LAWRENCE
LIVERMORE
NATIONAL
LABORATORY

LLNL-JRNL-402812

Nuclear Electric Dipole Moment of 3He

I. Stetcu, C. P.Liu, J. L. Friar, A. C. Hayes, Petr Navratil

April 9, 2008

Physics Letters B

Nuclear Electric Dipole Moment of ^{3}He

I. Stetcu [a], C.P. Liu [a,1], J. L. Friar [a], A. C. Hayes [a], P. Navrátil [b]

[a] *Theoretical Division, Los Alamos National Laboratory, Los Alamos, New Mexico 87545, U.S.A*

[b] *Lawrence Livermore National Laboratory, P.O. Box 808, L-414, California 94551, U.S.A*

Abstract

A permanent electric dipole moment (EDM) of a physical system would require time-reversal (T) violation, which is equivalent to charge-conjugation-parity (CP) violation by CPT invariance. Experimental programs are currently pushing the limits on EDMs in atoms, nuclei, and the neutron to regimes of fundamental theoretical interest. Nuclear EDMs can be studied at ion storage rings with sensitivities that may be competitive with atomic and neutron measurements. Here we calculate the magnitude of the CP-violating EDM of ^{3}He and the expected sensitivity of such a measurement to the underlying CP-violating interactions. Assuming that the coupling constants are of comparable magnitude for π-, ρ-, and ω-exchanges, we find that the pion-exchange contribution dominates. Finally, our results suggest that a measurement of the ^{3}He EDM is complementary to the planned neutron and deuteron experiments, and could provide a powerful constraint for the theoretical models of the pion-nucleon P,T-violating interaction.

1. Introduction

A permanent electric dipole moment (EDM) of a physical system would indicate direct violation of time-reversal (T) and parity (P) and thus CP violation through the CPT invariance. Presently there are several experimental programs pushing the limits on EDMs in atoms, nuclei, and the neutron to regimes of fundamental interest. The Standard Model (SM) predicts values for the EDMs of these systems that are too small to be detected in the foreseeable future, and hence a measured nonzero EDM in any of these systems is an unambiguous signal for a new source of CP violation and for physics beyond the SM. A new experimental scheme [1,2,3,4,5] for measuring EDMs of nuclei (stripped of their atomic electrons) in a magnetic storage ring suggests that the EDM of deuteron could be measured to an accuracy of better that $10^{-27}\,e\,$cm [4]. Unlike searches for CP-violating moments of the nucleus through measurements of atomic EDMs, a measurement for a stripped nucleus would not suffer from a suppression of the signal through atomic Schiff screening [6]. For this reason, the latter could represent about an order of magnitude better sensitivity to the underlying CP-violating interaction than the present limit on the neutron EDM, d_n [2]. Measurements using stripped nuclei in a magnetic storage ring are best suited to nuclei with small magnetic anomaly, making ^{3}He an ideal candidate for a high precision measurement. Here we examine the nuclear structure issues determining the EDM of ^{3}He and calculate the matrix elements of the relevant operators using the no-core shell model [7,8] and Podolsky's method for implementing second-order perturbation theory [9,10].

2. Sources of Nuclear *P*-,*T*-violation

A nuclear EDM consists of contributions from the following sources: (i) the intrinsic EDMs of the proton and neutron, d_p and d_n; (ii) the polarization effect caused by the P-,T-violating ($\not P \not T$) nuclear interaction, $H_{\not P \not T}$; (iii) the $\not P \not T$ meson-exchange charge operator appropriate for $H_{\not P \not T}$.

The contribution due to nucleon EDMs, $D^{(1)}$, which is purely one-body, can be easily evaluated by taking the matrix element

$$D^{(1)} = \langle\psi| \sum_{i=1}^{A} \frac{1}{2}\,\left[(d_p + d_n) + (d_p - d_n)\,\tau_i^z\right]\,\sigma_i^z|\psi\rangle\ , \quad (1)$$

where $|\psi\rangle$ is the nuclear state that has the maximal magnetic quantum number. In the particular case of interest in this paper, $|\psi\rangle = |0\rangle$ is the ground state of ^{3}He obtained by the diagonalization of the P,T-conserving interaction.

In perturbation theory, $H_{\not P \not T}$ induces a parity admixture to the nuclear state

$$\widetilde{|0\rangle} = \sum_{n\neq 0} \frac{1}{E_0 - E_n}\ |n\rangle\,\langle n|H_{\not P \not T}|0\rangle\ , \quad (2)$$

where $|n\rangle$ are eigenstates of energy E_n and opposite parity from $|0\rangle$, calculated with the P-,T-conserving Hamiltonian.

[1] Present Address: Department of Physics, University of Wisconsin-Madison, 1150 University Avenue, Madison, WI 53706-1390.

Hence, the polarization contribution $D^{(pol)}$ can be simply calculated as

$$D^{(pol)} = \langle 0|\ \hat{D}_z \widetilde{|0\rangle} + \text{c.c.}\,, \tag{3}$$

where

$$\hat{D}_z = \frac{e}{2} \sum_{i=1}^{A} (1+\tau_i^z)\, z_i \tag{4}$$

is the usual dipole operator projected in the z-direction.

The contribution due to exchange charge, $D^{(ex)}$, is typically at the order of $(v/c)^2$, and explicitly evaluated to be just a few percent of the polarization contribution for the deuteron case [11]; therefore, we ignore it and approximate the full two-body contribution, $D^{(2)}$, solely by the polarization term

$$D^{(2)} = D^{(pol)} + D^{(ex)} \cong D^{(pol)}\,. \tag{5}$$

The interaction $H_{\not P \not T}$ is conventionally formulated in a one-meson-exchange model. Including mesons with mass lower than 1 GeV, i.e., π, η, ρ, and ω, the full interaction is given as (see Refs. [11,12,13,14,15,16]):

$$\begin{aligned}
&H_{\not P \not T}(\boldsymbol{r}) \\
&= \frac{1}{2\,m_N} \Big\{ \boldsymbol{\sigma}_- \cdot \boldsymbol{\nabla} \left(\bar{G}^0_\eta\, y_\eta(r) - \bar{G}^0_\omega\, y_\omega(r) \right) \\
&+ \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2\, \boldsymbol{\sigma}_- \cdot \boldsymbol{\nabla} \left(\bar{G}^0_\pi\, y_\pi(r) - \bar{G}^0_\rho\, y_\rho(r) \right) \\
&+ \frac{\tau^z_+}{2}\, \boldsymbol{\sigma}_- \cdot \boldsymbol{\nabla} \left(\bar{G}^1_\pi\, y_\pi(r) - \bar{G}^1_\eta\, y_\eta(r) - \bar{G}^1_\rho\, y_\rho(r) - \bar{G}^1_\omega\, y_\omega(r) \right) \\
&+ \frac{\tau^z_-}{2}\, \boldsymbol{\sigma}_+ \cdot \boldsymbol{\nabla} \left(\bar{G}^1_\pi\, y_\pi(r) + \bar{G}^1_\eta\, y_\eta(r) + \bar{G}^1_\rho\, y_\rho(r) - \bar{G}^1_\omega\, y_\omega(r) \right) \\
&+ (3\,\tau^z_1\,\tau^z_2 - \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2)\, \boldsymbol{\sigma}_- \cdot \boldsymbol{\nabla} \left(\bar{G}^2_\pi\, y_\pi(r) - \bar{G}^2_\rho\, y_\rho(r) \right) \Big\}\,,
\end{aligned} \tag{6}$$

where $\bar{G}^T_x$ is defined as the product of a $\not P \not T$ meson–nucleon coupling $\bar{g}^T_x$, with T referring to the isospin, and its associated strong one, g_{xNN}. For example, $\bar{G}^0_\pi = g_{\pi NN}\, \bar{g}^0_\pi$ and $y_x(r) = e^{-m_x\, r}/(4\,\pi\, r)$ is the Yukawa function with a range determined by the mass of the exchanged x-meson.

In this work we mainly concentrate on pion exchange, which is long-ranged and therefore makes the dominant contribution [17] for the weakly bound ^{3}He nucleus. We nevertheless also include ρ- and ω-exchanges in our calculation. While a $H_{\not P \not T}$ in effective-field-theory (EFT) framework is still under development [18], the phenomenological ρ- and ω-exchanges should give some idea about the short-range part of $H_{\not P \not T}$ in EFT, which at the lowest order also anticipates five low-energy constants in order to characterize all possible $\not P \not T$ S–P transitions. [2]

By means of reliable hadronic calculations one should be able, in principle, to express the nucleon EDMs (and the $\not P \not T$ meson–nucleon couplings, as well) in terms of the underlying $\not P \not T$ parameters, including the QCD $\bar{\theta}$ term, etc., at the quark-gluon level. On the other hand, using the baryon-meson picture one can in fact estimate the nucleon EDMs through loop diagrams (see, for example, Refs. [21,22,23]) by using the same chiral Lagrangian that leads to $H_{\not P \not T}$ via meson exchange. For the later part of this paper we will rely on this picture whenever a nucleon EDM is to be expressed in terms of the $\not P \not T$ meson–nucleon couplings.

3. ^{3}He in the *ab initio* No-Core Shell Model

We solve the three-body problem in an *ab initio* no-core shell model (NCSM) framework [7,8]. The ground-state wave function is obtained by a direct diagonalization of an effective Hamiltonian in a truncated harmonic oscillator (HO) basis constructed in relative coordinates, as described in Ref. [24]. High-precision two-nucleon (NN) interactions, such as the local Argonne v_{18} [25,26] and the non-local charge-dependent (CD) Bonn potential [27] interactions, are used to derive an effective interaction in each model space via a unitary transformation [28,29,30] in a two-body cluster approximation. The Coulomb interaction between protons is also taken into account.

In addition to the phenomenological NN interaction models cited above, we consider two- and three-body interactions derived from EFT. In a recent work [31] using the *ab initio* NCSM, the presently available NN potential at N^{3}LO [32] and the three-nucleon (NNN) interaction at N^{2}LO [33,34] have been applied to the calculation of various properties of s- and p-shell nuclei. In that study a preferred choice of the two NNN low energy constants, c_D and c_E, was found (and the fundamental importance of the chiral EFT NNN interaction was demonstrated) by reproducing the structure of mid-p-shell nuclei. (Note that these interactions are fitted only for a momentum cutoff of 500 MeV, and therefore we are not able at this time to demonstrate a running of the observables with the cutoff.) This Hamiltonian was then used to calculate microscopically the photo-absorption cross section of ^{4}He [35], while the full technical details on the local chiral EFT NNN interaction that was used were given in Ref. [36]. We use an identical Hamiltonian in the present work, and we compare its predictions against the phenomenological potentials.

In the NCSM the basis states are constructed using HO wave functions. Hence, all the calculations involve two parameters: the HO frequency Ω and N_{max}, the number of oscillator quanta included in the calculation. At large enough N_{max}, the results become independent of the frequency, although the rate of convergence depends on Ω. Thus, for short-range operators, one can expect a faster convergence for larger values of Ω, as the characteristic length of the HO is $b = 1/\sqrt{m_N\,\Omega}$. The convergence also depends upon the P-,T-conserving interaction, H_0, used to solve the three-body problem. Thus the results obtained with Argonne v_{18} show the slowest convergence, because the NN interaction has a more strongly repulsive core than the interactions obtained from EFT, which have faster convergence rates.

[2] For more justification, see the established development in applying EFT to the parity-violating (but time-reversal-conserving) sector, Refs. [19,20].

The nucleonic contribution $D^{(1)}$ in Eq. (1) involves only H_0, and is easily calculated once the three-body problem is solved. We therefore concentrate on the part involving $H_{\not P \not T}$.

Equation (2) suggests that in order to calculate the dipole moment one needs to obtain high-accuracy excited states of ^{3}He in the continuum, an extremely difficult task in a NCSM framework, where the basis states are constructed using bound-state wave functions. The most straightforward technique for evaluating Eq. (2) is to use Podolsky's method [9,10], in which $\widetilde{|0\rangle}$ is obtained as the solution of the Schrödinger equation with an inhomogeneous term

$$(E_0 - H_0)\,\widetilde{|0\rangle} = H_{\not P \not T}\;|0\rangle\;. \tag{7}$$

The exceptionally nice feature of this method is that continuum states do not have to be *explicitly* calculated (they are, of course, implicitly included). In this sense the technique is a relatively simple extension of bound-state methods, which have been well studied and are robust. Moreover, in this approach the convergence of the EDM reduces to a large degree to the issue of the convergence of the ground state.

We express the solution of Eq. (7) as a superposition of a handful of vectors generated using the Lanczos algorithm [37,38]. Indeed, one can show that if we start with the inhomogeneous part of Eq. (7) as the starting Lanczos vector $|v_1\rangle = H_{\not P \not T}\,|0\rangle$, the solution becomes [39]

$$\widetilde{|0\rangle} \approx \sum_k g_k(E_0)\,|v_k\rangle\,, \tag{8}$$

where the summation over the index k runs over a finite and usually small number of iterations. The coefficients $g_k(E)$ are easily obtained using finite continued fractions [40].

We alter this approach in practice for efficiency reasons. Because Eq. (3) is symmetrical in $\hat{D}_z$ and $H_{\not P \not T}$, we are free to choose $|v_1\rangle = \hat{D}_z\,|0\rangle$ as the starting vector. This allows us to isolate the two isospin contributions for $H_{\not P \not T}$ in each run. Once we compute a second vector, $|v\rangle = H^\dagger_{\not P \not T}\,|0\rangle$, the polarization contribution to the EDM is finally evaluated as

$$D^{(pol)} = 2\sum_k g_k(E_0)\,\langle v|v_k\rangle\,. \tag{9}$$

(We have verified that the altered approach gives the same results as the original one.) As a particular test case we have considered the electric polarizability

$$\alpha_E = \frac{1}{2\pi^2}\int d\omega\,\frac{\sigma(\omega)}{\omega^2} = 2\,\alpha\sum_n \frac{\langle 0|\hat{D}_z|n\rangle\langle n|\hat{D}_z|0\rangle}{E_n - E_0} \tag{10}$$

(where α is the fine structure constant), which reduces Eq. (9) to $\alpha_E = -2\,\alpha\,g_1(E_0)\,\langle v_1|v_1\rangle$. We estimate that the electric polarizability of the ^{3}He nucleus is 0.183 fm^3 for the Argonne v_{18} potential, compared with 0.159 fm^3 reported in Ref. [41] for the same interaction. The 15% discrepancy is most likely the result of a difference in the theoretical approaches, as the result reported in Ref. [41] involves a matching of the ground-state energy to experiment (i.e., 7.72 MeV), although the calculation gives 6.88 MeV binding [42] in the absence of three-body forces (our converged binding energy for Argonne v_{18} is 6.92 MeV). Since the electric polarizability scales roughly with the inverse of the square of the binding energy, the discrepancy between the two results is reasonable. Moreover, we have made the additional check of the Levinger-Bethe sum rule [43], which in the case of tritium relates the total dipole strength to the charge radius, and we found it to be satisfied in all model spaces to a precision better than 10^{-5}. Finally, the ^{3}He polarizability calculated using the two- and three-body chiral interactions is 0.148 fm^3, compared with 0.145 fm^3 with Argonne v_{18} and Urbana IX [41] two- and three-body forces. In both cases excellent agreement with the experimental binding energy is achieved.

In a consistent approach the same transformation used to obtain the effective interaction should be used to construct the effective operators in truncated spaces. While this has been done in the past for general one- and two-body operators [44,45], such an approach is very cumbersome for the present investigation because both the dipole transition operator and $H_{\not P \not T}$ change the parity of the states. We have therefore chosen not to renormalize the operators involved, except for the P-,T-conserving Hamiltonian. This problem is largely overcome by the fact that long-range operators (like the dipole) have been found to be insensitive to the renormalization in the two-body cluster approximation [44,45], which is the level of truncation for the effective interaction. We have to point out that since $H_{\not P \not T}$ has short range, one can expect that the renormalization of this operator would improve the convergence pattern, especially for small frequencies. As with all operators, the effect of the renormalization decreases as the size of the model space increases, so that in large model spaces (like the ones in the present calculation) this effect can be safely neglected and good convergence of observables achieved.

4. Results and Discussions

We start the discussion of our results with the one-body contribution to the EDM of ^{3}He. In Table 1, we summarize the isoscalar and isovector contribution to $D^{(1)}$, decomposed into their respective coupling constants ($d_p + d_n$ for isoscalar, and $d_p - d_n$ for isovector).

Table 1
The nucleonic contribution (in e fm) to the ^{3}He EDM for different potential models. We decompose our results into isoscalar ($d_p + d_n$) and isovector ($d_p - d_n$) contributions.

	CD Bonn	v_{18}	EFT	
			NN	NN+NNN
$d_p + d_n$	0.430	0.415	0.437	0.433
$d_p - d_n$	−0.467	−0.462	−0.468	−0.468

All interactions give similar results, with only the Argonne v_{18} result deviating more significantly from the others, albeit by less than 6%. Table 1 also allows us an indirect check of our calculation. Since ^{3}He is mostly an $S_{1/2}$ state, its magnetic moment can be approximated similarly to Eq. (1), with the proton and neutron dipole moments d_p and d_n replaced by g_p and g_n, respectively, with $g_p = 2.79$ and $g_n = -1.91$ being the proton and neutron gyromagnetic ratios. We estimate that the magnetic moment is $0.90\,g_n - 0.04\,g_p = -1.83$ for chiral two- and three-body interactions, which compares reasonably well with the experimental value of $-2.12762485(7)$; the 15% difference can be explained by the relatively small admixture of S' and D states and the missing meson-exchange currents. Furthermore, the results agree quite well with a simple single-particle model, in which ^{3}He is described by a valence neutron hole. In such a simplified description, the magnetic and one-body EDM are g_n and d_n, respectively, in close agreement with $0.90\,g_n - 0.04\,g_p$ and $0.90\,d_n - 0.04\,d_p$, obtained from Table 1.

In Fig. 1 we present for four HO frequencies the running with N_{max} of the EDM induced by the pion-exchange part of $H_{\not P \not T}$. Two- and three-body EFT interactions have been used for this calculation, in order to obtain an accurate description of the ground-state of the three-body system. For the nuclear EDM we mix two types of operators: $H_{\not P \not T}$, which is short range, and $\hat{D}_z$, which is long range. The convergence pattern is therefore not as straightforward as discussed for the electric polarizability. However, the short-range part dominates the convergence pattern, and we thus observe faster convergence for larger frequencies (smaller HO parameter length). This behavior is opposite to the convergence in the case of the electric polarizability discussed above. While we do not show the convergence of α_E, we found faster convergence for smaller frequencies. Nevertheless, just as in Fig. 1, the results become independent of the frequency at large N_{max}. Note in the insert the convergence behavior of the ground-state energy, which converges to the experimental value already at $N_{max} \approx 22$ for most frequencies presented in the figure.

Similar convergence patterns can be observed for the other meson exchanges as well as other potential models. In Table 2 we summarize these results.

For pion exchange, all potential models give basically the same result, as the long-range part ($r \gtrsim 1/m_\pi$) of the ^{3}He wave function shows negligible model dependence. It is interesting to note the effect of the three-body force by comparing the results with and without NNN interactions. When only the NN EFT interaction is used, the binding energy is underestimated by about 500 keV. Therefore, since the ground-state energy is in the denominator of Eq. (2), one could naively expect that introducing the three-body forces (which increase the binding) decreases $D^{(pol)}$. Instead we obtain nearly the same results for both isoscalar and isovector contributions. This implies that the NNN interaction reshuffles the strength to compensate for the change in binding energy, most likely at low energies. This is not a surprise, because it was already found previously

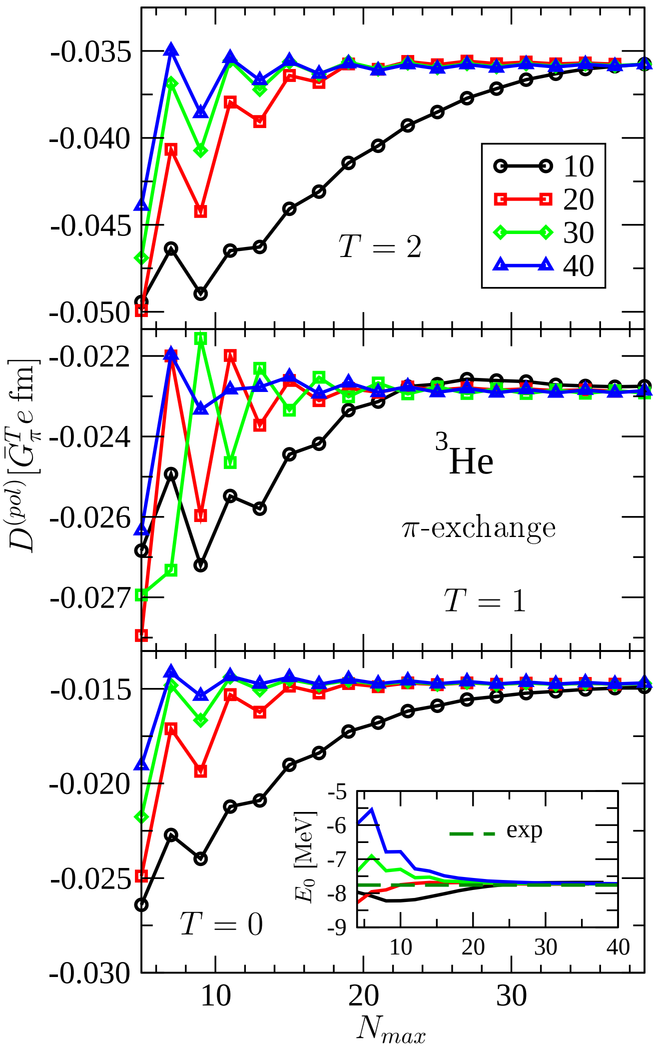

Figure 1. Isoscalar, isovector, and isotensor pion-exchange contributions to the $H_{\not P \not T}$-induced EDM in ^{3}He. We show four different frequencies in each case: $\Omega = 10$ MeV (circles), $\Omega = 20$ MeV (squares), $\Omega = 30$ MeV (diamonds), and $\Omega = 40$ MeV (triangles). In the insert, we present the convergence of the ground-state energy, which in the limit of large N_{max} approaches the experimental value (dashed line). Both NN and NNN EFT interactions have been used for diagonalization.

that the main effect of the three-body forces for the dipole response is an attenuation of the peak region at low energies both in the three- [41] and four-body [35] systems.

For ρ- and ω-exchanges one immediately sees that their corresponding coefficients are at most 10% of the pion-exchange ones, because only the short-range wave function ($r \lesssim 1/m_{\rho,\omega}$) contribute substantially. Also because of their short-range sensitivity, the model dependence becomes more pronounced, which can be as large as 50% for some channels. While a detailed explanation for the model dependence is too intricate to be disentangled, one can roughly see the general trend that the calculation using Argonne v_{18} gives consistently smaller results than ones using CD Bonn and chiral EFT, as Argonne v_{18} has a harder core than the other two. The best way to discuss the short-range interaction is using the language of EFT, and at the low-

Table 2
The polarization contribution to ^{3}He EDM (in units of e fm), decomposed in $\bar{G}_x^T$, where x stands for π, ρ, or ω meson exchanges.

	π				ρ				ω			
	CD Bonn	v_{18}	EFT		CD Bonn	v_{18}	EFT		CD Bonn	v_{18}	EFT	
			NN	NN+NNN			NN	NN+NNN			NN	NN+NNN
$\bar{G}_x^0$	−0.013	−0.012	−0.015	−0.015	0.0012	0.0006	0.0012	0.0013	−0.0008	−0.0005	−0.0009	−0.0007
$\bar{G}_x^1$	−0.022	−0.022	−0.023	−0.023	−0.0011	−0.0009	−0.0013	−0.0012	0.0011	0.0011	0.0017	0.0018
$\bar{G}_x^2$	−0.035	−0.034	−0.037	−0.036	0.0019	0.0015	0.0028	0.0027	-	-	-	-

est order one expects five low-energy constants to be determined phenomenologically. In fact, one can simulate this framework by treating $\bar{G}_\rho^{0,1,2}$ and $\bar{G}_\omega^{0,1}$ as free parameters in the one-meson-exchange picture. As it is very unlikely that we will have enough precise EDM measurements to determine these parameters in the near future, we assume naturalness, which requires that all $\bar{G}_{\pi,\rho,\omega}$ have similar magnitude. By this assumption, the pion-exchange has the dominant contributions to nuclear EDMs, with the heavy-meson exchanges (the short-range interaction) giving roughly a 10% correction to the pion-exchange contribution.

Assuming the dominance of pion exchange, $D^{(2)}$ has an almost model-independent expression

$$D^{(2)} \approx (-0.015\,\bar{G}_\pi^0 - 0.023\,\bar{G}_\pi^1 - 0.036\,\bar{G}_\pi^2)\,e\,\mathrm{fm}. \qquad (11)$$

The single-nucleon EDMs can be estimated using the non-analytic term that results from calculating the one-pion loop diagram, which dominates in the chiral limit (see, for example, [11,21,22,23])

$$d_{p,n} \approx \mp \frac{e}{4\,\pi^2\,m_N}\,(\bar{G}_\pi^0 - \bar{G}_\pi^2)\,\ln\left(\frac{m_N}{m_\pi}\right)\,. \qquad (12)$$

Folding this result into $D^{(1)}$ and using the physical nucleon and pion masses, $\ln(m_N/m_\pi) \approx 1.90$, we get

$$D^{(1)} \approx 0.009\,(\bar{G}_\pi^0 - \bar{G}_\pi^2)\,e\,\mathrm{fm}\,. \qquad (13)$$

Therefore, the total EDM of ^{3}He is estimated to be

$$\begin{aligned} D &= D^{(1)} + D^{(2)} \\ &= (-0.006\,\bar{G}_\pi^0 - 0.023\,\bar{G}_\pi^1 - 0.045\bar{G}_\pi^2)\,e\,\mathrm{fm}\,. \end{aligned} \qquad (14)$$

Expressing the EDMs of the neutron and deuteron also by Eq. (12) and pion-exchange dominance, one can see from Table 3 that an EDM measurement in ^{3}He is very complementary to the former two. Supposing that similar sensitivities can be reached in these three measurements, the $\not{P}\not{T}$ pion–nucleon coupling constants could be well-constrained if the assumption of pion-exchange dominance is not too far off.

5. Summary

In this paper, we have calculated the nuclear EDM of ^{3}He, which arises from the intrinsic electric EDMs of nucleons and the P-,T-violating nucleon-nucleon interaction.

Table 3
The EDMs (in units of e fm) of neutron, deuteron, and ^{3}He decomposed into contributions proportional to $\bar{G}_\pi^{0,1,2}$, while assuming the dominance of pion-exchange forces in $H_{\not{P}\not{T}}$ and estimating nucleon EDMs via pion loops.

	$\bar{G}_\pi^0$	$\bar{G}_\pi^1$	$\bar{G}_\pi^2$
neutron	0.010	0.000	−0.010
deuteron	0.000	0.015	0.000
^{3}He	−0.006	−0.023	−0.045

Several potential models for the P-,T-conserving nuclear interaction (including the latest-generation NN and NNN chiral EFT forces) have been used in order to obtain the solution to the nuclear three-body problem. The results obtained with these potential models agree within 2% in the $\not{P}\not{T}$ pion-exchange sector. Though larger spreads in $\not{P}\not{T}$ ρ- and ω-exchanges are found (as the results sensitively depend on the wave functions at short range), we expect them to be non-essential as pion-exchange dominates the observable – unless the $\not{P}\not{T}$ parameters associated with heavy-meson exchanges are unnaturally much larger than the ones with pion exchange. We further demonstrate that a measurement of ^{3}He EDM would be complementary to those of the neutron and deuteron, and in combination they can be used to put stringent constraints on the three P-,T-violating pion–nucleon coupling constants. We therefore strongly encourage experimentalists to consider such a ^{3}He measurement in a storage ring, in addition to the existing deuteron proposal [4].

Acknowledgements

I.S. thanks W. Leidemann, S. Quaglioni and S. Bacca for useful discussions. The work of I.S., C.P.L., A.C.H. and J.L.F. was performed under the auspices of the U. S. DOE. Prepared by LLNL under Contract DE-AC52-07NA27344. Support from the U.S. DOE/SC/NP (Work Proposal No. SCW0498), and from the U. S. Department of Energy Grant DE-FG02-87ER40371 is acknowledged.

References

[1] I. B. Khriplovich, Phys. Lett. B 444 (1998) 98.
[2] I. B. Khriplovich, Hyperfine Interact. 127 (2000) 365.
[3] F. J. M. Farley, et al., Phys. Rev. Lett. 93 (2004) 052001.

[4] Y. K. Semertzidis, et al., AIP Conf. Proc. 698 (2004) 200.
[5] Y. F. Orlov, W. M. Morse, Y. K. Semertzidis, Phys. Rev. Lett. 96 (2006) 214802.
[6] L. I. Schiff, Phys. Rev. 132 (1963) 2194.
[7] P. Navrátil, J. P. Vary, B. R. Barrett, Phys. Rev. Lett. 84 (2000) 5728.
[8] P. Navrátil, J. P. Vary, B. R. Barrett, Phys. Rev. C 62 (2000) 054311.
[9] B. Podolsky, Proc. Nat. Acad. Sci. U.S.A. 14 (1928) 253.
[10] J. L. Friar, G. L. Payne, Phys. Rev. C 72 (2005) 014002.
[11] C.-P. Liu, R. G. E. Timmermans, Phys. Rev. C 70 (2004) 055501.
[12] G. Barton, Nuovo Cimento 19 (1961) 512.
[13] W. C. Haxton, E. M. Henley, Phys. Rev. Lett. 51 (1983) 1937.
[14] P. Herczeg, in: N. R. Roberson, C. R. Gould, J. D. Bowman (Eds.), Tests of Time Reversal Invariance in Neutron Physics, World Scientific, Singapore, 1987, pp. 24–53.
[15] V. P. Gudkov, X.-G. He, B. H. J. McKellar, Phys. Rev. C 47 (1993) 2365.
[16] I. S. Towner, A. C. Hayes, Phys. Rev. C 49 (1994) 2391.
[17] J. L. Friar, G. L. Payne, Few-Body Systems 19 (1995) 203.
[18] M. J. Ramsey-Musolf, U. van Kolck, private communication.
[19] S.-L. Zhu, C. M. Maekawa, B. R. Holstein, M. J. Ramsey-Musolf, U. van Kolck, Nucl. Phys. A 748 (2005) 435.
[20] C.-P. Liu, Phys. Rev. C 75 (2007) 065501.
[21] G. Barton, E. D. White, Phys. Rev. 184 (1969) 1660–1662.
[22] R. J. Crewther, P. Di Vecchia, G. Veneziano, E. Witten, Phys. Lett. 88B (1979) 123, 91B, 487(E) (1980).
[23] X.-G. He, B. H. J. McKellar, S. Pakvasa, Int. J. Mod. Phys. A 4 (1989) 5011, 6, 1063(E) (1991).
[24] P. Navrátil, G. P. Kamuntavicius, B. R. Barrett, Phys. Rev. C 61 (2000) 044001.
[25] R. B. Wiringa, V. G. J. Stocks, R. Schiavilla, Phys. Rev. C 51 (1995) 38.
[26] S. Pieper, R. B. Wiringa, Annu. Rev. Nucl. Part. Sci. 51 (2001) 53.
[27] R. Machleidt, F. Sammarruca, Y. Song, Phys. Rev. C 53 (1996) 1483.
[28] S. Okubo, Prog. Theor. Phys. 12 (1954) 603.
[29] J. Da Providencia, C. M. Shakin, Ann. of Phys. 30 (1964) 95.
[30] K. Suzuki, S. Lee, Prog. Theor. Phys. 64 (1980) 2091.
[31] P. P. Navrátil, V. G. Gueorguiev, J. P. Vary, W. E. Ormand, A. Nogga, Phys. Rev. Lett. 99 (2007) 042501.
[32] D. R. Entem, R. Machleidt, Phys. Rev. C (2003) 041001(R).
[33] U. van Kolck, Phys. Rev. C 49 (1994) 2932.
[34] E. Epelbaum, A. Nogga, W. Glöckle, H. Kamada, U.-G. Meissner, H. Witala, Phys. Rev. C 66 (2002) 064001.
[35] S. Quaglioni, P. Navrátil, Phys. Lett. B652 (2007) 370.
[36] P. Navrátil, Few Body Syst. 41 (2007) 117.
[37] C. Lanczos, J. Res. Natl. Bur. Stand. 45 (1950) 255.
[38] J. H. Wilkinson, The Algebraic Eigenvalue Problem, Clarendon, Oxford, 1997.
[39] R. Haydock, J. Phys. A 7 (1974) 2120.
[40] M. A. Marchisio, N. Barnea, W. Leidemann, G. Orlandini, Few Body Syst. 33 (2003) 259.
[41] W. Leidemann, Few Body Syst. Suppl. 14 (2003) 313.
[42] W. Leidemann, private communication.
[43] J. S. Levinger, H. A. Bethe, Phys. Rev. 78 (1950) 115.
[44] I. Stetcu, B. R. Barrett, P. Navrátil, J. P. Vary, Phys. Rev. C 71 (2005) 044325.
[45] I. Stetcu, B. R. Barrett, P. Navrátil, J. P. Vary, Phys. Rev. C 73 (2006) 037307.