt et al., Phys. Rev. C 26, 55 (1982).
[24] J.R. Comfort and B.C. Karp, Phys. Rev. C 21, 2162 (1980).
[25] J. Gao et al., Phys. Rev. Lett. 84, 3265 (2000).
[26] J. M. Udíñas, J. A. Caballero, E. Moya de Guerra, J. E. Amaro, and T. W. Donnelly, Phys. Rev. Lett. 83, 5451 (1999).
[27] J.M. Udías, J.A. Caballero, E. Moya de Guerra, J.R. Vignote, and A. Escuderos, Phys. Rev. C 64, 024614 (2001).
[28] J.I. Johanson and H.S. Sherif, Phys. Rev. C 59, 3481 (1999).
[29] S. Boffi, C. Giusti, F.D. Pacati, and M. Radici, Nucl. Phys. A 518, 639 (1990).
[30] M.M. Giannini and G. Ricco, Ann. Phys. (N.Y.) 102, (1976).