Hadronic parity violation in few-body systems

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Abstract. Recent interest, both from experimental and theoretical point of view, on hadronic parity violation is reviewed, with particular emphasis on an effective theory description. After discussing the minimal form of the parity-violating $NN$ contact effective Lagrangian, we concentrate on the calculation of the neutron spin rotation in $\vec{n} - d$ scattering at zero energy. We find that this observable is sensitive to the long-range component of the parity-violating $NN$ potential due to the pion exchange, and that it is expected to be one order of magnitude larger than in $\vec{n} - p$ scattering.

1. Introduction

The understanding of weak interaction, while satisfactory at the level of quarks and leptons, is complicated in the hadronic sector by the non-perturbative character of QCD in this regime. At the nuclear level, weak interactions can only be studied by analyzing small effects induced by characteristic properties of this interaction, like flavour changing or parity violation.

There remain indeed many puzzles in the domain of hadronic weak interactions (HWI) \cite{1}, which are yet to be explained: among the most prominent ones are the origin of the $\Delta I = 1/2$ rule in the non-leptonic decays of kaons and hyperons, largely dominated by transition amplitudes with $\Delta I = 1/2$ \cite{2}, the difficulties in describing simultaneously the $S$- and $P$-waves of hyperon non-leptonic decays \cite{3}, and the anomalously large parity-violating asymmetries in the radiative hyperon decays \cite{4}. All these problems involve the strange quark and might be ascribed to a lack of convergence of the SU(3) chiral series, for several possible reasons (see e.g. Ref. \cite{5}). In order to isolate the genuine properties of HWI, without being much affected by the presence of dynamical strange quarks, it is therefore interesting to focus on the $\Delta S = 0$ component of HWI, by examining the hadronic parity violation (HPV). In this respect, few-nucleon systems provide a clean theoretical laboratory, due to the possibility of performing accurate \textit{ab initio} calculations. Improved experimental techniques allow now to measure on such systems the tiny effects of HPV, which are of the order of $G_F m_\pi^2 \sim O(10^{-7})$.

The main observables under experimental consideration are: i) the longitudinal analyzing power in $\vec{p}p$ scattering; ii) the asymmetry $A_\gamma$ of the emitted photon in the radiative capture $\vec{n}p \rightarrow d\gamma$ with respect to the initial neutron polarization, with a sensitivity of $10^{-8}$ expected at the NPDGamma experiment \cite{6}; iii) the asymmetry $A_t'$ in $\vec{n}d \rightarrow t\gamma$, proposed as a follow up...
of the latter experiment at the Spallation Neutron Source (SNS), Oak Ridge; \textit{iv}) neutron spin rotation in $\vec{np}$, $\vec{nd}$ and $\vec{n}\alpha$, also planned at the SNS.

In the following, we will concentrate on the calculation of the neutron spin rotation in $\vec{n} - d$ scattering at zero energy \cite{7}. In section 2 we discuss the two models for the PV $NN$ potential that we have used, with particular emphasis on the leading-order effective Lagrangian, exhibited in its minimal form \cite{8} (see also Ref. \cite{9}); the neutron spin rotation observable is briefly described in section 3, while section 4 contains our results.

2. Models of parity-violating $NN$ interaction

The traditional framework for the analysis of HPV has been, for almost three decades, the meson-exchange parity-violating $NN$ potential of Desplanques, Donoghue and Holstein \cite{10}, which incorporates $\pi$, $\rho$ and $\omega$ exchange. It is parametrized by 7 parity-violating meson-nucleon couplings, to be determined from experiment. These couplings have rather broad allowed ranges, and moreover, the consistency between different experiments is not entirely satisfactory \cite{11}. A more modern point of view has been advocated recently in Ref. \cite{12}, by recasting HPV in the language of effective field theory (EFT). This has the advantage of being a model-independent approach, based on the separation of the physical scales involved in the problem: subleading corrections can in principle be evaluated systematically in a $p/\Lambda_H$ expansion, $p$ being the typical momentum and $\Lambda_H$ the hadronic scale, characteristic of hadrons whose mass is not protected by chiral symmetry. It is then ideally suited for a situation in which experimental data are scarce. The effective theory is formulated in terms of nucleons and pions only, transforming non-linearly under the broken chiral symmetry of QCD \cite{13}. In addition to the (strong) parity-conserving couplings of chiral perturbation theory, one has to consider the induced PV couplings from weak interactions. Virtual $W$ and $Z$ exchanges induce (at the quark level) four-quark, current-current operators in the effective hamiltonian, which have definite transformation rules under the broken chiral symmetry of strong interactions. Hadronic operators, built in terms of pions and nucleons, are therefore introduced in the effective theory, with the same transformation properties \cite{14}, yielding a PV effective Lagrangian, whose infinite number of terms is ordered according to the power counting of chiral perturbation theory. The effect of heavier particles is encoded in the value of the low-energy coupling constants. At the leading order, the relevant Lagrangian is the sum of two pieces,

$$L^{PV} = L^{PV}_{\pi N} + L^{PV}_{NN},$$

(1)

the first one containing the interaction of pions with a single nucleon, the second one containing 2-nucleon contact vertices. At energies much smaller than the pion mass, even pions can be integrated out, and the pion mass be taken as the ultraviolet cutoff $\Lambda$. The “pionless effective field theory” ($\#EFT$) is then expressed in terms of nucleons only, with interactions described by the effective Lagrangian $L^{PV}_{NN}$. In this case however, the value of the low-energy constants will be changed, compared to Eq. (1), since in $\#EFT$ they include the effects of the integrated-out pions.

In the following we will focus our attention on the contact Lagrangian $L^{PV}_{NN}$. It has been written in Ref. \cite{12} as consisting of 12 operators,

$$L^{PV}_{NN} = \sum_{i=1}^{6} (C_i O_i + \tilde{C}_i \tilde{O}_i),$$

(2)
Table 1. Components of the PV NN potential corresponding to the DDH model and to the \#EFT.

| n | $c_n^{\text{DDH}}$ | $f_n^{\text{DDH}}(r)$ | $c_n^{\text{EFT}}$ | $f_n^{\text{EFT}}(r)$ | $O_{ij}^{(n)}$ |
|---|----------------|---------------------|----------------|-------------------|----------------|
| 1 | $+g_n \hbar^{1/2}$ | $f_x(r)$ | $2\mu^2 \chi C_6$ | $f_\mu(r)$ | $(\tau_i \times \tau_j)_z (\sigma_i + \sigma_j) \cdot \mathbf{X}_{ij,+}^{(1)}$ |
| 2 | $-\frac{g_n \hbar^{1/2}}{\alpha m}$ | $f_\rho(r)$ | 0 | 0 | $(\tau_i \cdot \tau_j) (\sigma_i - \sigma_j) \cdot \mathbf{X}_{ij,+}^{(2)}$ |
| 3 | $-\frac{g_n \hbar^{1/2}(1+\kappa_\rho)}{m} f_\rho(r)$ | 0 | 0 | $(\tau_i \cdot \tau_j) (\sigma_i \times \sigma_j) \cdot \mathbf{X}_{ij,-}^{(3)}$ |
| 4 | $-\frac{g_n \hbar^{1/2}}{2m}$ | $f_\rho(r)$ | $\frac{\mu^2}{\chi} (C_2 + C_4)$ | $f_\mu(r)$ | $(\tau_i + \tau_j)_z (\sigma_i - \sigma_j) \cdot \mathbf{X}_{ij,+}^{(4)}$ |
| 5 | $-\frac{g_n \hbar^{1/2}(1+\kappa_\rho)}{2m} f_\rho(r)$ | 0 | 0 | $(\tau_i + \tau_j)_z (\sigma_i \times \sigma_j) \cdot \mathbf{X}_{ij,-}^{(5)}$ |
| 6 | $-\frac{g_n \hbar^{1/2}}{2\sqrt{5} m}$ | $f_\rho(r)$ | $-\frac{2\mu^2}{\chi} C_5$ | $f_\mu(r)$ | $(3 \tau_i \cdot \tau_j)_z (\sigma_i - \sigma_j) \cdot \mathbf{X}_{ij,+}^{(6)}$ |
| 7 | $-\frac{g_n \hbar^{1/2}(1+\kappa_\rho)}{2\sqrt{5} m} f_\rho(r)$ | 0 | 0 | $(3 \tau_i \cdot \tau_j)_z (\sigma_i \times \sigma_j) \cdot \mathbf{X}_{ij,-}^{(7)}$ |
| 8 | $-\frac{g_n \hbar^{1/2}(1+\kappa_\omega)}{m} f_\omega(r)$ | $\frac{2\mu^2}{\chi} C_1$ | $f_\mu(r)$ | $(\sigma_i - \sigma_j) \cdot \mathbf{X}_{ij,+}^{(8)}$ |
| 9 | $-\frac{g_n \hbar^{1/2}(1+\kappa_\omega)}{m} f_\omega(r)$ | $\frac{2\mu^2}{\chi} C_1$ | $f_\mu(r)$ | $(\sigma_i \times \sigma_j) \cdot \mathbf{X}_{ij,-}^{(9)}$ |
| 10 | $-\frac{g_n \hbar^{1/2}}{2m} f_\omega(r)$ | 0 | 0 | $(\tau_i + \tau_j)_z (\sigma_i - \sigma_j) \cdot \mathbf{X}_{ij,+}^{(10)}$ |
| 11 | $-\frac{g_n \hbar^{1/2}(1+\kappa_\omega)}{m} f_\omega(r)$ | 0 | 0 | $(\tau_i + \tau_j)_z (\sigma_i \times \sigma_j) \cdot \mathbf{X}_{ij,-}^{(11)}$ |
| 12 | $-\frac{g_n \hbar^{1/2}}{2m} f_\rho(r)$ | 0 | 0 | $(\tau_i - \tau_j)_z (\sigma_i + \sigma_j) \cdot \mathbf{X}_{ij,+}^{(12)}$ |
| 13 | $-\frac{g_n \hbar^{1/2}}{2m} f_\rho(r)$ | 0 | 0 | $(\tau_i \cdot \tau_j)_z (\sigma_i + \sigma_j) \cdot \mathbf{X}_{ij,-}^{(13)}$ |

where

\begin{align}
O_1 &= \bar{\psi}\gamma^\mu \psi \bar{\gamma}_\mu \gamma_5 \psi, \\
O_2 &= \bar{\psi}\gamma^\mu \psi \bar{\gamma}_\tau \gamma_5 \psi, \\
O_3 &= \bar{\psi}\tau^a \gamma^\mu \psi \bar{\gamma}_\mu \gamma_5 \psi, \\
O_4 &= \bar{\psi}\gamma^\mu \psi \bar{\gamma}_\mu \gamma_5 \psi, \\
O_5 &= \mathcal{I}_{ab} \bar{\psi}\gamma^\mu \psi \bar{\gamma}_b \gamma_5 \psi, \\
O_6 &= i\epsilon^{abc} \bar{\psi}\tau^a \gamma^\mu \psi \bar{\gamma}_b \gamma_5 \psi,
\end{align}

the nucleon field $\psi$ is a doublet in isospin space and $\mathcal{I} = \text{diag}(1,1,-2)$. However, this set of operators is redundant: by using fields’ equations of motion and Fierz rearrangements, the following relations may be shown to hold [8],

\begin{align}
O_1 &= O_3, \\
O_2 &= O_4 = 2O_6, \\
O_2 + O_4 &= m(O_2 + O_4), \\
O_2 - O_4 &= -2mO_6 - O_6,
\end{align}

so that the number of independent operators can be reduced to 6. One can then make use of the fact that the operators $O_6$ and $\tilde{O}_6$ give rise, in the leading order of the non-relativistic reduction, to the same structure, as already observed in Ref. [12]. Therefore, the number of independent operators can be further reduced to 5 up to order $O(Q)$, and the minimal PV
two-nucleon non-relativistic contact Lagrangian may be taken to assume the form

$$\mathcal{L}_{NN}^{PV} = \frac{1}{\Lambda^2_N} \left\{ C_1 (N\uparrow \bar{\sigma} N \cdot N\uparrow i \nabla N - N\uparrow N N\uparrow i \nabla \cdot \bar{\sigma} N) \\
- \bar{C}_1 \epsilon_{ijk} N\uparrow \sigma^i N \nabla j (N\uparrow \sigma^k N) \\
- (C_2 + C_4) \epsilon_{ijk} [N\uparrow \tau_3 \sigma^i N \nabla j (N\uparrow \sigma^k N) + N\uparrow \sigma^i N \nabla j (N\uparrow \tau_3 \sigma^k N)] \\
- \bar{C}_3 \epsilon_{ijk} \tau^i \sigma^j \nabla (N\uparrow \tau^b \sigma^k N) \\
+ C_6 \epsilon^{abk} \nabla (N\uparrow \tau^a N) \cdot N\uparrow \tau^b \bar{\sigma} N \right\}, \quad (5)$$

where the notations for the coupling constants have been chosen so as to conform to Ref. [12]. It is worth noting that the reduction of the number of independent operators down to five has no practical consequences at the present stage of phenomenological analyses since only five combinations of low-energy constants are relevant at low energies. This was already noticed in Refs. [12, 1] on the basis of the observation that, since only $S$- and $P$-wave amplitudes are important in this regime, several operators give rise to identical matrix elements (see also the discussion in Ref. [15]). Nevertheless, in view of a description of nuclear parity violation with the chiral effective theory, it is important to use a truly minimal set of operators.

The resulting PV $NN$ potential consists of 5 different components, as opposed to the 13 components of the DDH model. It is convenient to express both potential in a unified notation,

$$v_{ij}^\alpha = \sum_{n=1}^{13} c_n^\alpha O_{ij}^{(n)}; \quad \alpha = \text{DDH or EFT.} \quad (6)$$

The parameters $c_n^\text{EFT}$ are in one-to-one correspondence with the 5 low-energy constants of the effective Lagrangian, while the $c_n^\text{DDH}$ are expressed in terms of 7 unknown PV meson-nucleon couplings ($h_1^\pi, h_0^{\pi,1,2}, h_1^\rho$, and $h_0^\rho$) as displayed in table 1 together with the corresponding space-spin-isospin structures. The vector spatial operators are defined as

$$X_{ij,+}^{(n)} \equiv [p_{ij}, f_n(r_{ij})]_+, \quad X_{ij,-}^{(n)} \equiv i[p_{ij}, f_n(r_{ij})]_-; \quad (7)$$

where $[...]_-$ $([...]_+)$ stands for the commutator (anticommutator) , and $p_{ij} = (p_i - p_j)/2$ for the relative momentum of nucleons $i$ and $j$. In the DDH model the radial functions $f_x(r)$ (with $x = \pi, \rho, \omega$) are Yukawa functions modified by the inclusion of monopole form factors,

$$f_x(r) = \frac{1}{4\pi r} \left\{ e^{-m_x r} - e^{-A_x r} \left[ 1 + \frac{A_x r}{2} \left( 1 - \frac{m_x^2}{A_x^2} \right) \right] \right\}, \quad (8)$$

while in the EFT there is a single cutoff function $f_\mu(r)$, taken as a Yukawa function of range $1/\mu$,

$$f_\mu(r) = \frac{1}{4\pi r} e^{-\mu r}. \quad (9)$$

3. **Neutron spin rotation in $\bar{n}d$ scattering**

Due to the presence of PV interactions (as exemplified in figure 3 by a $\sigma \cdot p$ interaction), neutrons polarized in the transverse direction (e.g. along $x$, with $|x| \propto |+\rangle + |-\rangle$), while traveling along $z$, accumulate different phases along $|+\rangle$ and $|-\rangle$, and the spin $\sigma$ is then viewed to rotate by an angle $\phi$ proportional to the forward scattering amplitude with the deuteron,

$$\frac{1}{\rho} \frac{d\phi}{dz} = \frac{1}{3v_{rel}} \sum_{m_n, m_d} \epsilon_{m_n} \langle p\bar{z}; m_n, m_d | v^{PV} | p\bar{z}; m_n, m_d \rangle_{\text{out}}. \quad (10)$$
Here $\rho$ is the deuteron density, $v_{PV}$ is the PV $NN$ potential, $p = p\hat{z}$ is the relative momentum, taken along the spin-quantization axis, $v_{rel} = p/\mu$ is the relative velocity, $\mu$ being the $n - d$ reduced mass, and $m_d$ ($m_n$) the deuteron (neutron) spin projection. The phase factor $\epsilon_{m_n} = \pm 1$ depending on whether $m_n = \pm 1/2$. The scattering states are calculated from the AV18 potential [16], supplemented by the three-nucleon interaction Urbana IX [17], by means of the hyperspherical harmonics method [18]. They are expanded in partial waves as [19]

$$|p; \tilde{m}_n, \tilde{m}_d \rangle_{\text{in/out}} = 4\pi \sum_{LSJ} \epsilon_{m_n} \langle \frac{1}{2} m_n, 1 m_d | SS_z \rangle_{\text{in}} \langle SS_z; L0| J J_z \rangle_{\text{in}} Y^*_L(\hat{z}) |p; LS; J, J_z \rangle_{\text{in/out}}. \quad (11)$$

For very low-energy neutrons, only the $S$ to $P$ transitions are relevant, and the spin rotation observable is written, using the time-reversal transformation properties of the ingoing and outgoing states, as

$$\frac{1}{\rho} \frac{d\phi}{dz} = \frac{8\pi}{\sqrt{3} v_{rel} \sum_{m_n} \sum_{S J} \epsilon_{m_n} \langle \frac{1}{2} m_n, 1 m_d | JJ_z \rangle_{\text{in}} \langle \frac{1}{2} m_n, 1 m_d | SJ_z \rangle_{\text{in}} \langle SJ_z, 10 | JJ_z \rangle_{\text{in}} \langle SS_z; L0| J J_z \rangle_{\text{in}} \langle p; 1 S; J, J_z | v_{PV} | p; 0 J; J, J_z \rangle_{\text{out}}.$$  

$$\text{Im} \left[ \sum_{m_n} \sum_{S J} \epsilon_{m_n} \langle \frac{1}{2} m_n, 1 m_d | JJ_z \rangle_{\text{in}} \langle \frac{1}{2} m_n, 1 m_d | SJ_z \rangle_{\text{in}} \langle SJ_z, 10 | JJ_z \rangle_{\text{in}} \langle SS_z; L0| J J_z \rangle_{\text{in}} \langle p; 1 S; J, J_z | v_{PV} | p; 0 J; J, J_z \rangle_{\text{out}} \right]. \quad (12)$$

The PV potential connects the doublet $S$-wave state $^2S_{1/2}$ to both the doublet $^2P_{1/2}$ and quartet $^4P_{1/2}$ $P$-wave states, and similarly for $J = 3/2$.

### 4. Results

In the calculation of the spin rotation observable, we do not consider any definite value for the coupling constants. Instead, we express the result as a sum

$$\frac{1}{\rho} \frac{d\phi}{d\rho} = \sum_{n=1}^{13} c_n^\rho I_n^\rho \quad (13)$$

where the contribution from the single components of the PV $NN$ potential are computed separately. Details of the calculation can be found in Ref. [7]. The numerical procedure has also been tested for the case of gaussian-like wave functions, for which the calculations can be done almost analytically. We have selected two sets of ultraviolet cutoffs, as indicated in table 2, for both the DDH model and the $\pi$EFT, in order to test the cutoff dependence.

The results for the individual contributions $I_n$ as defined in Eq. (13) are displayed in table 3 for the DDH model with and without cutoff. The effect of the inclusion of the three-nucleon interaction is also shown. The coefficients $I_{6,7}^{\text{DDH}}$ vanish because of isospin selection rules, while the results for $I_{13}^{\text{DDH}}$ are not reported, because the coupling constant $h_3^{J_1}$ that multiplies this contribution has been estimated to be much smaller [20] than the “best” values estimated for...
integrations are not shown, but are typically at the 1-2% level.

Table 2. Different sets of values, in GeV units, for the short-range cutoffs $\Lambda_x$ and $\mu$ used in the DDH model and in the $^\pi$EFT. The masses $m_x$, $m_\rho$, and $m_\omega$ are taken respectively as 0.138, 0.771, and 0.783 in units of GeV. The row with $\Lambda_x = \infty$ corresponds to point-like couplings.

|       | $\Lambda_\pi$ | $\Lambda_\rho$ | $\Lambda_\omega$ | $\mu$ |
|-------|---------------|---------------|-----------------|-------|
| DDH-I | 1.72          | 1.31          | 1.50            | EFT-I | 0.138 |
| DDH-II| $\infty$      | $\infty$      | 1.50            |       | EFT-II| 1.0   |

Table 3. The coefficients $l_n^{DDH}$ (in fm) in Eq. (13) corresponding to the DDH PV potential in combination with the AV18 and AV18/UIX strong-interaction potentials for the two different sets of cutoff parameters given in table 2. The statistical errors associated with the Monte Carlo integrations are not shown, but are typically at the 1-2% level.

| $n$   | AV18   | AV18/UIX | AV18   | AV18/UIX |
|-------|--------|----------|--------|----------|
| 1     | 0.256E + 03 | 0.270E + 03 | 0.257E + 03 | 0.274E + 03 |
| 2     | -0.444E + 01 | -0.691E + 01 | -0.719E + 01 | -0.118E + 02 |
| 3     | 0.444E + 01  | 0.401E + 01  | 0.732E + 01  | 0.761E + 01  |
| 4     | -0.231E + 01 | -0.881E + 00 | -0.332E + 01 | -0.148E + 01 |
| 5     | -0.247E + 01 | -0.122E + 01 | -0.387E + 01 | -0.234E + 01 |
| 6     | 0.420E + 01  | 0.362E + 01  | 0.543E + 01  | 0.516E + 01  |
| 7     | 0.111E + 01  | -0.117E + 00 | 0.136E + 01  | -0.189E + 00 |
| 8     | -0.253E + 01 | -0.991E + 00 | -0.314E + 01 | -0.141E + 01 |
| 9     | -0.280E + 01 | -0.144E + 01 | -0.369E + 01 | -0.225E + 01 |
| 10    | 0.316E + 01  | 0.339E + 01  | 0.455E + 01  | 0.546E + 01  |

the other PV meson-nucleon couplings in the DDH model [10]. As it appears, the matrix element of the pion-range component of the DDH potential is larger by two orders of magnitude than the components of shorter range induced by $\rho$ and $\omega$ exchanges. Moreover, the long-range component is quite independent of the cutoff, and it is not much affected by the inclusion of the three-nucleon force. On the contrary, the other individual contributions due to vector meson exchanges exhibit considerable dependence on the strong-interaction input and on the cutoff. This model dependence, though, has little impact on the neutron spin rotation observable. This is also shown in table 4 where the contributions from the long- and short-range components to this observable are indicated separately, for the cutoff version of the DDH model (DDH-I): the long-range component accounts for almost the 90% of the entire contribution. For the sake of comparison, two different set of strong- and weak-interaction coupling constants are used, corresponding to those reported in tables I and II of Ref. [21]: the “best” values are taken from Ref. [10], while the “adjusted” set is the one found in Ref. [22] to reproduce precise measurements of the longitudinal asymmetry in $\vec{n} - p$ scattering. The PV pion-nucleon coupling $h_\pi^\perp$ is the same for both sets. Interestingly enough, the neutron spin rotation in deuterium is predicted to be one order of magnitude larger than in the case of $\vec{n} - p$ [21], which should make the measurement easier, in principle.

In table 5 we show the individual contributions $I_n^{EFT}$ from the PV $NN$ potential derived in $^\pi$EFT, corresponding to $n = 1, 4, 8$ and 9. The term $n = 6$ is isotensor, and has vanishing matrix
Table 4. Spin rotation for $\vec{n} - d$ scattering in units of $10^{-7}$ rad/cm at zero energy, obtained for the DDH-I model assuming a liquid deuterium density of $\rho = 0.4 \times 10^{23}$ atoms cm$^{-3}$. The columns labeled DDH-best (DDH-adj) list the spin rotation predicted using the “best” (“adjusted”) values for the weak-interaction coupling constants, as reported in table II (I) of Ref. [21]. The row labeled “$\pi$” (“$\rho$-$\omega$”) lists the results obtained by including only the pion ($\rho$ and $\omega$ mesons), while that labeled “TOT” gives the total contributions.

|            | AV18, AV18/UIX |
|------------|----------------|
| DDH - best | DDH - adj      |
| $\pi$      | 0.457          |
| $\rho$-$\omega$ | 0.063      |
| TOT        | 0.520          |

Table 5. Same as in table 3 but for the pionless EFT PV potential. Note that there are no potential components with $O^{(n)}_{ij}$ with $n=2, 3, 5, 7, 10$ and 11, and that the coefficient $I_{n=6}^{EFT}$ vanishes because of isospin selection rules.

|            | EFT-I          | EFT-II         |
|------------|----------------|----------------|
|            | AV18/UIX       | AV18/UIX       |
| $n$        | AV18           | AV18           |
| 1          | $0.257E + 03$  | $0.274E + 03$  |
| 4          | $-0.154E + 03$ | $-0.508E + 02$ |
| 8          | $0.260E + 03$  | $0.189E + 03$  |
| 9          | $0.244E + 00$  | $-0.359E + 02$ |

elements. Contrary to the DDH model, all contributions are of the same order of magnitude and correspond to radial functions with a same range $\mu$. Of course, they depend significantly on the value of the cutoff, chosen as $\mu = m_{\pi}$ (EFT-I), as appropriate in $\pi$EFT, or as $\mu = 1$ GeV (EFT-II), as appropriate in the pionful version of the EFT. In the latter case, the leading-order component of the PV $NN$ potential has the same form as the pion-exchange term in the DDH model. Also notice that, while the $n = 1$ contribution is weakly dependent on the strong-interaction input, the same is not true for the other coefficients. However, we did not attempt to establish predictions for the spin rotation observable in $\pi$EFT, since at the present stage the value of the low-energy constants is essentially unknown.

5. Conclusions
The EFT approach, in conjunction with precise ab initio techniques for handling the few-body problem, allows to investigate on a firm theoretical basis the hadronic parity violation. Future and on-going experimental work will hopefully allow to pin down the PV low-energy constants, and shed light on long-standing problems of non-perturbative QCD. We have constructed the relevant leading-order $NN$ contact Lagrangian in its minimal form, which is an essential ingredient in both the pionful and pionless versions of the EFT. We have computed the neutron spin rotation in liquid deuterium at zero energy, and parametrized it in terms of the low-energy constants or the parameters of the DDH model. Within the DDH model, this observable is dominated by the contribution of the long-range part of the PV potential.
due to pion exchange, and this contribution is almost model-independent. The predicted value, \( d\phi/dz \sim 0.5 \times 10^{-7} \text{ rad/cm}^{-1} \), is an order of magnitude larger than expected in \( \vec{n} - p \) scattering [21]. An experimental investigation of this observable could therefore provide a further constraint, complementary to that coming from measurements of the photon asymmetry in \( \vec{n} - p \) radiative capture, on the strength of the long-range component of the hadronic weak interaction.

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