Relativistic Effects and the Role of Heavy Meson Exchange in Deuteron Photodisintegration\textsuperscript{∗†}

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Relativistic effects and the role of heavy meson exchange in deuteron photodisintegration are studied systematically for photon energies below the pion production threshold. In a $\left(p/M\right)$-expansion, all leading order relativistic one-body and $\pi$-exchange as well as all static heavy meson exchange currents consistent with the Bonn OBEPQ model are included. In addition, one- and two-body boost effects have been investigated. Sizeable effects from the various two-body contributions beyond $\pi$-exchange have been found in almost every observable considered, i.e., differential cross section and single polarization observables.

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I. INTRODUCTION

Electromagnetic deuteron disintegration is one of the basic processes in order to study various aspects of the strong interaction of nucleons in nuclei. For example, a large number of experimental and theoretical papers have clarified the role of pion degrees of freedom as manifest in meson exchange currents (MEC) and the importance of relativistic contributions \cite{1}. However, the role of heavy meson exchange (with exception of the $\rho$ meson), which give an important contribution to the $NN$ interaction remained largely unclear. One of the reasons lies in the Siegert theorem which provides via the Siegert operators in conjunction with the Siegert hypothesis a convenient calculational tool to include implicitly the major part of MEC in the electric multipoles \cite{2}.

But admittedly such a procedure overshadows the underlying physics, and only “patches” an inconsistent calculation. A calculation that uses a set of consistent electromagnetic (e.m.) operators with respect to the hadronic interaction model should always be preferred. Such a consistent treatment has been given some time ago allowing only pions to interact with nucleons \cite{3}. In particular, the importance of consistency of the leading order relativistic contributions has been pointed out in this work. However, the extension to a realistic interaction model was missing. Only recently, Levchuk has presented a nonrelativistic calculation of deuteron photodisintegration for the Bonn OBEPR model where all heavy meson exchange currents were included explicitly so that Siegert operators were not needed \cite{4}. Unfortunately the results are presented in such a way, that the specific contributions from heavy meson exchange are not evident. Furthermore relativistic contributions have been neglected, which however become increasingly important at higher energies. As a side remark, we would like to mention another earlier calculation based on the Bonn potential models with a consistent pion exchange current but contributions of heavier mesons were included via the Siegert operators only \cite{5,6}.

More recently, we have extended the work of \cite{3} in order to investigate this question for deuteron electrodisintegration \cite{7} taking as interaction model the Bonn OBEPQ versions \cite{8,9}. As general result we found that the $\rho$ meson gives the most important heavy meson contribution whereas the influence of $\eta, \omega, \sigma, \delta, \text{and } \gamma\pi\rho/\omega$ is much smaller, in some observables completely negligible, in particular, near the quasifree kinematics. However, it is a priori not clear that this conclusion will be valid also for photodisintegration because of the fixed energy-momentum transfer relation. Therefore, we want to provide with the present work the missing study of the influence of heavy meson exchange on deuteron photodisintegration with inclusion of competing relativistic effects in the one-body and pion exchange sector.

The calculational framework is the same as in \cite{8} and will be very briefly reviewed in Sect. \ref{II}. The results are presented and discussed in Section \ref{III} restricting ourselves to the differential cross section $d\sigma/d\Omega$, and all single

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polarization observables, i.e., the photon asymmetry $\Sigma_l$, the target asymmetries $T_{IM}, IM \in \{11, 20, 21, 22\}$, and the proton and neutron polarization components $P_y(p)$ and $P_y(n)$. Section IV gives a short summary and outlook.

II. THEORETICAL FRAMEWORK

The calculation of the photodisintegration process is based on the equations-of-motion method for the derivation of the hadronic interaction model and the corresponding electromagnetic current operators. It is described in detail in [3] where for the first time a consistent treatment including leading order relativistic contributions had been presented for a pure one-pion-exchange model. As mentioned above, we have extended this work to the realistic Bonn OBEPQ versions and first applied it to photodisintegration [3] where further details can be found. In particular, all explicit expressions for the electromagnetic operators, which can be derived in the equation-of-motion method [3] or in the unitarily equivalent $S$-matrix approach [10], are listed in the Appendix of [3], including in addition the $\gamma\pi\rho$ and $\gamma\pi\omega$-currents and the currents involving $\Delta$-isobars. These dissociation and isobar currents introduce additional uncertainties and model dependence and will not be considered in the present study.

When calculating electromagnetic properties of hadronic systems one must, of course, use a set of electromagnetic operators that is consistent with the underlying hadronic interaction model, as is demanded by the requirement of gauge invariance. In order to construct such a consistent set of operators one should use the same theoretical basis as for the hadronic interaction, i.e., for a one-boson-exchange potential (OBEP) one should calculate all corresponding meson exchange current (MEC) operators consistently. However, it is not sufficient to simply take the same meson coupling constants and cutoffs as used in the potential model. Especially for the pion, care must be taken, because there are several sources of unitary freedom in constructing the pionic operators, leading to unitary parameters in the corresponding current expressions [3,10]. These should be chosen consistently with the OBE potential and for this reason we will briefly discuss them now.

First of all, in view of the unitary equivalence of pseudoscalar (ps) and pseudovector (pv) pion nucleon couplings, one introduces a mixing parameter $\mu$, that allows arbitrary mixing of the two coupling types, where $\mu = 0$ corresponds to pure ps and $\mu = 1$ to pure pv coupling. Secondly, it is well known, that this equivalence breaks down in the presence of the electromagnetic interaction. Then one has the choice between a ps-pv equivalent theory or one which is not, i.e. one can retain or leave out the equivalence breaking term. Clearly, chiral symmetry prefers the pv coupling. Thus it is customary to multiply the equivalence breaking term with a new parameter $\nu$ that switches this term off ($\nu = 0$) or on ($\nu = 1$) and introduces as additional parameter

$$\gamma = \mu + \nu,$$

so that $\gamma = 1$ corresponds to a chiral invariant interaction theory whereas $\gamma = 0$ violates chiral symmetry.

Another freedom arises in the $p/M$ expansion for the nonrelativistic reduction, the so-called Barnhill freedom described by the Barnhill parameter $c$ [11], which can be incorporated in the definition of a parameter

$$\tilde{\mu} = \mu + c(1 - \mu).$$

Note that only for ps coupling one has a dependence on the Barnhill parameter.

Finally, another unitary parameter stems from retardation. Although the OBEPQ versions are static potentials, it is not possible to construct a gauge invariant set of electromagnetic operators that are purely static [10]. This is due to the non-local nature of the e.m. operators when one leaves the nonrelativistic limit, e.g., the charge density associated with the pion in flight fulfills the gauge condition with the other retarded operators [3]. We have generated the retarded potential through a Taylor expansion of the pion propagator keeping the leading term only

$$V_{ret}(\vec{k}) = V_0(\vec{k})\Delta(\vec{k}^2)k_0^2,$$

where $V_0(\vec{k})$ is the static, nonrelativistic potential, $k_0$ the energy transfer at the vertex, and

$$\Delta(\vec{k}^2) = \frac{1}{m^2 + \vec{k}^2}$$

the static meson propagator. Certainly, this approximation is valid only below the pion production threshold. Here one has again the freedom to express the energy transfer of the pion by the energy transfers of the individual nucleons parametrized by a retardation parameter $\nu_{ret}$.
\[ k_0^2 = -k_0^{(1)} k_0^{(2)} + \frac{1 - \nu_{ret}}{2} (k_0^{(1)} + k_0^{(2)})^2 \]
\[ = \frac{1}{4M^2} \left( \bar{k} \cdot \bar{Q}_1 \bar{k} \cdot \bar{Q}_2 + \frac{1 - \nu_{ret}}{2} \left( \bar{k} \cdot (\bar{Q}_1 - \bar{Q}_2) \right)^2 \right), \]

(5)

where \( k_0^{(i)} \) denotes the energy transfer of nucleon \( "^i\)n. This freedom can be exploited to eliminate the retarded potential in the center-of-mass (c.m.) frame by the choice \( \nu_{ret} = \frac{1}{2} \). The retarded e.m. operators must then be constructed consistent with this choice.

With respect to the Bonn OBEPQs, one must be aware that these were constructed from the three dimensional Blankenbecler-Sugar reduction of the Bethe-Salpeter equation, thus yielding the nonrelativistic form of the kinetic energy operator, while still being a fully relativistic potential. The e.m. operators, constructed within the time-ordered perturbation theory or within an equation-of-motion approach [3], contain naturally relativistic one-body operators.

The connection between potential operators, calculated within these two different approaches, can be made by the “Coester” transformation [10]. This leads to the conclusion that the operators in [10] are consistent with the Bonn potentials for the choice \( \tilde{\mu} = -1 \), as was found in [3].

For the exchange of heavy mesons, the operators given in [10] are consistent with the Bonn potentials due to the simplicity of the corresponding Hamiltonians in leading order. For this reason, no additional unitary parameters appear. They should show up in higher order terms, which however, can safely be neglected due to the large meson masses.

For the \( \rho \) MEC we had distinguished in [3] between a “Pauli” current, corresponding to the \( \rho \) contribution proportional to \((1 + \kappa_\rho)^2\) which can be generated from the \( \pi \) MEC by substituting terms of the form \( \vec{\sigma} \cdot \vec{a} \) by \( \vec{\sigma} \times \vec{a} \), and the “Dirac” current for the remaining operators. This distinction will also be used later in the discussion.

With respect to the \( \sigma \) meson, we would like to mention that the Bonn potentials OBEPQ (A,B,C) need different meson parameters for the isospin \( T = 0 \) and \( T = 1 \) channels. Introducing the isospin projection operators, the potential (and the e.m. operators) can be viewed as a superposition of an isoscalar and isovector scalar meson for the two sets of parameters, yielding effectively four scalar mesons

\[ V^\sigma = \frac{1}{4} (1 - \vec{r}_1 \cdot \vec{r}_2) V^{\sigma_0} + \frac{1}{4} (3 + \vec{r}_1 \cdot \vec{r}_2) V^{\sigma_1}. \]

(6)

So, strictly speaking, one has to take into account artificial isovector currents introduced through this procedure, which obviously tend to cancel each other. A simpler approximation, i.e., comparing the “pure” \( \sigma \) meson exchanges of the two parameter sets, also leads to the conclusion that this “anomaly” has no visible effect [3].

Finally, for a consistent treatment of leading order relativistic contributions one has to include the wave function boost which takes into account the fact that initial and final hadronic states move with different velocities, i.e., their rest frames to which the intrinsic motion refers are different. A convenient method for treating this frame dependence of the intrinsic motion is to introduce a unitary transformation generated by a boost operator \( \chi(\vec{P}) \)

\[ | \vec{P}, \vec{p} \rangle = | \vec{P}_{c.m.} \rangle \otimes e^{-i\chi(\vec{P})} | \vec{p}_{int} \rangle, \]

(7)

where \( | \vec{p}_{int} \rangle \) describes the intrinsic wave function in the rest frame [12]. Instead of transforming the wave functions themselves, the boost effect is incorporated into the operators by

\[ e^{i\chi \Omega} e^{-i\chi} \approx \Omega + i [\chi, \Omega], \]

(8)

where in the commutator only the nonrelativistic part of \( \Omega \) has to be considered.

The operator \( \chi \) can be separated into a kinematic, interaction independent part \( \chi_0 \) and an interaction dependent part \( \chi_V \)

\[ \chi = \chi_0 + \chi_V. \]

(9)

For the two nucleon system one has [12]

\[ \chi_0 = - \left( \frac{\vec{r} \cdot \vec{P} (\vec{r} \cdot \vec{P})}{16M^2} \right) + \text{h.c.} + \frac{((\vec{\sigma}_1 - \vec{\sigma}_2) \times \vec{p}) \cdot \vec{P}}{8M^2}, \]

(10)

whereas a nonvanishing, interaction dependent boost operator exists only for pseudoscalar meson exchange \((\pi, \eta)\) in the case of pseudoscalar coupling [13,14]

\[ \chi_\pi V = - (\vec{r}_1 \cdot \vec{r}_2) \frac{i}{8M} \left( \frac{\pi_{NN}}{2M} \right)^2 (1 - \hat{\mu}) \int \frac{d^3k}{(2\pi)^3} e^{ik \cdot r} \Delta(\vec{k}) \vec{\sigma}_1 \cdot \vec{P} \vec{\sigma}_2 \cdot \vec{k} + (1 \leftrightarrow 2). \]

(11)

Therefore, for the Bonn potentials, the potential dependent boost appears with \((1 - \hat{\mu}) = 2\).
III. RESULTS AND DISCUSSION

We have calculated the unpolarized differential cross section \( \frac{d\sigma}{d\Omega} \) and all single polarization observables, i.e., photon asymmetry \( \Sigma^i \), the target asymmetries \( T_{11}, T_{20}, T_{21}, T_{22} \), and the final nucleon polarization \( P_y \) for proton and neutron. Their formal expressions in terms of the basic T-matrix elements

\[
T_{sm,\lambda m_d} = \pi \sqrt{\frac{k}{q}} \alpha E_d (s m_s | j_\lambda (q) | m_d)
\]

are given in [12], where \( \alpha \) denotes the fine structure constant, \( k \) the asymptotic relative momentum of the outgoing nucleons in the c.m. frame, \( q \) the photon momentum, \( E_d \) the energy of the initial deuteron, and \( j_\lambda \) the nucleon current in a spherical basis.

For the explicit calculation we have chosen four representative photon energies \( E_\gamma = 4.5, 40, 100, \) and \( 140 \) MeV covering the region between the maximum of the total cross section and pion production threshold. In order to distinguish the different influences from pion, rho, and other heavy meson exchanges, we show their effects in separate panels for each observable and each energy. In addition, we show an overview and the potential model dependence with respect to the versions A, B, and C of the Bonn OBEPQ. Thus each observable is represented by a figure consisting of four columns, one for each energy, and five rows for the overview, \( \pi \) exchange, \( \rho \) exchange, additional heavy meson exchange, and potential model dependence. The notation of the curves is the same for all observables.

In the first row we present an overview of the following effects: the nonrelativistic one-body current (long-dashed curve), the relativistic one-body current (dash-dotted curve), to this added the nonrelativistic \( \pi \) MEC (dotted curve), and finally the total result including all heavy meson exchanges (full curve).

The next row shows the contributions from \( \pi \) MEC, starting from the relativistic one-body current (dash-dotted curve of the first row) to which first the nonrelativistic \( \pi \) MEC is added (dotted curve) and then the relativistic \( \pi \) MEC -- including retardation corrections -- (dashed curve).

The third row displays the effects of the \( \rho \) MEC: here we start from the relativistic \( \pi \) MEC (dashed curve of the previous row) and include first the Pauli \( \rho \) MEC (short-dashed curve) and then the Dirac \( \rho \) MEC (long-dashed-dotted curve).

The effects of the various heavy meson exchanges are presented in the fourth row. To the relativistic one-body current plus relativistic \( \pi \) MEC and full \( \rho \) MEC (long-dash-dotted curve, as in the third row), we add consecutively \( \delta \) MEC (dotted curve), \( \omega \) MEC (short-dashed curve), \( \sigma \) MEC (dashed curve), and finally the \( \eta \) MEC (full curve).

The last row shows the potential dependence of the observable with respect to the different versions of the Bonn OBEPQ potential, where the full curve represents the version B, the short-dashed version A, and the dotted version C.

Now we will discuss in detail the various observables starting with the differential cross section in Fig. 1. The overview shows that in the maximum of the total cross section, at 4.5 MeV this observable is dominated by the nonrelativistic one-body current while only the nonrelativistic \( \pi \) MEC gives a 10 percent enhancement. Obviously relativistic effects and heavy mesons are negligible as well as the potential model dependence.

At the next higher energy (40 MeV) the nonrelativistic \( \pi \) MEC becomes comparable to the one-body current. All other contributions give a small overall reduction, somewhat more pronounced in forward and backward direction. However, if one looks at the separate contributions, one notices a subtle destructive interference of different larger effects. First, relativistic \( \pi \) MEC gives a slight reduction in the maximum but leaves the forward and backward directions almost unchanged. Next, from \( \rho \) one sees a strong forward and backward reduction from the Pauli current, whereas the Dirac contribution mainly leads to a sizeable enhancement in the maximum which, however, is largely cancelled by the additional heavy mesons. Finally, one finds a small model dependence of a few percent.

Considering now the two higher energies (100 and 140 MeV) one readily notices a dramatic increase of relativistic effects. First a sizeable reduction appears from the relativistic one-body current showing the well known effect of diminishing the differential cross section at forward and backward angles which comes mainly from the dominant spin-orbit (SO) current [16]. The further reduction from the remaining contributions shows again in detail a strong destructive interference. In fact, first the relativistic \( \pi \) MEC surprisingly enhances the cross section in forward direction while then the \( \rho \) Pauli current results in a drastic reduction at both extreme angles. The Dirac contribution gives again an overall enhancement but of smaller size. This effect of the Dirac \( \rho \) current is somehow surprising, because it is roughly of the same size as that of the Pauli current whereas from the size of the coupling constants one would have expected a suppression by a factor of about 50. The additional heavy mesons beyond the \( \rho \) show a much smaller influence. Their individual contributions are surprisingly big (up to \( \sim 5\% \)), in particular compared to their role in the parametrization of the \( NVV \) force and their importance in the electrodisintegration of the deuteron. Most prominent is the effect of the \( \delta \) meson leading to a reduction of the differential cross section. However, looking at the
overall result, the heavy meson exchanges tend to cancel each other. With respect to the potential model dependence, one sees now a larger variation, in particular also at forward and backward angles which increases with energy.

The photon asymmetry $\Sigma^l$ in Fig. 2 is very sensitive to two-body effects, as is long known. At 4.5 MeV only the nonrelativistic one-body current contributes to $\Sigma^l$ and no potential model dependence appears. Then at 40 MeV the nonrelativistic $\pi$ MEC becomes sizeable as well as the Pauli $\rho$ current. All other effects, relativistic one-body and $\pi$ MEC, Dirac $\rho$ and additional heavy meson effects are very small, as is the potential model dependence. At higher energies the relativistic one-body current and the relativistic $\pi$ MEC becomes important, too. The first leads to a sizeable reduction of the photon asymmetry, the latter to a smaller increase. The $\rho$ MEC increases the photon asymmetry, of which the Pauli current is the most dominant part while the Dirac current is comparably small, although its size increases with the photon energy. The influence of the various heavy meson exchanges are much more pronounced than in the differential cross section, mainly coming from the $\delta$ MEC. But again the various heavy mesons tend to interfere destructively. Also the potential dependence is quite large for 100 MeV and 140 MeV, where the OBEPQ version A yields the biggest asymmetry, version B intermediate values, and version C the lowest photon asymmetry. Thus one might be tempted to single out one potential against others by comparison with experimental data. However, one has to be careful in such a comparison, because before drawing definite conclusions as to which model should be preferred, one has to study in detail the remaining theoretical uncertainties due to the strength of the dissociation and isobar currents. Here, the additional independent measurement of the unpolarized cross section and the vector target asymmetry $T_{11}$ could help in fixing the respective strengths of these contributions.

The next observable, the tensor target asymmetry $T_{20}$ in Fig. 4 shows sizeable effects from the various currents mainly for the regions around 0° and 180°, except at 4.5 MeV where only a very small influence from the nonrelativistic $\pi$ MEC is seen. For the higher energies, the largest effect stems from the nonrelativistic $\pi$ MEC. In addition, at 100 and 140 MeV also the relativistic one-body current becomes equally important, particularly strong in the backward direction of increasing size. This observable reacts only slightly when the relativistic $\pi$ MEC is added, and the same is true for the $\rho$ current. The additional heavy mesons appear more interesting. The $\omega$ and $\sigma$ meson exchanges lead to a visible reduction at forward and backward angles. Finally, $T_{20}$ is stable against a change of the potential version. This is valid for all tensor target asymmetries.

The next observable, the tensor target asymmetry $T_{21}$ in Fig. 5 shows the largest sensitivity to the nonrelativistic $\pi$ meson exchange at the low energy of 4.5 MeV, although the observable is small in absolute size. Also for the higher energies the only relevant contribution besides the one-body part comes from the nonrelativistic $\pi$ MEC, but the relative importance diminishes with increasing energy. Otherwise, $T_{21}$ is a very stable observable with respect to all remaining contributions and to a potential variation.

The last observable, the tensor target asymmetry $T_{22}$ in Fig. 6 is like $T_{21}$ also sensitive to two-body effects of which the nonrelativistic $\pi$ MEC is dominating again. It leads for photon energies of 40, 100, and 140 MeV to a drastic increase of the asymmetry. But this observable shows almost no influence from relativistic contributions, neither from the relativistic one-body nor from the relativistic $\pi$ MEC. Also the $\rho$ current shows an increasing contribution with increasing energy, mainly from the Pauli part. The additional heavy meson exchange becomes significant, too, with the largest part from $\omega$ MEC, leading to a slight overall decrease. The variation with the potential model is negligible.

Finally, we show the outgoing proton and neutron polarizations $P_p(p)$ and $P_p(n)$ in Figs. 7 and 8, respectively. They show very similar behaviour with respect to the different contributions, except for the two higher energies where one notes some differences. At 4.5 MeV the polarization is dominated by the nonrelativistic one-body current. There is only a tiny contribution from the nonrelativistic $\pi$ MEC, and all other effects are completely negligible. The relativistic one-body current shows some increasing influence with increasing energy, but the most important contribution is here again the nonrelativistic $\pi$ MEC whereas the sensitivity to the relativistic $\pi$ MEC is very small. The $\rho$ contribution stems mainly from the Pauli $\rho$ MEC, leading to a reduction at backward angles. The additional heavy meson exchanges lead to a small but increasing effect for higher photon energies. The potential dependence becomes sizeable at 100 and 140 MeV, in particular for the proton polarization, and also increases with energy.

Finally, we would like to mention that boost contributions have turned out to be totally negligible. This is true for both the one-body boost operators, see also [17], as well as for the contributions of the two-body boost currents. It is in contrast to deuteron electrodisintegration where the one-body boosts cannot be neglected.
IV. SUMMARY AND OUTLOOK

Summarizing our results, we may state that relativistic contributions beyond the leading spin orbit current are important and should be included in a consistent manner within a realistic hadronic interaction model. In comparison with the work of Ref. [3], where a pure pionic model had been used, we see that the relativistic effects are roughly of the same size. For the differential cross section one notices that the relativistic pionic current in the present work has a different angular distribution than in [3], whereas the effects in the photon asymmetry are very similar. These differences stem from the different potential models and were expected.

Of the heavy mesons beyond the pion, the $\rho$ meson is the most important one confirming earlier investigations. We also find that in many cases the Pauli current is dominating its contribution justifying to some extent the neglect of the Dirac current. However, we have found some observables for which the Dirac current gives contributions comparable to the Pauli current. With respect to the other heavy mesons beyond the $\rho$ meson, we have found that each individual meson shows a sizeable influence on the various observables of deuteron photodisintegration. But taken together, we find in many of the observables considered strong destructive interference between them ($\delta$, $\omega$, $\sigma$, and $\eta$) so that their overall effect is quite small. Nonetheless, there are certain polarization observables, like the tensor target asymmetry $T_{20}$, that obviously demand the incorporation of the heavy meson contributions.

We have not attempted to make a comparison with experimental data because, first of all, our main interest was the study of the relative importance of relativistic effects for a realistic $NN$ interaction model compared to the current contributions from heavy meson exchange. Secondly, since at low and medium energies the results for cross sections and polarization observables agree essentially with those of other potential models, one can expect the same kind of agreement with experimental data for these energies, whereas at higher energies (100 and 140 MeV) isobar contributions, particularly from the $\Delta$ resonance, have to be included for a meaningful comparison. Furthermore, above pion production threshold, the approximative treatment of retardation in the present work breaks down and the correct incorporation of pion retardation in the $NN$- and $N\Delta$-sector becomes important [10]. These effects are definitively much more crucial than the effects of heavy meson exchange. Also relativistic contributions to the excitation of the $\Delta$ resonance itself might show some sizeable influence. These questions will be investigated in future work.

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FIG. 1. The differential cross section for various laboratory photon energies. Notation of the curves: (1) overview: nonrelativistic one-body current (long-dashed); relativistic one-body current (dash-dotted); nonrelativistic $\pi$ MEC added (dotted); total result (full); (2) $\pi$ meson: relativistic one-body current (dash-dotted); nonrelativistic $\pi$ MEC added (dotted); relativistic $\pi$ MEC including retardation (dashed); (3) $\rho$ meson: relativistic one-body plus complete $\pi$ MEC (dashed); Pauli MEC (short-dashed); Dirac MEC (long-dashed-dotted); (4) heavy meson: relativistic one-body current plus complete $\pi$ and full $\rho$ MEC (long-dash-dotted); $\delta$ MEC (dotted); $\omega$ MEC (short-dashed); $\sigma$ MEC (dashed); $\eta$ MEC (full); (5) potential: OBEPQ version B (full); version A (short-dashed); version C (dotted).
FIG. 2. The photon asymmetry $\Sigma^1$. Notation of the curves as in Fig. 1.
FIG. 3. The vector target asymmetry $T_{11}$. Notation of the curves as in Fig. 1.
FIG. 4. The tensor target asymmetry $T_{20}$. Notation of the curves as in Fig. [1]
FIG. 5. The tensor target asymmetry $T_{21}$. Notation of the curves as in Fig. 1.
FIG. 6. The tensor target asymmetry $T_{22}$. Notation of the curves as in Fig. [1]
FIG. 7. The polarization $P_y$ of the outgoing proton. Notation of the curves as in Fig. 1.
FIG. 8. The polarization $P_y$ of the outgoing neutron. Notation of the curves as in Fig. [1].