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Nucleon-nucleon potential: Drift effects

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In the rest frame of a many-body system, used in the calculation of its static and scattering properties, the center of mass of a two-body subsystem is allowed to drift. We show, in a model-independent way, that drift corrections to the nucleon-nucleon potential are relatively large and arise from both one- and two-pion exchange processes. As far as chiral symmetry is concerned, corrections to these processes begin, respectively, at $\mathcal{O}(q^2)$ and $\mathcal{O}(q^4)$. The two-pion exchange interaction also yields a new spin structure, which promotes the presence of P waves in trinuclei and is associated with profile functions that do not coincide with either central or spin-orbit ones. In principle, the new spin terms should be smaller than the $\mathcal{O}(q^3)$ spin-orbit components. However, in the isospin-even channel, a large contribution defies this expectation and gives rise to the prediction of important drift effects.

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I. INTRODUCTION

This work is motivated by a private question posed by Alejandro Kievsky some years ago, concerning the possibility of novel forms of spin dependence in the nucleon-nucleon interaction, when one is not in the center-of-mass frame of the two-body system. In the study of static and scattering properties of many-body nuclei, calculations are performed in the rest frame of the larger system and a two-body subsystem is allowed to drift. This picture led him to introduce a phenomenological three-body force [1], which improved the description of the N-d vector analyzing power A_y.

Nowadays, the outer layers of the *NN* interaction, represented by one-pion and two-pion exchange potentials (OPEP and TPEP, respectively), are set in solid foundations owing to the use of chiral symmetry. Nuclear processes are dominated by the light quarks u and d, and one is not far from the massless limit, in which QCD becomes invariant under both isospin and chiral SU(2) × SU(2) transformations. Chiral symmetry is realized in the Nambu-Goldstone mode and the QCD vacuum can bear collective excitations, identified as pions. A suitably formulated chiral perturbation theory (ChPT) allows deviations from the massless limit to be treated systematically [2]. As low-energy QCD calculations are prevented by its non-Abelian character, in practice one works with chiral effective theories, in which elementary nucleons interact by exchanging pions.

In chiral perturbation, one uses a typical scale $q \ll 1$ GeV, set by either pion four-momenta or nucleon three-momenta. The leading term [3] in the *NN* interaction is the OPEP, at $\mathcal{O}(q^0)$. The TPEP begins at $\mathcal{O}(q^2)$ and two independent expansions up to $\mathcal{O}(q^4)$ are presently available. One of them, based on heavy-baryon ChPT [4], uses nonrelativistic Lagrangians from the very beginning and the inverse of the nucleon mass as an expansion parameter. The other one, produced recently by our group [5,6], is based on relativistic expressions, written in terms of observable coefficients and covariant loop integrals. The use of a relativistic language frees one from particular reference frames and allows a straightforward treatment of two-body interactions in which the center of mass is able to move. Here, we rely on our previous work in order to derive the drift contributions to the *NN* potential. For the sake of definiteness, we stay in the realm of three-body nuclei, but results can be easily generalized to larger systems.

Our presentation is organized as follows. In Sec. II, we review the dynamical role of two-body interactions in trinuclei. This sets the stage for the derivation of drift interactions, which is performed in Sec. III. Results are summarized in Sec. IV, whereas technical issues, concerning kinematics and spin operators, are left to appendices.

II. DYNAMICS

The interactions of a three-nucleon system in momentum space are represented by the operator W, defined by [7,8]

$$\langle \boldsymbol{p}_1', \boldsymbol{p}_2', \boldsymbol{p}_3' | \hat{W} | \boldsymbol{p}_1, \boldsymbol{p}_2, \boldsymbol{p}_3 \rangle = -(2\pi)^3 \delta^3(\boldsymbol{p}_1' + \boldsymbol{p}_2' + \boldsymbol{p}_3' - \boldsymbol{p}_1 - \boldsymbol{p}_2 - \boldsymbol{p}_3) \times \bar{t}_3(\boldsymbol{p}_1', \boldsymbol{p}_2', \boldsymbol{p}_3', \boldsymbol{p}_1, \boldsymbol{p}_2, \boldsymbol{p}_3), \quad (1)$$

where $\bar{t}_3$ is the proper part of the nonrelativistic three-body transition matrix. In configuration space, the position of nucleon i is described by $\boldsymbol{r}_i$ and one uses the Jacobi variables

$$\boldsymbol{R} = (\boldsymbol{r}_1 + \boldsymbol{r}_2 + \boldsymbol{r}_3)/3, \quad \boldsymbol{r} = \boldsymbol{r}_2 - \boldsymbol{r}_1, \quad \boldsymbol{\rho} = (2\boldsymbol{r}_3 - \boldsymbol{r}_1 - \boldsymbol{r}_2)/\sqrt{3}, \quad (2)$$

which correspond to

$$\boldsymbol{p}_1 = \frac{\boldsymbol{P}}{3} - \boldsymbol{p}_r - \frac{\boldsymbol{p}_\rho}{\sqrt{3}}, \quad \boldsymbol{p}_2 = \frac{\boldsymbol{P}}{3} + \boldsymbol{p}_r - \frac{\boldsymbol{p}_\rho}{\sqrt{3}}, \quad \boldsymbol{p}_3 = \frac{\boldsymbol{P}}{3} + \frac{2\boldsymbol{p}_\rho}{\sqrt{3}}. \quad (3)$$

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The Schrödinger equation for the internal degrees of freedom is obtained by using $\boldsymbol{P} = \boldsymbol{P}' = 0$ and given by

$$\left[-\frac{1}{m}\nabla^2_{r'} - \frac{1}{m}\nabla^2_{\rho'} - \epsilon\right]\psi(\boldsymbol{r}', \boldsymbol{\rho}') = -\left[\frac{\sqrt{3}}{2}\right]^3 \int d\boldsymbol{r}\, d\boldsymbol{\rho} \times W(\boldsymbol{r}', \boldsymbol{\rho}'; \boldsymbol{r}, \boldsymbol{\rho})\psi(\boldsymbol{r}, \boldsymbol{\rho}), \tag{4}$$

with

$$W(\boldsymbol{r}', \boldsymbol{\rho}'; \boldsymbol{r}, \boldsymbol{\rho}) = -\frac{1}{(2\pi)^{12}}\left[\frac{2}{\sqrt{3}}\right]^6 \int d\boldsymbol{Q}_r d\boldsymbol{Q}_\rho d\boldsymbol{q}_r d\boldsymbol{q}_\rho \times e^{i[\boldsymbol{Q}_r\cdot(\boldsymbol{r}'-\boldsymbol{r})+\boldsymbol{Q}_\rho\cdot(\boldsymbol{\rho}'-\boldsymbol{\rho})+\boldsymbol{q}_r\cdot(\boldsymbol{r}'+\boldsymbol{r})/2+\boldsymbol{Q}_\rho\cdot(\boldsymbol{\rho}'+\boldsymbol{\rho})/2]} \times \bar{t}_3(\boldsymbol{Q}_r, \boldsymbol{Q}_\rho, \boldsymbol{q}_r, \boldsymbol{q}_\rho), \tag{5}$$

$\boldsymbol{Q}_i = (\boldsymbol{p}'_i + \boldsymbol{p}_i)/2$ and $\boldsymbol{q}_i = (\boldsymbol{p}'_i - \boldsymbol{p}_i)$, for $i = (r, \rho)$.

In this work we are interested in describing two-body interactions between nucleons 1 and 2 and note that the conservation of $\boldsymbol{p}_3$ implies $\boldsymbol{q}_\rho = 0$. We write

$$\bar{t}_3(\boldsymbol{p}'_1, \boldsymbol{p}'_2, \boldsymbol{p}'_3; \boldsymbol{p}_1, \boldsymbol{p}_2, \boldsymbol{p}_3) = (2\pi)^3\left[\frac{\sqrt{3}}{2}\right]^3 \delta^3(\boldsymbol{q}_\rho) \times \bar{t}_2(\boldsymbol{Q}_\rho, \boldsymbol{Q}_r, \boldsymbol{q}_r), \tag{6}$$

where $\bar{t}_2$ is the two-body t matrix, and the corresponding potential becomes

$$W_2(\boldsymbol{r}', \boldsymbol{\rho}'; \boldsymbol{r}, \boldsymbol{\rho}) = -\frac{1}{(2\pi)^9}\left[\frac{2}{\sqrt{3}}\right]^3 \int d\boldsymbol{Q}_r\, d\boldsymbol{Q}_\rho\, d\boldsymbol{q}_r \times e^{-i[\boldsymbol{Q}_r\cdot(\boldsymbol{r}'-\boldsymbol{r})+\boldsymbol{Q}_\rho\cdot(\boldsymbol{\rho}'-\boldsymbol{\rho})+\boldsymbol{q}_r\cdot(\boldsymbol{r}'+\boldsymbol{r})/2]} \times \bar{t}_2(\boldsymbol{Q}_r, \boldsymbol{Q}_\rho, \boldsymbol{q}_r). \tag{7}$$

In isospin space, the amplitude $\bar{t}_2$ reads

$$\bar{t}_2 = t^+ + \boldsymbol{\tau}^{(1)}\cdot\boldsymbol{\tau}^{(2)} t^-. \tag{8}$$

The usual spin decomposition is obtained by going to the center-of-mass (c.m.) frame of the two-body system, where one finds

$$t_2^\pm]_{\text{c.m.}} = t_C^\pm + \frac{\boldsymbol{\Omega}_{LS}}{m^2}t_{LS}^\pm + \frac{\boldsymbol{\Omega}_{SS}}{m^2}t_{SS}^\pm + \frac{\boldsymbol{\Omega}_T}{m^2}t_T^\pm + \frac{\boldsymbol{\Omega}_Q}{m^4}t_Q^\pm, \tag{9}$$

with two-component operators defined by

$$\boldsymbol{\Omega}_{LS} = i(\boldsymbol{\sigma}^{(1)} + \boldsymbol{\sigma}^{(2)})\cdot\boldsymbol{q}_r\times\boldsymbol{Q}_r/2, \tag{10}$$
$$\boldsymbol{\Omega}_{SS} = \boldsymbol{q}_r^2\boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{\sigma}^{(2)}, \tag{11}$$
$$\boldsymbol{\Omega}_T = -\boldsymbol{q}_r^2(3\boldsymbol{\sigma}^{(1)}\cdot\hat{\boldsymbol{q}}_r\,\boldsymbol{\sigma}^{(2)}\cdot\hat{\boldsymbol{q}}_r - \boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{\sigma}^{(2)}), \tag{12}$$
$$\boldsymbol{\Omega}_Q = 4\boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{q}_r\times\boldsymbol{Q}_r\,\boldsymbol{\sigma}^{(2)}\cdot\boldsymbol{q}_r\times\boldsymbol{Q}_r. \tag{13}$$

In this formulation, the two-body interaction does not depend on $\boldsymbol{Q}_\rho$ and is completely decoupled from the larger system it is immersed in. The Fourier transform of this result produces the configuration space potential, given by

$$W_2(\boldsymbol{r}', \boldsymbol{\rho}'; \boldsymbol{r}, \boldsymbol{\rho}) = \delta^3(\boldsymbol{r}'-\boldsymbol{r})\delta^3(\boldsymbol{\rho}'-\boldsymbol{\rho})\left[\frac{2}{\sqrt{3}}\right]^3 V(\boldsymbol{r})^\pm]_{\text{c.m.}}, \quad V(r)^\pm]_{\text{c.m.}} = V_C^\pm + V_{LS}^\pm\Omega_{LS} + V_{SS}^\pm\Omega_{SS} + V_T^\pm\Omega_T, \tag{14}$$

where we have kept only local and spin-orbit contributions and the spin operators read

$$\Omega_{LS} = \boldsymbol{L}\cdot(\boldsymbol{\sigma}^{(1)} + \boldsymbol{\sigma}^{(2)})/2, \tag{15}$$
$$\Omega_{SS} = \boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{\sigma}^{(2)}, \tag{16}$$
$$\Omega_T = 3\boldsymbol{\sigma}^{(1)}\cdot\hat{\boldsymbol{r}}\boldsymbol{\sigma}^{(2)}\cdot\hat{\boldsymbol{r}} - \boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{\sigma}^{(2)}. \tag{17}$$

The radial functions are given by

$$V_C^\pm(r) = U_C^\pm(x), \tag{18}$$
$$V_{LS}^\pm(r) = \frac{\mu^2}{m^2}\frac{1}{x}\frac{d}{dx}U_{LS}^\pm(x), \tag{19}$$
$$V_{SS}^\pm(r) = -\frac{\mu^2}{m^2}\left[\frac{d^2}{dx^2} + \frac{2}{x}\frac{d}{dx}\right]U_{SS}^\pm(x), \tag{20}$$
$$V_T^\pm(r) = \frac{\mu^2}{m^2}\left[\frac{d^2}{dx^2} - \frac{1}{x}\frac{d}{dx}\right]U_T^\pm(x), \tag{21}$$

with $x = \mu r$ and

$$U_I^\pm(x) = -\int\frac{d^3q}{(2\pi)^3}e^{i\boldsymbol{q}\cdot\boldsymbol{r}}t_I^\pm(q), \quad I = \{C, LS, SS, T\}. \tag{22}$$

As we discuss below, the fact that the two-body c.m. is allowed to drift gives rise to extra interaction operators in the potential.

III. DRIFT TERMS

Corrections to the NN potential owing to the motion of the c.m. are derived by evaluating $\mathcal{T}$, the covariant t-matrix for the on-shell process $N(p_1)N(p_2) \to N(p'_1)N(p'_2)$, and writing the result in terms of two-component spinors, using the expressions of Appendix B. This gives rise to an amplitude expanded in terms of Pauli spin operators. Dividing it by the factor $4mE$ present in the relativistic normalization, one obtains the amplitude $\bar{t}_2$, which is to be fed into Eq. (7). In this work we concentrate on contributions from processes due to the exchanges of one and two pions.

The transformation of a t-matrix into a potential to be used in a dynamical equation is not trivial and depends on a number of important conventions. These range from the very nature of the equation adopted to tacit assumptions concerning the off-shell behavior of the potential. The latter class of effects appears as corrections to leading-order effects and was discussed in a comprehensive paper by Friar [9]. Here we stick to the conventions used long ago by Partovi and Lomon [10] and da Rocha and Robilotta [11].

A. OPEP

The covariant amplitude for on-shell nucleons reads

$$\mathcal{T} = \boldsymbol{\tau}^{(1)}\cdot\boldsymbol{\tau}^{(2)}\frac{g_A^2 m^2}{f_\pi^2}\frac{1}{q^2 - \mu^2}[\bar{u}\gamma_5 u]^{(1)}[\bar{u}\gamma_5 u]^{(2)}, \tag{23}$$

where g_A, f_π, μ, and m are, respectively, the axial and pion decay constants and the pion and nucleon masses. Using Eq. (A8) for the momentum q and Eq. (B4) for the spinor matrix element, one finds the two-component amplitude

$$
\begin{aligned}
\mathcal{T} = \boldsymbol{\tau}^{(1)} \cdot \boldsymbol{\tau}^{(2)} \frac{g_A^2 m^2}{f_\pi^2} & \frac{1}{(\boldsymbol{q}_r^2+\mu^2) - 4(\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho)^2/3\mathcal{E}^2} \\
& \times \mathcal{N}^2\{[m+\mathcal{E}/2+2\boldsymbol{Q}_r\cdot\boldsymbol{Q}_\rho/\sqrt{3}\mathcal{E}]\boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{q}_r \\
& -2\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho\boldsymbol{\sigma}^{(1)}(\boldsymbol{Q}_r+\boldsymbol{Q}_\rho/\sqrt{3})/\sqrt{3}\mathcal{E}\}\{[m+\mathcal{E}/2 \\
& -2\boldsymbol{Q}_r\cdot\boldsymbol{Q}_\rho/\sqrt{3}\mathcal{E}]\boldsymbol{\sigma}^{(2)}\cdot\boldsymbol{q}_r \\
& +2\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho\boldsymbol{\sigma}^{(2)}\cdot(\boldsymbol{Q}_r-\boldsymbol{Q}_\rho/\sqrt{3})/\sqrt{3}\mathcal{E}\},
\end{aligned} \tag{24}
$$

where $\mathcal{N}^2$ is a normalization factor,

$$
\begin{aligned}
\mathcal{N}^2 = \{&[(m+\mathcal{E}/2)^2+[4(\boldsymbol{Q}_r\cdot\boldsymbol{Q}_\rho)^2-(\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho)^2]/3\mathcal{E}^2]^2 \\
&-16(m+\mathcal{E}/2)^2(\boldsymbol{Q}_r\cdot\boldsymbol{Q}_\rho)^2/3\mathcal{E}^2\}^{-1/2}
\end{aligned} \tag{25}
$$

and $\mathcal{E}$ is the total energy of the two-nucleon system, determined by the condition

$$
\begin{aligned}
\mathcal{E}^4 - 4\big(m^2+\boldsymbol{q}_r^2+4\boldsymbol{Q}_r^2+\boldsymbol{Q}_\rho^2/3\big)\mathcal{E}^2 + (4/3)[(\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho)^2 \\
+4(\boldsymbol{Q}_r\cdot\boldsymbol{Q}_\rho)^2] = 0.
\end{aligned} \tag{26}
$$

This t-matrix is fully relativistic and contains no approximations. All its terms involving the variable $\boldsymbol{Q}_\rho$ vanish in the rest frame of the two-body system and therefore can be interpreted as drift effects. With the provisos discussed in Ref. [9], it could already be used as input into a momentum-space dynamical equation. Alternatively, in the framework of ChPT, one might wish to rewrite it as a power series, truncated at a given order.

In configuration space, the variables $\boldsymbol{Q}_r$ and $\boldsymbol{Q}_\rho$ correspond to nonlocal operators and are usually associated with gradients acting on the wave function. To restrict the corresponding complications to a minimum, we remain in the limited scope of Eq. (4) and keep only terms linear in these momenta. This procedure is referred to as the *linear gradient approximation*.

Within this approximation, the OPEP retains its usual local form, given by

$$
t_{SS}^{-} = -t_T^{-} \overset{\text{lga}}{\simeq} \frac{g_A^2 m^2}{12 f_\pi^2}\frac{1}{\boldsymbol{q}_r^2+\mu^2} + \text{local corrections}. \tag{27}
$$

B. TPEP

Quite generally, for each isospin channel, the spin content of the TPEP is given by [5]

$$
\begin{aligned}
\mathcal{T}^{\pm} = [\bar{u}u]^{(1)}[\bar{u}u]^{(2)}(\mathcal{I}_{DD}^{\pm}) - \frac{i}{2m}[\bar{u}u]^{(1)}[\bar{u}\sigma_{\mu\lambda}(p'-p)^\mu u]^{(2)} \\
\times(\mathcal{I}_{DB}^{\pm})^\lambda - \frac{i}{2m}[\bar{u}\sigma_{\mu\lambda}(p'-p)^\mu u]^{(1)}[\bar{u}u]^{(2)}(\mathcal{I}_{BD}^{\pm})^\lambda \\
-\frac{1}{4m^2}[\bar{u}\sigma_{\mu\lambda}(p'-p)^\mu u]^{(1)}[\bar{u}\sigma_{\nu\rho}(p'-p)^\nu u]^{(2)}(\mathcal{I}_{BB}^{\pm})^{\lambda\rho},
\end{aligned} \tag{28}
$$

where the functions $\mathcal{I}$ involve loop integrals and have a Lorentz structure realized in terms of the kinematical variables W, z, and q, defined in Appendix A. Terms proportional to q do not contribute for on-shell nucleons and we have

$$
(\mathcal{I}_{DB}^{\pm})^\lambda = \frac{W^\lambda}{2m}\mathcal{I}_{DB}^{(w)\pm} + \frac{z^\lambda}{2m}\mathcal{I}_{DB}^{(z)\pm}, \tag{29}
$$

$$
(\mathcal{I}_{BD}^{\pm})^\lambda = \frac{W^\lambda}{2m}\mathcal{I}_{DB}^{(w)\pm} - \frac{z^\lambda}{2m}\mathcal{I}_{DB}^{(z)\pm}, \tag{30}
$$

$$
(\mathcal{I}_{BB}^{\pm})^{\lambda\rho} = g^{\lambda\rho}\mathcal{I}_{BB}^{(g)\pm} + \frac{W^\lambda W^\rho}{4m^2}\mathcal{I}_{BB}^{(w)\pm} + \frac{z^\lambda z^\rho}{4m^2}\mathcal{I}_{BB}^{(z)\pm}. \tag{31}
$$

The amplitudes $\mathcal{I}$ were explicitly calculated in Ref. [5], as functions of the invariants W^2, z^2, and q^2, and the two-pion exchange interaction is described by

$$
\begin{aligned}
\mathcal{T}^{\pm} = [\bar{u}u]^{(1)}[\bar{u}u]^{(2)}\left[\mathcal{I}_{DD}^{\pm} + \frac{q^2}{2m^2}\mathcal{I}_{DB}^{(w)\pm} + \frac{q^4}{16m^4}\mathcal{I}_{BB}^{(w)\pm}\right] \\
-\frac{i}{2m}\{[\bar{u}u]^{(1)}[\bar{u}\sigma_{\mu\lambda}(p'-p)^\mu u]^{(2)} \\
-[\bar{u}\sigma_{\mu\lambda}(p'-p)^\mu u]^{(1)}[\bar{u}u]^{(2)}\}\frac{z^\lambda}{2m}\left[\mathcal{I}_{DB}^{(w)\pm}+\mathcal{I}_{DB}^{(z)\pm}\right. \\
\left.+\frac{q^2}{4m^2}\mathcal{I}_{BB}^{(w)\pm}\right] - \frac{1}{4m^2}[\bar{u}\sigma_{\mu\lambda}(p'-p)^\mu u]^{(1)} \\
\times[\bar{u}\sigma_{\nu\rho}(p'-p)^\nu u]^{(2)} \\
\times\left[g^{\lambda\rho}\mathcal{I}_{BB}^{(g)\pm} + \frac{z^\lambda z^\rho}{4m^2}\big(-\mathcal{I}_{BB}^{(w)\pm}+\mathcal{I}_{BB}^{(z)\pm}\big)\right].
\end{aligned} \tag{32}
$$

This result can be recast in a form similar to Eq. (24), by using the spinor matrix elements given in Appendix B. One finds

$$
\begin{aligned}
\mathcal{T}^{\pm} = \mathcal{N}^2\big\{\mathcal{I}_{DD}^{\pm}\big[2m(m+\mathcal{E}/2) - 2(\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho)^2/3\mathcal{E}^2 + \boldsymbol{q}_r^2/2 \\
- i\boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{q}_r\times\big(\boldsymbol{Q}_r+\boldsymbol{Q}_\rho/\sqrt{3}\big)\big] \\
\times\big[2m(m+\mathcal{E}/2) - 2(\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho)^2/3\mathcal{E}^2 + \boldsymbol{q}_r^2/2 \\
- i\boldsymbol{\sigma}^{(2)}\cdot\boldsymbol{q}_r\big(\boldsymbol{Q}_r-\boldsymbol{Q}_\rho/\sqrt{3}\big)\big] + \cdots\big\},
\end{aligned} \tag{33}
$$

and its full drift content becomes explicit. However, for the sake of simplicity, we remain in the framework of the linear gradient approximation. Using Eqs. (A6)–(A8), one learns that, in this case, the variables W^2, z^2, and q^2 do not depart from their c.m. values and the only sources of drift corrections are the spin functions. The results of Appendix B yield the following nonrelativistic amplitude:

$$
t_2^{\pm} = t_2^{\pm}]_{\text{c.m.}} + \frac{\boldsymbol{\Omega}_D}{m^2}t_D^{\pm}, \tag{34}
$$

where the drift operator $\boldsymbol{\Omega}_D$ is given by

$$
\boldsymbol{\Omega}_D = i(\boldsymbol{\sigma}^{(1)}-\boldsymbol{\sigma}^{(2)})\cdot\boldsymbol{q}_r\times\boldsymbol{Q}_\rho/2\sqrt{3} \tag{35}
$$

and the profile functions read

$$
\begin{aligned}
t_D^{\pm} = -\frac{m}{e}\left\{\frac{4m^2}{\lambda^2}\left(1+\frac{\boldsymbol{q}_r^2}{\lambda^2}\right)\left[\mathcal{I}_{DD}^{\pm} - \frac{\boldsymbol{q}_r^2}{2m^2}\mathcal{I}_{DB}^{(w)\pm}\right.\right. \\
\left.\left.+\frac{\boldsymbol{q}_r^4}{16m^4}\mathcal{I}_{BB}^{(w)\pm}\right] + \frac{e\boldsymbol{q}_r^2}{m\lambda^2}\mathcal{I}_{BB}^{(g)\pm}\right\},
\end{aligned} \tag{36}
$$

with $\lambda^2 = 4m(e+m)$ and $e = \sqrt{m^2+\boldsymbol{q}_r^2+4\boldsymbol{Q}_r^2}$. This result is fully model independent, since it springs directly from Lorentz covariance and is constrained just by the linear gradient approximation. The profile functions $t_D^{\pm}$ do not coincide with any other components of the TPEP, which, in the same approximation, are given by [5]

$$t_C^{\pm} = \frac{m}{e}\left\{\left(1+\frac{\boldsymbol{q}_r^2}{\lambda^2}\right)^2\left[\mathcal{I}_{DD}^{\pm} - \frac{\boldsymbol{q}_r^2}{2m^2}\mathcal{I}_{DB}^{(w)\pm} + \frac{\boldsymbol{q}_r^4}{16m^4}\mathcal{I}_{BB}^{(w)\pm}\right] + \frac{\boldsymbol{q}_r^4}{16m^4}\mathcal{I}_{BB}^{(g)\pm}\right\}, \tag{37}$$

$$t_{LS}^{\pm} = \frac{m}{e}\left\{\left(1+\frac{\boldsymbol{q}_r^2}{\lambda^2}\right)\left[-\frac{4m^2}{\lambda^2}\mathcal{I}_{DD}^{\pm} + \left(1+\frac{2\boldsymbol{q}_r^2}{\lambda^2}\right)\mathcal{I}_{DB}^{(w)\pm} + \mathcal{I}_{DB}^{(z)\pm} - \frac{\boldsymbol{q}_r^2}{4m^2}\left(1+\frac{\boldsymbol{q}_r^2}{\lambda^2}\right)\mathcal{I}_{BB}^{(w)\pm} - \frac{\boldsymbol{q}_r^2}{4m^2}\left(1+\frac{4m^2}{\lambda^2}\right)\right] \times \mathcal{I}_{BB}^{(g)\pm}\right\}, \tag{38}$$

$$t_T^{\pm} = t_{SS}^{\pm}/2 = \frac{m}{e}\left\{-\frac{1}{12}\mathcal{I}_{BB}^{(g)\pm}\right\}. \tag{39}$$

We consider here the expansion of the TPEP to $\mathcal{O}(q^4)$, using Eqs. (9) and (34), which requires $t_C^{\pm}\to\mathcal{O}(q^4)$ and $\{t_{LS}^{\pm}, t_T^{\pm}, t_{SS}^{\pm}, t_D^{\pm}\}\to\mathcal{O}(q^2)$. The expansion of the various profile functions is performed using the results $\{\mathcal{I}_{BB}^{(g)+}, \mathcal{I}_{BB}^{(w)+}, \mathcal{I}_{BB}^{(z)+}\}\to\mathcal{O}(q^0)$, $\{\mathcal{I}_{DB}^{(w)+}, \mathcal{I}_{DB}^{(z)+}, \mathcal{I}_{DB}^{(w)-}, \mathcal{I}_{DB}^{(z)-}, \mathcal{I}_{BB}^{(g)-}\}\to\mathcal{O}(q^1)$, $\{\mathcal{I}_{DD}^{-}\}\to\mathcal{O}(q^2)$, $\{\mathcal{I}_{DD}^{+}\}\to\mathcal{O}(q^3)$, and $\{\mathcal{I}_{BB}^{(w)-}, \mathcal{I}_{BB}^{(z)-}\}\sim 0$, and one finds

$$t_D^{+} = -\frac{m}{e}\left\{\frac{\boldsymbol{q}_r^2}{8m^2}\mathcal{I}_{BB}^{(g)+} + \frac{1}{2}\left[\mathcal{I}_{DD}^{+} - \frac{\boldsymbol{q}_r^2}{2m^2}\mathcal{I}_{DB}^{(w)+}\right]\right\} \times \to \{\mathcal{O}(q^2) + [\mathcal{O}(q^3)]\}, \tag{40}$$

$$t_D^{-} = -\frac{m}{e}\left\{\frac{1}{2}\mathcal{I}_{DD}^{-}\right\} \to \{\mathcal{O}(q^2)\}. \tag{41}$$

In the expression for t_D^{+}, the term within square brackets is $\mathcal{O}(q^3)$. Nevertheless, we have kept it, for it is anomalously large. Considering comparable terms in Eq. (38), one writes

$$t_C^{+} = \frac{m}{e}\left\{\mathcal{I}_{DD}^{+} - \frac{\boldsymbol{q}_r^2}{2m^2}\mathcal{I}_{DB}^{(w)+}\right\} \to \{\mathcal{O}(q^3)\}, \tag{42}$$

$$t_{LS}^{+} = \frac{m}{e}\left\{\mathcal{I}_{DB}^{(w)+} + \mathcal{I}_{DB}^{(z)+} - \left[\frac{\boldsymbol{q}_r^2}{4m^2}\left(\mathcal{I}_{BB}^{(w)+} + \frac{3}{2}\mathcal{I}_{BB}^{(g)+}\right)\right]\right\} \to \{\mathcal{O}(q) + [\mathcal{O}(q^2)]\}, \tag{43}$$

$$t_C^{-} = \frac{m}{e}\{\mathcal{I}_{DD}^{-}\} \to \{\mathcal{O}(q^2)\}, \tag{44}$$

$$t_{LS}^{-} = \frac{m}{e}\left\{\mathcal{I}_{DB}^{(w)-} + \mathcal{I}_{DB}^{(z)-} - \left[\frac{1}{2}\mathcal{I}_{DD}^{-}\right]\right\} \to \{\mathcal{O}(q) + [\mathcal{O}(q^2)]\}. \tag{45}$$

These results show that the drift potential has little affinity with the spin-orbit term and, at the chiral order considered, can be written as

$$t_D^{+} = \frac{3\boldsymbol{q}_r^2}{4m^2}t_{SS}^{+} - \frac{1}{2}t_C^{+}, \tag{46}$$

$$t_D^{-} = -\frac{1}{2}t_C^{-}. \tag{47}$$

The Fourier transform of Eq. (34) yields the configuration-space structure

$$V(r)^{\pm} = V(r)^{\pm}\big]_{\text{c.m.}} + V_D^{\pm}\Omega_D, \tag{48}$$

with

$$\Omega_D = \frac{1}{4\sqrt{3}}(\boldsymbol{\sigma}^{(1)} - \boldsymbol{\sigma}^{(2)})\cdot\boldsymbol{r}\times(-i\overleftrightarrow{\nabla}_\rho), \tag{49}$$

$$V_D^{\pm}(r) = \frac{\mu^2}{m^2}\frac{1}{x}\frac{d}{dx}U_D^{\pm}(x), \tag{50}$$

$$U_D^{\pm}(x) = -\int\frac{d\boldsymbol{q}_r}{(2\pi)^3}e^{i\boldsymbol{q}_r\cdot\boldsymbol{r}}t_D^{\pm}(q_r). \tag{51}$$

The presence of the operator

$$\overleftrightarrow{\nabla}_\rho = \overrightarrow{\nabla}_\rho - \overleftarrow{\nabla}_\rho \tag{52}$$

in Eq. (49) ensures that results are symmetric under the exchange of initial and final states. Using results (46) and (47), one has

$$U_D^{+} = \tfrac{3}{4}V_{SS}^{+} - \tfrac{1}{2}V_C^{+}, \tag{53}$$

$$U_D^{-} = -\tfrac{1}{2}V_C^{-}. \tag{54}$$

In Figs. 1 and 2 we display the profile functions for the drift and spin-orbit potentials, derived from our $\mathcal{O}(q^4)$ expansion of the TPEP [5,6]. These results do not include short-range effects and cannot be trusted for $r < 1$ fm. We recall that both components of the force are multiplied by $\mathcal{O}(q^2)$ spin operators and hence we need to keep just $\mathcal{O}(q^2)$ terms in V_D. As shown in Eqs. (40)–(45), in principle one should have $V_{LS}^{\pm}\sim\mathcal{O}(q) > V_D^{\pm}\sim\mathcal{O}(q^2)$. These expectations are confirmed in the figures, provided one uses the $\mathcal{O}(q^2)$ dotted curve for V_D^{+}. However, when the $\mathcal{O}(q^3)$ term associated with the central potential is kept, one has a complete subversion of the expected chiral hierarchy, associated with the prediction of a rather large effect in the isospin-even channel.

To produce a feeling for the role of drift interactions in trinuclei, we note that their ground states contain S, P, and D waves, but they are heavily dominated by the principal S component, which is fully symmetric under the exchange of nucleon coordinates. Using the notation of Ref. [12], we write

$$|S\rangle = S(\boldsymbol{r},\boldsymbol{\rho})\Gamma_{1/2i}^{1/2\mu}, \tag{55}$$

where

$$\Gamma_{1/2i}^{1/2\mu} = \frac{1}{\sqrt{2}}\left[|m^-\mu\rangle_S|m^+i\rangle_I - |m^+\mu\rangle_S|m^-i\rangle_I\right] \tag{56}$$

is the totally antisymmetric spin-isospin $=(1/2, 1/2)$ wave function with third components μ and i, whereas $|m^+\rangle$ and $|m^-\rangle$ represent, respectively, even and odd mixed-symmetry states under permutation of particles 1 and 2. The leading term of the function $S(\boldsymbol{r},\boldsymbol{\rho})$ is known [13] to depend just on the hyper-radius $\xi\equiv\sqrt{\boldsymbol{r}^2+\boldsymbol{\rho}^2}$ and hence the most important coupling introduced by the drift potential is associated with

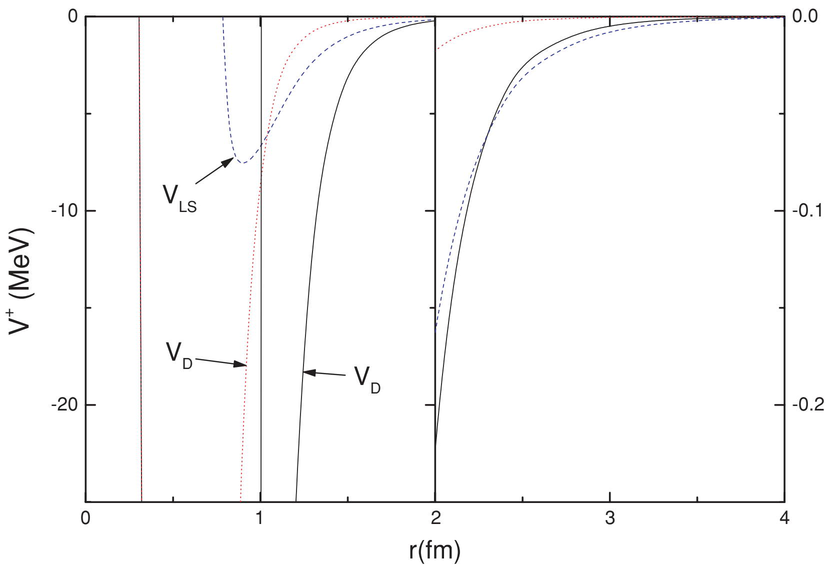

FIG. 1. (Color online) Isospin-even drift (full and dotted lines) and spin-orbit (dashed line) potentials; the dotted line is $\mathcal{O}(q^2)$; the full one is $\mathcal{O}(q^2) + \mathcal{O}(q^3)$.

the structure

$$\Omega_D|S\rangle \sim \Omega_D S(\xi)\Gamma^{1/2\mu}_{1/2i} = \frac{2\pi r\rho}{3\sqrt{3}\xi}\frac{\partial S(\xi)}{\partial\xi}\{\{-[[Y_1(\hat{\boldsymbol{r}})\otimes \times Y_1(\hat{\boldsymbol{\rho}})]_1 \otimes |m^+\rangle_S]^{\mu}_{1/2} + \sqrt{2}[[Y_1(\hat{\boldsymbol{r}})\otimes Y_1(\hat{\boldsymbol{\rho}})]_1 \otimes |s\rangle_S]^{\mu}_{1/2}\} \times |m^+ i\rangle_I + [[Y_1(\hat{\boldsymbol{r}})\otimes Y_1(\hat{\boldsymbol{\rho}})]_1 \otimes \times |m^-\rangle_S]^{\mu}_{1/2}|m^- i\rangle_I\}, \tag{57}$$

$|s\rangle_S$ being the spin 3/2 state. This result indicates that the drift potential enhances the role of P waves in trinuclei, as one might have guessed directly from Eq. (49).

IV. SUMMARY

In nuclei containing three or more nucleons, the center of mass of a two-body subsystem is allowed to drift. This kind of movement does affect the forms of both one- and two-pion exchange contributions and gives rise to important nonlocal corrections to the potential. As interactions of this type are difficult to deal with in configuration space, we have restricted ourselves to the simplest possible nonlocal operators, proportional to single gradients acting on the wave function, that arise in two-pion processes. Using a relativistic chiral expansion of the two-pion exchange NN potential to $\mathcal{O}(q^4)$ derived previously, we have shown, in a model-independent way, that the profile functions of the drift corrections do not coincide with any of its components. The spin dependence of the drift term is implemented by the operator

$$\Omega_D = \frac{1}{4\sqrt{3}}(\boldsymbol{\sigma}^{(1)} - \boldsymbol{\sigma}^{(2)})\cdot \boldsymbol{r}\times(-i\overleftrightarrow{\nabla}_\rho),$$

where $\boldsymbol{r}$ and $\boldsymbol{\rho}$ are Jacobi coordinates associated with two and three bodies. This structure promotes couplings between S and P waves, enhancing the role of the latter in trinuclei.

As far as chiral symmetry is concerned, drift corrections begin at $\mathcal{O}(q^4)$ and, in principle, should be smaller than spin-orbit terms, which begin at $\mathcal{O}(q^3)$. However, in the isospin-

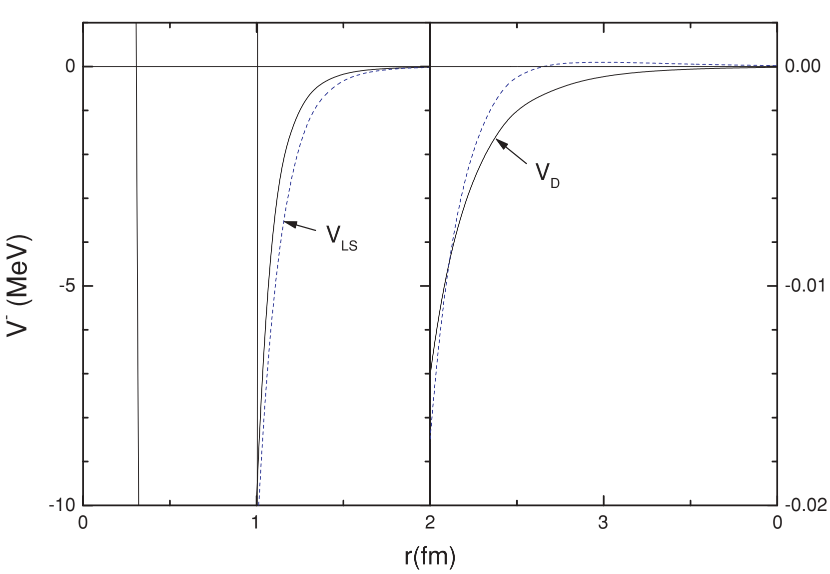

FIG. 2. (Color online) Isospin-odd drift (full line) and spin-orbit (dashed line) potentials.

even channel, the same dynamical contribution that makes its $\mathcal{O}(q^3)$ central component larger than the $\mathcal{O}(q^2)$ odd counterpart subverts the expected chiral hierarchy and gives rise to the prediction of important drift effects.

Note added in proof. I thank Prof. A. Gal for drawing my attention, in August 2006, to a paper by Close and Dalitz [14], in which a drift operator identical with that derived here was used in the framework of the quark model. Please also see the work by Kim [15] quoted in their paper.

ACKNOWLEDGMENTS

I would like to thank Alejandro Kievsky for stressing the importance of having interactions that are symmetric under exchanges between initial and final states and Renato Higa for supplying his numerical profile functions for the chiral two-pion exchange potential.

APPENDIX A: KINEMATICS

The conventions used here are the same as in Ref. [5]. The initial and final nucleon momenta are denoted by p and p' and we define the variables

$$W = p_1 + p_2 = p'_1 + p'_2, \tag{A1}$$

$$z = [(p_1 + p'_1) - (p_2 + p'_2)]/2, \tag{A2}$$

$$q = p'_1 - p_1 = p_2 - p'_2. \tag{A3}$$

The interacting nucleons are assumed to be on shell and the following constraints hold:

$$m^2 = (W^2 + z^2 + q^2)/4, \tag{A4}$$

$$W \cdot z = W \cdot q = z \cdot q = 0. \tag{A5}$$

Using the Jacobi variables defined in Eq. (3), one has

$$W = (\mathcal{E}, -2\boldsymbol{Q}_\rho/\sqrt{3}), \tag{A6}$$

$$z = (4\boldsymbol{Q}_r \cdot \boldsymbol{Q}_\rho/\mathcal{E}\sqrt{3}, -2\boldsymbol{Q}_r), \tag{A7}$$

$$q = (2\boldsymbol{q}_r \cdot \boldsymbol{Q}_\rho/\mathcal{E}\sqrt{3}, -\boldsymbol{q}_r), \tag{A8}$$

where $\mathcal{E}$ is the total energy of the two-body system. If there were no drift, this energy would be written in terms of the single-particle c.m. energy e as

$$\mathcal{E}_{\text{c.m.}} = 2e = 2\sqrt{m^2 + \boldsymbol{q}_r^2 + 4\boldsymbol{Q}_r^2}. \tag{A9}$$

Explicit calculation yields

$$\mathcal{E}^2 + (4/3)\left[(\boldsymbol{q}_r \cdot \boldsymbol{Q}_\rho)^2/\mathcal{E}^2 + 4(\boldsymbol{Q}_r \cdot \boldsymbol{Q}_\rho)^2/\mathcal{E}^2 - \boldsymbol{Q}_\rho^2\right] = 4e^2 \tag{A10}$$

and hence, in the linear gradient approximation,

$$\mathcal{E} \overset{\text{lga}}{\simeq} 2e. \tag{A11}$$

APPENDIX B: SPIN OPERATORS

We present here the changes induced in the spin operators owing the drift of the two-body c.m. With our conventions, we write

$$[\bar{u}\Gamma u]^{(i)} = \left\{\mathcal{N}\chi^\dagger\left[E'+m, -\boldsymbol{\sigma}\cdot\boldsymbol{p}'\right]\begin{bmatrix} \cdot & \cdot \\ \cdot & \cdot \end{bmatrix}\begin{bmatrix} E+m \\ \boldsymbol{\sigma}\cdot\boldsymbol{p}\end{bmatrix}\chi\right\}^{(i)}, \tag{B1}$$

$$\mathcal{N} = 1/\sqrt{(E'+m)(E+m)}, \tag{B2}$$

for a generic Dirac matrix Γ. We display results for nucleon 1 and those corresponding to nucleon 2 are obtained by making $\boldsymbol{q}_r \to -\boldsymbol{q}_r$ and $\boldsymbol{Q}_r \to \boldsymbol{Q}_r$. For the normalization, one has

$$\begin{aligned}\mathcal{N} = \big[&(m+\mathcal{E}/2)^2 + 4(m+\mathcal{E}/2)\boldsymbol{Q}_r\cdot\boldsymbol{Q}_\rho/\sqrt{3}\mathcal{E} \\ &+ [4(\boldsymbol{Q}_r\cdot\boldsymbol{Q}_\rho)^2 - (\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho)^2]/3\mathcal{E}^2\big]^{-1/2} \overset{\text{lga}}{\simeq} 1/(m+e),\end{aligned} \tag{B3}$$

where the last equality corresponds to the linear gradient approximation.

The OPEP, Eq. (23), is based on the function

$$\begin{aligned}[\bar{u}\gamma_5 u]^{(1)} &= \mathcal{N}\chi^\dagger\{[m+\mathcal{E}/2 + 2\boldsymbol{Q}_r\cdot\boldsymbol{Q}_\rho/\sqrt{3}\mathcal{E}]\boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{q}_r \\ &\quad - 2\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho\boldsymbol{\sigma}^{(1)}\cdot(\boldsymbol{Q}_r + \boldsymbol{Q}_\rho/\sqrt{3})/\sqrt{3}\mathcal{E}\}\chi \\ &\quad \times \overset{\text{lga}}{\simeq} [\bar{u}\gamma_5 u]^{(1)}_{\text{c.m.}} \\ &= \chi^\dagger[\boldsymbol{\sigma}^{(1)}\cdot\boldsymbol{q}_r]\chi.\end{aligned} \tag{B4}$$

The expression for the TPEP is given by Eq. (31) and employs the operators

$$\begin{aligned}[\bar{u}(\boldsymbol{p}')u(\boldsymbol{p})]^{(1)} &= \{\mathcal{N}\chi^\dagger[2m(m+\mathcal{E}/2) - 2(\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho)^2/3\mathcal{E}^2 \\ &\quad + \boldsymbol{q}_r^2/2 - i\boldsymbol{\sigma}\cdot\boldsymbol{q}_r\times(\boldsymbol{Q}_r + \boldsymbol{Q}_\rho/\sqrt{3})]\chi\}^{(1)} \\ &\overset{\text{lga}}{\simeq} [\bar{u}(\boldsymbol{p}')u(\boldsymbol{p})]^{(1)}_{\text{c.m.}} - \left\{\chi^\dagger\left[\frac{i}{(e+m)}\boldsymbol{\sigma}\cdot\boldsymbol{q}_r\right.\right. \\ &\quad \left.\left.\times\,\boldsymbol{Q}_\rho/\sqrt{3}\right]\chi\right\}^{(1)},\end{aligned} \tag{B5}$$

$$\begin{aligned}&\left[\frac{i}{2m}\bar{u}(\boldsymbol{p}')\sigma_{\mu 0}(p'-p)^\mu u(\boldsymbol{p})\right]^{(1)} \\ &= \{(\mathcal{N}/2m)\chi^\dagger[(m+\mathcal{E}/2)[\boldsymbol{q}_r^2 - 2i\boldsymbol{\sigma}\cdot\boldsymbol{q}_r(\boldsymbol{Q}_r + \boldsymbol{Q}_\rho/\sqrt{3})] \\ &\quad - 2(\boldsymbol{q}_r\cdot\boldsymbol{Q}_\rho)\boldsymbol{q}_r\cdot(\boldsymbol{Q}_r + \boldsymbol{Q}_\rho/\sqrt{3})/\sqrt{3}\mathcal{E}]\chi\}^{(1)} \\ &\overset{\text{lga}}{\simeq} \left[\frac{i}{2m}\bar{u}(\boldsymbol{p}')\sigma_{\mu 0}(p'-p)^\mu u(\boldsymbol{p})\right]^{(1)}_{\text{c.m.}} \\ &\quad - \left\{\chi^\dagger\left[\frac{i}{m}\boldsymbol{\sigma}\cdot\boldsymbol{q}_r\boldsymbol{Q}_\rho/\sqrt{3}\right]\chi\right\}^{(1)},\end{aligned} \tag{B6}$$

$$\left[\frac{i}{2m}\bar{u}(\boldsymbol{p}')\sigma_{\mu j}(p'-p)^{\mu}u(\boldsymbol{p})\right]^{(1)}$$
$$= \{\mathcal{N}\chi^{\dagger}[(m+\mathcal{E}/2)i\boldsymbol{\sigma}\times\boldsymbol{q}_r + \big[-\boldsymbol{q}_r^2 + 2i\boldsymbol{\sigma}\cdot\boldsymbol{q}_r\big(\boldsymbol{Q}_r$$
$$+\boldsymbol{Q}_{\rho}/\sqrt{3}\big)4(\boldsymbol{q}_r\cdot\boldsymbol{Q}_{\rho})^2/3\mathcal{E}^2\big]\big(\boldsymbol{Q}_r+\boldsymbol{Q}_{\rho}/\sqrt{3}\big)/2m$$
$$-(\boldsymbol{q}_r\cdot\boldsymbol{Q}_{\rho})\big[\boldsymbol{q}_r + 2i\boldsymbol{\sigma}\big(\boldsymbol{Q}_r+\boldsymbol{Q}_{\rho}/\sqrt{3}\big)\big]/\sqrt{3}\mathcal{E}]\chi\}^{(1)}$$
$$\overset{\text{lga}}{\simeq}\left[\frac{i}{2m}\bar{u}(\boldsymbol{p}')\sigma_{\mu j}(p'-p)^{\mu}u(\boldsymbol{p})\right]^{(1)}_{\text{c.m.}} - \left\{\chi^{\dagger}\frac{1}{2(e+m)}\right.$$
$$\left.\times\big[\boldsymbol{q}_r^2\boldsymbol{Q}_{\rho}/\sqrt{3}m + (\boldsymbol{q}_r\cdot\boldsymbol{Q}_{\rho})\boldsymbol{q}_r/\sqrt{3}e\big]_j\chi\right\}^{(1)}. \quad \text{(B7)}$$

These results allow one to write

$$\big\{[\bar{u}u]^{(1)}[\bar{u}u]^{(2)}\big\}\overset{\text{lga}}{\simeq}\{\cdots\}_{\text{c.m.}} - \left[\frac{4m}{(e+m)} + \frac{\boldsymbol{q}_r^2}{(e+m)^2}\right]\boldsymbol{\Omega}_D, \quad \text{(B8)}$$

$$\left\{-\frac{i}{2m}[\bar{u}u]^{(1)}[\bar{u}\sigma_{\mu\lambda}(p'-p)^{\mu}u]^{(2)} - (1\leftrightarrow 2)\right\}\frac{z^{\lambda}}{2m} \overset{\text{lga}}{\simeq}\{\cdots\}_{\text{c.m.}}\frac{z^{\lambda}}{2m}, \quad \text{(B9)}$$

$$\left\{-\frac{1}{4m^2}[\bar{u}\sigma_{\mu\lambda}(p'-p)^{\mu}u]^{(1)}[\bar{u}\sigma_{\nu\rho}(p'-p)^{\nu}u]^{(2)}\right\}g^{\lambda\rho} \overset{\text{lga}}{\simeq}\{\cdots\}_{\text{c.m.}}\,g^{\lambda\rho} - \frac{e\boldsymbol{q}_r^2}{m^2(e+m)}\boldsymbol{\Omega}_D, \quad \text{(B10)}$$

$$\left\{-\frac{1}{4m^2}[\bar{u}\sigma_{\mu\lambda}(p'-p)^{\mu}u]^{(1)}[\bar{u}\sigma_{\nu\rho}(p'-p)^{\nu}u]^{(2)}\right\}\frac{z^{\lambda}z^{\rho}}{4m^2} \overset{\text{lga}}{\simeq}\{\cdots\}_{\text{c.m.}}\frac{z^{\lambda}z^{\rho}}{4m^2}, \quad \text{(B11)}$$

where the functions $\{\cdots\}_{\text{c.m.}}$ are given by Eqs. (A32)–(A35) of Ref. [5] and the two-component spin operators $\boldsymbol{\Omega}$ are defined as

$$\boldsymbol{\Omega}_D = i(\boldsymbol{\sigma}^{(1)} - \boldsymbol{\sigma}^{(2)})\cdot\boldsymbol{q}_r\times\boldsymbol{Q}_{\rho}/2\sqrt{3}. \quad \text{(B12)}$$

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