Letter

Renormalization of *CP*-violating nuclear forces

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Electric dipole moments of nuclei, diamagnetic atoms, and certain molecules are induced by *CP*-violating nuclear forces. Naive dimensional analysis predicts these forces to be dominated by long-range one-pion-exchange processes with short-range forces entering only at next-to-next-to-leading order in the chiral expansion. Based on renormalization arguments we argue that a consistent picture of *CP*-violating nuclear forces requires a leading-order short-distance operator contributing to $^1S_0$ - $^3P_0$ transitions due to the attractive and singular nature of the strong tensor force in the $^3P_0$ channel. The short-distance operator leads to $O(1)$ corrections to static and oscillating, relevant for axion searches, electric dipole moments. We discuss strategies how the finite part of the associated low-energy constant can be determined in the case of *CP* violation from the QCD $\bar{\theta}$ term by the connection to charge-symmetry violation in nuclear systems.

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Introduction. Electric dipole moments (EDMs) of nuclei, atoms, and molecules are excellent probes of new sources of *CP* violation [1,2]. *CP* violation in the fermion mixing matrices of the standard model (SM) only induce EDMs through multiple electroweak loops and lead to immeasurably small values for EDMs [3,4], implying that any nonzero measurement is either due to the so-far undiscovered QCD $\bar{\theta}$ term or from beyond-the-SM (BSM) sources of *CP* violation. Current experimental EDM limits [5–7] set strong constraints on BSM models with additional *CP*-violating phases, such as supersymmetry, leptoquarks, or left-right symmetric models, and scenarios of electroweak baryogenesis [8]. In the framework of the SM effective field theory (SMEFT), EDM limits constrain a large set of *CP*-odd dimension-six operators at the multi-TeV scale, well above limits from collider experiments [9].

The interpretation of EDM experiments requires care. It is a nontrivial task to connect EDMs of complex objects, such as nuclei to the underlying *CP*-violating source at the quark level. Recent years have seen significant theoretical improvements towards model-independent first-principles calculations of EDMs from a combination of lattice QCD [10–12], chiral EFT (χEFT) [13–15], and nuclear calculations [16–20]. The chain of logic is roughly as follows: The SMEFT framework allows for the derivation of a general set of dimension-four (the $\bar{\theta}$ term) and dimension-six *CP*-violating operators involving light quarks, gluons, and photons. χEFT, the low-energy EFT of QCD, is used to construct the corresponding *CP*-violating interactions among the relevant low-energy degree of freedoms: pions, nucleons, and photons. Each chiral interaction comes with a low-energy constant (LEC) that encodes the nonperturbative QCD dynamics, ideally calculated from lattice QCD (LQCD). EDMs can be then be calculated in terms of the LECs in the chiral Lagrangian.

The χEFT framework provides an expansion of hadronic and nuclear amplitudes in terms of p/Λ_χ where $p \approx k_F \approx m_\pi \approx O(100\ \text{MeV})$ and $\Lambda_\chi \approx 4\pi F_\pi \approx O(1\ \text{GeV})$ [21–23], where $F_\pi \simeq 92.4$ MeV is the pion decay constant. The electric dipole form factors of nucleons were calculated up to next-to-next-to-leading order (N^{2}LO) [24–27]. Nuclear EDMs require the derivation of *CP*-violating forces and currents. The *CP*-odd nucleon-nucleon (*NN*) potential was calculated up to N^{2}LO [20,28] and used to calculate nuclear and atomic EDMs [16–20].

The derivation of the *CP*-odd *NN* potential of Refs. [20,28] is based on Weinberg's power-counting scheme [29]. In this scheme, the *CP*-odd potential arises from one-pion-exchange (OPE) diagrams, whose LECs can be fixed from processes involving just nucleons and pions (only in principle as πN-scattering experiments are not sufficiently accurate). Chiral symmetry does not forbid purely nuclear short-distance interactions with LECs that can only be fixed in nuclear systems. Indeed, in the *CP*-conserving potential the LO potential consists of OPE diagrams and two nonderivative contact interactions in $^1S_0$ and $^3S_1$ waves. In the *CP*-violating case, *NN* interactions require, at least, one space-time derivative, and Weinberg's power-counting scheme predicts short-distance operators to enter at N^{2}LO. This is welcome news: It implies that nuclear EDMs can be calculated in terms of only a few LECs and ratios of EDMs can be used to pinpoint the underlying *CP*-violating source [30].

Weinberg's power-counting scheme is based on naive dimensional analysis (NDA) of the *NN* LECs [31] which is not always reliable for nuclear physics. NDA does not in all cases lead to order-by-order renormalized nuclear amplitudes [32,33]. This is most clear in partial waves where OPE is attractive and nonperturbative, such as the $^3P_0$ channel where phase shifts show oscillatory limit-cycle-like cutoff dependence [34] that cannot be renormalized at LO in Weinberg's scheme, although there is some debate about this [35–37]. The same problem affects external currents inserted in *NN*-scattering states [38,39]. Here we investigate *CP*-violating OPE potentials and use cutoff dependence of observables as a diagnostic tool to demonstrate that a LO short-distance operator for $^1S_0$ - $^3P_0$ transitions is required. This directly affects the interpretation of EDM experiments and other time-reversal-odd observables, such as magnetic quadrupole moments or neutron-nucleus scattering.

An axion dark matter (DM) field can induce oscillating EDMs that can be probed in experiments [40]. Such DM axions would constitute a coherent oscillating classical field $a \approx (a_0/f_a)\cos \omega t$ where the angular frequency is given by $\omega \simeq m_a$ where m_a is the axion mass, a_0 is the local amplitude of the axion DM field, and f_a is the axion decay constant [41,42]. The hadronic and nuclear matrix elements that connect static EDMs to the $\bar\theta$ term are identical to those that connect oscillating EDMs to the axion field. Here we focus on the former, but all expressions below are applicable to DM searches for axions by replacing $\bar\theta$ by $(a_0/f_a)\cos \omega t$.

Setup of the calculation. We first consider the QCD $\bar\theta$ term. The relevant Lagrangian is given by [43,44]

$$\mathcal{L} = \bar q i \not{D} q - \bar q(\mathcal{M} - i\gamma_5 m_\star \bar\theta)q, \tag{1}$$

where $q = (ud)^T$ denotes the quark field, $\mathcal{D}_\mu$ denotes the color and electromagnetic covariant derivative, and $\mathcal{M} = \mathrm{diag}(m_u, m_d)$ denotes the quark mass matrix $m_\star = m_u m_d/(m_u + m_d)$. The chiral Lagrangian can be constructed with well-known methods [45], and the leading *CP*-even and *CP*-odd pion-nucleon interactions are

$$\mathcal{L}_{\pi N} = -\frac{g_A}{2F_\pi}\nabla\vec\pi\cdot\bar N\vec\tau\boldsymbol{\sigma}N + \bar g_0 \bar N\vec\pi\cdot\vec\tau N + \cdots, \tag{2}$$

in terms of the nonrelativistic nucleon doublet $N = (pn)^T$ and the pion triplet $\vec\pi$, $g_A \simeq 1.27$ is the nucleon axial coupling, and $\bar g_0 = O(m_\star\bar\theta/F_\pi)$ is a *CP*-odd LEC. The dots denote interactions involving more pions. The QCD $\bar\theta$ term is related by a chiral rotation to the isospin-breaking component of the quark masses [24], giving a precise determination of $\bar g_0$ [46],

$$\bar g_0 = \frac{\delta m_N^{\rm str}(1-\varepsilon^2)}{4F_\pi\varepsilon}\bar\theta = -(14.7 \pm 2.3)\times 10^{-3}\bar\theta, \tag{3}$$

where $\delta m_N^{\rm str}$ is the quark-mass-induced part of the proton-neutron mass splitting that has been calculated with LQCD [47] and $\varepsilon = (m_u - m_d)/(m_u + m_d)$. The value of $\bar g_0$ agrees with a LQCD extraction [12].

From the interactions in Eq. (2) we calculate the OPE *NN* potentials,

$$V_{{\rm str},\pi} = -\frac{1}{(2\pi)^3}\left(\frac{g_A}{2F_\pi}\right)^2 \vec\tau_1\cdot\vec\tau_2\frac{(\boldsymbol{\sigma}_1\cdot\mathbf{q})(\boldsymbol{\sigma}_2\cdot\mathbf{q})}{\mathbf{q}^2+m_\pi^2},$$
$$V_{\bar g_0} = -\frac{1}{(2\pi)^3}\frac{g_A\bar g_0}{2F_\pi}\vec\tau_1\cdot\vec\tau_2\frac{i(\boldsymbol{\sigma}_1-\boldsymbol{\sigma}_2)\cdot\mathbf{q}}{\mathbf{q}^2+m_\pi^2}, \tag{4}$$

where $\mathbf{q} = \mathbf{p} - \mathbf{p}'$ is the momentum transfer between in- and outgoing nucleon pairs with relative momenta $\mathbf{p}$ and $\mathbf{p}'$, respectively ($|\mathbf{p}| = p$ and $|\mathbf{p}'| = p'$), and m_π denotes the pion mass. We include the *CP*-even *NN* interactions,

$$V_{\rm str,sd} = \frac{1}{(2\pi)^3}\left(C_sP_s + C_tP_t + \frac{1}{4}pp'C_PP_p\right), \tag{5}$$

where $P_{s,t,p}$ project, respectively, on the $^1S_0$, $^3S_1$, and $^3P_0$ waves. In Weinberg's power counting the *S*-wave contact terms appear at LO whereas the *P*-wave counterterm enters at N^{2}LO. To obtain the strong *NN*-scattering wave functions we solve a Lippmann-Schwinger (LS) equation,

$$T_{\rm str} = V_{\rm str} + V_{\rm str}G_0T_{\rm str}, \quad G_0 = (E - p^2/m_N + i\varepsilon)^{-1}, \tag{6}$$

with $V_{\rm str} = (V_{{\rm str},\pi} + V_{\rm str,sd})f_\Lambda(p,p')$, where $f_\Lambda(p,p')$ is a regulator function,

$$f_\Lambda(p,p') = e^{-(\mathbf{p}/\Lambda)^4}e^{-(\mathbf{p}'/\Lambda)^4}, \tag{7}$$

in terms of a momentum space cutoff Λ. The LS equation is solved numerically for a wide range of Λ to ensure that observables are cutoff independent.

We briefly discuss results for waves with total angular momentum $j = 0, 1$. Solving the LS equation for just the strong OPE potential leads to $^1S_0$ and $^3S_1$ - $^3D_1$ phase shifts and mixing angles that are cutoff dependent, see Fig. 1(a). In Weinberg's power counting, this is resolved by including the short-distance counterterms C_s and C_t acting in $^1S_0$ and $^3S_1$ waves. Fitting the LECs to reproduce the strong phases shifts at $E_{CM} = 5$ MeV, the phase shifts become cutoff independent for a wide range of energies as exemplified in Fig. 1(b).

In the $^3P_1$ and $^1P_1$ waves, the strong OPE potential lead to cutoff independent phase shifts, see Fig. 2(a) that at low energies agree well with experimental data. In the $^3P_0$ channel, however, the phase shifts arising from OPE are strongly cutoff dependent and undergo a dramatic limit-cycle-like behavior, see Fig. 2. There does not appear to be a counterterm that can absorb this regulator dependence in Weinberg's power counting. Following Ref. [34], we promote the $^3P_0$ counterterm with LEC C_P in Eq. (5) to LO and fit C_P to the $^3P_0$ phase shift at $E_{\rm CM} = 5$ MeV. With this modified power counting the phase shifts become cutoff independent, see Fig. 2(a). The regulator dependence of C_s, C_t, and C_P are given in Fig. 2(b). The LECs $C_{s,t,P}$ show significant Λ dependence, which is of no concern as they are not observable. All results agree with Refs. [34,48].

We now insert the *CP*-odd potential $V_{\bar g_0}$ which causes $^1S_0$ - $^3P_0$ and $^3S_1$ - $^1P_1$ transitions. We treat $V_{\bar g_0}$ to very good accuracy in perturbation theory and write

$$T_{\bar g_0} = V_{\bar g_0} + V_{\bar g_0}G_0T_{\rm str} + T_{\rm str}G_0V_{\bar g_0} + T_{\rm str}G_0V_{\bar g_0}G_0T_{\rm str}. \tag{8}$$

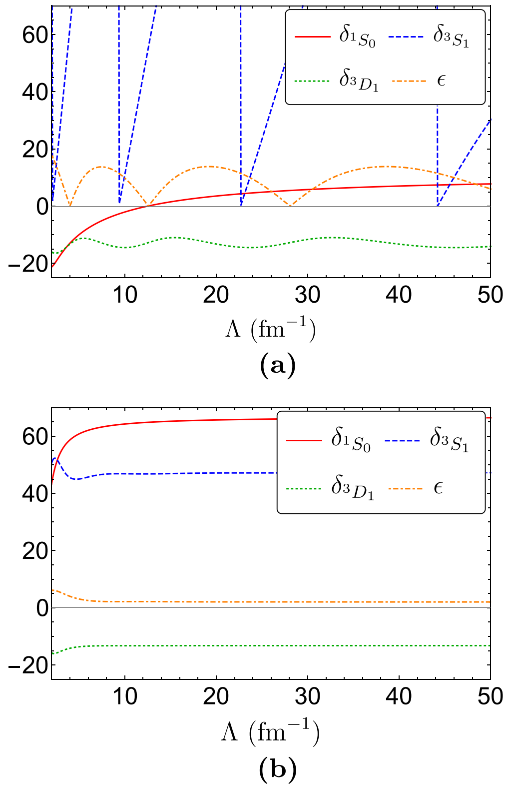

FIG. 1. Phase shifts and mixing angle for the ${}^1S_0$ and ${}^3S_1$ - ${}^3D_1$ channels at $E_{\rm CM} = 50$ MeV from just the strong OPE potential (a) and after renormalization (b) as function of the regulator Λ.

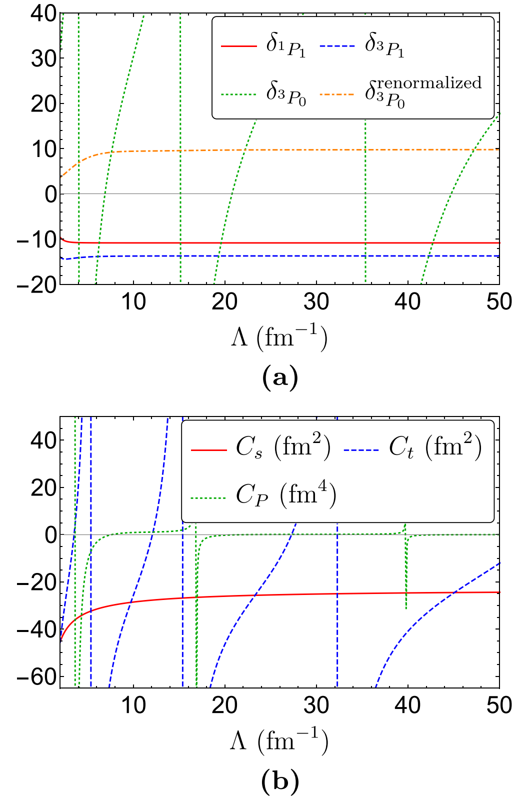

FIG. 2. (a) ${}^3P_0$, ${}^3P_1$, and ${}^1P_1$ phase shifts at $E_{\rm CM} = 50$ MeV as the function of the regulator Λ. The green (dotted) line denotes the ${}^3P_0$ phase shifts in Weinberg's power counting where no counterterm is available for renormalization. The orange (dashed-dotted) is the result after promoting C_P to LO. (b) Low-energy constants C_s, C_t, and C_P as function of Λ.

The on-shell scattering matrix $T = T_{\rm str} + T_{\bar g_0}$ is related to the S matrix,

$$S(E_{\rm CM}) = 1 - i\pi m_N^{3/2} E_{\rm CM}^{1/2} T(p = p' = \sqrt{E_{\rm CM} m_N}), \tag{9}$$

where m_N is the nucleon mass. For $j = 0$ we parametrize the S matrix by

$$S_{j=0} = \begin{pmatrix} e^{2i\delta_{{}^1S_0}} & \epsilon^0_{\rm SP} e^{i[\delta_{{}^1S_0}+\delta_{{}^3P_0}]} \\ -\epsilon^0_{\rm SP} e^{i[\delta_{{}^1S_0}+\delta_{{}^3P_0}]} & e^{2i\delta_{{}^3P_0}} \end{pmatrix}, \tag{10}$$

where $\epsilon^0_{\rm SP} \approx \bar\theta$ denotes the small ${}^1S_0$ - ${}^3P_0$ mixing angle. The $j = 1$ channel is more complicated because of strong ${}^3S_1$ - ${}^3D_1$ mixing, and for simplicity, we expand in the small S-D mixing angle ϵ. Up to $O(\epsilon^3)$,

$$S_{j=1} = \begin{pmatrix} e^{2i\delta_{{}^3S_1}} \cos 2\epsilon & ie^{i[\delta_{{}^3S_1}+\delta_{{}^3D_1}]} \sin 2\epsilon & x_{\rm SP} \\ ie^{i[\delta_{{}^3S_1}+\delta_{{}^3D_1}]} \sin 2\epsilon & e^{2i\delta_{{}^3D_1}} \cos 2\epsilon & x_{\rm DP} \\ -x_{\rm SP} & -x_{\rm DP} & e^{2i\delta_{{}^1P_1}} \end{pmatrix},$$

$$x_{\rm SP} = \left[\epsilon^1_{\rm SP} + i\epsilon\epsilon_{\rm DP}\right] e^{i[\delta_{{}^3S_1}+\delta_{{}^1P_1}]},$$

$$x_{\rm DP} = \left[\epsilon_{\rm DP} + i\epsilon\epsilon^1_{\rm SP}\right] e^{i[\delta_{{}^3D_1}+\delta_{{}^1P_1}]}, \tag{11}$$

in terms of two *CP*-odd mixing angles $\epsilon^1_{\rm SP}$ and $\epsilon_{\rm DP}$. S is antisymmetric in the S-P and P-D elements due to time-reversal violation. The *CP*-odd mixing angles $\epsilon^{0,1}_{\rm SP}$ and $\epsilon_{\rm DP}$ are observable in, for example, spin rotation of polarized ultracold neutrons on a polarized hydrogen target [49], but it is unlikely that these experiments can reach a sensitivity that is competitive with EDM experiments, although neutron transmission experiments using heavy target nuclei might be up to the task [50,51]. Nuclear EDMs can be written as linear combinations of the mixing angles in addition to contributions from CP-odd electromagnetic currents, such as constituent nucleon EDMs.

The CP-odd mixing angles are observable and should be independent of the value of Λ up to NLO corrections. We find that this is the case for $\epsilon^1_{\rm SP}$ and $\epsilon_{\rm DP}$ which quickly converge as shown in Fig. 3(a). However, $\epsilon^0_{\rm SP}$ shows an oscillatory behavior and even changes sign as the function of Λ. There is no sign of convergence whatsoever. We have checked that no regulator dependence appears for any $j = 2$ transition after renormalizing the strong $j = 2$ scattering states. The difference between the behavior of ${}^1S_0$ - ${}^3P_0$ and ${}^3\{S, D\}_1$ - ${}^1P_1$ arises from the absence of a strong counterterm in the ${}^1P_1$ channel. The observed regulator dependence arises from divergences in diagrams contributing to $T_{\bar g_0}$ with topology of the left diagram in Fig. 4, where $V_{\bar g_0}$ is dressed on both sides by a strong short-distance interaction (an infinite number of LO diagrams are generated by adding additional strong interactions on either side). At LO this only occurs for ${}^1S_0$ - ${}^3P_0$ transitions. In χEFT calculations using Weinberg's power counting, P-wave counterterms appear at N^{2}LO but are iterated to all orders in the solution of the LS equation [52]. Divergent diagrams with the topology of Fig. 4 reappear, and the *CP*-odd

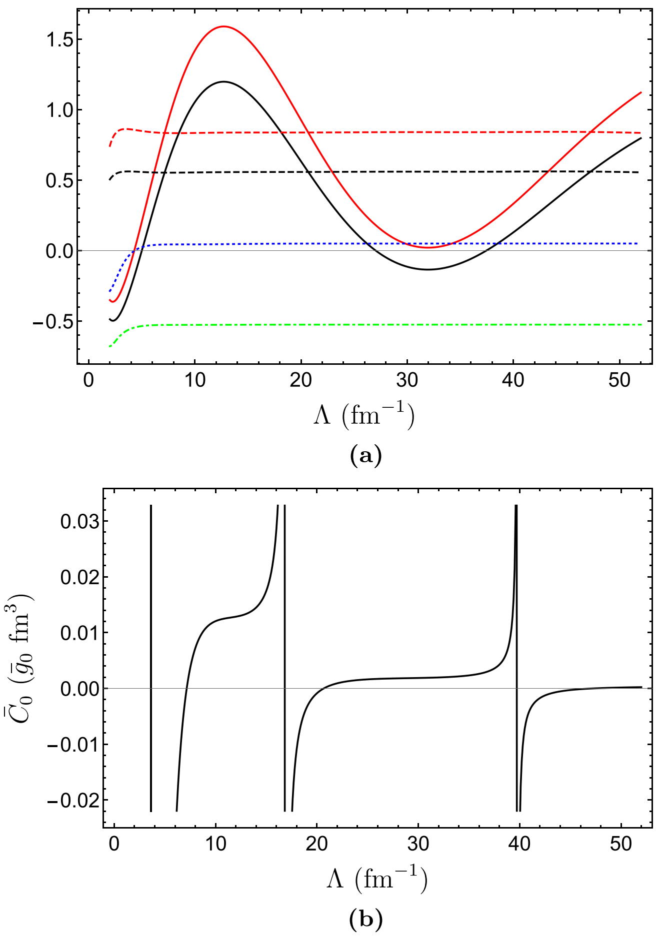

FIG. 3. (a) $\epsilon^0_{\rm SP}$ in units of $\bar{g}_0$ as function of Λ before (solid) and after (dashed) promoting $\bar{C}_0$ to LO for $E_{\rm CM} = \{25, 50\}$ MeV in, respectively, black and red. The blue (dotted) and green (dashed-dot) lines denote, respectively, $\epsilon^1_{\rm SP}$ and $\epsilon_{\rm DP}$ for $E_{\rm CM} = 50$ MeV. (b) $\bar{C}_0$ as a function of Λ.

transitions become regulator dependent. This might be hard to see numerically as regulators are only varied in a tiny window around $\Lambda = 500$ MeV [18,20].

The need for a counterterm. The observation that $\epsilon^0_{\rm SP}$ is cutoff dependent implies that *CP*-odd observables that depend on ${}^1S_0$ - ${}^3P_0$ mixing cannot be directly calculated from $\bar{g}_0$, and, thus, $\bar{\theta}$ via Eq. (3). An observable that shows regulator dependence in an EFT calculation indicates there must be an associated counterterm that encapsulates missing short-distance physics and absorbs the divergence. In the present context, such counterterms are provided by short-range *CP*-odd *NN* interactions, see the right diagram of Fig. 4 of the form [14,15]

$$\mathcal{L}_{NN} = \bar{C}_0\left[\bar{N}\boldsymbol{\sigma}N\cdot\nabla(\bar{N}N) + \frac{1}{3}\bar{N}\vec{\tau}\boldsymbol{\sigma}N\cdot\nabla(\bar{N}\vec{\tau}N)\right], \quad (12)$$

which projects on ${}^1S_0$ - ${}^3P_0$. $\bar{C}_0$ is a LEC that depends on Λ in such a way to make $\epsilon^0_{\rm SP}$ Λ independent. NDA suggests $\bar{C}_0 = O[m_\star\bar{\theta}/(F_\pi^2\Lambda_\chi^2)]$ and a N ^{2}LO contribution, but renormalization enhances $\bar{C}_0$ to LO.

We now show that promoting $\bar{C}_0$ to LO indeed renormalizes the ${}^1S_0$ - ${}^3P_0$ transition. We fit $\bar{C}_0$ at a specific kinematical point to a fictitious measurement of $\epsilon^0_{\rm SP}$, picking $\epsilon^0_{\rm SP,fit} = 0.01\bar{g}_0$ at $E_{\rm CM} = 5$ MeV for concreteness. The regulator dependence of $\bar{C}_0$ is shown in Fig. 3(b) and shows a limit-cycle-like behavior driven by C_P. The resulting $\epsilon^0_{\rm SP}$ is regulator independent for a wide range of energies as depicted by the dashed lines in Fig. 3(a). Although this method accounts for the regulator-dependent part of short-distance contributions and renormalizes the *CP*-odd amplitude, it cannot account for possible finite contributions from $\bar{C}_0$. That is, the results in Fig. 3(a) can shift up or down (they remain flat) if we were to pick different values for $\epsilon^0_{\rm SP,fit}$. The best way to obtain the total short-distance contribution is by fitting to a measurement of $\epsilon^0_{\rm SP}$. This is at present not possible, and even if there was data it would not be satisfactory. We would like to use such data to extract a value of $\bar{\theta}$.

Fixing the value of the short-distance LEC. We discuss two possible methods to obtain a value for $\bar{C}_0$ in the absence of data. The first one is to perform a LQCD calculation of $NN \to NN$ scattering in the presence of a nonzero $\bar{\theta}$ background. There have been significant developments in calculations of nucleon EDMs arising from the $\bar{\theta}$ term by applications of the gradient flow [12,53]. The same techniques could be used to study four-point functions in a $\bar{\theta}$ vacuum. A major challenge will be to control the signal to noise. Already for *CP*-conserving $NN \to NN$ processes, signal-to-noise considerations demand pion masses well above the physical point [54]. Going to smaller pion masses is even more daunting for the $\bar{\theta}$ term as the signal scales as $\approx\bar{\theta}m_\pi^2$. If such LQCD calculations are possible, we can obtain $\bar{C}_0$ from a matching calculation of χEFT to lattice data after taking the appropriate continuum and infinite-volume limits.

On a shorter timescale a promising approach is to apply chiral-symmetry relations between the $\bar{\theta}$ term and the quark masses similar to the relation between $\bar{g}_0$ and $\delta m_N^{\rm str}$ in Eq. (3). Using $SU(2)_L \times SU(2)_R$ χEFT, the operators in Eq. (12) arise from

$$\mathcal{L}_{NN} = -\frac{iC_0}{8}{\rm Tr}[\chi_-]\left[\bar{N}\boldsymbol{\sigma}N\cdot\nabla(\bar{N}N) + \frac{1}{3}\bar{N}\vec{\tau}\boldsymbol{\sigma}N\cdot\nabla(\bar{N}\vec{\tau}N)\right], \quad (13)$$

where $\chi_- = u^\dagger\chi u^\dagger - u\chi^\dagger u$, $u = \exp[i\vec{\tau}\cdot\vec{\pi}/(2F_\pi)]$, $\chi = 2B(\mathcal{M} + im_\star\bar{\theta})$, and $B = -\langle\bar{q}q\rangle/F_\pi^2$. Expanding the trace gives $\bar{C}_0 = (Bm_\star\bar{\theta})C_0$ and a relation to the *CP*-conserving,

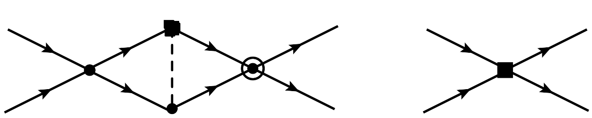

FIG. 4. Left: Diagram contributing to the regulator dependence of $\epsilon^0_{\rm SP}$. Solid (dashed) lines denote nucleons (pions). The square denotes $\bar{g}_0$ whereas the circles denote the g_A or C_s vertices. The circled circle denotes C_P. Right: short-distance contribution proportional to $\bar{C}_0$.

isospin-breaking $NN\pi$ operators [28],

$$\mathcal{L}_{NN,\pi} = \frac{C_0 B(m_d - m_u)}{2}\frac{\pi_0}{F_\pi}\left[\bar{N}\boldsymbol{\sigma}N\cdot\nabla(\bar{N}N) + \frac{1}{3}\bar{N}\vec{\tau}\boldsymbol{\sigma}N\cdot\nabla(\bar{N}\vec{\tau}N)\right]. \quad (14)$$

These operators contribute to charge-symmetry-breaking (CSB) in $NN \to NN\pi$ processes [55–58]. A LO contribution to this CSB process arises from the $N\pi\pi$ vertex related to $\delta m_N^{\rm str}$ by chiral symmetry,

$$\mathcal{L}_{\rm CSB} = -\frac{\delta m_N^{\rm str}}{4F_\pi^2}\bar{N}\vec{\tau}\cdot\vec{\pi}\pi_0 N. \quad (15)$$

The contact operator in Eq. (14) contributes at N^{2}LO in Weinberg's counting. At the pion threshold, the transition operator for the process $^1S_0$ - $^3P_0 + \pi$ due to Eq. (15) is of the same form as $V_{\bar{g}_0}$. As such, the regulator dependence seen in Fig. 3 appears, and C_0 must be promoted to LO for renormalization. Unfortunately the simplest process where CSB data are available $pn \to d\pi^0$ is not sensitive to C_0 due to the isosinglet nature of the deuteron. This motivates an investigation of $dd \to \alpha\pi^0$ using renormalized χEFT to fit C_0 to CSB data [59] to directly obtain $\bar{C}_0 = (Bm_\star\bar{\theta})C_0$.

Other sources of CP or P violation. At the dimension-six level there appear other *CP*-odd sources. For the present discussion the most relevant operators are quark chromo-EDMs and chiral-breaking four-quark operators, which are induced in a wide range of BSM models [30,60]. In addition to the isoscalar $\bar{g}_0$ term in Eq. (2), the LO *CP*-odd chiral Lagrangian contains an isovector term,

$$\mathcal{L}_{\pi N} = \bar{g}_1\bar{N}\pi_0 N, \quad (16)$$

whereas a potential isotensor term is subleading [14]. In combination with the strong g_A vertex, an OPE involving $\bar{g}_1$ causes $^1S_0$ - $^3P_0$ and $^3S_1$ - $^3P_1$ transitions. Strong $^3P_1$ interactions arise solely from OPE, and the divergent diagrams in Fig. 4 do not appear. We expect no regulator dependence for $^3S_1$ - $^3P_1$ transitions, which is confirmed by explicit calculations. The $j = 0$ transition, up to an isospin factor, shows the same regulator dependence as the $\bar{g}_0$ case, and, thus, a LO isospin-breaking counterterm is needed. The associated operator takes the form

$$\mathcal{L}_{NN} = \bar{C}_1[\bar{N}\tau^3\boldsymbol{\sigma}N\cdot\nabla(\bar{N}N) + \bar{N}\boldsymbol{\sigma}N\cdot\nabla(\bar{N}\tau^3 N)], \quad (17)$$

which projects unto $^1S_0$ - $^3P_0$, but only for the neutron-neutron and proton-proton cases. The simplest EDM that depends on $\bar{g}_1$ is the deuteron EDM [61] targeted in storage-ring experiments [62]. Due to the isosinglet nature of the deuteron, its EDM only depends on $^3S_1$ - $^3P_1$ transitions which do not require a counterterm. There is no such selection rule for more complex EDMs, such as ^{3}He, ^{199}Hg, or ^{225}Ra [17–20,63,64], and $\bar{C}_1$ must be included at LO.

The finiteness of $^3S_1$ - $^3P_1$ transitions is relevant for the field of hadronic parity (P) violation [65]. The LO P odd, but *CP* even, chiral Lagrangian induced by P-odd four-quark operators contains a single πN term [66], usually parametrized as $(h_\pi/\sqrt{2})\bar{N}(\vec{\pi}\times\vec{\tau})^3 N$ that in combination with g_A leads to $^3S_1$ - $^3P_1$ transitions [67,68]. We have checked explicitly that no regulator dependence appears and no counterterms are needed. The value of h_π recently determined from P-violating asymmetries in $\vec{n}p \to d\gamma$ [69] can, thus, be directly applied in calculations of other P-odd observables.

Conclusion. We have argued the need for a leading-order short-range *CP*-violating counterterm in $^1S_0$ - $^3P_0$ transitions that affects calculations of EDMs and *CP* violation in nucleon-nucleon and neutron-nucleus scatterings at the $O(1)$ level. This directly affects the interpretation of experimental limits, and hopefully future signals, in terms of the QCD $\bar{\theta}$ term and other *CP*-odd sources, and the interpretation of axion DM searches via oscillating EDMs. For *CP* violation from the $\bar{\theta}$ term, we have proposed strategies to obtain the value of the associated low-energy constant $\bar{C}_0$ from existing data on charge-symmetry breaking in few-body systems. We hope our results stimulate determinations of $\bar{C}_0$ using lattice QCD, analyses of CSB data, and calculations of the impact of the short-range operator on observables of experimental interest, such as (oscillating) EDMs, magnetic quadrupole moments, and time-reversal-odd scattering observables.

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