Parity-violating polarization in $np \to d\gamma$ with a pionless effective field theory

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We consider the two-nucleon weak interaction with a pionless effective field theory. Dibaryon fields are introduced to ensure fast convergence of the perturbative expansion. Weak interactions are accounted for with the parity-violating dibaryon-nucleon-nucleon vertices, which contain unknown weak coupling constants. We apply the model to the calculation of a parity-violating observable in the radiative neutron capture on a proton at threshold. Result is obtained up to the linear order of the weak dibaryon-nucleon-nucleon coupling constants. We compare our result to the ones obtained from other approaches, and discuss investigation of the weak interaction in few-body systems.

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I. INTRODUCTION

Weak nucleon-nucleon (NN) interaction has recently been formulated in the framework of effective field theory (EFT) [1], where parity-violating (PV) weak NN potentials were obtained up to next-to-next-to leading order in both pionful and pionless theories. Weak potentials thus obtained were subsequently employed in the calculation of PV observables in the two-nucleon systems [2–4], and the results demonstrate that the perturbative expansion converges reasonably, and thus the EFT is a working tool for the description of weak NN interaction at low energies.

In this work, we consider a pionless theory with dibaryon fields, which are auxiliary fields describing the two-nucleon states. Introducing a dibaryon field for the deuteron, we can take into account the effective range contribution ($\gamma\rho_d \sim 0.4$) to the deuteron propagator up to infinite order, and it consequently makes the convergence of the theory better than the pionless EFT that does not have dibaryon fields. Since the scattering lengths and the effective ranges are large in the two-nucleon S states, summation of effective range contribution to infinite order in the dibaryon formalism is especially useful for the two-nucleon systems that are dominated by S states. Parity-violating vertex in the pionless theory in Ref. [1] consists of the multiplication of a two-nucleon field in an S state with one in a P state. Given a rule to map a two-nucleon state to the corresponding dibaryon field, we can obtain the PV Lagrangian that describes the weak NN interaction in terms of PV dibaryon-nucleon-nucleon or dibaryon-dibaryon vertices. In this work, we obtain the PV Lagrangian with dibaryon fields by transforming the two-nucleon fields in the S states to the corresponding dibaryon ones, while the P states are represented in terms of the two-nucleon fields. Then the weak NN interaction is described with the PV dibaryon-nucleon-nucleon (dNN) vertices, which have unknown weak dNN coupling constants.

With the weak interaction thus obtained, we calculate the PV polarization (P_γ) in $np \to d\gamma$ at threshold. Parity-violating polarization has been calculated with the weak one-meson exchange potentials (conventionally referred to as DDH potential [5]), and with a few strong interaction models [6]. The results in Ref. [6] show substantial dependence on the strong interaction model, and are dominated by the ρ- and ω-meson exchange terms in the DDH potential. In the EFT, ρ, ω, and heavier mesons are integrated out because their masses are large scales in low-energy few-body processes, and their contributions are embedded in the low-energy constants in the NN contact terms. Because the PV polarization in $np \to d\gamma$ is dominated by the heavy-meson terms in the DDH potential, if it is considered in the EFT, only the contact terms are relevant. Therefore, pionless theories provide a natural framework for the investigation of the problem in the context of EFT. With the weak interaction described by the PV dNN vertices, we calculate P_γ and obtain the result up to linear order of the unknown weak dNN coupling constants. We compare the result with the ones obtained from other approaches, and discuss the problems for the weak interaction at low energies in the few-body system.

We outline the paper as the following. In Sec. II, we present the parity-conserving and the parity-violating Lagrangians that are relevant to the calculation in the work. In Sec. III, we obtain the PV polarization in the unpolarized neutron capture on a proton at threshold with the Lagrangians obtained in Sec. II, and discuss the result. We summarize the paper and discuss extension of the investigation to few-body systems in Sec. IV.

II. EFFECTIVE LAGRANGIAN

The parity-conserving part of the Lagrangian includes strong and electromagnetic (EM) interactions. Parity-conserving terms with dibaryon fields can be written as

$$\mathcal{L}_{\rm PC} = \mathcal{L}_N + \mathcal{L}_s + \mathcal{L}_t + \mathcal{L}_{st}, \tag{1}$$

where $\mathcal{L}_N$, $\mathcal{L}_s$, and $\mathcal{L}_t$ represent the strong interactions for the nucleon, dibaryon in the $^1S_0$ state, and dibaryon in the $^3S_1$ state, respectively, and $\mathcal{L}_{st}$ the EM transition between the $^1S_0$

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and the $^3S_1$ states. Terms relevant to this work read

$$\mathcal{L}_N = N^\dagger\left(iv\cdot D + \frac{1}{2m_N}\{(v\cdot D)^2 - D^2\}\right)N, \quad (2)$$

$$\mathcal{L}_s = \sigma_s s_a^\dagger\left\{iv\cdot D + \frac{1}{4m_N}[(v\cdot D)^2 - D^2] + \Delta_s\right\}s_a - y_s\left\{s_a^\dagger\left[N^T P_a^{(^1S_0)}N\right] + \text{H.c.}\right\}, \quad (3)$$

$$\mathcal{L}_t = \sigma_t t_i^\dagger\left\{iv\cdot D + \frac{1}{4m_N}[(v\cdot D)^2 - D^2] + \Delta_t\right\}t_i - y_t\left\{t_i^\dagger\left[N^T P_i^{(^3S_1)}N\right] + \text{H.c.}\right\}, \quad (4)$$

$$\mathcal{L}_{st} = \frac{L_1}{m_N\sqrt{r_0\rho_d}}[t_i^\dagger s_3 B_i + \text{H.c.}], \quad (5)$$

where the projection operators for the $^1S_0$ and the $^3S_1$ states are defined, respectively, as

$$P_a^{(^1S_0)} = \frac{1}{\sqrt{8}}\sigma_2\tau_2\tau_a, \quad (6)$$

$$P_i^{(^3S_1)} = \frac{1}{\sqrt{8}}\sigma_2\sigma_i\tau_2. \quad (7)$$

In the nonrelativisitic limit, we assume $v^2 = 1$ for the velocity vector v_μ, and the covariant derivative is defined as $D_\mu = \partial_\mu - i\mathcal{V}_\mu^{\text{ext}}$, where $\mathcal{V}_\mu^{\text{ext}}$ is the external vector field. Dibaryon fields in the $^1S_0$ and the $^3S_1$ states are denoted by s_a and t_i, respectively, and B_i is the magnetic field given by $\vec{B} = \nabla\times\vec{\mathcal{V}}^{\text{ext}}$. σ_s and σ_t are sign factors having a value -1, and Δ_s and Δ_t are the mass difference between the dibaryon and the two-nucleon states, $\Delta_{s,t} = m_{s,t} - 2m_N$. Low-energy constants y_s and y_t are the strong dNN coupling constants, which are determined from the empirical values of the effective ranges in the $^1S_0$ and $^3S_1$ states, respectively. We obtain $y_s = \frac{2}{m_N}\sqrt{\frac{2\pi}{r_0}}$ and $y_t = \frac{2}{m_N}\sqrt{\frac{2\pi}{\rho_d}}$, where r_0 is the effective range in the $^1S_0$ state and ρ_d is the effective range for the deuteron. Low-energy constant L_1 is a photon-dibaryon-dibaryon coupling constant for the M1 transition, and it will be determined from experiments.

Parity-violating terms for the two-nucleon weak interactions can be written as

$$\mathcal{L}_{\text{PV}} = \sum_{\Delta I}\mathcal{L}_{\text{PV}}^{\Delta I}, \quad (8)$$

where ΔI denotes the isospin change at the PV vertex. Because $\Delta(L+S+I)$ is even and ΔL is odd at a PV vertex, $\Delta(S+I) = 1$ for the two-nucleon system, and thus we have

$$\mathcal{L}_{\text{PV}} = \mathcal{L}_{\text{PV}}^0 + \mathcal{L}_{\text{PV}}^1. \quad (9)$$

Because the total angular momentum is conserved in the interaction, parity admixtures allowed by the PV vertices for the lowest orbital states are $^1S_0 \leftrightarrow {}^3P_0$, and $^3S_1 \leftrightarrow {}^1P_1$ due to $\mathcal{L}_{\text{PV}}^0$, and $^3S_1 \leftrightarrow {}^3P_1$ due to $\mathcal{L}_{\text{PV}}^1$. As a result, nonrelativistic P-odd and T-even Lagrangian for the neutron-proton system with $\Delta I = 0$ can be written as

$$\mathcal{L}_{\text{PV}}^0 = \frac{h_{\text{dNN}}^{0s}}{2\sqrt{2\rho_d r_0}m_N^{5/2}}s_3^\dagger N^T\sigma_2\sigma_i\tau_2\tau_3\frac{i}{2}(\overleftarrow{\nabla} - \overrightarrow{\nabla})_i N + \text{H.c.} \quad (10)$$

$$+ \frac{h_{\text{dNN}}^{0t}}{2\sqrt{2}\rho_d m_N^{5/2}}t_i^\dagger N^T\sigma_2\tau_2\frac{i}{2}(\overleftarrow{\nabla} - \overrightarrow{\nabla})_i N + \text{H.c.}, \quad (11)$$

where h_{dNN}^{0s} and h_{dNN}^{0t} denote the weak dNN coupling constants. Spin-isospin operator $\sigma_2\sigma_i\tau_2\tau_a$ in Eq. (10) projects the two-nucleon system to the $^3P_0$ state, and thus the PV vertex in the equation produces $^3P_0$ mixture in the $^1S_0$ state. Similarly, the operator $\sigma_2\tau_2$ in Eq. (11) is the projection operator for the $^1P_1$ state, and thus the term mixes the $^1P_1$ state to the $^3S_1$ state. For the $\Delta I = 1$ part, we have $^3P_1$ admixture to the $^3S_1$ state, so the Lagrangian reads

$$\mathcal{L}_{\text{PV}}^1 = i\frac{h_{\text{dNN}}^1}{2\sqrt{2}\rho_d m_N^{5/2}}\epsilon_{ijk}t_i^\dagger N^T\sigma_2\sigma_j\tau_2\tau_3\frac{i}{2}(\overleftarrow{\nabla} - \overrightarrow{\nabla})_k N + \text{H.c.} \quad (12)$$

Lagrangians given by Eqs. (10), (11), (12) account for the weak interactions between a neutron and a proton in the pionless theory with dibaryon fields.

III. RESULT AND DISCUSSION

In the pionless theory, expansion parameters are Q/m_π or Q/Λ, where Q is a small momentum, m_π the pion mass, and Λ a symmetry breaking scale. Because the scattering lengths and effective ranges in the $^1S_0$ and the $^3S_1$ states are large, we count their inverse as small scales, that is, $(\gamma, 1/a_s, 1/a_t, 1/r_0, 1/\rho_d) \sim Q$. $a_{s(t)}$ is the scattering length in the $^1S_0(^3S_1)$ state, and $\gamma = \sqrt{m_N B}$, where B is the binding energy of the deuteron. Propagators for the nucleon and the dibaryon are counted as $1/Q^2$ and a loop integral contributes an order of Q^5.

Feynman diagrams for the PV amplitude at leading order (Q^0) are depicted in Fig. 1. Single straight and wavy lines represent the nucleon and the photon fields, respectively. Double line with solid circle denotes the dressed dibaryon field. Because $y \propto r_0^{-1/2} \sim Q^{1/2}$, the order of a nucleon loop inserted in the dibaryon propagator is Q^0, which does not change the order of the dibaryon propagator. Therefore, we have to sum nucleon loops in the dibaryon propagators up to infinite order, and the result of the summation is called the dressed dibaryon field. Small dots at the dibaryon-nucleon-nucleon vertices denote the strong dNN vertices proportional to y_s or y_t, and the circles with a cross represent the weak dNN vertices proportional to h_{dNN}^{0s}, h_{dNN}^{0t}, and h_{dNN}^1. For the photon-nucleon-nucleon coupling in Figs. 1(a)–1(c), we employ the vertex function of the convection current given by

$$i\Gamma_{\text{VNN}}(E1) = \frac{i}{2m_N}(1+\tau_3)\frac{1}{2}(\vec{p} + \vec{p}\,')\cdot\vec{\epsilon}_\gamma^{\,*}, \quad (13)$$

where $\vec{p}$ and $\vec{p}\,'$ are momenta for the incoming and outgoing nucleons, respectively, and $\vec{\epsilon}_\gamma^{\,*}$ is the polarization of the outgoing photons. For the weak photon-dibaryon-nucleon-nucleon vertices in Figs. 1(d)–1(f), we assume minimal coupling to the weak dNN vertex,

$$\vec{\nabla} \to \vec{\nabla} - i\frac{e}{2}(1+\tau_3)\vec{\mathcal{V}}^{\text{ext}}, \quad (14)$$

where $\vec{\mathcal{V}}^{\text{ext}}$ denotes the external photon field.

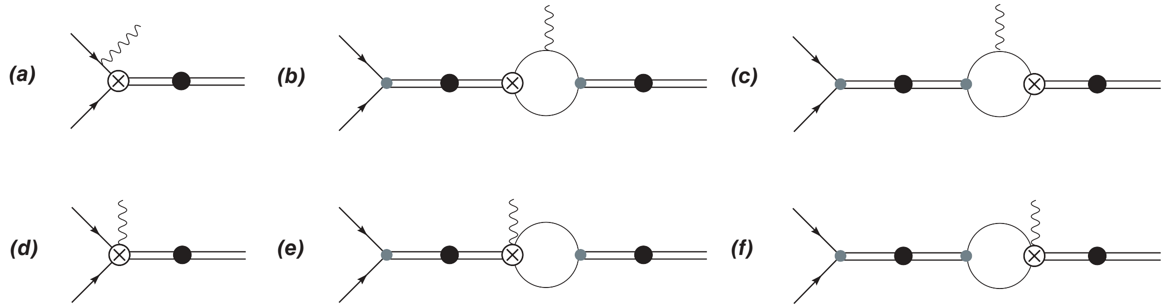

FIG. 1. Leading order (Q^0) PV diagrams for np capture. Single straight line denotes a nucleon, wavy line a photon, and a double line with a solid circle stands for dressed dibaryon propagator. Circle with a cross represents a PV dNN vertex.

Parity-violating polarization P_γ in $np \to d\gamma$ is defined as

$$P_\gamma = \frac{\sigma_+ - \sigma_-}{\sigma_+ + \sigma_-}, \tag{15}$$

where σ_+ and σ_- are the total cross sections for the photons with right and left helicities, respectively. P_γ was measured in 70s, and the result reads $P_\gamma = (1.8 \pm 1.8) \times 10^{-7}$ [8]. Because of the difficulty to measure the polarization of outgoing photons in $np \to d\gamma$, modern facilities with intense laser beams are more interested in measuring PV asymmetry in $d\vec{\gamma} \to np$. Parity-violating asymmetry in $d\vec{\gamma} \to np$ has been recently calculated with the DDH potential up to 10 MeV above threshold [9,10]. Magnitude of the asymmetry is maximum at threshold and it decreases very quickly as the energy increases. Since P_γ is equal to the PV asymmetry in $d\vec{\gamma} \to np$ at threshold because of detailed balance, measurement of the asymmetry at threshold can be directly related to P_γ.

Transition amplitude for the neutron-proton capture process can be written as

$$iM_{np} = [Y\vec{\epsilon}_d^{\,*} \cdot (\hat{k} \times \vec{\epsilon}_\gamma^{\,*}) - iZ\vec{\epsilon}_d^{\,*} \cdot \vec{\epsilon}_\gamma^{\,*}]N^T P_3^{(^1S_0)} N, \tag{16}$$

where Y and Z are the parity-conserving and parity-violating terms, respectively. At threshold parity-conserving amplitude Y is dominated by the M1 transition, and we take the result in Ref. [11],

$$Y = \frac{\sqrt{2\pi}}{m_N^2}\sqrt{\frac{\gamma}{1-\gamma\rho_d}}[(1+\kappa_V)(1-\gamma a_s) - \gamma^2 a_s L_1], \tag{17}$$

where $\kappa_V = 3.706$, $\gamma = \sqrt{m_N B} = 45.7$ MeV, $\rho_d = 1.764$ fm, and $a_s = -23.732$ fm. Factor $1/(1-\gamma\rho_d)$ in the square root originates from the infinite sum of the nucleon loop in the dibaryon propagator. Low-energy constant L_1 is fitted to the np capture cross section at threshold, $\sigma_{\rm exp} = 334.2 \pm 0.5$ mb, and we obtain $L_1 = -4.427 \pm 0.015$ fm [11]. Same as the parity-conserving amplitude, transition also occurs from the initial $^1S_0$ state to the final $^3S_1$ one for the parity-violating amplitude Z. Each diagram in Fig. 1 gives the PV amplitude,

$$Z_a = -\frac{1}{3}\frac{h^{0t}_{\rm dNN}}{m_N^2\sqrt{m_N\rho_d}}\sqrt{\frac{\gamma}{1-\gamma\rho_d}}\frac{p^2}{\gamma^2+p^2}, \tag{18}$$

$$Z_b = -\frac{1}{3}\frac{h^{0s}_{\rm dNN}}{m_N^2\sqrt{m_N\rho_d}}\sqrt{\frac{\gamma}{1-\gamma\rho_d}}\frac{1}{\frac{1}{a_s}-\frac{1}{2}r_0p^2+ip}\frac{\gamma^3+ip^3}{\gamma^2+p^2}, \tag{19}$$

$$Z_c = -\frac{1}{3}\frac{h^{0t}_{\rm dNN}}{m_N^2\sqrt{m_N\rho_d}}\sqrt{\frac{\gamma}{1-\gamma\rho_d}}\frac{1}{\frac{1}{a_s}-\frac{1}{2}r_0p^2+ip}\frac{\gamma^3+ip^3}{\gamma^2+p^2}, \tag{20}$$

$$Z_d = \frac{1}{2}\frac{h^{0t}_{\rm dNN}}{m_N^2\sqrt{m_N\rho_d}}\sqrt{\frac{\gamma}{1-\gamma\rho_d}}, \tag{21}$$

$$Z_e = \frac{1}{2}\frac{h^{0s}_{\rm dNN}}{m_N^2\sqrt{m_N\rho_d}}\sqrt{\frac{\gamma}{1-\gamma\rho_d}}\frac{\gamma}{\frac{1}{a_s}-\frac{1}{2}r_0p^2+ip}, \tag{22}$$

$$Z_f = -\frac{1}{2}\frac{h^{0t}_{\rm dNN}}{m_N^2\sqrt{m_N\rho_d}}\sqrt{\frac{\gamma}{1-\gamma\rho_d}}\frac{ip}{\frac{1}{a_s}-\frac{1}{2}r_0p^2+ip}, \tag{23}$$

where $r_0 = 2.70$ fm. Taking the limit $p \to 0$ at threshold, we obtain the net PV amplitude,

$$Z = \frac{1}{m_N^2\sqrt{m_N\rho_d}}\sqrt{\frac{\gamma}{1-\gamma\rho_d}} \times \left[h^{0t}_{\rm dNN}\left(\frac{1}{2}-\frac{1}{3}\gamma a_s\right)+\frac{1}{6}h^{0s}_{\rm dNN}\gamma a_s\right], \tag{24}$$

and the PV polarization P_γ at leading order reads

$$\begin{aligned} P_\gamma &= -2\frac{\mathrm{Re}(YZ^*)}{|Y|^2} \\ &= -\sqrt{\frac{2}{\pi m_N\rho_d}}\frac{\left(\frac{1}{2}-\frac{1}{3}\gamma a_s\right)h^{0t}_{\rm dNN}+\frac{1}{6}\gamma a_s h^{0s}_{\rm dNN}}{(1+\kappa_V)(1-\gamma a_s)-\gamma^2 a_s L_1} \\ &= -\left(2.59h^{0t}_{\rm dNN} - 1.01h^{0s}_{\rm dNN}\right)\times 10^{-2}. \end{aligned} \tag{25}$$

Parity-violating polarization depends on two weak coupling constants $h^{0t}_{\rm dNN}$ and $h^{0s}_{\rm dNN}$, and thus we cannot determine them uniquely from a single measurement of P_γ at threshold. In order to determine them unambiguously, we need more data for P_γ at energies above threshold, or measurement of observables that are dependent on $h^{0t}_{\rm dNN}$ and/or $h^{0s}_{\rm dNN}$. We will discuss this matter in the next section.

We now try to compare our result to the one obtained with a pionless EFT without dibaryon fields [2]. The comparison may be viable if we have relations for the weak coupling constants in the two theories. Lagrangians for the strong interactions at leading order in the pionless theory are given as

$$\mathcal{L} = N^\dagger\left(i\partial_0 + \frac{\nabla^2}{2m_N}\right)N - C_0^{(^1S_0)}(N^T P^{(^1S_0)}N)^\dagger(N^T P^{(^1S_0)}N) - C_0^{(^3S_1)}(N^T P^{(^3S_1)}N)^\dagger(N^T P^{(^3S_1)}N) + \cdots, \quad (26)$$

where $C_0^{(^1S_0)}$ and $C_0^{(^3S_1)}$ are the strong coupling constants for the $^1S_0$ and the $^3S_1$ states, respectively. Comparing the strong Lagrangians given by Eq. (26) to those with dibaryon fields, Eqs. (3) and (4) in this work, we obtain the rules to transform the two-nucleon field to the dibaryon one as

$$N^T P_a^{(^1S_0)}N \to \frac{y_s}{C_0^{(^1S_0)}}s_a, \quad N^T P_i^{(^3S_1)}N \to \frac{y_t}{C_0^{(^3S_1)}}t_i. \quad (27)$$

Substituting the transformations given by Eq. (27) into the pionless weak Lagrangians in Ref. [7], and comparing them with the weak Lagrangians in Eqs. (10)–(12) in this work, we can have the relations for the weak coupling constants in the two theories. Employing the strong coupling constants in the power divergence subtraction scheme,

$$\frac{1}{C_0^{(^1S_0)}} = \frac{m_N}{4\pi}\left(\frac{1}{a_0} - \mu\right), \quad \frac{1}{C_0^{(^3S_1)}} = \frac{m_N}{4\pi}\left(\gamma - \tfrac{1}{2}\gamma^2\rho_d - \mu\right), \quad (28)$$

where μ is the renormalization point, we obtain the relations for the weak coupling constants as

$$h_{\rm dNN}^{0s} = 16\sqrt{\frac{\rho_d m_N}{2\pi}}m_N^2\left(\frac{1}{a_0} - \mu\right)\left(\mathcal{C}_{\Delta I=0}^{(^1S_0-^3P_0)} - 2\mathcal{C}_{\Delta I=2}^{(^1S_0-^3P_0)}\right), \quad (29)$$

$$h_{\rm dNN}^{0t} = 16\sqrt{\frac{\rho_d m_N}{2\pi}}m_N^2\left(\gamma - \tfrac{1}{2}\gamma^2\rho_d - \mu\right)\mathcal{C}^{(^3S_1-^1P_1)}. \quad (30)$$

Inserting Eqs. (29) and (30) to the PV amplitude in Eq. (24) and assuming $\mu = m_\pi$, we obtain

$$Z \propto \mathcal{C}^{(^3S_1-^1P_1)} - 0.56\left(\mathcal{C}_{\Delta I=0}^{(^1S_0-^3P_0)} - 2\mathcal{C}_{\Delta I=2}^{(^1S_0-^3P_0)}\right). \quad (31)$$

Isospin change is zero at the vertex denoted by $\mathcal{C}^{(^3S_1-^1P_1)}$ (i.e., $\Delta I = 0$). Assuming a rough relation $\mathcal{C}^{(^3S_1-^1P_1)} \sim \mathcal{C}_{\Delta I=0}^{(^1S_0-^3P_0)}$, we obtain $Z \sim 0.44\mathcal{C}_{\Delta I=0}^{(^1S_0-^3P_0)} + 1.12\mathcal{C}_{\Delta I=2}^{(^1S_0-^3P_0)}$, where the ratio of the coefficient for $\mathcal{C}_{\Delta I=0}^{(^1S_0-^3P_0)}$ ($\Delta I = 0$) to that for $\mathcal{C}_{\Delta I=2}^{(^1S_0-^3P_0)}$ ($\Delta I = 2$) is approximately 1 : 2.5. Parity-violating polarization has been calculated in Ref. [2], where strong interaction is described by Argonne v18 model (Av18) [12], weak interaction by the pionless EFT, and the EM operator is obtained by using the Siegert theorem. The result in Ref. [2] is represented in terms of Danilov parameters. Transforming the Danilov parameters to the PV low-energy constants in the pionless theory, P_γ in Ref. [2] is written as

$$P_\gamma(\text{hybrid}) = (-0.25C_1 + 2.14C_3 + 4.18C_5) \times 10^{-3}, \quad (32)$$

where C_1 and C_3 correspond to $\Delta I = 0$ vertices and C_5 to $\Delta I = 2$ one. Assuming $C_3 \sim C_1$ and comparing the coefficients for C_1 ($\Delta I = 0$) to that for C_5 ($\Delta I = 2$) in Eq. (32), we obtain a ratio 1 : 2.3, which is similar to the ratio from our result. A similar value of the ratio for $\Delta I = 0$ to $\Delta I = 2$ was also obtained from the calculation that employed DDH potential for the weak interaction, Av18 for the strong interaction, and EM operator with the Siegert theorem [6].

IV. CONCLUSION

We have calculated the parity-violating polarization in the radiative neutron capture on a proton at threshold with a new framework, pionless EFT with dibaryon fields. Two-nucleon weak interactions are described with the parity-violating dibaryon-nucleon-nucleon vertices, whose coupling constants are not determined yet. Parity-violating polarization is obtained as a linear combination of two weak dNN coupling constants. If the weak dNN coupling constants are determined precisely from either experiments or theories, we may be able to understand and predict weak phenomena in the two-nucleon system at low energies with the pionless EFT with dibaryon fields.

There are five unknown weak coupling constants in the pionless theory [13] and therefore we need at least five data in order to determine them unambiguously. P_γ may be one of them. Recently PV longitudinal asymmetries in $\vec{p}p$, $\vec{n}p$, and $\vec{n}n$ scattering have been calculated with a pionless EFT [7]. Longitudinal asymmetry in $\vec{p}p$ depends on three PV coupling constants $\mathcal{C}_{\Delta I=0,1,2}^{(^1S_0-^3P_0)}$, and thus the datum at 13.6 MeV [14] provides a constraint for the weak coupling constants. Parity-violating asymmetry in $\vec{n}p \to d\gamma$ and the anapole moment of the deuteron have been calculated in both conventional and EFT approaches [15–19], and the results turn out to be dominated by $h_{\rm dNN}^1$ (or the weak pion-nucleon coupling constant h_π^1). Measurement of the PV asymmetry in $\vec{n}p \to d\gamma$ at SNS is expected to provide an important constraint with which we can determine the value of weak coupling constant $h_{\rm dNN}^1$ (or h_π^1). Turning to the possibilities in the three-body system, one can find a recent calculation of the weak effect in the spin rotation in $\vec{n}d$ scattering [20]. DDH potential was employed for the weak interaction, and the result is dominated by h_π^1. Measurement of the asymmetry in $\vec{n}p$ and the spin rotation in $\vec{n}d$ will provide a double-check for h_π^1 (or $h_{\rm dNN}^1$). To determine the remaining weak coupling constants in the pionless EFT, we need calculations and measurements for as many observables as possible.

Among the PV observables in the few-body system, PV asymmetry in $\vec{n}d \to t\gamma$ at threshold provides an interesting probe. The PV asymmetry in the process was measured at ILL [21], and the result was

$$A_\gamma^t = (4.2 \pm 3.8) \times 10^{-6}.$$

Theoretical calculation of A_γ^t in Ref. [22] adopted DDH potential for the weak interaction, and examined the dependence on the strong interaction models such as de Tourreil-Sprung (TS) and Reid soft core (RSC). The results are interesting in the following aspects. First, dependence on the strong interaction model is non-negligible; TS model gives a result $A_\gamma^t(\text{TS}) = 0.81 \times 10^{-6}$ while RSC gives $A_\gamma^t(\text{RSC}) = 0.61 \times 10^{-6}$. Second, isoscalar, isovector, and isotensor parts of the weak interaction in the DDH potential give similar contribution to A_γ^t (i.e., 0.40, 0.45, and −0.04 in units of 10^{-6}, respectively) with the TS model. This means that the contributions from the π-, ρ-, and ω- exchanges in the DDH potential are similar. Non-negligible dependence on the strong interaction model, and the significant contribution from the heavy mesons in the weak interaction are the features common with the PV polarization in $np \to d\gamma$ [6]. In this problem, therefore, pionless EFT will provide us with a natural and systematic way to parametrize the parity mixing in the few-body system due to weak interactions. Pionless EFT with dibaryon fields has recently been applied to the EM transitions in $nd \to t\gamma$ [23,24], and the results agree to data to a good accuracy. Now the pionless EFT with dibaryon fields well accounts for the strong and the electromagnetic interactions in the two- and three-nucleon systems. With the formalism established thus far, therefore, we may be able to make a self-contained prediction for the PV asymmetry in $\vec{n}d \to t\gamma$.

Parity violation in the few-body system may show us effects that are not accessible in the two-nucleon systems. For instance, strong three-nucleon force can give non-negligible correction to the one- and two-body contributions to the PV observables. There has been no consideration on the weak three-nucleon force, but we have recently obtained nonzero components of the weak three-nucleon force in a preliminary calculation [25]. Two- and three-body PV meson-exchange currents may be important issues, too. We expect that the EFT will play a vital role in the understanding of the weak interaction in the few-body system.

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