Editors' Suggestion

Standard Model Prediction for Paramagnetic Electric Dipole Moments

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Standard model *CP* violation associated with the phase of the Cabibbo-Kobayashi-Maskawa quark mixing matrix is known to give small answers for the electric dipole moment (EDM) observables. Moreover, predictions for the EDMs of neutrons and diamagnetic atoms suffer from considerable uncertainties. We point out that the *CP*-violating observables associated with the electron spin (paramagnetic EDMs) are dominated by the combination of the electroweak penguin diagrams and $\Delta I = 1/2$ weak transitions in the baryon sector, and are calculable within chiral perturbation theory. The predicted size of the semileptonic operator C_S is 7×10^{-16}, which corresponds to the *equivalent* electron EDM $d_e^{\mathrm{eq}} = 1.0 \times 10^{-35}\ e\,\mathrm{cm}$. While still far from the current observational limits, this result is 3 orders of magnitude larger than previously believed.

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Introduction.—The searches for the electric dipole moments (EDMs) of elementary particles [1–4] represent an important way of probing the TeV scale new physics [5–7]. Recent breakthrough sensitivity to *CP* violation connected to electron spin (which we will refer to as "paramagnetic EDMs") [3] established a new limit on the linear combination of the electron EDM d_e and semi-leptonic nucleon-electron $\bar{N}N\bar{e}i\gamma_5 e$ operators, commonly parametrized by a C_S coefficient. The rapid progress of the last decade, as well as some additional hopes for increased accuracy (see, e.g., [8–10]), makes one revisit the standard model (SM) sources of *CP* violation, and the expected size of the paramagnetic EDMs in the SM.

The SM has two sources of *CP* violation. The first source, undetected thus far, corresponds to the nonperturbative effects parametrized by the QCD vacuum angle θ. Recently it has been shown [11] that paramagnetic EDMs are dominated by the two-photon exchange mechanism, and the leading chiral behavior of the hadronic part of the diagram is given by the t-channel exchange by π^0, η. *CP* violation due to θ comes through the $\pi^0(\eta)\bar{N}N$ coupling. The result, in combination with the experimental bound [3], sets the independent limit on $|\theta| < 3 \times 10^{-8}$, which is still subdominant to the limit provided by $d_n(\theta)$.

The second source of the SM *CP* violation is the celebrated Kobayashi-Maskawa (KM) phase δ_{KM} [12], which is now observed to rather good accuracy in a plethora of flavor transitions in B and K mesons. Observations are often matched by rather precise theoretical predictions, starting from [13]. The predictions of EDM-like observables induced by δ_{KM} thus far can be summarized by two adjectives: small and uncertain. The suppression comes from the necessity to involve at least two W bosons and multiple loops [14–16] involving all three generations of quarks. As a result, short distance contributions to quark EDMs do not exceed $10^{-33}\ e\,\mathrm{cm}$ level [17]. At the same time, it is clear that long-distance nonperturbative contributions, typically described as a combination of two transitions changing strangeness by one unit, $\Delta S = \pm 1$, dominate d_n and nucleon-nucleon forces [18–22]. A more recent estimate [23] places d_n in the ballpark of few $\times 10^{-32}\ e\,\mathrm{cm}$ with a wide order-of-magnitude expected range. It is fair to say that magnitudes of d_n and nucleon-nucleon forces (that feed into the nuclear-spin-dependent atomic EDMs) cannot be accurately predicted at this point.

What is the size of paramagnetic EDMs induced by δ_{KM}? Recent estimates of d_e [24] (dominated again by long-distance effects) converge at the tiniest value of $\sim 6 \times 10^{-40}\ e\,\mathrm{cm}$, presumably with considerable uncertainties corresponding to hadronic modeling of quark loops. This result is subdominant to the C_S estimate due to the two-photon exchange mechanism in combination with $\Delta S = \pm 1$ transitions [25], which corresponds to equivalent d_e of $\sim 10^{-38}\ e\,\mathrm{cm}$. To introduce useful notations, this is the $\mathrm{EW}^2\mathrm{EM}^2$ order effect, where EW and EM stand for electroweak and electromagnetic, respectively.

In this Letter, we demonstrate that the dominant contribution to paramagnetic EDMs associated with the KM

CP violation is given by the semileptonic C_S induced in EW^3 order. It has an unambiguous answer in the flavor $SU(3)$ chiral limit and is calculable to $\sim 30\%$ accuracy, which can be further improved. Remarkably, the result reaches the level of $\sim 10^{-35}\ e\,\mathrm{cm}$ in terms of the d_e equivalent, which is 3 orders of magnitude larger than previously believed [25].

Our starting point is the expression for the *equivalent* d_e that follows from atomic and molecular theory, and defines the linear combination of two Wilson coefficients constrained by the most precise paramagnetic EDM measurements performed with the ThO molecule:

$$d_e^{\mathrm{equiv}} = d_e + C_S \times 1.5 \times 10^{-20}\ e\,\mathrm{cm}, \tag{1}$$

where e is the positron charge. Current experimental limit [3] stands as $|d_e^{\mathrm{equiv}}| < 1.1 \times 10^{-29}\ e\,\mathrm{cm}$. As per convention, C_S is defined with the Fermi constant factored out, and γ_5 corresponds to the $\frac{1}{2}\gamma_\mu(1-\gamma_5)$ definition of the left-handed current

$$\mathcal{L}_{eN} = C_S \frac{G_F}{\sqrt{2}} (\bar{e} i\gamma_5 e)(\bar{p}p + \bar{n}n). \tag{2}$$

Our goal is to calculate $C_S(\delta_{\mathrm{KM}})$.

Leading chiral order C_S calculation.—Because of the conservation of the electron chirality in the SM, it is clear that $C_S \propto m_e$. This in turn rules out single-photon exchange (EM penguin) as origin of $m_e \bar{e} i\gamma_5 e$, and one would need either a two-photon mechanism [11,25] or the EW penguin Z-boson exchange and W-box diagram. The most crucial property of EW penguins is that, although they are formally of the second order in weak interactions, their size is enhanced by the heavy top, so that the result scales as $G_F^2 m_t^2$. EW penguins (as is well known, EW penguins must also include W-box diagrams, and we include both) induce $B_{s,d} \to \mu^+\mu^-$ decays, and dominate the dispersive part of $K_L \to \mu^+\mu^-$ amplitude. Dropping the vector part of the lepton current (as not leading to $m_e \bar{e} i\gamma_5 e$), and integrating out heavy W, Z, t particles, one can concisely write the semileptonic operator as

$$\mathcal{L}_{\mathrm{EWP}} = -\mathcal{P}_{\mathrm{EW}} \times \bar{e}\gamma_\mu\gamma_5 e \times \bar{s}\gamma^\mu(1-\gamma_5)d + (\mathrm{H.c.}), \tag{3}$$

where

$$\mathcal{P}_{\mathrm{EW}} = \frac{G_F}{\sqrt{2}} \times V_{ts}^* V_{td} \times \frac{\alpha_{\mathrm{EM}}(m_Z)}{4\pi \sin^2\theta_W} I(x_t), \tag{4}$$

and the loop function is given by [26]

$$I(x_t) = \frac{3}{4}\left(\frac{x_t}{x_t-1}\right)^2 \log x_t + \frac{1}{4}x_t - \frac{3}{4}\frac{x_t}{x_t-1}, \qquad x_t = \frac{m_t^2}{m_W^2}. \tag{5}$$

These results are well established, and unlike the case of four-quark operators, the subsequent renormalization group evolution of Eq. (3) introduces only small corrections (see, e.g., [27,28]). This is because the QCD evolution is trivial (apart from small threshold corrections at m_W) due to the partially conserved nature of the quark current, and QED evolution is small $\propto \alpha_{\mathrm{EM}}/\pi$.

The most convenient representation of the Cabibbo-Kobayashi-Maskawa (CKM) matrix is when δ_{KM} enters mostly in V_{td}. It enters the imaginary part of $\mathcal{P}_{\mathrm{EW}}$ and couples the axial vector current of leptons to the $\bar{s}\gamma_\mu(1-\gamma_5)d - \bar{d}\gamma_\mu(1-\gamma_5)s$ quark current. This current can create or annihilate CP-even combination of the neutral kaons that (in neglection of small ϵ_K) can be identified with K_S field. Same operator in the muon channel induces $K_S \to \mu^+\mu^-$ meson decay [29,30]. Within chiral perturbation theory, the axial vector current of leptons is treated as an external left-handed current, which gives rise to

$$\mathcal{L}_{Uee} = -\frac{if_0^2}{2}\mathcal{P}_{\mathrm{EW}} \times \bar{e}\gamma_\mu\gamma_5 e \times \mathrm{Tr}[h^\dagger(\partial^\mu U)U^\dagger] + (\mathrm{H.c.}), \tag{6}$$

where U is the exponential of the meson octet M, $U = \exp[2iMf_0^{-1}]$, in our convention it transforms as $U' = LUR^\dagger$, and $h_{ij} = \delta_{i2}\delta_{j3}$. At linear order, this leads to $\partial_\mu K \times \bar{e}\gamma^\mu\gamma_5 e$, and upon application of the equation of motion for electrons we arrive at

$$\mathcal{L}_{Kee} = -2\sqrt{2} f_0 m_e \bar{e} i\gamma_5 e (K_S \times \mathrm{Im}\mathcal{P}_{\mathrm{EW}} + K_L \times \mathrm{Re}\mathcal{P}_{\mathrm{EW}}). \tag{7}$$

In this expression, f_0 is the meson coupling constant, which in the $SU(3)$ symmetric limit is equal to $\simeq 134$ MeV, and we follow Ref. [31] conventions. Subsequent m_s-dependent corrections renormalize this coupling to $f_0 \to f_K \simeq 160$ MeV. While other s-quark containing resonances may also contribute, the neutral kaon exchange, Fig. 1, will give the only m_s^{-1}-enhanced contribution in the chiral limit.

We now need to find out how the neutral kaons couple to the nucleon scalar densities, $\bar{p}p$ and $\bar{n}n$, that occur due to $\Delta S = \pm 1$ transitions in the EW^1 order. Instead of attempting such calculation from first principles (see, e.g., [32]) we will use flavor $SU(3)$ relations and connect this coupling to the s-wave amplitudes of hadronic decays of strange hyperons, following [31]. It is well known that empirical $\Delta I = 1/2$ rule holds for hyperon decays, and the leading order $SU(3)$ relations fit s-wave amplitudes with $O(10\%)$ accuracy. It is strongly suspected that these amplitudes are indeed induced by strong penguins (SPs), although this assumption is not crucial for us. With that, one can write the two types of couplings consistent with $(8_L, 1_R)$ transformation properties:

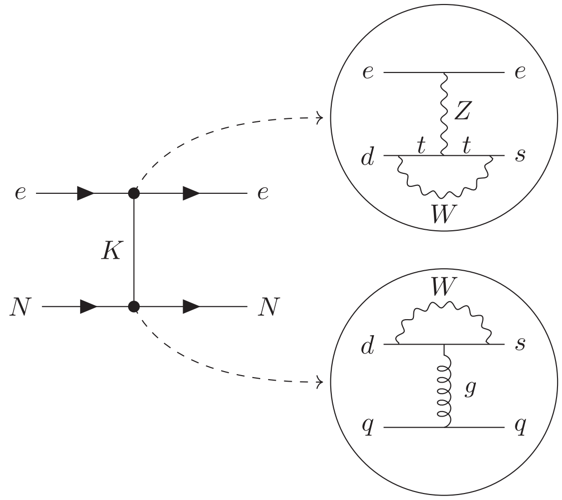

FIG. 1. EW^3 order diagram that dominates in the chiral limit. The top vertex is the CP-odd, P-even $K_S\bar{e}i\gamma_5 e$ generated in EW^2 order, and the bottom vertex is CP-even, P-odd $K_S\bar{N}N$ coupling generated at EW^1 order.

$$\mathcal{L}_{\rm SP} = -a\mathrm{Tr}(\bar{B}\{\xi^\dagger h\xi, B\}) - b\mathrm{Tr}(\bar{B}[\xi^\dagger h\xi, B]) + (\mathrm{H.c.}). \quad (8)$$

In this expression, B is the baryon octet matrix, and $\xi = \exp[iMf_0^{-1}]$. Assuming a and b to be real, and taking $f_0 = f_\pi$, they are fit by [31] to be [33]

$$a = 0.56 G_F f_\pi \times [m_{\pi^+}]^2; \qquad b = -1.42 G_F f_\pi \times [m_{\pi^+}]^2. \quad (9)$$

Brackets over m_{π^+} indicate that these are numerical values taken, 139.5 MeV, rather than $m_u + m_d$-proportional theoretical quantity m_π. These values can be easily found via the least square fit to the nonleptonic s-wave amplitude, which also indicates 10% theoretical accuracy of this fit. In the assumption of a and b being real, only the K_S meson couples to nucleons, $2^{1/2}f_0^{-1}[(b-a)\bar{p}p + 2b\bar{n}n]K_S$, which will provide the dominant contribution. This type of coupling breaks P but respects CP symmetry. Restoring the CKM factors, one can also include much subdominant coupling to K_L so that we have

$$\mathcal{L}_{KNN} \simeq -\frac{\sqrt{2}G_F \times [m_{\pi^+}]^2 f_\pi}{|V_{ud}V_{us}|f_0} \times 2.84(0.7\bar{p}p + \bar{n}n) \times (\mathrm{Re}(V_{ud}^* V_{us})K_S + \mathrm{Im}(V_{ud}^* V_{us})K_L). \quad (10)$$

At the last step, we integrate out the K mesons as shown in Fig. 1. Adopting it for a nucleus containing $A = Z + N$ nucleons, one arrives at a straightforward prediction for the $\delta_{\rm KM}$-induced size of the electron-nucleon interaction:

$$C_S \simeq \mathcal{J} \times \frac{N + 0.7Z}{A} \times \frac{13[m_{\pi^+}]^2 f_\pi m_e G_F}{m_K^2} \times \frac{\alpha_{\rm EM} I(x_t)}{\pi \sin\theta_W^2}, \quad (11)$$

where $\mathcal{J}$ is the rephasing invariant combination of the CKM angles,

$$\mathcal{J} = \mathrm{Im}(V_{ts}^* V_{td} V_{ud}^* V_{us}) \simeq 3.1 \times 10^{-5}, \quad (12)$$

which carries about $\sim 6\%$ uncertainty. Notice that the f_0 factor in the numerator of Eq. (7) cancels against f_0 in the denominator of Eq. (10), and this cancellation would persist even one changes f_0 for f_K.

The overall scaling of this formula in the chiral limit and at large x_t is

$$G_F C_S \propto \mathcal{J} G_F^3 m_t^2 m_e m_s^{-1} \Lambda_{\rm hadr}^2, \quad (13)$$

where $\Lambda_{\rm hadr}$ is a typical hadronic energy or momentum scale. Notice that this is far more singular behavior with m_q of a light quark than that arising in the chiral-loop-induced expressions for d_n. Also notice that the K_S exchange dominates for any conventional parametrization of the CKM matrix, and the role of K_L exchange is to add small pieces of the amplitude that take $\mathrm{Re}(V_{ud}V_{us}^*)\mathrm{Im}(V_{ts}V_{td}^*)$, arising from K_S exchange, to full $\mathcal{J}$. Substituting all SM parameters, we obtain the following leading order (LO) result:

$$C_S(\mathrm{LO}) \simeq 5 \times 10^{-16}. \quad (14)$$

In order to estimate accuracy of the LO $\sim O(m_s^{-1})$ result, one could try to evaluate the next-to-leading-order (NLO) corrections in the expansion over small m_s. These corrections can be divided into two groups: (1) corrections to the $K\bar{N}N$ vertex at $m_s \log m_s$ order and (2) diagrams that do not reduce to the t-channel K-meson exchange. Type (1) corrections involve essentially the same diagrams as those appearing in the corresponding corrections to the s-wave hyperon decays [31,35]. The analysis of Ref. [35] showed that when the loop corrections are included with the tree-level a and b parameters and the total theoretical result is fit to experimental data, one notices that the tree-level values for a and b come out smaller than in Eq. (9), while the *total* result is rather close to the tree-level fit for a, b. This comes mostly from the renormalization of the meson and baryon wave functions. The lesson from this is that the corrections of type (1) for KNN weak coupling are expected to mirror results of Ref. [35] for s-wave amplitudes, and therefore would not deviate substantially from Eq. (10).

We then estimate type (2) corrections. It turns out that they parametrically dominate over other types of corrections, as the baryon pole diagrams, Fig. 2, contribute. The m_s scaling of these corrections is set by the ratio of the loop integral, proportional to m_K (at $m_K^2 \gg m_\pi^2$ limit), divided by mass splitting Δm_B in the baryon octet, e.g., $m_\Lambda - m_n$. This quantity scales as $m_s^{-1/2}$ and therefore these baryon pole diagrams dominate the NLO contributions in the chiral limit. They are fully calculable (i.e., do not depend on

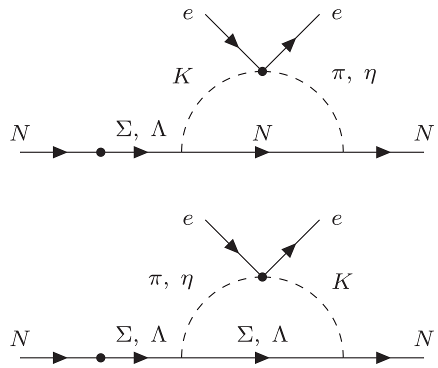

FIG. 2. The baryon pole diagrams that contribute to C_S at the NLO level in the chiral limit. The left vertex is the nucleon-hyperon mixing induced by Eq. (8), while the top vertex is induced by Eq. (6). The vertices without black dots are the strong interaction with the coupling constants D and F. The diagrams with the nucleon-hyperon mixing on the right side give the same amount of contribution.

unknown counterterms), and the results for these corrections are

$$\frac{C_{S,\mathrm{NLO}}(p)}{C_{S,\mathrm{LO}}(p)} = \frac{m_K^3(0.77D^2 + 2.7DF - 2.3F^2)}{24\pi f_0^2(m_{\Sigma^+} - m_p)} \quad (15)$$

$$\frac{C_{S,\mathrm{NLO}}(n)}{C_{S,\mathrm{LO}}(n)} = \frac{m_K^3}{24\pi f_0^2}\left[\frac{(a/b+3)}{2\sqrt{6}(m_\Lambda - m_n)}(-0.44D^2 + 3.2DF + 1.3F^2) + \frac{a/b-1}{2\sqrt{2}(m_{\Sigma^0} - m_n)}(-0.53D^2 - 1.9DF + 1.6F^2)\right]. \quad (16)$$

It has been obtained using heavy baryon chiral perturbation theory, and D, F are the coupling constants characterizing the strength of the $SU(3)$-invariant baryon-meson strong interaction, with $F = 0.46$, $D = 0.8$ typically used [31]. Since the dominant contribution comes from loops with $K - \pi$ transition, it is appropriate to take $f_0^2 \simeq f_\pi f_K$. Using these numbers, we discover that NLO corrections interfere constructively with LO, and give 30% correction for the proton, and 40% for the neutron, correspondingly. Combining LO and NLO, we arrive at our final result:

$$C_S(\mathrm{LO} + \mathrm{NLO}) \simeq 6.9 \times 10^{-16} \Rightarrow d_e^{\mathrm{equiv}} \simeq 1.0 \times 10^{-35}\ e\,\mathrm{cm}. \quad (17)$$

The size of the NLO corrections also allows us to estimate the accuracy of this computation as $O(30\%)$.

As stated in the introduction, this result is much larger than previously believed, and exceeds any contributions of d_e into d_e^{equiv} by at least 4 orders of magnitude. The enhancement of C_S at EW^3 order compared to $\mathrm{EW}^2\mathrm{EM}^2$ can be roughly ascribed to $\alpha_W/\alpha_{\mathrm{EM}}^2 \sim O(10^3)$. We note that, although translating C_S to d_e^{equiv} depends on the atoms or molecules that one considers (ThO above), this dependence is mild and d_e^{equiv} is within the same ballpark if we instead consider, e.g., Tl or YbF [25].

Stepping away from chiral expansion, one can formulate the necessary hadronic matrix element that will be required to generate C_S in combination with the dominant $\mathrm{Im}\mathcal{P}_{\mathrm{EW}}$ channel of Eq. (3). The corresponding d-to-s transitions need to be taken in the first order, EW^1, that break P and C separately but conserve CP:

$$\langle N|i(\bar{s}\gamma_\mu(1-\gamma_5)d - \bar{d}\gamma_\mu(1-\gamma_5)s)|N\rangle_{\mathrm{EW}^1} = \frac{f_S}{m_N} i q_\mu \bar{N}N + \frac{f_T}{m_N} q_\nu \bar{N}\sigma_{\mu\nu}\gamma_5 N. \quad (18)$$

In this formula, q_μ stands for the momentum transfer. It turns out that there are only two form factors on the right-hand side of this expression that have the same CP properties as the left-hand side. Moreover, f_T in combination with $\mathrm{Im}\mathcal{P}_{\mathrm{EW}}$ leads to CP-odd P-even interactions that do not induce EDMs. Therefore only f_S form factor (that sometimes is called induced scalar) at $q^2 \to 0$ is relevant. We have provided the first two terms in the chiral expansion of f_S for neutrons and protons, so that effectively $f_S \propto a(b) \times m_N m_K^{-2} + \cdots$. While we use chiral perturbation theory, in principle, calculation of f_S can be attempted using lattice QCD methods.

Finally, we note that other semileptonic operators such as $\bar{e}e\bar{N}i\gamma_5 N$ that lead to nuclear-spin-dependent effects are not generated the same way at EW^3 order and therefore will be suppressed compared to Eq. (11).

Conclusions.—We have shown that δ_{KM} induces the CP-odd electron nucleon interaction at the level much larger than previous estimates [25]. The main mechanism is not a two-photon exchange, $\mathrm{EW}^2\mathrm{EM}^2$, between electron and the nucleus, but the combination of a weak nonleptonic EW^1 transition with the semileptonic EW^2 electroweak penguin. Although the result is still small, it is not unthinkable that the progress in sensitivity to paramagnetic EDMs may reach the level of d_e^{equiv} in the future. Indeed, some novel proposals [8] envision that statistical sensitivity to paramagnetic EDMs can be brought down to $d_e \sim O(10^{-35}$–$10^{-37})$ e cm.

It is not surprising that the C_S operator can be predicted, at least in the $SU(3)$ chiral expansion, rather precisely. This clearly distinguishes our C_S calculation from $d_n(\delta_{\mathrm{KM}})$ estimate that carries an order of magnitude uncertainty with unclear prospects for improvement. In contrast, the only significant source of uncertainty in C_S is in the induced scalar form factor, Eq. (18), that can be improved in the future with the use of lattice QCD methods.

Even if one takes chiral $SU(3)$ expansion skeptically, it is clear that unique m_s^{-1} (LO) and $m_s^{-1/2}$ (NLO) contributions to C_S identified in our Letter would not be cancelled —unless completely accidentally—by other contributions, mirroring a similar argument of [36] made for $d_n(\theta)$. Therefore, 10^{-35} e cm should be adopted as the robust δ_{KM}-induced SM benchmark value for all experiments attempting the search of d_e using electron spins in heavy atoms and molecules. It also allows for establishing the *maximum* sensitivity to *CP*-violating new physics via d_e. Taking a one-loop perturbative scaling, $d_e \propto (\alpha/\pi) m_e \Lambda_{\mathrm{NP}}^{-2}$, and equating it to $d_e^{\mathrm{equiv}}(\delta_{\mathrm{KM}})$, one arrives at the maximum scale that is possible to probe with paramagnetic EDMs: $\Lambda_{\mathrm{NP}}^{\max} \sim 5 \times 10^7$ GeV. Notice, however, that in models with no chiral m_e suppression of d_e and/or tree-level C_S generation by new physics, the ultimate scale can be larger [37,38].

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