Two parity violating asymmetries from $np \to d\gamma$ in pionless effective field theories

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Abstract

We consider parity-violating observables from the processes $\bar{n}p \to d\gamma$ and $np \to d^\gamma$. We perform calculations using pionless effective field theory both with and without explicit dibaryon fields. After combining these results with ones we have already obtained on parity-violating asymmetries in $\bar{N}N$ scattering, experimental input would in principle allow the extraction of all five parameters occurring at leading order in the parity-violating Lagrangian.

Keywords: Parity violation, effective field theory

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I. INTRODUCTION

Low energy hadronic parity violation is of current interest both because of recent theoretical developments and because of the many and varied experiments underway or proposed to study the problem \[1, 2, 3, 4, 5, 6\]. In particular, parity violation in the two nucleon sector remains an open problem. In a previous paper \[7\] we addressed the parity violating asymmetries in $\bar{N}N$ scattering. Here we consider processes involving photons, in particular polarized and unpolarized radiative neutron capture on protons. Experiments to measure asymmetries in $np \to d\gamma$ are difficult, but are important enough that the experimental community continues to push for better limits. With new results expected in the next several years, it is timely to revisit the theoretical problem.

The parity-violating (PV) component of hadronic interactions is caused by weak interactions of quarks contained in the hadrons. Because of the relative strength of the weak interaction, its manifestation is highly suppressed compared to the strong interaction. Therefore, we consider observables that would be zero without the presence of parity-violating effects. The reaction $np \to d\gamma$, with suitable polarizations of the neutron or photon, allows access to two different asymmetries: the photon asymmetry $A_\gamma$ in polarized neutron capture, $\bar{n}p \to d\gamma$, and the circular photon polarization $P_\gamma$ in unpolarized capture, $np \to d\gamma$. In fact, due to the difficulty of measuring final state photon circular polarization, the inverse reaction of deuteron photo-disintegration $\gamma d \to np$ may be more experimentally feasible. The asymmetry from this inverse reaction is equal to $P_\gamma$ for exactly reversed kinematics.

Weak interactions are well understood in the context of the standard model, but we are interested here in weak manifestations in hadrons at energies where QCD is not perturbative. Therefore, we turn to effective field theories (EFTs) that allow for a perturbative treatment in quantities other than the strong coupling constant. While it is true that investigating hadronic parity violation necessarily involves complications from nonperturbative QCD, this fact can be used as an opportunity for probing nonperturbative QCD phenomena in hadrons. With nucleons, photons, and the deuteron as physical degrees of freedom we form the set of leading-order operators that obey the symmetries of QCD, but allow for parity violation. Because the processes we are interested in occur at energies well below where the pion is dynamical, we use an EFT in which the pion is integrated out, EFT($\not{\pi}$), and its physics is encoded in low energy constants (LECs).

Traditionally, hadronic parity violation has been studied using either the so-called Danilov amplitudes \[8\] or one-boson-exchange models \[9, 10, 11\]. While the one-boson-exchange models, in particular the one of Ref. \[11\], have been the standard for analyzing experiments, some possible inconsistencies have emerged (see e.g. Fig. 5 in Ref. \[12\]). The effective field theory treatment of parity-violating hadronic interactions as performed in this paper allows for a systematic and model-independent study of few-body low-energy phenomena. As can be seen from the Lagrangian in the following section, the EFT($\not{\pi}$) approach is more closely related to the Danilov amplitudes than the boson-exchange models. The use of EFTs to study parity violation goes back more than a decade (see e.g. \[13, 14, 15\]), with a comprehensive formulation of both pionless and pionful theories given in Ref. \[16\]. The asymmetry from $\bar{n}p \to d\gamma$ has been calculated previously in an EFT that included pions \[14\] and using the dibaryon formalism \[17\].

In the next section we reiterate the five independent PV operators that occur at leading order in EFT($\not{\pi}$). One contributes to the $\bar{n}p \to d\gamma$ process and three to $np \to d\gamma$. Other
linear combinations are accessible through asymmetries in $\bar{N}N$ scattering. However, the power counting in the $NN$ system is still an open question for some observables. It is clear that the $NN$ scattering lengths $a$ are anomalously large, and that it is necessary to resum a polynomial series in $a$. It is less clear how to treat the $NN$ effective range $r$. For some processes, particularly those involving the deuteron, a much improved description of data is found by treating $r$ as large, or resumming the series in $r$ (see Ref. [18], motivated in part by results in Ref. [19]). But that may not be true for all processes. Therefore, in Sec. III, we present calculations of $A_\gamma$ and $P_\gamma$ both in the non-dibaryon formalism, where $r$ is considered to be of "natural" size and terms involving $r$ are treated as higher order, as well as in the dibaryon formalism, where $r$ is treated as anomalously large and the resummed nucleon bubbles are encoded into a dynamical "dibaryon" field. Stating our results in both languages will provide the flexibility of developing a common language with other calculations done with more nucleons and/or using the non-dibaryon formalism.¹

Due to the difficulties inherent in measuring parity violation in the $NN$ system, information on $A_\gamma$ and $P_\gamma$ is sparse. Currently the asymmetry $A_\gamma$ from $\bar{n}p \rightarrow d\gamma$ is consistent with zero [21, 22], but an ongoing experimental effort is expected to improve the current value [5, 23]. Results for the photon polarization $P_\gamma$ from $np \rightarrow d\gamma$ are also consistent with zero [24]. As discussed above, the asymmetry from the inverse reaction is equal to $P_\gamma$ (for suitable kinematics) and, while current results are again consistent with zero [22, 25], there has been recent interest in performing this measurement [1, 2].

II. LAGRANGIANS

In EFT($\pi\not\pi$) nucleons interact through contact interactions. The leading order operators contain the minimum number of necessary derivatives. The parity-conserving (PC) Lagrangian is given by
\[
\mathcal{L}_{PC} = N^\dagger (iD_0 + \frac{\vec{D}^2}{2M})N + \frac{e}{2M} N^\dagger (\kappa_0 + \tau_3 \kappa_1) \sigma \cdot B N
- C^{(1S_0)}_0 (N^T P^{(1S_0)}_a N)^\dagger (N^T P^{(1S_0)}_a N) - C^{(3S_1)}_0 (N^T P^{(3S_1)}_i N)^\dagger (N^T P^{(3S_1)}_i N) + \ldots ,
\]
(1)

with the normalized projection operators [26]
\[
P^{(1S_0)}_a = \frac{1}{\sqrt{8}} \tau_2 \tau_a \sigma_2 , \quad P^{(3S_1)}_i = \frac{1}{\sqrt{8}} \tau_2 \sigma_2 \sigma_i .
\]
(2)
The $\sigma_i$ and $\tau_a$ are SU(2) Pauli matrices in spin and isospin space, respectively, $D_\mu N$ is the nucleon covariant derivative,
\[
D_\mu N = \partial_\mu N + ie \frac{1 + \tau_3}{2} A_\mu N ,
\]
(3)
and $\kappa_0$ and $\kappa_1$ are the isoscalar and isovector nucleon magnetic moments. Eq. (1) shows only the leading-order interaction terms. In the power counting for EFT($\pi\not\pi$) contributions

¹ While writing up this work we became aware of contemporary results on $P_\gamma$ in the dibaryon formalism from Shin, Ando, and Hyun [20].
from other terms are suppressed by powers of \(Q/m_\pi\), where \(Q \sim p \sim 1/a\); \(p\) is the relative \(NN\) momentum \([27, 28, 29]\) and \(a\) is the scattering length in either the \(^1S_0\) or \(^3S_1\) channel. Using the power-divergence subtraction (PDS) scheme \([28, 30]\) to renormalize loop diagrams the low-energy constants (LECs) are given by

\[
C_0^{(1S_0)} = \frac{4\pi}{M} \frac{1}{a^{(1S_0)} - \mu}, \tag{4}
\]

\[
C_0^{(3S_1)} = \frac{4\pi}{M} \frac{1}{a^{(3S_1)} - \mu}. \tag{5}
\]

Here, \(a^{(1S_0)}/a^{(3S_1)}\) are the scattering lengths in the \(^1S_0\) and \(^3S_1\) channel, respectively, and \(\mu\) is the subtraction point.

Since the following calculations contain a deuteron in the final state one has to choose an interpolating field for the deuteron. We follow Ref. \([7]\) and use \(D_i = N^iP_i^{(3S_1)}N\).

For the leading-order PV Lagrangian we use the partial wave notation of Ref. \([7]\):

\[
\mathcal{L}_{\text{PV}} = - \left[ C^{(3S_1-1P_1)} (N^T\sigma_2 \bar{\tau}_2 N)^\dagger \left( N^T\sigma_2 \bar{\tau}_2 i \bar{D} \right) N \right] + C^{(1S_0-3P_0)} (N^T\sigma_2 \bar{\tau}_2 N)^\dagger \left( N^T\sigma_2 \bar{\tau}_2 i \bar{D} \right) \tag{6}
\]

where \(a \mathcal{O} \bar{D} b = a \mathcal{O} \bar{D} b - (\bar{D}a)\mathcal{O}b\) with \(\mathcal{O}\) some spin-isospin-operator, and

\[
\mathcal{I} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix}.
\]

As shown in Ref. \([7]\), this form of the Lagrangian is equivalent to the one given in Ref. \([31]\). (Note, however, the different placement of the gauged derivative from the ungauged derivatives in the operators given in Ref. \([7]\). While the ungauged derivative is unaffected by isospin Pauli matrices, it is important to maintain consistent placement for the gauged derivative.)

The terms in Eqs. (5) and (6) are considered to be of leading order if we assume that the effective ranges \(r\) in the \(^1S_0\) and \(^3S_1\) channel are “natural”, i.e. terms like \(r/a\) and \(rp\) are numerically suppressed. In an alternative power counting the effective ranges are considered large and have to be resummed to all orders. This is most conveniently achieved by use of dynamical dibaryon fields \([18, 32, 33]\). The PC dibaryon Lagrangian is given by \([18]\)

\[
\mathcal{L}_{\text{PC}}^d = N^i(iD_0 + \bar{D}^2/2M)N - t_1^i \left( i\partial_0 + \bar{D}^2/4M - \Delta^{(3S_1)} \right) t_i - g^{(3S_1)} \left[ t_i^\dagger N^TP_i^{(3S_1)}N + \text{h.c.} \right]
\]

\[
- s_1^\dagger \left( i\partial_0 + \bar{D}^2/4M - \Delta^{(1S_0)} \right) s_1 - g^{(1S_0)} \left[ s_1^\dagger N^TP_a^{(1S_0)}N + \text{h.c.} \right]. \tag{7}
\]
where \( t_i \) and \( s_a \) are the dibaryon fields in the \( ^3S_1 \) and \( ^1S_0 \) channel, respectively. The couplings can be determined either by integrating out the dibaryon fields or by reproducing the effective range expansion of the \( NN \) scattering amplitude. This leads to (suppressing channel subscripts)

\[
g^2 = \frac{8\pi}{M^2 r}, \quad \Delta = \frac{2}{M r} \left( \frac{1}{a} - \mu \right). \tag{8}\]

Note that these procedures only fix the magnitude of \( g \) and not its sign. Since the dibaryon-\( NN \) couplings are of leading order, insertions of nucleon loops in the dibaryon propagator are not suppressed. This means that the leading-order dibaryon propagator gets dressed by an infinite series of nucleon loop insertions (e.g., [18]). The dressed propagator is given by

\[
S_d(E) = \frac{4\pi}{M g^2 \mu + \frac{4\pi}{M g^2} \Delta - \frac{4\pi}{M g^2} E + i\sqrt{M E}}, \tag{9}\]

where we have again dropped the partial wave superscripts. For the dibaryon calculations the choice of the deuteron interpolating field is simply the dibaryon field \( t_i \) [18].

Parts of the PV dibaryon Lagrangian were given in Refs. [17, 34]. The complete Lagrangian is given by

\[
\mathcal{L}^d_{PV} = - \left[ g^{(3S_1-1P_1)} t_i^\dagger \left( N^T \sigma_2 \tau_2 i \bar{D}_i N \right) + g^{(1S_0-3P_0)} s_a^\dagger \left( N^T \sigma_2 \hat{\sigma} \cdot \tau_2 \tau_i D N \right) + g^{(1S_0-3P_0)} \epsilon^{3ab} (s_a)^\dagger \left( N^T \sigma_2 \hat{\sigma} \cdot \tau_2 \tau^b D N \right) + g^{(1S_0-3P_0)} \hat{T}^{ab} (s_a)^\dagger \left( N^T \sigma_2 \hat{\sigma} \cdot \tau_2 \tau^b i \bar{D} j N \right) + g^{(3S_1-3P_1)} \epsilon^{ijk} (t^i)^\dagger \left( N^T \sigma_2 \sigma^k \tau_2 \tau_3 \bar{D} j N \right) \right] + h.c. \tag{10}\]

By performing the path integral over the dibaryon fields in the action we can relate the couplings in the two formalisms:

\[
g^{(X-Y)} = \sqrt{8 \Delta^{(X)}} g^{(X)} C^{(X-Y)} = \sqrt{\frac{\pi}{r(X)}} \frac{8}{M} C^{(X-Y)} \tag{11}\]

For example,

\[
g^{(1S_0-3P_0)}_{(\Delta I=0)} = \sqrt{8 \Delta^{(1S_0)}} g^{(1S_0)} C^{(1S_0-3P_0)}_{(\Delta I=0)} \tag{12}\]

### III. RESULTS

The invariant amplitude for \( np \rightarrow d\gamma \) can be parameterized as [14]

\[
\mathcal{M} = e_X N^T \tau_2 \sigma_2 \left[ \sigma \cdot q \ e_d^* \cdot e_d^* - \sigma \cdot e_d^* \cdot q \cdot e_d^* \right] N + i e_Y e^{ijk} e_d^{*j} q^{*i} e_d^* \left( N^T \tau_2 \sigma_2 N \right) + i e_W e^{ijk} e_d^{*j} e_d^* \left( N^T \tau_3 \sigma_2 N \right) + e V e_d^* \cdot e_d^* \left( N^T \tau_2 \sigma_2 N \right) + \ldots \tag{13}\]
FIG. 1: Leading-order diagrams contributing to the parity-conserving amplitude $Y$. Solid lines denote nucleons, wavy lines denote photons. The large solid circle stands for the resummed $NN$ scattering, the gray oval for the deuteron interpolating field, and the small open circle for a coupling to the nucleon magnetic moment.

where the ellipsis stands for terms that are not needed in our calculation. $\epsilon_d$ and $\epsilon_\gamma$ are the polarization vectors of the deuteron and photon, respectively, $q$ is the outgoing photon momentum and $e > 0$. The amplitudes $X$ and $Y$ are parity-conserving, while $V$ and $W$ are parity-violating. For $Y$ and $V$ the initial $NN$ state is in a relative $^1S_0$ wave, while for $X$ and $W$ it is in a $^3S_1$ wave. At leading order $Y$ and $W$ contribute to the photon asymmetry in $\bar{n}p \to d\gamma$, while the circular polarization in $np \to d\gamma$ stems from interference between $Y$ and $V$. Below we discuss the two processes.²

A. Photon asymmetry in $\bar{n}p \to d\gamma$

The photon asymmetry $A_\gamma$ for $\bar{n}p \to d\gamma$ at threshold is defined by

$$\frac{1}{\Gamma} \frac{d\Gamma}{d\cos \theta} = 1 + A_\gamma \cos \theta,$$

with $\Gamma$ the $np \to d\gamma$ width and $\theta$ the angle between the neutron polarization and the outgoing photon momentum. The polarization of the neutron leads to interference between the PC amplitudes $X$ and $Y$ and the PV amplitude $W$. At leading order $X = 0$ and the asymmetry $A_\gamma$ is given in terms of the amplitudes by

$$A_\gamma = -2 \frac{M}{\gamma^2} \frac{\text{Re}[Y^* W]}{|Y|^2},$$

where $\gamma = \sqrt{MB}$ is the deuteron momentum with $B$ the deuteron binding energy.

The leading-order diagrams contributing to the PC amplitude $Y$ are shown in Fig. 1, yielding

$$Y = -\frac{2}{M} \sqrt{\frac{\pi}{\gamma^3}} \kappa_1 \left(1 - \gamma a(1S_0)\right).$$

² Some of these results have been previously presented by the authors ³⁵. However, while writing up this work we became aware of a recent result from Shin, Ando, and Hyun ²⁰ on $np \to d\gamma$. 
FIG. 2: Leading-order diagrams contributing to the parity-violating amplitudes $W$ and $V$. Solid lines denote nucleons, wavy lines denote photons. The large solid circle stands for the resummed $NN$ scattering, the gray oval for the deuteron interpolating field. The solid black square is the PV operator. Photons are minimally coupled.

The diagrams shown in Fig. 2 contribute to the PV amplitude $W$ for the initial $np$ state in a $^3S_1$ partial wave. Using $1/a(^3S_1) = \gamma$ (valid at this order), the result for $W$ is

$$W = \frac{8}{3} \sqrt{\frac{\gamma}{\pi}} (\gamma - \mu) C(^3S_1-^3P_1),$$

which, using Eq. (5) and again $1/a(^3S_1) = \gamma$, can be written as

$$W = \frac{32\pi}{3M} \sqrt{\frac{\gamma}{\pi}} \frac{C(^3S_1-^3P_1)}{C_0(^3S_1)}.$$  

This gives the asymmetry

$$A_\gamma = \frac{32}{3} \frac{M}{\kappa_1 (1 - \gamma a(^3S_0))} \frac{C(^3S_1-^3P_1)}{C_0(^3S_1)}.$$  

Note the appearance of the ratio $\frac{C(^3S_1-^3P_1)}{C_0(^3S_1)}$. $A_\gamma$ is a physical quantity and must be independent of the subtraction point $\mu$. $C_0(^3S_1)$ has the $\mu$ dependence shown in Eq. (5), so the $\mu$ dependence of $C(^3S_1-^3P_1)$ must have the same form

$$C(^3S_1-^3P_1) \sim \frac{1}{a(^3S_1) - \mu}.$$  

1
FIG. 3: Leading-order diagrams contributing to the parity-conserving amplitude $Y$ in the dibaryon formalism. Solid lines denote nucleons, thick solid lines denote dressed dibaryons and wavy lines denote photons. The gray oval stands for the deuteron interpolating field, and the small open circle for a coupling to the nucleon magnetic moment.

This echoes the results discussed in Ref. [7].

The diagrams that need to be considered when using the dibaryon Lagrangians of Eqs. (7) and (10) are shown in Figs. 3 and 4. In this approach the result for the PC amplitude $Y$ is given by

$$Y^d = \frac{2}{M} \sqrt{\frac{\pi}{3}} \frac{1}{\sqrt{1 - \gamma r^{(3S_1)}}} \kappa_1 \left( 1 - \gamma a^{(1S_0)} \right),$$

(21)

which, when expanded in $r^{(3S_1)}$, reproduces the result of Eq. (16) up to a factor of $-1$. Since the amplitude itself is not an observable this sign difference is of no significance and could be absorbed by a field redefinition $t_i \rightarrow -t_i$.

For the PV amplitude $W$ we obtain

$$W^d = -2 \sqrt{\frac{\gamma r^{(3S_1)}}{1 - \gamma r^{(3S_1)}}} \left( 1 - \frac{1}{3} \gamma a^{(3S_1)} \right) g^{(3S_1-3P_1)}.$$

(22)

Expanding in $r^{(3S_1)}$ and using Eq. (11), as well as $\gamma = 1/a^{(3S_1)}$ at leading order, we reproduce the result of Eqs. (17) and (18) up to a factor of $-1$ as discussed above.\(^3\)

The asymmetry is given by

$$A_\gamma = 2M^2 \sqrt{\frac{\gamma r^{(3S_1)}}{\pi}} \frac{1 - \gamma a^{(3S_1)}}{\kappa_1 (1 - \gamma a^{(1S_0)})} g^{(3S_1-3P_1)},$$

(23)

which exactly reproduces Eq. (19) at leading order in $r^{(3S_1)}$.

### B. Circular polarization in $np \rightarrow d\gamma$

The photon circular polarization in $np \rightarrow d\gamma$ is defined by

$$P_\gamma = \frac{\sigma_+ - \sigma_-}{\sigma_+ + \sigma_-},$$

(24)

\(^3\) We note that our values for $W^d$ and $Y^d$ disagree with those in Ref. [17] (with $L_1 = 0$ since it is higher order) by a factor of $-2$ and $2$, respectively.
FIG. 4: Leading-order diagrams contributing to the parity-violating amplitudes $W$ and $V$ in the dibaryon formalism. Solid lines denote nucleons, thick solid lines denote dressed dibaryons and wavy lines denote photons. The gray oval stands for the deuteron interpolating field. The solid black box is the PV operator. The photons are minimally coupled.

where $\sigma_{+/-}$ is the total cross section for photons with positive/negative helicity. The polarization again stems from interference between PC and PV amplitudes and, up to the order to which we are working, is given by

$$P_\gamma = \frac{2 M \text{Re}[Y^*V]}{\gamma^2 |Y|^2},$$

(25)

where we have again used $X = 0$ to this order. The expression for $Y$ is already given in Eq. (15). The amplitude $V$ is calculated from the diagrams in Fig. 2 with the $np$ initial state in a $^1S_0$ partial wave. We find

$$V = 4 \sqrt{\frac{\gamma}{\pi}} \left[ \left(1 - \frac{2}{3} \gamma a^{(1S_0)} \right) \left(\gamma - \mu\right) C^{(3S_1-1P_1)} + \frac{\gamma a^{(1S_0)}}{3} \left(\frac{1}{a^{(1S_0)}} - \mu\right) \left( C^{(1S_0-3P_0)}_{(\Delta I=0)} - 2 C^{(1S_0-3P_0)}_{(\Delta I=2)} \right) \right],$$

(26)

or, using Eqs. (14) and (15),

$$V = \frac{16 \pi}{M} \sqrt{\frac{\gamma}{\pi}} \left[ \left(1 - \frac{2}{3} \gamma a^{(1S_0)} \right) \frac{C^{(3S_1-1P_1)}}{C_0^{(3S_1)}} + \frac{1}{3} \gamma a^{(1S_0)} \frac{C^{(1S_0-3P_0)}_{(\Delta I=0)}}{C_0^{(1S_0)}} - 2 \frac{C^{(1S_0-3P_0)}_{(\Delta I=2)}}{C_0^{(1S_0)}} \right].$$

(27)
The photon polarization $P_\gamma$ is given by

$$P_\gamma = -16 \frac{M}{\kappa_1 (1 - \gamma a^{(1S_0)})} \left[ \left( 1 - \frac{2}{3} \gamma a^{(1S_0)} \right) \frac{C^{(3S_1-1P_1)}}{C_0^{(3S_1)}} + \frac{\gamma a^{(1S_0)}}{3} \frac{C^{(3S_0-3P_0)}}{C_0^{(1S_0)}} - 2C^{(3S_0-3P_0)} \right].$$

In the dibaryon formalism we obtain

$$V^d = -2 \sqrt{\frac{\gamma a^{(3S_1)}}{1 - \gamma a^{(3S_1)}}} \left[ \left( 1 - \frac{2}{3} \gamma a^{(1S_0)} \right) g^{(3S_1-1P_1)} + \frac{\gamma a^{(1S_0)}}{3} \sqrt{\frac{\pi}{r^{(3S_1)}}} \left( g^{(3S_0-3P_0)} - 2g^{(1S_0-3P_0)} \right) \right],$$

which leads to the polarization

$$P_\gamma = -2 \sqrt{\frac{\gamma a^{(1S_0)}}{1 - \gamma a^{(1S_0)}}} \left[ \left( 1 - \frac{2}{3} \gamma a^{(1S_0)} \right) g^{(3S_1-1P_1)} + \frac{\gamma a^{(1S_0)}}{3} \sqrt{\frac{\pi}{r^{(3S_1)}}} \left( g^{(3S_0-3P_0)} - 2g^{(1S_0-3P_0)} \right) \right].$$

Using the relations for the PV dibaryon couplings in Eq. (11) we see that this reproduces the result of Eq. (28).

**IV. CONCLUSION**

We obtained results for the photon asymmetry $A_\gamma$ in polarized neutron capture and for the photon polarization $P_\gamma$ in unpolarized capture. We provide results expressed in both the dibaryon and the non-dibaryon coefficient language so that our results can be used in conjunction with additional calculations performed in either operator set. In Eq. (11) we present the complete leading-order, parity-violating dibaryon Lagrangian required for these calculations. Our results will allow the extraction of two of the PV parameters once the experiments on $np \rightarrow d\gamma$ become available.

Measurements are available on parity violation in more complicated nuclear (and atomic) systems, but of course these are more difficult for theorists to treat systematically, and to date it is not clear if the experimental results have a consistent theoretical interpretation. A solid understanding of the two-nucleon sector is likely a necessary prerequisite to understanding the many-nucleon parity violating observables. Fortunately, EFTs allow the extraction of parameters from two nucleon observables that can then be consistently used in calculations on more complicated systems.

To obtain an experimental determination of the five parity violating parameters appearing at leading order, at least five experiments will be required. The two asymmetries from $np \rightarrow d\gamma$ provide two of them. $\bar{n}p \rightarrow d\gamma$ provides $C^{(3S_1-3P_1)}$ (or in the dibaryon language, $g^{(3S_1-3P_1)}$) while $np \rightarrow d\gamma$ yields a linear combination of $C^{(3S_1-1P_1)}$, $C^{(3S_0-3P_0)}$, and $C^{(1S_0-3P_0)}$ (or the corresponding dibaryon coefficients). Asymmetries from $\bar{n}n$, $\bar{n}p$, and $\bar{p}p$ scattering would yield three more linear combinations, including a dependence on $C^{(3S_0-3P_0)}$, but experimental results on $\bar{n}p$ and $\bar{n}n$ scattering are unlikely in the near future. To obtain further constraints on the parameters would require extending the treatment to few-body systems.
From a theoretical perspective a three-body calculation is the natural extension of the current program, in particular considering recent experimental interest in the reaction $\bar{n}d \rightarrow t\gamma$ [36].

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[1] Ch. Sinclair et al., “Letter-of-Intent 00-002 for PAC 17: Study of the Parity Nonconserving Force Between Nucleons Through Deutron Photodisintegration”.
[2] B. Wojtsekhowski and W.T.H. van Oers, “Summary of the Working Group Meeting on Parity Violation in Deuteron Photodisintegration with Circularly Polarized Photons,” 13-14 April, 2000, Jefferson Lab.
[3] D.M. Markoff, J. Res. Natl Inst. Stan. Tech. 110, 209 (2005).
[4] E. Stiliaris, Eur. Phys. J. A 24S2, 175 (2005).
[5] B. Lauss et al., AIP Conf. Proc. 842, 790 (2006) [arXiv:nucl-ex/0601004).
[6] C. D. Bass et al., [arXiv:0905.0395 [nucl-ex].
[7] D. R. Phillips, M. R. Schindler and R. P. Springer, Nucl. Phys. A 822, 1 (2009) [arXiv:0812.2073 [nucl-th]]
[8] G. S. Danilov, Phys. Lett. 18, 40 (1965); Phys. Lett. B35, 579 (1971). Sov. J. Nucl. Phys. 14, 443 (1972).
[9] F. C. Michel, Phys. Rev. 133, B329 (1964).
[10] B. Desplanques and J. H. Missimer, Nucl. Phys. A 300, 286 (1978).
[11] B. Desplanques, J. F. Donoghue and B. R. Holstein, Annals Phys. 124, 449 (1980).
[12] W. C. Haxton, [arXiv:0802.2984 [nucl-th].
[13] D. B. Kaplan and M. J. Savage, Nucl. Phys. A 556, 653 (1993) [Erratum-ibid. A 570, 833 (1994 ERRAT,A580,679.1994)].
[14] D. B. Kaplan, M. J. Savage, R. P. Springer and M. B. Wise, Phys. Lett. B 449, 1 (1999) [arXiv:nucl-th/9907081].
[15] M. J. Savage and R. P. Springer, Nucl. Phys. A 644, 235 (1998) [Erratum-ibid. A 657, 457 (1999)] [arXiv:nucl-th/9807014].
[16] S. L. Zhu, C. M. Maekawa, B. R. Holstein, M. J. Ramsey-Musolf and U. van Kolck, Nucl. Phys. A 748, 435 (2005) [arXiv:nucl-th/0407087].
[17] M. J. Savage, Nucl. Phys. A 695, 365 (2001) [arXiv:nucl-th/0012043].
[18] S. R. Beane and M. J. Savage, Nucl. Phys. A 694, 511 (2001) [arXiv:nucl-th/0011067].
[19] D. R. Phillips, G. Rupak and M. J. Savage, Phys. Lett. B 473, 209 (2000) [arXiv:nucl-th/9908054].
[20] J. W. Shin, S. Ando and C. H. Hyun, [arXiv:0907.3995 [nucl-th].
[21] J. F. Cavaignac, B. Vignon and R. Wilson, Phys. Lett. B 67 (1977) 148.
[22] J. Alberi et al., Can. J. Phys. 66 (1988) 542.
[23] P.-N. Seo, private communication, on NPDgamma experiment at ORNL/SNS.
[24] V. A. Knyaz'kov et al., Nucl. Phys. A 417, 209 (1984).
[25] E. D. Earle et al., Can. J. Phys. 66 (1988) 534.
[26] D. B. Kaplan, M. J. Savage and M. B. Wise, Phys. Rev. C 59, 617 (1999)
  [arXiv:nucl-th/9804032].
[27] U. van Kolck, arXiv:hep-ph/9711222; Nucl. Phys. A 645, 273 (1999) [arXiv:nucl-th/9808007].
[28] D. B. Kaplan, M. J. Savage and M. B. Wise, Phys. Lett. B 424, 390 (1998)
  [arXiv:nucl-th/9801034].
[29] J. Gegelia, Phys. Lett. B 429, 227 (1998).
[30] D. B. Kaplan, M. J. Savage and M. B. Wise, Nucl. Phys. B 534, 329 (1998)
  [arXiv:nucl-th/9802075].
[31] L. Girlanda, Phys. Rev. C 77, 067001 (2008) [arXiv:0804.0772 [nucl-th]].
[32] D. B. Kaplan, Nucl. Phys. B 494, 471 (1997) [arXiv:nucl-th/9610052].
[33] P. F. Bedaque and H. W. Griesshammer, Nucl. Phys. A 671, 357 (2000)
  [arXiv:nucl-th/9907077].
[34] C. H. Hyun, J. W. Shin and S. I. Ando, Mod. Phys. Lett. A 24, 827 (2009) [arXiv:0809.4892
  [nucl-th]].
[35] R. P. Springer, Duke High Energy Physics Seminar, 4 November 2008; M. R. Schindler, North
  Carolina State University Theory Talk, 24 April 2009; M. R. Schindler, 4th International
  Workshop “From Parity Violation to Hadronic Structure and more...”, 22-26 June 2009, Bar
  Harbor, http://web.mit.edu/pavi09/talks/Schindler_pavi09.pdf.
[36] C. B. Crawford, 4th International Workshop “From Parity Violation to Hadronic Structure and more...”, 22-26 June 2009, Bar Harbor,
  http://web.mit.edu/pavi09/talks/Crawford_neutron_cap_pv.pdf.