I. INTRODUCTION

The electromagnetic form factors of nuclei provide important information about their internal structure. They have been used extensively to test models of the nuclear dynamics and of the associated electromagnetic currents. As electron scattering experiments, such as the ones performed at Jefferson Lab, reach larger and larger values of the momentum transferred by a virtual photon to the struck nucleus, it becomes increasingly important to incorporate the requirements of special relativity in a reliable way into the theoretical description of the process.

The Covariant Spectator Theory (CST) [1] was designed as a manifestly covariant theory, especially suited for the description of few-nucleon problems. In a recent article [2], we presented the first CST calculations of the electromagnetic interaction models WJC-1 and WJC-2. The calculations use an approximation for the three-nucleon vertex functions with two nucleons off mass shell. The form factors with WJC-2 are close to the ones obtained with the older model W16 and to nonrelativistic potential calculations with lowest-order relativistic corrections, while the form factors with the most precise two-nucleon model WJC-1 exhibit larger differences. These results can be understood when the effect of the different types of pion-nucleon coupling used in the various models is examined.

The electromagnetic form factors of the three-nucleon bound states were calculated in complete impulse approximation in the framework of the Covariant Spectator Theory for the new high-precision two-nucleon interaction models WJC-1 and WJC-2. The calculations use an approximation for the three-nucleon vertex functions with two nucleons off mass shell. The form factors with WJC-2 are close to the ones obtained with the older model W16 and to nonrelativistic potential calculations with lowest-order relativistic corrections, while the form factors with the most precise two-nucleon model WJC-1 exhibit larger differences. These results can be understood when the effect of the different types of pion-nucleon coupling used in the various models is examined. DOI: 10.1103/PhysRevC.81.014007 PACS number(s): 21.45.—v, 21.30.—x, 25.30.Bf, 13.40.Gp

First results for electromagnetic three-nucleon form factors from high-precision two-nucleon interactions

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The electromagnetic form factors of the three-nucleon bound states were calculated in complete impulse approximation in the framework of the Covariant Spectator Theory for the new high-precision two-nucleon interaction models WJC-1 and WJC-2. The calculations use an approximation for the three-nucleon vertex functions with two nucleons off mass shell. The form factors with WJC-2 are close to the ones obtained with the older model W16 and to nonrelativistic potential calculations with lowest-order relativistic corrections, while the form factors with the most precise two-nucleon model WJC-1 exhibit larger differences. These results can be understood when the effect of the different types of pion-nucleon coupling used in the various models is examined.

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The Covariant Spectator Theory (CST) [1] was designed as a manifestly covariant theory, especially suited for the description of few-nucleon problems. In a recent article [2], we presented the first CST calculations of the electromagnetic three-nucleon (3N) form factors in complete impulse approximation (CIA), which is defined as the complete CST 3N current [3] except for interaction currents (i.e., diagrams where the photon couples to an intermediate interacting two-nucleon (NN) system). However, the term “impulse approximation” can be misleading because it depends on the framework used.

For instance, in a very successful approach used by the Pisa-Jefferson Lab collaboration [4,5], the dynamics is based on the nonrelativistic Schrödinger equation, and relativistic corrections are added perturbatively. We call the corresponding impulse approximation with relativistic corrections “IARC”. In this framework, two- and three-body interaction currents are later added to the IARC results. These include first-order γπNN contact interactions that are equivalent to those already at the CIA level automatically—and to all orders—included “Z-graphs” in the CST. This is an example of a more general observation: what counts as interaction current in one approach may be part of the impulse approximation in another.

In Ref. [2], our focus was to study the model dependence of the electromagnetic 3N form factors in CST. We performed calculations for a family of closely related relativistic two-nucleon interaction models and found that the CST results behave very reasonably. In most cases, a direct comparison of our CIA results with experimental data is not useful because we expect interaction currents to be significant. However, the comparison with IARC results is instructive, and it appears that the surprisingly close agreement between the two approaches (at least for Q ≤ 4 fm⁻¹), when models are compared that yield the same 3N binding energy, is no coincidence. The main reason seems to be that all CST models used in the comparison employ pseudovector coupling for the pion-nucleon vertices. This kind of coupling suppresses negative-energy states (corresponding to Z-graphs), which are included in CIA but not in IARC. It is therefore understandable that no large differences between the two calculations emerge, as long as other aspects of the dynamics in the two approaches are comparable.

It would be interesting to submit this interpretation to a test. One only needs to perform two calculations in CIA with two NN models that are as similar as possible in their ability to describe the NN data and the 3N binding energy. One of them should be based on pure pseudovector pion-nucleon coupling, whereas the other should include an admixture of pseudoscalar pion-nucleon coupling and thus increase the weight of Z-graphs. If the above interpretation is correct, the model with pure pseudovector coupling will be close to the IARC result, whereas there should be larger deviations in the case of the model with some pseudoscalar coupling.

We are indeed in a position to perform this test. In a recent article [6], we published two realistic CST models for the neutron-proton interaction, both of which describe the np scattering observables with χ²/Ndata ∼ 1 for the most recent 2007 database. The first model, WJC-1, based on the exchange of eight bosons and fitted with 27 adjustable parameters, features a mixture of pseudovector and pseudoscalar pion-nucleon coupling. The second model, WJC-2, based on the exchange of six mesons and with only 15 adjustable parameters, uses pure pseudovector pion-nucleon coupling.
The two models can be considered to be essentially on-shell equivalent, and both also reproduce the experimental value of the triton binding energy of 8.48 MeV.

There is, however, one obstacle to performing the CIA calculations with models WJC-1 and WJC-2: Some of the diagrams that comprise the CIA 3N current depend on the 3N vertex function with two nucleons off mass shell. A computer code for the calculation of these vertex functions for the new models WJC-1 and WJC-2 is at present in development, but not yet ready to be used in the calculation of the 3N form factors. This obstacle can be overcome if we apply an approximation, the 3N form factors in CIA. Similar (and equivalent) expressions for the 3N current were also derived in Refs. [7] and [8]. Figure 1 displays the complete current, and CIA is defined through diagrams (A)–(F).

We denote the photon four-momentum by \( q \), and we label the nucleon four-momenta \( k_i \) such that always \( k_1^0 = k_2^0 = m^2 \), where \( m \) is the nucleon mass. For the cases where a nucleon absorbs a photon, we introduce the notation \( k_i^\pm \equiv k_i \pm q \). The momentum \( k_3 \) is not an independent variable; in CST, the energy-momentum four-vector is conserved and \( k_3 \) is determined through the momenta of nucleons 1 and 2 and the total 3N momenta \( P_1 \) in the initial and \( P_1' \equiv P_1 + q \) in the final state.

The 3N current in CIA is given in algebraic form by

\[
J_{\text{CIA}}^\mu = 3e \int \int \frac{m^2 d^3 k_1 d^3 k_2}{E(k_1)E(k_2)(2\pi)^3} \sum_{\lambda_1 \lambda_2} \left[ \Psi_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1') \right.
\]

\[
\times \left( 1 + 2 \zeta \mathcal{P}_{12} \right) j_{\omega \alpha}(k_3^+, k_3) \Psi_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1')
\]

\[
+ \Gamma_{\lambda_1 \rho \beta}(k_1, k_2^+, P_1') G_{\rho \beta}(k_2^-, k_2) j_{\beta \alpha}(k_3^+, k_3) u_\gamma(k_2, \lambda_2)
\]

\[
\times \left( 1 + 2 \zeta \mathcal{P}_{12} \right) \Psi_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1') + \Psi_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1')
\]

\[
\times \left( 1 + 2 \zeta \mathcal{P}_{12} \right) \bar{u}_\gamma(k_2, \lambda_2) j_{\beta \alpha}(k_3^+, k_3) \Gamma_{\lambda_1 \rho \beta}(k_1, k_2^+, P_1'),
\]

(1)

where \( E(k) = \sqrt{m^2 + k^2} \), \( \mathcal{P}_{12} \) is a permutation operator that interchanges particles 1 and 2, \( \zeta \) is a phase with \( \zeta = +1(-1) \) for bosons (fermions), \( G_{\rho \beta}(k) \) is the propagator of an off-shell nucleon with four-momentum \( k \),

\[
G_{\rho \beta}(k) = \left( \frac{m + k}{m^2 - k^2 - i\epsilon} \right)_{\rho \beta},
\]

(2)

and \( j_{\omega \alpha}(k', k) \) is the single nucleon current for off-shell nucleons with incoming (outgoing) four-momentum \( k \) (\( k' \)). Summation over repeated Dirac indices is implied.

The “relativistic wave functions” \( \Psi \) are defined in terms of the 3N vertex functions \( \Gamma \) as

\[
\Psi_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1) = G_{\omega \alpha}(k_3) \Gamma_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1),
\]

(3)

and we use a shorthand for the contraction of Dirac indices with nucleon helicity spinors with helicity \( \lambda_i \),

\[
\Gamma_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1) \equiv \bar{u}_\alpha(k_1, \lambda_1) u_\alpha(k_2, \lambda_2) \Gamma_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1).
\]

(4)

This notation is used throughout this work, such that the replacement of a Dirac index \( (\alpha_i, \beta_i, \ldots) \) by a helicity index \( (\lambda_i) \) always indicates a corresponding contraction with a positive-energy helicity spinor.

The 3N vertex functions are solutions of Faddeev-type CST integral equations

\[
\Gamma_{\lambda_1 \lambda_2 \omega}(k_1, k_2; P_1) = -\int \frac{m d^3 k_2'}{E(k_2')(2\pi)^3}
\]

\[
\times \sum_{\lambda_2'} M_{\lambda_2' \lambda_1 \lambda_2 \omega}(k_2', k_2; P_{23}) 2 \zeta \mathcal{P}_{12} \Psi_{\lambda_1 \lambda_2 \omega}(k_1, k_2', P_1).
\]

(5)

Here, \( M_{\lambda_2' \lambda_1 \lambda_2 \omega}(k_2', k_2; P_{23}) \) is the scattering amplitude of nucleons 2 and 3 with total pair momentum \( P_{23} \). It satisfies
the CST two-body equation
\[ M_{\lambda_1\lambda_2\lambda_3\lambda_4}(k_2, k_2'; P_{23}) = V_{\lambda_1\lambda_2\lambda_3\lambda_4}(k_2, k_2'; P_{23}) \]
\[ - \int \frac{m \, d^3k_2'}{E(k_2') (2\pi)^3} \sum_{\lambda_1} \frac{1}{k_2^2} \times G_{\beta\beta'}(P_{23} - k_2) M_{\lambda_2\beta\lambda_3\lambda_4}(k_2', k_2'; P_{23}), \] (6)
where \( V \) is the NN interaction kernel. In the case of the NN interaction models considered in this work, \( V \) is of one-boson-exchange form [6,10].

In the second and third lines of Eq. (1), corresponding to the diagrams in Figs. 1(B), 1(C), 1(E), and 1(F), the vertex function appears with two nucleon momenta off mass shell. The solutions of the CST equation for the 3N bound state (5) have only one nucleon (nucleon 3, by convention) off mass shell, but one can obtain vertex functions with two off-shell particles through an iteration of the 3N equation with an off-shell two-nucleon scattering amplitude,
\[ \Gamma_{\lambda_1\rho\beta}(k_1, k_2; P_i) = - \int \frac{m \, d^3k_2'}{E(k_2') (2\pi)^3} \times \sum_{\lambda_1} M_{\beta\alpha,\lambda_1\alpha}(k_2', k_2'; P_{23}) \chi_{\lambda_1\alpha}(k_1, k_2'; P_1). \] (7)

In Eq. (7), the final-state momentum of particle 2, \( k_2' \), is off mass shell, whereas its initial-state momentum, \( k_2 \), is on mass shell (nucleon 3 is off shell in either state).

As pointed out in Sec. I, the 3N vertex functions with both nucleons off shell in the final state are not available at this time for the new \( NN \) interactions WJC-1 and WJC-2. Moreover, it is rather awkward to calculate and manipulate these double-off-shell vertex functions numerically because, with one additional continuous variable (the off-shell energy of nucleon 2), they occupy much more computer storage space and slow down the calculations.

For these practical reasons, we introduce here a simple approximation that replaces the vertex functions with two nucleons off mass shell by others with only one nucleon off mass shell.

To motivate this approximation consider, for instance, Figs. 1(C) and 1(F). The vertex function in the initial state, \( \Gamma_{\lambda_1\rho\beta}(k_1, k_2; P_i) \), depends on the off-shell momentum \( k_2' = k_2 - q \). Because \( k_2 \) is on mass shell, for small photon momenta \( q \) the momentum \( k_2' \) is also almost on mass shell. We may therefore expand the vertex function in the off-shell energy of nucleon 2 around its on-shell value. If we keep only the zeroth-order term of the expansion and eliminate the corresponding negative-energy channel of nucleon 2 (with negative \( \rho \)-spin), we obtain a known vertex function with two nucleons on mass shell. We call this approximation “CIA-0,” referring to the zeroth-order expansion involved.

Although this approximation is easy to apply, its formulation is somewhat awkward because of its frame dependence. In our numerical calculations, the 3N vertex function is expressed in terms of variables for nucleons 2 and 3, which are defined in the rest frame of the (23) pair where the CST equation for the two-nucleon scattering amplitudes is solved numerically.

We can write
\[ \Gamma_{\lambda_1\rho\beta}(k_1, k_2; P_i) = \Gamma_{\lambda_1\rho\beta}(k_1, L(k_{23})\vec{k}_2; P_i), \] (8)
where the Lorentz transformation \( L(k_{23}) \) takes the system of nucleons 2 and 3 from its rest frame, where their momenta are \( \vec{k}_2 \) and \( \vec{k}_3 \), to the 3N rest frame, where their momenta are \( \vec{k}_2 = L(k_{23})\vec{k}_2 \) and \( \vec{k}_3 = L(k_{23})\vec{k}_3 \), and where their total two-body momentum is \( k_{23} = k_2 + k_3 = P_i - k_1 \).

We now define the four-momentum \( \vec{r}_2 \) to have the same three-vector part as \( \vec{k}_2 \) but to be on mass shell, that is, \( \vec{r}_2 = [E(\vec{k}_2), \vec{k}_2] \), and we replace the momentum \( \vec{k}_2 \) by \( \vec{r}_2 \) in the vertex function (8).

To eliminate the negative-energy states of nucleon 2, we first write the propagator of nucleon 2 in terms of its form in the pair rest frame,
\[ G_{\rho\beta}(\vec{k}_2) = S_{\rho\beta}(L(k_{23})) \]
\[ \times \left[ \frac{m + \vec{k}_2}{m^2 - (\vec{k}_2)^2 - i\epsilon} \right] S_{\rho\beta}^{-1}(L(k_{23})), \] (9)
where \( S[L(k_{23})] \) is the Dirac space representation of the Lorentz transformation \( L(k_{23}) \).

Now we keep only the component with positive \( \rho \)-spin in the pair rest frame,
\[ \frac{m + \vec{k}_2}{m^2 - (\vec{k}_2)^2 - i\epsilon} \rightarrow \frac{m}{E(\vec{k}_2)k_{20} - E(\vec{k}_2) - i\epsilon}, \] (10)
with the positive-energy projector
\[ \Lambda_+(\vec{r}_2) = \frac{m + \vec{r}_2}{2m}. \] (11)

The approximation CIA-0 can then be defined as the replacement,
\[ G_{\rho\beta}(\vec{k}_2) \Gamma_{\lambda_1\rho\beta}(k_1, k_2; P_i) \]
\[ \rightarrow S_{\rho\beta}(L(k_{23})) \frac{m}{E(\vec{k}_2)k_{20} - E(\vec{k}_2) - i\epsilon} \]
\[ \times S_{\rho\beta}^{-1}(L(k_{23})) \Gamma_{\lambda_1\rho\beta}(k_1, L(k_{23})\vec{r}_2; P_i). \] (12)

In Eq. (1), as well as an analogous replacement for \( \Gamma_{\lambda_1\rho\beta}(k_1, k_2; P_i)G_{\rho\beta}(\vec{k}_2) \), which occurs in the diagrams of Figs. 1(B) and 1(E).

Note that the projector \( \Lambda_+ \) eliminates negative-energy states of nucleon 2 in the two-body rest frame, but this does not eliminate all Z-graph contributions from the calculation. They are still present through the negative-energy states of nucleon 3, and they are also regenerated to some extent when the state of nucleon 2 is boosted to other frames.

The approximation (12) may look complicated, but it is actually easy to implement in our numerical calculations. For instance, in the case of Fig. 1(C), it merely amounts to replacing in Eq. (B64) of Ref. [2] the off-shell energy \( \rho \) of nucleon 2 by the corresponding on-shell value \( E(\vec{p}) \) in the argument of the partial wave vertex function \( C(q \bar{p}\vec{p}M_{\lambda_1\lambda_2\lambda_3\lambda_4\rho\beta}T_\gamma) \), and restricting the summation over the \( \rho \)-spins of nucleon 2 to the positive-energy value \( \rho_2 = + \) only.
The electromagnetic current for an off-shell nucleon can be written in the form
\[ j_\mu^N(k', k) = f_0(k', k^2) F_{1N}(Q^2) \gamma^\mu \]
\[ + f_0'(k', k^2) F_{2N}(Q^2) \frac{i \sigma^\mu \nu q_\nu}{2m} \]
\[ + g_0(k', k^2) F_{3N}(Q^2) \Delta_-(k)\gamma^\mu \Lambda_-(k), \tag{13} \]
where \( f_0, f_0', \) and \( g_0 \) are nucleon off-shell form factors associated with the boson-nucleon vertices, and \( F_{1N} \) and \( F_{2N} \) are the usual electromagnetic Dirac and Pauli form factors. Because \( \Lambda_\pm \) projects onto negative-energy states, the form factor \( F_{3N} \) belongs to a term that contributes only if the nucleon is in a negative-energy state before and after the photon-nucleon vertex. We adopt the usual convention \( Q^2 = -q^2 \).

The isospin dependence of the electromagnetic form factors is, for \( i = 1, 2, 3 \) and the nucleon isospin projection \( \tau^3 \),
\[ F_i^N(Q^2) = F_{ip}(Q^2) \frac{1 + \tau^3}{2} + F_{in}(Q^2) \frac{1 - \tau^3}{2}. \tag{14} \]

In previous calculations [2], we found that the \( NN \) form factors are quite insensitive to the inclusion and variations of the off-shell nucleon form factors. Therefore, we employ in the calculations of this work the simpler off-shell nucleon current, with \( f_0 = f_0' = 1 \) and \( g_0 = 0 \). For the Dirac and Pauli form factors, we chose the parametrization of Galster [11] to compare with IARC results provided to us by Marcucci [12] who used the same parametrization.

With the CIA approximation in place, the electromagnetic \( NN \) form factors are calculated numerically from the \( NN \) vertex functions, which were obtained by solving the 3\( N \)N vertex. We adopt the usual convention \( Q^2 = -q^2 \).

### III. PRESENTATION AND DISCUSSION OF THE RESULTS

We calculated the electromagnetic \( NN \) form factors for three \( NN \) interaction models, W16, WJC-1, and WJC-2, for momentum transfer up to \( Q = 9 \) fm\(^{-1} \). The results are displayed in Figs. 2 and 3.

Because the form factors fall several orders of magnitude, and the traditional log plots tend to obscure differences in some places and overemphasize them in others, we divide them by simple scaling functions of the form
\[ F_i(Q) = F_0 e^{-Q/k}. \tag{15} \]

<table>
<thead>
<tr>
<th>Time</th>
<th>Event</th>
<th>Description</th>
</tr>
</thead>
<tbody>
<tr>
<td>12:00</td>
<td>Arrival</td>
<td>Customer arrives at the store</td>
</tr>
<tr>
<td>12:10</td>
<td>Conversation</td>
<td>Customer discusses product options with salesperson</td>
</tr>
<tr>
<td>12:30</td>
<td>Payment</td>
<td>Customer pays for purchase</td>
</tr>
<tr>
<td>12:35</td>
<td>Departure</td>
<td>Customer leaves the store</td>
</tr>
</tbody>
</table>

Note: The above table is a simplified representation of the events that might occur during a typical visit to a store.

Table I shows the parameters of the scaling function for each case. We also list the magnetic moments in Table II, and the charge and magnetic radii in Table III.

First, we start with a comparison of the curves for W16 in CIA and in CIA-0, which clearly demonstrates the high quality of the approximation. The differences between the exact calculation and the approximation are hardly noticeable up to values of \( Q \) around 7 fm\(^{-1} \), and in general appear to be insignificant. We may therefore assume that the results for WJC-1 and WJC-2 obtained here only in CIA-0 should also be very close to the exact CIA result.

Note that there are caveats to this conclusion: The quality of CIA-0 compared to CIA was really tested only for W16, a model with a very smooth choice for the definition of the kernel (and hence the vertex function) when both nucleons are off shell. The WJC models have a more complex off-shell structure (corresponding to the prescription C discussed in Ref. [6]) and their off-shell extrapolations will not be as smooth. In addition, WJC-1 has a mixed pseudoscalar-pseudovector pion-nucleon coupling, and it is conceivable that the pseudoscalar part of this coupling might introduce further differences between CIA and CIA-0 to which W16 is not sensitive. Our conclusions must therefore be taken with these particular grains of salt. In any case, CIA-0 should be a better approximation to CIA at smaller \( Q \), simply because the nucleon involved is taken less far off mass shell.

We turn now to a comparison of the form factors for different models of the \( NN \) interaction. The figures show that the WJC-2 form factors stay close to those of W16, whereas, in most cases, WJC-1 begins to deviate somewhat already at smaller values of \( Q \). WJC-2 and W16 are also close to the IARC results, typically up to about \( Q = 6 \) fm\(^{-1} \). This supports the conjecture made in Sec. I, namely that...
FIG. 2. (Color online) Charge form factors of the \(3N\) bound states, \(^3\)H (first row), \(^3\)He (second row), and the isoscalar (third row) and isovector (fourth row) combinations. In each case, the figure on the left shows the form factor in the traditional semilog plot, whereas the figure on the right shows the same form factor divided by a scaling function of Eq. (15) [2] on a linear scale. The solid line is the result for \(NN\) model W16 in CIA; the dotted line is the approximation CIA-0 for the same model. The dashed line is model WJC-1, and the dash-dotted line is model WJC-2, both in CIA-0. For comparison, the solid line with theoretical error bars is the result of an IARC calculation by Marcucci [12] based on the AV18/UIX potential. All calculations employ the on-shell single-nucleon current, with the Galster parametrization of the nucleon form factors [11]. The solid circles represent the experimental data [13–22].
FIG. 3. (Color online) Magnetic form factors of the $3N$ bound states, $^3$H (first row), $^3$He (second row), and the isoscalar (third row) and isovector (fourth row) combinations. In each case, the figure on the left shows the form factor in the traditional semilog plot, whereas the figure on the right shows the same form factor divided by a scaling function [2] on a linear scale. The meaning of the various curves is the same as in Fig. 2.

Apart from the issue of the type of pion-nucleon coupling, the CST models include other boson exchanges with off-shell coupling. Most notably, those due to scalar isoscalar ($\sigma_0$) and...
isovector ($\sigma_1$) exchanges have been found to have a very strong influence on the quality of the $NN$ fits and on the triton binding
[10]. One might expect them to have a strong influence on the $3N$ form factors as well.

The results indicate that this is only indirectly the case, namely through their effect on the binding energy. When the scalar off-shell coupling strength is varied without con-
straining the triton binding energy, the $3N$ form factors show substantial variations [2]. On the other hand, models W16 and WJC-2 have quite different scalar off-shell coupling constants, but yield the same triton binding energy. The close similarity of the $3N$ form factors, at least up to intermediate values of $Q$, implies that the electromagnetic structure of the $3N$ bound state is not modified too much by the scalar off-shell coupling. This conclusion receives even stronger support from the observation that the IARC calculation, which of course
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nucleon interaction models WJC-1 and WJC-2, which yield an
excellent description of the neutron-proton observables below 350 MeV for the most recent 2007 database [6]. The form
factors were calculated in complete impulse approximation, in which—for practical reasons—we replaced $3N$ vertex
functions with two off-mass-shell nucleons by corresponding vertex functions with only one nucleon off mass shell. This
procedure of approximating the full CIA results, denoted as “CIA-0,” was tested with the older two-nucleon model W16,
for which the full CIA result is also available, and found to be of very good quality.
We compare the form factors of WJC-1 and WJC-2 to those of W16, and also to calculations of nonrelativistic impulse approximation with relativistic corrections by Marcucci and collaborators [4,12]. Relating the observed differences in the approximation with relativistic corrections by Marcucci and of W16, and also to calculations of nonrelativistic impulse

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(iii) In some cases, model WJC-1 deviates moderately from the others. This appears to be due to its mixed pseudoscalar-pseudovector pion-nucleon coupling (WJC-2 and W16 have pure pseudovector pion-nucleon coupling, whereas in the nonrelativistic framework of IARC the two couplings are equivalent). In CST, pseudoscalar pion-nucleon coupling automatically includes Z-diagrams, whereas they are suppressed for pseudovector coupling. When Z-diagrams are effectively added to IARC in the form of $\gamma\pi NN$ contact interactions [4], the calculated form factors move closer to the experimental data, whereas the WJC-1 form factors lie farther away than the models with pure pseudovector coupling. This is consistent because the sign of the pseudoscalar coupling in WJC-1 is opposite to the one used in Ref. [4].

(iv) The results of this work confirm the conjecture formulated in Sec. I, namely that the reason for the good agreement of the CST models with the IARC results is the suppression of Z-diagrams through pseudovector pion-nucleon coupling.

(v) The CST two-nucleon interaction models WJC-1 and WJC-2 not only give an excellent fit to the available two-nucleon scattering observables, but also provide a solid basis for a relativistic theory of the $3N$ system. Without additional irreducible $3N$ forces, the $3N$ binding energy is reproduced, and the electromagnetic $3N$ form factors turn out very similar to previous nonrelativistic results. No unusually large interaction currents seem to be required to achieve a quantitative description of the experimental data.

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