

## Electric dipole moment of $^{13}\text{C}$

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We calculate for the first time the electric dipole moment (EDM) of  $^{13}\text{C}$  generated by the isovector charge conjugation-parity (CP)-odd pion exchange nuclear force in the  $\alpha$ -cluster model, which describes well the structures of low-lying states of the  $^{13}\text{C}$  nucleus. The linear dependence of the EDM of  $^{13}\text{C}$  on the neutron EDM and the isovector CP-odd nuclear coupling is found to be  $d_{^{13}\text{C}} = -0.33d_n - 0.0020\bar{G}_\pi^{(1)}$ . The linear enhancement factor of the CP-odd nuclear coupling is smaller than that of the deuteron, due to the difference of the structure between the  $1/2^-$  state and the opposite-parity ( $1/2^+$ ) states. We clarify the role of the structure played in the enhancement of the EDM. This result provides good guiding principles to search for other nuclei with large enhancement factor. We also mention the role of the EDM of  $^{13}\text{C}$  in determining the new physics beyond the standard model.

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### I. INTRODUCTION

One of the most important challenges in particle physics and cosmology is to explain the baryon number asymmetry of the universe. In the standard model, it is, however, known that the charge conjugation-parity (CP) violation is not sufficient to realize the currently observed matter abundance [1–5]. This actually gives a strong motivation to search for new physics beyond the standard model with large CP violation.

The most promising experimental observable to probe the CP violation of new physics is the *electric dipole moment* (EDM) [6–16]. The EDM is measurable in many systems, and many experimental measurements have been done so far, including the neutron [17], atoms [18–22], molecules [23,24], etc. Recently, new technology using storage rings is making it possible to measure the EDM of charged particles [25–33]. The experimental measurement of the nuclear EDM is also developed and is expected to be realized in the near future.

The study of the nuclear EDM has several notable advantages. The first advantage is the absence of electrons which screen the EDM of the nucleus, as dictated by the theorem of Schiff [34]. The second one is the small contribution from the standard model CP violation [35], which makes it an experimentally clean observable [36–41]. The final reason is the potential enhancement of the EDM by the nuclear many-body effect [42–44].

Due to those advantages, theoretical evaluations of the nuclear EDM have been extensively done so far [45–64]. The EDM of nuclei was first estimated in the simple core+valence model, and an enhancement of the CP violation was suggested [42–44]. For light nuclei, the applicability of this model is, however, questionable, since other correlations such as the cluster structure [65–67], which are thought to be important

in light nuclear systems, were not taken into account. In a recent work, the EDM of light nuclei was evaluated in the cluster model, and it was found that the cluster structure may enhance the nuclear EDM [63]. There the EDMs of  $^6\text{Li}$  and  $^9\text{Be}$  were calculated within the framework of  $\alpha + p + n$  and  $2\alpha + n$  three-body systems, respectively. The cluster model was successfully applied in the calculation of the  $^6\text{Li}$  and  $^9\text{Be}$  nuclei, and the result suggested an enhancement compared to the deuteron EDM, thanks to the cluster structure. This fact motivates us to theoretically evaluate the EDM of other unknown light nuclei with the expectation to find sensitive and experimentally advantageous nuclear systems.

The study of the EDM of multiple hadronic systems is not only essential in finding experimentally sensitive observables, but it is also absolutely necessary for constraining the hadron-level CP-violating interaction with many unknown couplings. Recent investigations of hadronic CP violation in the chiral effective field theory is indicating that some CP-violating interactions such as the three-pion interaction or the contact CP-odd nuclear force also receive contribution from the quark- and gluon-level CP violations in the leading order [55,62,68]. As the determination of unknown couplings requires at least the same number of experimental values of different observables, further study of the EDM of light nuclei is mandatory to unveil the CP violation beyond the standard model.

As mentioned, the EDM of light nuclei has been studied up to the three-body system in the  $\alpha$ -cluster model. The next step is then to study four-body systems. In this work, we therefore evaluate the EDM of  $^{13}\text{C}$ , which is the simplest one. The  $^{13}\text{C}$  nucleus was recently studied within the framework of  $3\alpha + n$  four-body cluster model [69]. The four-body problem can be solved by using the Gaussian expansion method [70,71], which was applied in many systems of a wide range of physical hierarchy [69,72–89]. In the study of  $^{13}\text{C}$ , its structure up to the excitation energy 15 MeV was successfully described, and

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it was argued that this nucleus does not have a simple shell structure, so a significant change of the nuclear EDM from the simple core+valence model prediction is expected. In this work, we will therefore calculate the nuclear EDM of  $^{13}\text{C}$  in the four-body cluster model using the Gaussian expansion method and investigate the effect of the cluster correlations on the EDM.

The plan of this paper is as follows. In the next section, we first give a brief overview of the calculation of the structure of  $^{13}\text{C}$ . In Sec. III, the formulation of the EDM and the CP-odd nuclear force is given. We then show the result of the calculation of the EDM of  $^{13}\text{C}$  in the Gaussian expansion method and analyze it. In Sec. V, we discuss the role of  $^{13}\text{C}$  EDM in the determination of a hadron-level CP violation. In Sec. VI, we show the prospects for the discovery of new physics beyond the standard model by measuring the EDM of  $^{13}\text{C}$ . The final section is devoted to the summary.

## II. STRUCTURE OF $^{13}\text{C}$ IN THE FOUR-BODY CLUSTER MODEL

Here we review the setup of the calculation of  $^{13}\text{C}$  and the CP-even nuclear force used within the framework of the  $3\alpha + N$  four-body cluster model [69]. The total wave function of  $^{13}\text{C}$   $\tilde{\Psi}_J(A=13)$ , with the total angular momentum  $J$  and total isospin  $T = \frac{1}{2}$  in the  $3\alpha + N$  cluster model, is expressed by the product of the internal wave functions of  $\alpha$  clusters  $\phi(\alpha)$  and the relative wave function  $\Psi_J(A=13)$  among the  $3\alpha$  clusters and the extra neutron,

$$\tilde{\Phi}_J(A=13) = \Phi_J(A=13)\phi(\alpha_1)\phi(\alpha_2)\phi(\alpha_3). \quad (1)$$

The relative wave function  $\Phi_J(A=13)$  is expanded in terms of the Gaussian basis as follows:

$$\Phi_J(A=13) = \sum_{p=1}^4 \sum_{c^{(p)}} \sum_{v^{(p)}} f_{c^{(p)}}^{(p)}(v^{(p)}) \Phi_{c^{(p)}}^{(p)}(v^{(p)}), \quad (2)$$

$$\Phi_{c^{(p)}}^{(p)}(v^{(p)}) = \mathcal{S}_\alpha \left[ [\varphi_{\ell_1^{(p)}}(\mathbf{R}_1^{(p)}, v_1^{(p)})][\varphi_{\ell_2^{(p)}}(\mathbf{R}_2^{(p)}, v_2^{(p)}) \right. \\ \left. \times \varphi_{\ell_3^{(p)}}(\mathbf{R}_3^{(p)}, v_3^{(p)}) \right]_{L_{23}^{(p)}} \left]_{L^{(p)}} \xi_{\frac{1}{2}}(N) \right]_J, \quad (3)$$

$$\langle u_F | \Phi_J(^{13}\text{C}) \rangle = 0, \quad (4)$$

where we assign the cluster number, 1, 2, and 3, to the three  $\alpha$  clusters (spin 0), and the number 4 to the extra nucleon (spin  $\frac{1}{2}$ ).  $\Phi_{c^{(p)}}^{(p)}[v^{(p)}]$  denotes the relative wave function with respect to the  $p$ th Jacobi coordinate system of the four-body  $3\alpha + N$  model (either of coordinate type of  $K$  or  $H$ ) shown in Fig. 1 of Ref. [69], where  $c^{(p)}$  represents the angular-momentum channel for the  $p$ th Jacobi coordinate system.  $\mathcal{S}_\alpha$  stands for the symmetrization operator acting on all  $\alpha$  particles obeying the Bose statistic, and  $\xi_{\frac{1}{2}}(N)$  is the spin function of the extra nucleon.  $\phi(\alpha)$  denotes the intrinsic wave function of the  $\alpha$  cluster with the  $(0s)^4$  shell-model configuration.  $v^{(p)}$  denotes the set of size parameters,  $v_1^{(p)}$ ,  $v_2^{(p)}$ , and  $v_3^{(p)}$ , of the normalized Gaussian function,  $\varphi_\ell(\mathbf{R}, v) = N_\ell(v) R^\ell \exp(-vR^2) Y_\ell(\hat{\mathbf{R}})$ , and  $v$  is taken to be of geometrical progression,

$$v_n = 1/b_n^2, \quad b_n = b_{\min} a^{n-1}, \quad n = 1 \sim n_{\max}. \quad (5)$$

It is noted that this prescription is found to be very useful in optimizing the ranges with a small number of free parameters  $(b_{\min}, a, n_{\max})$  with high accuracy [70,72]. In this work, the geometric series for each coordinate is made of seven terms, i.e.,  $7^3 = 343$  bases per angular -momentum channel. To converge the  $1/2_1^-$  and  $1/2_1^+$  states, we use 59 and 69 channels, respectively. The calculation of the EDM requires parity mixing, so 128 angular-momentum channels, i.e., 43 904 bases, are required.

The  $3\alpha + N$  Hamiltonian for  $\Phi_J(A=13)$  is presented as

$$\mathcal{H} = \sum_{i=1}^4 T_i - T_{\text{c.m.}} + \sum_{i<j=1}^3 V_{2\alpha}(i,j) + \sum_{i=1}^3 V_{\alpha N}(i,4) \\ + V_{3\alpha}(1,2,3) + \sum_{i<j=1}^3 V_{2\alpha N}(i,j,4) \\ + V_{3\alpha N}(1,2,3,4) + V_{\text{Pauli}}, \quad (6)$$

where  $T_i$ ,  $V_{2\alpha}$  ( $V_{\alpha N}$ ),  $V_{3\alpha}$  ( $V_{2\alpha N}$ ), and  $V_{3\alpha N}$  stand for the kinetic energy operator for the  $i$ th cluster,  $\alpha$ - $\alpha$  ( $\alpha$ - $N$ ) potential, three-body potential among the three  $\alpha$  particles (the two  $\alpha$  particles and extra nucleon), and the four-body potential, respectively. The center-of-mass kinetic energy ( $T_{\text{c.m.}}$ ) is subtracted from the Hamiltonian.

The effective  $\alpha$ - $\alpha$  potential  $V_{2\alpha}$  is constructed by the folding procedure from an effective two-nucleon force including the proton-proton Coulomb force. Here we take the Schmid-Wildermuth (SW) force [90] as the effective  $NN$  force. This folded  $\alpha$ - $\alpha$  potential reproduces nicely the  $\alpha$ - $\alpha$  scattering phase shifts and the energies of the  $^8\text{Be}$  ground-band state ( $J^\pi = 0^+ - 2^+ - 4^+$ ) within the framework of the  $2\alpha$  cluster model. As for the  $\alpha$ - $N$  potential, we use the Kanada-Kaneko potential [91], which reproduces the low-energy  $\alpha$ - $N$  scattering phase shifts. We introduce the phenomenological effective three-body forces,  $V_{3\alpha}$  and  $V_{2\alpha N}$ , the former depending on the total angular momentum of  $^{12}\text{C}$ . The  $3\alpha$  cluster model then describes well the structure of the low-lying states of  $^{12}\text{C}$  including the  $2_2^+$ ,  $0_3^+$ , and  $0_4^+$  states above the Hoyle state, which have quite recently been observed in experiments [92,93]. The structures of the low-lying states of  $^9\text{Be}$  and  $^9\text{B}$  are also well described by the  $2\alpha + N$  cluster model. The four-body interaction,  $V_{3\alpha N}$ , is given to reproduce the  $^{13}\text{C}$  ground-state energy.

Equation (4) represents the orthogonality condition that the total wave function (2) should be orthogonal to the Pauli-forbidden states of the  $3\alpha + N$  system,  $u_F$ 's, which are constructed from the Pauli forbidden states between two  $\alpha$  particles and those between  $\alpha$  particle and extra nucleon  $N$  [94–98]. The Pauli-forbidden states are removed by using the Pauli-blocking operator  $V_{\text{Pauli}}$  [99] in Eq. (6),

$$V_{\text{Pauli}} = \lim_{\lambda \rightarrow \infty} \lambda \sum_f |u_f\rangle \langle u_f|, \quad (7)$$

which rules out the Pauli-forbidden  $\alpha$ - $\alpha$  relative states ( $f = 0S, 1S, 0D$ ) and the Pauli-forbidden  $\alpha$ - $n$  relative state ( $f = 0S$ ) from the four-body  $3\alpha$ - $n$  wave function. In the present study, we take  $\lambda = 10^4$  MeV. The ground state of  $^{13}\text{C}$  with the

dominant shell-model-like configuration  $(0s)^4(0p)^9$ , then, can be properly described in the present  $3\alpha + N$  cluster model.

The equation of motion of  $^{13}\text{C}$  with the  $3\alpha + N$  four-body cluster model is obtained by the variational principle,

$$\delta[\langle \Phi_J(A=13) | \mathcal{H} - E | \Phi_J(A=13) \rangle] = 0, \quad (8)$$

where  $E$  denotes the eigenenergy of  $^{13}\text{C}$  measured from the  $3\alpha + n$  threshold. The energy  $E$  and the expansion coefficients  $f_{c^{(p)}}^{(p)}$  in the total wave function shown in Eq. (2) are determined by solving a secular equation derived from Eq. (8).

### III. CALCULATION OF THE NUCLEAR EDM

To induce the nuclear EDM, the existence of the P, CP-odd nucleon-level processes is required. In this work, we assume the following effective CP-odd Lagrangian [10,100,101]:

$$\begin{aligned} \mathcal{L}_{PT} = & -\frac{i}{2} \sum_{N=p,n} \bar{d}_N \bar{N} \sigma_{\mu\nu} \gamma_5 N F^{\mu\nu} \\ & + \sum_{N=p,n} \left[ \sum_{a=1}^3 \bar{g}_{\pi NN}^{(0)} \bar{N} \tau^a N \pi^a + \bar{g}_{\pi NN}^{(1)} \bar{N} N \pi^0 \right. \\ & \left. + \sum_{a=1}^3 \bar{g}_{\pi NN}^{(2)} (\bar{N} \tau^a N \pi^a - 3\bar{N} \tau^3 N \pi^0) \right]. \quad (9) \end{aligned}$$

Here the P, CP-odd coupling constants depend on QCD and elementary-level CP violation. In this work, we consider them as small and given. There are other CP-odd hadron-level effective interactions which contribute in the leading order of chiral perturbation theory, such as the three-pion interaction or the contact CP-odd  $NN$  interactions [68],

$$\begin{aligned} \mathcal{L}'_{PT} = & m_N \Delta_{3\pi} \pi^z \sum_{a=1}^3 \pi_a^2 + \bar{C}_1 \bar{N} N \partial_\mu (\bar{N} S^\mu N) \\ & + \sum_{a=1}^3 \bar{C}_2 \bar{N} \tau_a N \cdot \partial_\mu (\bar{N} S^\mu \tau_a N). \quad (10) \end{aligned}$$

The three-pion interaction (term with  $\Delta_{3\pi}$ ) is isovector, and the radiative correction is known to sizably contribute to  $\bar{g}_{\pi NN}^{(1)}$ . Here the contact interaction [terms with  $\bar{C}_1$  and  $\bar{C}_2$  of Eq. (10)] is isoscalar, and it receives a contribution from the  $\eta$  meson exchange. It is interesting, since it is also an important probe of the Weinberg operator [55]. In this work, we, however, do not consider it, since its effect suffers from large theoretical uncertainty in the nuclear-level calculation [58,59].

The physical nucleon EDMs  $d_n$  and  $d_p$  are not only due to the bare terms  $\bar{d}_n$  and  $\bar{d}_p$  but also receive contribution from the isoscalar CP-odd pion-nucleon interaction  $\bar{g}_{\pi NN}^{(0)}$ . In the leading order of chiral perturbation theory, it is given as [102]

$$d_N = \bar{d}_N - \tau_z \frac{e g_A \bar{g}_{\pi NN}^{(0)}}{4\pi^2 f_\pi} \ln \frac{\Lambda}{m_\pi}, \quad (11)$$

where  $\Lambda \approx 1$  GeV is the cutoff of the hadron-level effective theory and  $\tau_z = +1$  ( $-1$ ) for the proton (neutron). Here we neglect the effect of  $\bar{g}_{\pi NN}^{(1)}$  and  $\bar{g}_{\pi NN}^{(2)}$  which contributes at the higher order.

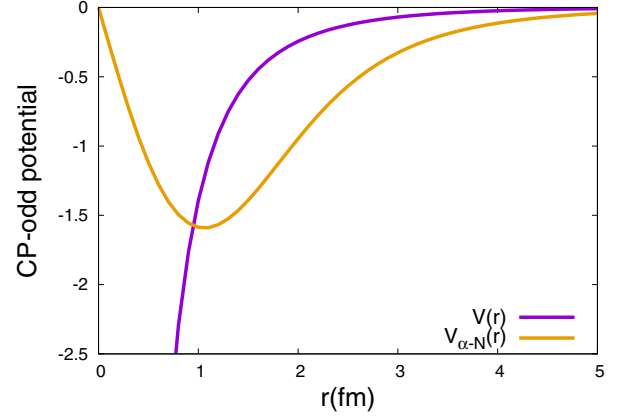


FIG. 1. The radial shape of the bare pion exchange CP-odd nuclear force  $V(r)$  and its folding potential  $V_{\alpha-N}(r)$ . The CP-odd coupling constant was factored out.

We now give the one-pion exchange CP-odd nuclear force. It has been studied and used in many previous works [48,50,54,56,59,63]. The CP-odd one-pion exchange nuclear force in the coordinate representation is [50,101,103,104]

$$\begin{aligned} H_{PT}^\pi = & \{ \bar{G}_\pi^{(0)} \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2 \boldsymbol{\sigma}_- + \frac{1}{2} \bar{G}_\pi^{(1)} (\tau_+^z \boldsymbol{\sigma}_- + \tau_-^z \boldsymbol{\sigma}_+) \\ & + \bar{G}_\pi^{(2)} (3\tau_1^z \tau_2^z - \boldsymbol{\tau}_1 \cdot \boldsymbol{\tau}_2) \boldsymbol{\sigma}_- \} \cdot \hat{\boldsymbol{r}} V(r), \quad (12) \end{aligned}$$

where  $\hat{\boldsymbol{r}} \equiv \frac{\boldsymbol{r}_1 - \boldsymbol{r}_2}{|\boldsymbol{r}_1 - \boldsymbol{r}_2|}$  is the unit vector. The spin and isospin notations are  $\boldsymbol{\sigma}_- \equiv \boldsymbol{\sigma}_1 - \boldsymbol{\sigma}_2$ ,  $\boldsymbol{\sigma}_+ \equiv \boldsymbol{\sigma}_1 + \boldsymbol{\sigma}_2$ ,  $\boldsymbol{\tau}_- \equiv \boldsymbol{\tau}_1 - \boldsymbol{\tau}_2$ , and  $\boldsymbol{\tau}_+ \equiv \boldsymbol{\tau}_1 + \boldsymbol{\tau}_2$ . The dimensionless CP-odd nuclear couplings  $\bar{G}_\pi^{(i)}$  ( $i = 0, 1, 2$ ) are given by

$$\bar{G}_\pi^{(0)} = -\frac{g_A m_N}{f_\pi} \bar{g}_{\pi NN}^{(0)}, \quad (13)$$

$$\bar{G}_\pi^{(1)} = -\frac{g_A m_N}{f_\pi} \bar{g}_{\pi NN}^{(1)}, \quad (14)$$

$$\bar{G}_\pi^{(2)} = \frac{g_A m_N}{f_\pi} \bar{g}_{\pi NN}^{(2)}, \quad (15)$$

in the leading order of chiral perturbation theory. The radial shape of the CP-odd  $NN$  potential is given by

$$V(r) = -\frac{m_\pi}{8\pi m_N} \frac{e^{-m_\pi r}}{r} \left( 1 + \frac{1}{m_\pi r} \right), \quad (16)$$

with the pion and nucleon masses  $m_\pi = 138$  MeV and  $m_N = 939$  MeV, respectively. The shape of the radial dependence of the CP-odd nuclear force is shown in Fig. 1.

To obtain the CP-odd  $N-\alpha$  potential, we use the folding by the density function of the  $\alpha$  cluster. It works as<sup>1</sup>

$$V_{\alpha-N}(R) \hat{\boldsymbol{R}} = \int d^3 \boldsymbol{R}' V(|\boldsymbol{R} - \boldsymbol{R}'|) \rho_\alpha(\boldsymbol{R}') \frac{\boldsymbol{R} - \boldsymbol{R}'}{|\boldsymbol{R} - \boldsymbol{R}'|}, \quad (17)$$

<sup>1</sup>This folding potential has been corrected from that of Ref. [63] by taking into account the directional dependence in the integral. That of Ref. [63] agrees with the present one only in the case where the nucleon and the  $\alpha$  cluster are distant from each other. Since  $^{13}\text{C}$  is not dilute, it is adequate to use this corrected folding potential.

where  $\rho_\alpha(R) = \frac{4}{\pi^{3/2}b^3} e^{-R^2/b^2}$  is the density function of the  $\alpha$  cluster with the center-of-mass effect removed, with  $b = 1.358\sqrt{\frac{4}{3}} \approx 1.57$  fm, and  $\mathbf{R}$  the  $N$ - $\alpha$  relative coordinate. The shape of the CP-odd  $\alpha$ - $N$  interaction is shown in Fig. 1. It is important to note that the folding cancels the isoscalar and isotensor CP-odd nuclear force, since the  $\alpha$  cluster has closed spin and isospin shells.

The nuclear EDM has two leading sources: (i) the intrinsic EDM of the constituent nucleons and (ii) the P, CP-odd  $NN$  interactions (CP-odd nuclear force).

Let us first see the contribution due to the intrinsic nucleon EDM. It is given by

$$\begin{aligned} d_A^{(\text{Nedm})} &= \sum_i^A d_i \langle \Phi_J(A) | \sigma_{iz} | \Phi_J(A) \rangle \\ &= \frac{1}{2} \sum_i^A [d_p \langle \Phi_J(A) | \sigma_{iz} (1 + \tau_i^3) | \Phi_J(A) \rangle \\ &\quad + d_n \langle \Phi_J(A) | \sigma_{iz} (1 - \tau_i^3) | \Phi_J(A) \rangle] \\ &\equiv C_A^{(0)}(d_p + d_n) + C_A^{(1)}(d_p - d_n), \end{aligned} \quad (18)$$

where  $|\Phi_J(A)\rangle$  is the polarized (in the  $z$  direction) nuclear wave function ( $A = {}^{13}\text{C}$ ) and  $\tau_i^3$  is the isospin Pauli matrix. From the above equation, we see that we just have to calculate the spin matrix elements  $C_A^{(0)} \equiv \frac{1}{2} \sum_i^A \langle \Phi_J(A) | \sigma_{iz} | \Phi_J(A) \rangle$  and  $C_A^{(1)} \equiv \frac{1}{2} \sum_i^A \langle \Phi_J(A) | \sigma_{iz} \tau_i^3 | \Phi_J(A) \rangle$  to obtain the effect of the nucleon EDM on the nuclear EDM. In this case, we do not need to calculate the mixing between parity-odd and parity-even states.

The nucleon EDM contribution to the nuclear EDM is in the general case not enhanced. If the spins of several nucleons are aligned inside the nucleus, the nuclear spin matrix elements may be enhanced. This is, however, not the case for nuclear ground states, due to the strong pairing correlation. The EDM of composite systems may also be enhanced through the polarization of the whole system by the EDM of the components. In heavy atoms, the polarization due to the electron EDM is known to be enhanced by the relativistic effect [105–107]. In nuclear systems, such enhancement is absent, since the system is nonrelativistic. Rather, it was recently suggested that the polarization of the nucleus due to the interaction between the EDM of the nucleon and the nuclear internal electric field can suppress the total intrinsic nucleon EDM contribution [108]. This phenomenon resembles the screening of the EDM of constituents in an electrically bound neutral system, first pointed out by Schiff [34]. In this work, we do not consider this effect. The nuclear spin matrix elements are also suppressed by the mixing of different angular-momentum configurations. In the general case, the nuclear wave function shares some portion of states which have different orbital angular momentum, so  $C_A^{(0)}$  and  $C_A^{(1)}$  are suppressed. The coefficient relating the nucleon EDM to the nuclear EDM  $\frac{1}{2}(C_A^{(0)} \pm C_A^{(1)})$  is therefore at most one for the majority of nuclei.

The evaluation of the nuclear polarization due to the P, CP-odd nuclear force is more complicated than the previous

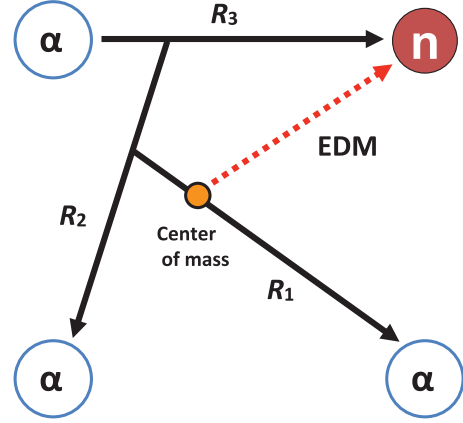


FIG. 2. Jacobi coordinates to express the EDM of  ${}^{13}\text{C}$ .

case, since it arises from the mixing between the parity-even and parity-odd states. In this work we assume that a P, CP-odd nuclear force with a small coupling constant exists. As this P, CP-odd coupling is small, the nuclear EDM should have a linear dependence on it. The polarization contribution of the P, CP-odd nuclear force to the nuclear EDM is given by

$$\begin{aligned} d_A^{(\text{pol})} &= \sum_{i=1}^A \frac{e}{2} \langle \tilde{\Phi}_{J=1/2} | (1 + \tau_i^3) r_{iz} | \tilde{\Phi}_{J=1/2} \rangle \\ &\approx \frac{e}{2} \sum_{i=1}^A \sum_{n \neq 0} \frac{1}{E_0 - E_n} \langle \Phi_{J^p=1/2_1^-} | (1 + \tau_i^3) r_{iz} \\ &\quad \times | \Phi_{J^p=1/2_2^+} \rangle \langle \Phi_{J^p=1/2_2^+} | H_{PT} | \Phi_{J^p=1/2_1^-} \rangle + (\text{c.c.}), \end{aligned} \quad (19)$$

where  $|\tilde{\Phi}_{J=1/2}\rangle$  is the polarized (in the  $z$  axis) nuclear wave function obtained by diagonalizing the Hamiltonian  $\mathcal{H} + H_{PT}$ ;  $r_{iz}$  is (the  $z$  component of) the position of the constituent nucleon in the nuclear center-of-mass frame. The second equality is the first-order perturbation in the P, CP-odd nuclear force  $H_{PT}$ , where  $|\Phi_{J^p=1/2_1^-}\rangle$  is the (polarized) nuclear wave function without opposite-parity states,  $|\Phi_{J^p=1/2_2^+}\rangle$  the (polarized) opposite-parity states, and  $E_n$  their corresponding energy.

The polarization operator of  ${}^{13}\text{C}$  in the  $\alpha$ -cluster model is given by

$$\begin{aligned} \sum_{i=1}^4 Q_i e r_i &= 2e r_1 + 2e r_2 + 2e r_3 + \frac{e}{2} (1 + \tau_4^z) r_4 \\ &= \frac{e}{2} \tau_4^z r_4 = e \tau_4^z \left[ -\frac{2}{13} \mathbf{R}_1 - \frac{2}{9} \mathbf{R}_2 + \frac{2}{5} \mathbf{R}_3 \right], \end{aligned} \quad (20)$$

where  $\mathbf{r}_i$  ( $i = 1, 2, 3$ ) denotes the coordinates of the  $\alpha$  clusters, and  $\mathbf{r}_4$  that of the nucleon, in the center-of-mass frame ( $4\mathbf{r}_1 + 4\mathbf{r}_2 + 4\mathbf{r}_3 + \mathbf{r}_4 = 0$ ). The last line is expressed in terms of the Jacobi coordinate (see Fig. 2). In this work, we do not consider the effect of meson exchange current [109–112].



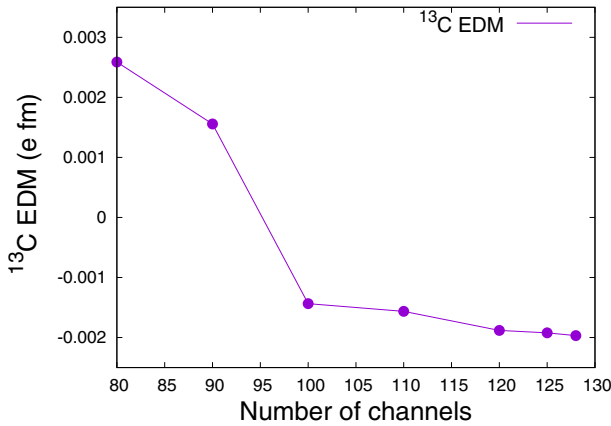


FIG. 3. Convergence of the EDM of  $^{13}\text{C}$  displayed in the function of the number of channels, with 343 bases for each. The CP-odd coupling constant  $\bar{G}_\pi^{(1)}$  was factored out.

#### IV. RESULT AND ANALYSIS

We first show the result of the evaluation of the intrinsic nucleon EDM contribution to the EDM of  $^{13}\text{C}$ . After calculation, we obtain the following value:

$$d_{^{13}\text{C}}^{(\text{Nedm})} = -0.33 d_n. \quad (21)$$

This result agrees with the formula of the EDM of  $^{13}\text{C}$  when we assume that the nucleus is a  $p$ -wave core+valence nucleon system [42],

$$d_{^{13}\text{C}}^{(\text{Nedm})} = -\frac{1}{3} d_n. \quad (22)$$

Here we repeat that the effect of the interaction between the nucleon EDM and the nuclear internal electric field [108] was not taken into account. This may suppress the nucleon EDM contribution to the nuclear EDM, although we do not evaluate it.

Let us try to see the consistency with the experimental data of the nuclear magnetic moment. If we assume the  $p$ -wave core+valence nucleon system with the valence neutron magnetic moment  $\mu_n = -1.9130$  is giving the entire contribution, then we have

$$\mu_{^{13}\text{C}} = 2\mu_n C_{^{13}\text{C}}^{(0)} = 0.63, \quad (23)$$

where  $C_{^{13}\text{C}}^{(0)}$  is defined in Eq. (18). This estimation is in reasonable agreement with the experimental value  $\mu_{^{13}\text{C}} = 0.702411(1)$ . Our result is therefore consistent with the experimental data of the nuclear magnetic moment.

Next, we show the result of the calculation of the polarization contribution. The effect of the CP-odd nuclear force to the nuclear EDM of  $^{13}\text{C}$  is

$$d_{^{13}\text{C}}^{(\text{pol})} = -0.0020 \bar{G}_\pi^{(1)} e \text{ fm}. \quad (24)$$

Note that only the isovector CP-odd nuclear force contributes to the  $^{13}\text{C}$  EDM in the  $\alpha$ -cluster model. We used  $7^3 \times 128 = 43\,904$  bases to converge the EDM of  $^{13}\text{C}$  (see Sec. II). The convergence of the nuclear EDM in the function of the angular-momentum channels is shown in Fig. 3. We see that the EDM of

$^{13}\text{C}$  is smaller than the result of the calculation for lighter nuclei in Ref. [63] [e.g., the deuteron EDM  $d_{^2\text{H}}^{(\text{pol})} = 0.0145 \bar{G}_\pi^{(1)} e \text{ fm}$ ]. This fact suggests that some suppression mechanisms of the EDM are relevant for  $^{13}\text{C}$ . We note that, although being smaller than other lighter nuclei, the EDM of  $^{13}\text{C}$  is only smaller than them by an order of magnitude. It is much larger than the nuclear EDM (not atomic EDM!) of  $^{129}\text{Xe}$  by several orders [64].

Let us try to analyze the physical mechanism which suppresses the EDM following Eq. (19). The first point is the distance between the valence nucleon and the core. The expectation value of the EDM becomes larger when the  $1/2_1^-$  and the opposite-parity  $1/2_1^+$  states have better overlap. Here good overlap points to the bra and ket wave functions which make large matrix elements of the CP-odd nuclear force and the EDM operator. Such a combination of wave functions must therefore be well related with a  $|\Delta L| = 1$  transition, with the constituents not distant from each other, since the CP-odd nuclear force is exponentially damping at long distance.

In looking at the  $1/2_1^-$  state, the root-mean-square radius of the core-valence distance of  $^{13}\text{C}$  is  $\langle \sqrt{r^2} \rangle_{1/2_1^-} = 2.81 \text{ fm}$ , whereas that of the  $1/2_1^+$  state is  $\langle \sqrt{r^2} \rangle_{1/2_1^+} = 3.95 \text{ fm}$  [69]. We emphasize that in this work we refer the ground and first excited states as the lowest and the next lowest energy states, respectively, obtained after diagonalizing the Hamiltonian including the CP-odd nuclear force for the system with angular momentum  $1/2$ , while the  $1/2_1^-$  and  $1/2_1^+$  states correspond to those unperturbed by the CP-odd nuclear force, without parity mixing. We see that  $^{13}\text{C}$  has a shell structure in the  $1/2_1^-$  state but a neutron halo structure in the  $1/2_1^+$  excited state. The bad overlap of the core-valence wave function may therefore suppress the EDM.

Another bad overlap between the ground and  $1/2_1^+$  states also exists. The core of  $^{13}\text{C}$ , i.e., the  $^{12}\text{C}$  subsystem, actually has a different structure between the two states. In Ref. [69], the spectroscopic factors for both states were calculated. There it was found that the core is dominated by the  $2^+$  state in the ground state. This is due to the strong  $LS$  attraction between the valence nucleon and the  $\alpha$  cluster which forces the alignment of their orbital angular momentum ( $p$ -wave) and the nucleon spin to form an angular-momentum  $3/2$  system. This fact, consequently, requires the core to have at least angular momentum 2 to be consistent with the total angular momentum  $1/2$  of the  $^{13}\text{C}$  nucleus. However, in the  $1/2_1^+$  state, the leading contribution is given by a  $0^+$  state. This makes a bad overlap between the cores of the  $1/2_1^-$  and  $1/2_1^+$  states and, consequently, suppresses the EDM.

The spectrum of  $^{13}\text{C}$  (not perturbed with CP-odd nuclear force) shows a  $1/2_1^+$  bound state at 3.1 MeV. In the study of the EDM of light nuclei, it is the first case where we can find the ground and first excited states both as bound states. In the leading order of perturbation (which works well in the present case), the EDM receives a contribution from the transition between the  $1/2_1^-$  and  $1/2_1^+$  states and that between the  $1/2_1^-$  state and the  $1/2$  continuum state [see Eq. (19) and Fig. 4]. It is interesting to compare those two effects. The EDM of the ground state (in the spectrum with the CP-odd nuclear force) of  $^{13}\text{C}$ , generated by the transition between  $1/2_1^-$  and  $1/2_1^+$

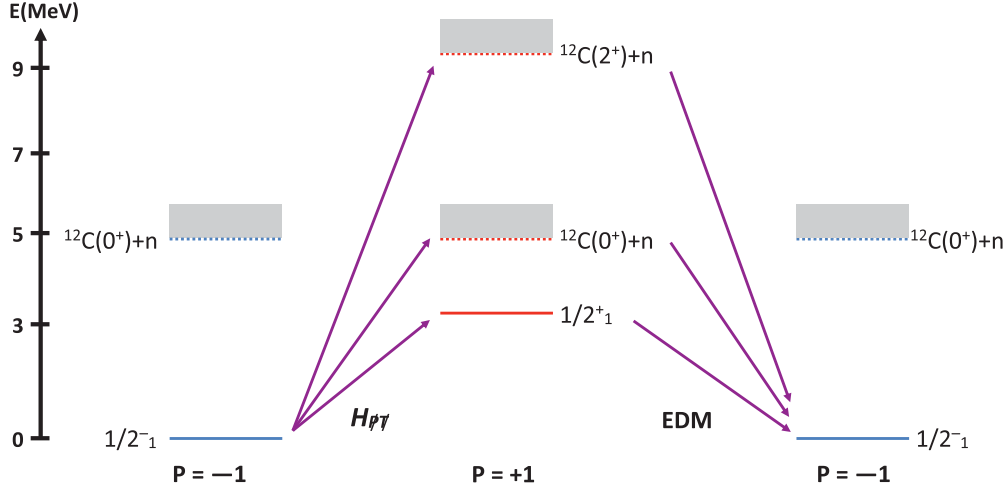


FIG. 4. Parity-odd transitions contributing to the EDM of  $^{13}\text{C}$  in the leading order of perturbation. The dashed lines represent the dissociation thresholds into  $^{12}\text{C}(0^+) + n$  or  $^{12}\text{C}(2^+) + n$  states. The gray bands correspond to the continuum states.

states (solved without CP-odd nuclear force), is

$$d_{^{13}\text{C}}^{\text{bound}} \approx \sum_i \frac{\langle \Phi_{J^P=1/2_1^-} | Q_i e r_i | \Phi_{J^P=1/2_1^+} \rangle}{E_{1/2_1^-} - E_{1/2_1^+}} \times \langle \Phi_{J^P=1/2_1^+} | H_{pT}^\pi | \Phi_{J^P=1/2_1^-} \rangle. \quad (25)$$

Our calculation gives

$$d_{^{13}\text{C}}^{\text{bound}} = 0.00025 \bar{G}_\pi^{(1)} e \text{ fm}. \quad (26)$$

The remaining part is the contribution from the transition to the continuum states and resonances, which is given by

$$d_{^{13}\text{C}}^{\text{rem}} = d_{^{13}\text{C}}^{(\text{pol})} - d_{^{13}\text{C}}^{\text{bound}} = -0.0022 \bar{G}_\pi^{(1)} e \text{ fm}. \quad (27)$$

In Ref. [69], five resonances have been identified for each parity. By calculating their contribution as in Eq. (25), it is found that the effect of resonances is not larger than 10% of  $d_{^{13}\text{C}}^{\text{rem}}$ . The transition to continuum is therefore dominant in the EDM of  $^{13}\text{C}$ .

How about the excited state (in the spectrum with the CP-odd nuclear force)? In our framework, the EDM of the first excited state (in the spectrum with the CP-odd nuclear force) can also be calculated. The EDM of the excited states is given by

$$d_{^{13}\text{C}^*}^{(\text{pol})} = 0.0244 \bar{G}_\pi^{(1)} e \text{ fm}. \quad (28)$$

We see that the EDM of the excited states is much larger in magnitude than the ground-state EDM (by about 10 times). Among them, the contribution from the bound-state transition is just

$$d_{^{13}\text{C}^*}^{\text{bound}} = -d_{^{13}\text{C}}^{\text{bound}} = -0.00025 \bar{G}_\pi^{(1)} e \text{ fm}. \quad (29)$$

For the remaining contribution, we, however, have

$$d_{^{13}\text{C}^*}^{\text{rem}} = d_{^{13}\text{C}^*}^{(\text{pol})} - d_{^{13}\text{C}^*}^{\text{bound}} = 0.0247 \bar{G}_\pi^{(1)} e \text{ fm}, \quad (30)$$

which is dominantly made by the transition to continuum. The contribution from resonances is less than 3%. We see that  $d_{^{13}\text{C}^*}^{\text{rem}}$  is much larger than that of the ground state. The

continuum contribution represents the majority for the EDM of excited  $^{13}\text{C}$ . These large values can be explained as follows. The analysis of the spectroscopic factor in Ref. [69] shows that the excited state ( $1/2_1^+$  in the absence of CP-odd nuclear force) has a dominant  $0^+$  state contribution for the core,  $^{12}\text{C}(0^+)$  [69]. The (bound) excited state is dominantly made of a halolike configuration of valence nucleon ( $s$  wave) and the  $^{12}\text{C}(0^+)$  core. This state should have a large overlap with the  $^{12}\text{C}(0^+)$  core+ $n$  ( $p$  wave) continuum state, for which the threshold opens just 1.8 MeV (for  $^{13}\text{C}$ ) above the energy level of the  $1/2_1^+$  state (4.9 MeV above the  $1/2_1^-$  state of  $^{13}\text{C}$ , see Fig. 4).

For the case of the ground state (with CP-odd nuclear force), the core has a dominant  $2^+$  configuration,  $^{12}\text{C}(2^+)$ . Therefore the EDM does not receive a large contribution from the continuum  $^{12}\text{C}(0^+) + n$  ( $s$  wave) due to the small overlap with this continuum. The closest threshold of continuum states that dominantly couple to the ground state,  $^{12}\text{C}(2^+) + n$ , is about 4 MeV above the  $^{12}\text{C}(0^+) + n$  threshold (about 9 MeV above the  $1/2_1^-$  state), since we need this energy to excite  $^{12}\text{C}(0^+)$  to the  $2_1^+$  state (see Fig. 4). From the expression of the perturbation (19), the EDM is inversely proportional to the energy difference. The EDM of the ground state, although having a good overlap with the  $^{12}\text{C}(2^+) + n$  continuum, is suppressed by the large energy difference 9 MeV and becomes smaller than the EDM of the excited state, for which the energy difference is only 1.8 MeV.

Our analysis of the EDM of  $^{13}\text{C}$  has qualitatively explained the suppression and the enhancement of the ground- and first-excited-state EDMs, respectively, by inspecting the energy levels and the overlap of opposite-parity states. This fact suggests that the nuclear EDM is very sensitive to the nuclear structure. As the structure of nuclei differs much between nuclei, its investigation also guides us in finding nuclei with a large EDM. As a nucleus which has a continuum state threshold near the ground state, we have the  $^7\text{Li}$  nucleus, which has a  $^3\text{H} + \alpha$  threshold at 2.47 MeV above the ground state. The  $^{19}\text{F}$  nucleus is also a very good candidate, as it has an opposite-parity first excited state at 110 keV above the ground state.

We also mention the theoretical uncertainty of our result. The intrinsic nucleon EDM contribution to the EDM of  $^{13}\text{C}$  is well explained in the valence nucleon picture, and the discrepancy of the calculated magnetic moment with the experimental data is about 10%. The error bar of our result [Eq. (21)] should be of the same order, like 10%.

The theoretical uncertainty due to the isovector CP-odd nuclear force [term with  $\bar{G}_\pi^{(1)}$ , Eq. (12)] is estimated from the energy spectrum calculated in detail in our previous work [69]. The energy levels are well reproduced, within about 30% for those above the  $^{12}\text{C}(0^+) + n$  threshold. As we showed in this section that the nuclear EDM does not receive much contribution from resonances, it is expected that it is less affected by the discrepancy of the energy-level scheme.

We also mention the theoretical uncertainty related to the  $E1$  transition. The observed energy spectrum of  $^{13}\text{C}$  contains a giant dipole resonance near 25 MeV which might sizably contribute to the EDM. The cluster model cannot fully describe it, so we have to consider this contribution as a systematic error. Here we shall show that it is not large. Since the giant dipole resonance is located at 25 MeV, its contribution to the EDM is damped by the large energy in the denominator [see Eq. (19)]. We also see from Eq. (19) that the EDM is generated not only by the isovector dipole operator but also by the isovector spin-flip operator of the CP-odd nuclear force [see Eq. (12)]. This latter induces an isovector spin-flip giant dipole resonance, which requires higher excitation energy than the giant dipole resonance. Unless those giant resonances are strongly coupled, they should not give important effects.

The  $^{12}\text{C} + n$  continuum states therefore play the most important role in the EDM of  $^{13}\text{C}$ . The  $3\alpha + n$  OCM (orthogonality condition model), which is a semimicroscopic cluster model, can simultaneously describe the  $^{12}\text{C} + n$  asymptotic behavior, as well as the  $^{12}\text{C} + n$  threshold energy, and the ground state of  $^{13}\text{C}$  with a shell-model-like structure, whereas the ordinary shell model or other mean-field approaches cannot. There the  $^{12}\text{C} + n$  continuum states should well be described, since the structure of  $^{12}\text{C}$  is accurately given in the  $3\alpha$  OCM [79,85,113]. Here we conservatively estimate the theoretical error bar as 50%, since the EDM of  $^{13}\text{C}$  is generated through a suppression mechanism.

## V. ROLE OF $^{13}\text{C}$ EDM IN THE DETERMINATION OF HADRON-LEVEL CP VIOLATION

Let us inspect the necessity to measure the EDM of  $^{13}\text{C}$ . Its dependence on the bare CP-odd couplings of Eq. (9), in the leading order of chiral perturbation theory, is given by

$$d_{13\text{C}} = -0.33\bar{d}_n + [-0.045\bar{g}_{\pi NN}^{(0)} - 0.025\bar{g}_{\pi NN}^{(1)}]e \text{ fm}. \quad (31)$$

Here we neglect the isotensor CP-odd pion-nucleon interaction [term with  $\bar{g}_{\pi NN}^{(2)}$  in Eq. (9)]. The dependence of  $d_{13\text{C}}$  on  $\bar{g}_{\pi NN}^{(0)}$  dominantly comes from the valence neutron EDM.

By using the relation between the neutron EDM and  $\bar{\theta}$ , obtained in the chiral effective field theory combined with lattice QCD data [62], the  $\theta$ -term contribution to  $^{13}\text{C}$  EDM is given by

$$d_{13\text{C}} = (9 \pm 5) \times 10^{-17} \bar{\theta} e \text{ c.m.} \quad (32)$$

This contribution is dominantly due to the valence neutron EDM. The dependence of the EDM of  $^{13}\text{C}$  on  $\bar{g}_{\pi NN}^{(0)}$  as well as on the  $\theta$  term is more pronounced than that of the deuteron and  $^6\text{Li}$ , since the effect of  $\bar{g}_{\pi NN}^{(0)}$  cancels in the leading order. We see below that even if the sensitivity of  $^{13}\text{C}$  on individual CP-odd interactions is smaller than other light nuclei (see Ref. [63]), it has its own speciality.

It is important to note that the experimental data of the EDM of a single system, even if a finite value is measured, are not sufficient to determine the new physics beyond the standard model. Obviously, the combination of the observations for several systems with linearly independent coefficients will absolutely be required to constrain multiple CP-odd couplings. Explicitly, there are four unknown CP-odd couplings [ $\bar{d}_0, \bar{d}_1, \bar{g}_{\pi NN}^{(0)}, \bar{g}_{\pi NN}^{(1)}$ ] in the hadron-level effective interaction we are considering [see Eq. (9)]. To fix them, at least four experimental measurements of EDM, using systems with linearly independent dependence on them, are required. We emphasize that those four unknown parameters do not form the complete set of leading CP-odd couplings, and additional CP-odd interactions such as the three-pion interaction or the contact interactions [see Eq. (10)] are known to contribute in the leading order of chiral perturbation theory [13,58,59,62,68]. To fix them, additional EDM experimental data are required. The experimental measurements of the EDM of lightest systems, namely the neutron, the proton, the deuteron, and  $^3\text{He}$ , are not sufficient to determine the new physics. In this circumstance, the experimental constraint from the EDM of  $^{13}\text{C}$  will play an important role.

As other potential candidates of nuclei to be measured, we have  $^6\text{Li}$  and  $^9\text{Be}$  [63]. The  $^{13}\text{C}$  and  $^9\text{Be}$  nuclei have similar dependence on the CP-odd couplings, since they probe  $d_n$  through the valence neutron, and  $\bar{g}_{\pi NN}^{(1)}$  via the polarization. It is maybe useful to note that the experimental manipulation of  $^9\text{Be}$  may have some disadvantages comparing with  $^{13}\text{C}$ , since the  $^9\text{Be}$  atom is unstable, and its raw component is toxic. The  $^{13}\text{C}$  nucleus is therefore a good candidate to probe the linear combination of  $d_n$  and  $\bar{g}_{\pi NN}^{(1)}$ .

In the analysis of new physics beyond the standard model, there is always some region of the parameter space which escapes the EDM constraints [114–118]. The measurement of  $^{13}\text{C}$  EDM may therefore be an important piece to determine or discriminate candidates of new physics. Moreover, the coefficient of the isovector nuclear force [ $\bar{g}_{\pi NN}^{(1)}$ ] and that for the intrinsic nucleon EDM are negative. Other light nuclei studied so far all have corresponding coefficients positive. The  $^{13}\text{C}$  nucleus has thus the possibility to provide a good check of the relative sign.

We should not forget to show the standard model contribution generated by the CP phase of the Cabibbo-Kobayashi-Maskawa matrix [41]:

$$d_{13\text{C}}^{(\text{SM})} = -3.3 \times 10^{-32} e \text{ c.m.} \quad (33)$$

We see that the effect of the standard model is small. Here we have only considered the polarization effect from the pion exchange CP-odd nuclear force. The nucleon EDM also contributes to the nuclear EDM [39,40], but its contribution should be small. The reason is the same as that for the

suppression of the effect of new physics candidates which contribute through the quark EDM.

## VI. PROSPECTS TO THE DISCOVERY OF NEW PHYSICS

Let us now see the prospect for the discovery of new physics beyond the standard model. The first typical CP-violating model to be inspected is the supersymmetric model [10,114–117,119–123]. If the EDM of  $^{13}\text{C}$  can be measured at the level of  $O(10^{-29})e\text{ cm}$  [33], then the mass scale of the supersymmetry breaking can be probed at the TeV scale (see Ref. [63]).

The class of models which generate the four-quark interaction or the Barr-Zee type diagram contributes through the combination  $\frac{\text{Im}(gg')}{m_{\tilde{\chi}_p}^2}$ . If the nuclear EDM is measured with the prospective sensitivity  $O(10^{-29})e\text{ cm}$ , then the latter combination of couplings and new particle mass can be probed at the level of  $O(0.1\text{ PeV})$ . Those new physics include the left-right symmetric model [57], the Higgs doublet models [124–126], supersymmetric models with  $R$ -parity violation [118,127–130], etc.

Models which contribute to the nuclear EDM only through the quark EDM are much less sensitive than those which generate the quark chromo-EDM. This fact is due to several physical reasons. The first reason is because the effect of the quark EDM to the nucleon EDM is suppressed by the nucleon tensor charge. The nucleon tensor charge is the linear coefficient which relates the quark EDM to the nucleon EDM, and the extraction from experimental data shows values smaller than 1 [131–134]. Recent lattice QCD analyses give also consistent data [135–143]. It is suggested that nucleon charges which give the quark spin in the nonrelativistic limit, such as the axial or tensor charges, are suppressed by the dynamical gluon dressing [144–146]. The second reason is the suppression of the Wilson coefficient of the quark EDM operator by the renormalization group evolution. By running this Wilson coefficient from the scale  $\mu = 1\text{ TeV}$  to the hadronic scale  $\mu = 1\text{ GeV}$ , it becomes less than 80% [147–149]. The third one is the suppression of the nucleon EDM contribution to the EDM of  $^{13}\text{C}$  by a factor of  $-\frac{1}{3}$ , due to the antiparallel valence nucleon spin and the core-valence orbital angular momentum. There may also be additional screening effect to the nucleon EDM, generated by the interaction between the nucleon EDM and the nuclear internal electric field [108]. For the case of the Barr-Zee type diagram, the fermion EDM may also be suppressed by the electromagnetic coupling and the fractional charge of the quarks,

relative to the strong coupling. For models which generate the quark electromagnetic EDM at leading order, it is therefore strongly recommended to probe them through the neutron EDM.

## VII. SUMMARY

In this work, we have calculated the EDM of  $^{13}\text{C}$  using the Gaussian expansion method. There we have assumed the  $\alpha$ -cluster model. As a result, we obtained that  $^{13}\text{C}$  has smaller enhancement factors than the deuteron,  $^3\text{He}$ ,  $^3\text{H}$ ,  $^6\text{Li}$ , and  $^9\text{Be}$ . The sensitivity of the EDM of  $^{13}\text{C}$  is nevertheless much larger than the EDM of the bare  $^{129}\text{Xe}$  nucleus, and this fact shows us that the EDM of light nuclei are promising observables. Its experimental measurement is absolutely required for determining multiple hadron-level CP-violating couplings.

By analyzing the suppression of the EDM of  $^{13}\text{C}$ , we have found a mechanism to enlarge the matrix elements of CP-odd operators between opposite-parity states, required to enhance the nuclear EDM. In the case of  $^{13}\text{C}$ , the bad overlap between the  $^{12}\text{C}$  core of the  $1/2_1^-$  and the  $1/2_1^+$  or the closest continuum states suppresses the EDM. This shows that the knowledge of the structure of nuclei is essential in giving the linear coefficient between the nuclear EDM and the CP-odd nuclear couplings.

The result of this work provides a very good guide to find nuclei with large EDM. In particular, we expect a large EDM for  $^7\text{Li}$  and  $^{19}\text{F}$ . This enhancement can potentially reach more than  $O(10)$  times than those so far known. The study of those nuclei will be the subject of future works.

Moreover, the determination of the new physics beyond the standard model cannot be achieved by only measuring the EDM of one sensitive system, and combining experimental data is essential. In particular, there are more than four types of CP-odd effective interactions contributing in the leading order of chiral perturbation theory, and the measurements of the EDM of light systems (neutron, proton, deuteron, and  $^3\text{He}$ ) are not sufficient to determine the hadron-level CP violation. The nuclear EDM of  $^{13}\text{C}$  can therefore play a significant role in driving the CP violation of new physics into a corner or discriminate among candidates.

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