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Parity violation in neutron capture on the proton: Determining the weak pion–nucleon coupling

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ABSTRACT

We investigate the parity-violating analyzing power in neutron capture on the proton at thermal energies in the framework of chiral effective field theory. By combining this analysis with a previous analysis of parity violation in proton–proton scattering, we are able to extract the size of the weak pion–nucleon coupling constant. The uncertainty is significant and dominated by the experimental error which is expected to be reduced soon.

Although parity violation (PV) induced by the weak interaction is well understood at the level of elementary quarks, its manifestation at the hadronic and nuclear level is not that clear. This holds particularly true for the strangeness-conserving part of the weak interaction which induces PV in hadronic and nuclear systems. The Standard Model predicts PV forces between nucleons. However, their forms and strengths are masked by the nonperturbative nature of QCD at low energies. Combined with the difficulty of doing experiments with sufficient accuracy to extract parity-violating signals, hadronic PV is one of the least tested parts of the Standard Model.

The understanding of low-energy strong interactions has increased tremendously by the use of effective field theories (EFTs). It has been realized that by writing down the most general interactions among the low-energy degrees of freedom that are consistent with the symmetries of QCD, one obtains an EFT, chiral perturbation theory (χPT), that is a low-energy equivalent of QCD. Each interaction term in the chiral Lagrangian comes with a coupling strength, or low-energy constant (LEC), which needs to be extracted from data or computed in lattice QCD. In contrast to low-energy QCD itself, χPT allows one to calculate observables in a perturbative framework with expansion parameter p/Λ_χ, where p is the momentum scale of the process and $\Lambda_\chi \sim 1$ GeV, the scale where the EFT breaks down. Although nuclear physics is intrinsically nonperturbative, the nucleon–nucleon (*NN*) potential can be calculated perturbatively within χPT. The resulting chiral potential is then iterated to all orders to calculate *NN*-scattering and bound state properties. This framework is usually called chiral nuclear EFT (for recent reviews, see Refs. [1,2]).

The success of chiral EFT in parity-conserving (PC) nuclear physics has led to an analogous program in the PV sector [3–7]. One starts with the four-quark operators that are induced when the heavy weak gauge bosons are integrated out. The next step entails constructing a PV chiral Lagrangian which contains all interaction terms that transform under chiral symmetry in the same way as the underlying four-quark operators. From the resulting chiral Lagrangian one then calculates the PV *NN* potential and electromagnetic current. In the final step the obtained PV potential and current are applied, in combination with the PC chiral potential and current, in calculations of nuclear processes. The PV LECs appearing in the PV chiral Lagrangian can be fitted to some data and other PV processes can then be predicted.

Although this sounds like a good strategy, in practice this procedure is complicated by the lack of data on PV processes. So far, hadronic PV has only been measured in a handful of experiments (see Refs. [8,9] for recent reviews). The longitudinal analyzing power (LAP), which would be zero in the limit of no PV, has been measured for proton–proton scattering at three different energies [10–12], for proton–alpha scattering only at a single

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energy [13,14], and recently for the first time a preliminary result has been reported for radiative neutron capture on the proton $\vec{n}p \to d\gamma$ at thermal energies [15]. Nonzero parity-violating signals have also been found in more complex systems, as exemplified by the radiative decay of the ^{19}F nucleus [16,17] and the anapole moment of the Cesium atom [18].

The first full chiral EFT analysis of PV nuclear forces has been done in Ref. [3] where it has been concluded that at leading order (LO) only a single interaction term appears:

$$\mathcal{L}_{\not P} = \frac{h_\pi}{\sqrt{2}} \bar{N}(\vec{\pi} \times \vec{\tau})^3 N \ , \tag{1}$$

written in terms of the pion isospin-triplet $\vec{\pi}$, the nucleon isospin-doublet $N = (p, n)^T$, and the weak pion–nucleon coupling constant h_π. The leading order PV potential arising from one-pion exchange takes the form

$$V_{\rm OPE} = -\frac{g_A h_\pi}{2\sqrt{2}F_\pi} i(\vec{\tau}_1 \times \vec{\tau}_2)^3 \frac{(\vec{\sigma}_1 + \vec{\sigma}_2)\cdot \vec{k}}{m_\pi^2 + \vec{k}^2}, \tag{2}$$

in terms of the momentum transfer $\vec{k} = \vec{p} - \vec{p}\,'$, where $\vec{p}$ and $\vec{p}\,'$ are the incoming and outgoing nucleon momenta in the center-of-mass frame, and $\vec{\sigma}_{1,2}$ and $\vec{\tau}_{1,2}$ the nucleon spin- and isospin-operators, respectively. $F_\pi = 92.4$ MeV denotes the pion decay constant, $m_\pi = 139.57$ MeV the charged pion mass, and $g_A = 1.29$ the nucleon axial-vector coupling constant taking into account the Goldberger–Treiman discrepancy, in order to represent the strong πNN-coupling.

Considering that there are no other terms at leading order, the one-pion exchange (OPE) potential can be expected to give the dominant contribution to PV in nuclear processes. Nevertheless, despite decades of experimental effort the existence of a long-range PV NN force has not been confirmed. This indicates that h_π could be smaller than expected from naive dimensional analysis which predicts $h_\pi \sim G_F F_\pi \Lambda_\chi \sim 10^{-6}$ (consistent with the often-used estimate $h_\pi = 4.6 \cdot 10^{-7}$ of Ref. [19]), with $G_F \simeq 1.67 \cdot 10^{-5}$ GeV^{-2} the Fermi coupling constant. In fact, the isovector nature of the weak pion–nucleon coupling already gives a natural suppression of $\sin^2\theta_W \sim 1/4$ [3,20], while a large-N_c analysis indicates that h_π is even further suppressed [20–22]. A first lattice QCD calculation gave $h_\pi \simeq 10^{-7}$ [23]. Finally, the absence of a PV signal in the γ-ray emission from ^{18}F leads to the bound $h_\pi \leq 1.3 \cdot 10^{-7}$ [24–26].

The evidence in favor of a small value of h_π is not conclusive. Large-N_c arguments can be misleading, especially for pionic interactions, while the lattice calculation did not include disconnected diagrams. The bound from ^{18}F depends on nuclear structure calculations of a relatively complicated nucleus and, despite being a careful work, might suffer from uncontrolled uncertainties. Finally, the Cesium anapole moment prefers a much larger value $h_\pi \simeq 10^{-6}$ although the involved uncertainties are also larger [27,28]. It seems that the only conclusive method of determining the size of h_π is through a fit to experiments using simple few-body processes which are theoretically much better under control. Unfortunately, only a few PV signals have been measured so far in such few-body processes. In recent work we investigated the data on $\vec{p}p$ scattering in a chiral EFT framework [6,29]. The main goal of this paper is to combine this analysis with the recent data on PV in radiative neutron capture on the proton $\vec{n}p \to d\gamma$ and extract a value of h_π. An analysis of PV in the inverse process $\vec{\gamma}d \to np$ within pionless EFT has recently been performed in Ref. [30].

Our task gets complicated by two things. First of all, the OPE potential in Eq. (2) changes the total isospin and does not contribute to $\vec{p}p$ scattering. The three data points still carry information on the size of h_π because the analyzing power does depend on h_π through the two-pion-exchange (TPE) diagrams [5,29,7]. The TPE diagrams appear at higher order in the chiral counting where additional contributions in the form of PV NN contact terms appear as well [5,31,32]. Secondly, although the PV OPE potential does contribute to $\vec{n}p \to d\gamma$ capture, if the coupling constant h_π is really as small as suggested, formally higher-order corrections can become relevant and need to be taken into account. Again such corrections appear as NN contact terms. We discuss these subleading terms in the PV potential and the current at a later stage.

The other ingredients required for the calculation of PV observables are the PC NN potential and the PC and PV electromagnetic currents. As in Ref. [29], we apply here the next-to-next-to-next-to-leading order (N^{3}LO) chiral EFT potential obtained in Ref. [33] and we refer the reader to this paper for all further details. The N^{3}LO potential exists for several values of the cut-off needed to regularize the scattering equation. Here, we regularize the PV potential in the same way as the PC potential via

$$V_{PV}(\vec{p},\, \vec{p}\,') \to e^{-p^6/\Lambda^6} V_{PV}(\vec{p},\, \vec{p}\,') e^{-p'^6/\Lambda^6} \ , \tag{3}$$

where three choices for $\Lambda = \{450, 550, 600\}$ MeV are applied, see Refs. [29,33]. TPE diagrams are regularized with a spectral cut-off $\Lambda_S = \{500, 600, 700\}$ MeV [33]. In recent work [34] an alternative regularization scheme (formulated in coordinate space) has been proposed which better preserves the long-range nature of pion-exchange terms in the potential. Considering the large experimental uncertainties in the field of nuclear parity violation, we do not expect drastic changes if the alternative regularization scheme is applied. Nevertheless, we will investigate this new scheme and its extension to N^{4}LO [35] in future work.

Within the chiral EFT power-counting rules the dominant PC current arises from the nucleon magnetic moments. At next-to-leading (NLO) order we encounter the one-body convection current,[1] which arises from gauging the nucleon kinetic energy term, and the leading OPE two-body currents. The total PC current up to NLO is then given by

$$\begin{aligned}\vec{J}_{PC} = & \sum_{j=1}^{2} \frac{e}{4m_N} \left\{ -\left[\mu_s + \mu_v \tau_j^3\right] i(\vec{\sigma}_j \times \vec{q}) + (1+\tau_j^3)(\vec{P}_j + \vec{P}_j^{\,\prime}) \right\} \\ & \times \delta^{(3)}[\vec{P}_j - \vec{P}_j^{\,\prime} - \vec{q}\,] \\ & + \frac{e g_A^2}{4F_\pi^2} i\left(\vec{\tau}_1 \times \vec{\tau}_2\right)^3 \left\{ 2\vec{k} \frac{\vec{\sigma}_1 \cdot (\vec{k} + \vec{q}/2)}{(\vec{k}+\vec{q}/2)^2 + m_\pi^2} \frac{\vec{\sigma}_2 \cdot (\vec{k} - \vec{q}/2)}{(\vec{k}-\vec{q}/2)^2 + m_\pi^2} \right. \\ & \left. - \vec{\sigma}_1 \frac{\vec{\sigma}_2 \cdot (\vec{k} - \vec{q}/2)}{(\vec{k}-\vec{q}/2)^2 + m_\pi^2} - \vec{\sigma}_2 \frac{\vec{\sigma}_1 \cdot (\vec{k} + \vec{q}/2)}{(\vec{k}+\vec{q}/2)^2 + m_\pi^2} \right\} \ , \end{aligned} \tag{4}$$

where $\mu_s = 0.88$ and $\mu_v = 4.72$ are the isoscalar and isovector nucleon magnetic moments. The momenta of the incoming and outgoing nucleon interacting with the photon (of outgoing momentum $\vec{q}$) are denoted by $\vec{P}_j$ and $\vec{P}_j^{\,\prime}$, respectively. The momenta carried by the intermediate pions are $\vec{k}+\vec{q}/2 = \vec{P}_1 - \vec{P}_1^{\,\prime}$ and $\vec{k}-\vec{q}/2 = \vec{P}_2^{\,\prime} - \vec{P}_2$. In contrast, the leading PV current is solely due to OPE diagrams where one of the pion-nucleon vertices is from Eq. (1)

[1] Here the power-counting rules of Ref. [33] are followed where recoil and relativistic corrections are relegated to higher order by counting $1/m_N \sim k/\Lambda_\chi^2$, where k is the typical momentum scale of the process. The magnetic moment operator is not a recoil correction and only scales as $1/m_N$ for conventional reasons. We thus treat $\mu_{s,v}/m_N \sim 1/\Lambda_\chi$ which is also justified by the large value of $\mu_v = 4.72$.