\[ e' \text{ from right-handed currents} \]

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Recent lattice determinations of direct CP violation in kaon decays, \( e' \), has been precisely measured over a decade ago\(^{1,2,3}\), the corresponding standard model (SM) predictions have not yet reached a comparable precision. Such a prediction is a challenging task that requires the calculation of nonperturbative matrix elements. Recently, these matrix elements have been determined using lattice QCD\(^4\), and suggest a \( 2 - 3 \sigma \) discrepancy between the SM and the experimental value. This discrepancy is in agreement with the results of Refs.\(^5,6,7\), while several analytic approaches find values for \( e' \) that are consistent with experiment\(^8,9,10\).

Assuming that this tension survives future improved lattice determinations, it is interesting to investigate possible explanations in terms of new physics. Several explanations in terms of vector-like quarks\(^{11}\), 331 models\(^{12}\), \( Z(t) \) couplings\(^{13}\), and supersymmetric scenarios\(^{14,15,16}\) have been discussed in the literature. Here we investigate a scenario involving a single gauge-invariant dimension-six operator\(^{17}\). This operator induces right-handed charged currents (RHCCs), which couple the \( W \) boson to right-handed quarks, and is given by

\[
\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \frac{2}{v^2} i \bar{\varphi}^3 D_\mu \varphi \bar{u}_R^i \gamma^\mu \xi_{ij} \bar{d}_R^j + \text{h.c.}, \quad \rightarrow \quad \mathcal{L}_{\text{SM}} + \frac{g}{\sqrt{2}} \left[ \xi_{ij} \bar{u}_R^i \gamma^\mu \bar{d}_R^j W^+_{\mu} \right] \left( 1 + \frac{h}{v} \right)^2 + \text{h.c.}, \quad (1)
\]

where \( \varphi \) is the Higgs doublet, \( v \approx 246 \) GeV is its vacuum expectation value (vev), and the covariant derivative is given by \( D_\mu = \partial_\mu - ig/2 \tau \cdot W_\mu - ig'/2B_\mu \), with \( g \) and \( g' \) the \( SU(2) \) and \( U(1)_Y \) gauge couplings. Finally, \( \xi_{ij} \) is a \( 3 \times 3 \) matrix in flavor space, whose elements are expected to scale as \( v^2/\Lambda^2 \), where \( \Lambda \) is the scale of new physics. This interaction is generated in left-right symmetric models\(^{18,19}\), which are based on the gauge group \( SU(2)_L \times SU(2)_R \times U(1)_{B-L} \). These models feature a new right-handed \( W_R \) boson which can undergo mass-mixing with the SM \( W \) boson. After integrating out the heavy \( W_R \) this mixing induces Eq. 1. Explicitly one has, \( \xi_{ij} \approx \frac{g^2}{2} \frac{\kappa \kappa'}{m_R^2} e^{i \alpha} (V_R)_{ij} \), where \( \kappa, \kappa' \sim v \) are the magnitudes of the vevs that break \( SU(2)_L \) and \( \alpha \) is the phase difference between them, \( m_R \) and \( g_R \) are mass of the \( W_R \) boson and its gauge coupling, while \( V_R \) is the right-handed analogue of the CKM matrix.

Here we do not restrict to a specific model and focus on two elements of the \( \xi \) matrix\(^a\), namely \( \xi_{ud} \) and \( \xi_{us} \). After integrating out the \( W \) boson these elements induce both \( \Delta S = 1 \) and

\(^a\)For the phenomenology of the remaining flavor structures see\(^{20}\).
\( \Delta S = 0 \) four-quark operators, which contribute to \( \epsilon' \) and hadronic EDMs, respectively. Both contributions depend on nonperturbative matrix elements. In the case of \( \epsilon' \) chiral symmetry allows one to relate the necessary matrix elements to those of SM operators calculated on the lattice. Although the EDM analysis depends on several nonperturbative quantities, not all of which are known, we argue that the same matrix elements allow one to estimate the leading contributions in this case as well.

We discuss the low-energy Lagrangian induced by Eq. 1 in section 2. The impact of this Lagrangian on \( \epsilon' \) and hadronic EDMs are derived in section 3 and 4, respectively. We discuss the resulting constraints and the possibility of a solution to the \( \epsilon' \) discrepancy in section 5.

2 Low-energy Lagrangian

2.1 Quark-level Lagrangian

After integrating out the heavy SM fields the couplings \( \xi_{ud} \) and \( \xi_{us} \) of Eq. 1 give rise to the following four-quark interactions

\[
\mathcal{L}_{LR} = -\sum_{i=1}^{2} \left( C_{iLR}^{ud} \mathcal{O}_{iLR}^{ud} + C_{iLR}^{us} \mathcal{O}_{iLR}^{us} + C_{iLR}^{us} \mathcal{O}_{iLR}^{ud} + C_{iLR}^{ud} \mathcal{O}_{iLR}^{us} + \text{h.c.} \right),
\]

(2)

where, \( \mathcal{O}_{iLR}^{ijlm} = \bar{q}^i \gamma^\mu P_L t^j \bar{q}^l \gamma^\mu P_R d^m \), and \( \mathcal{O}_{2LR}^{ij} = \bar{q}^i \gamma^\mu P_L t^j \bar{q}^l \gamma^\mu P_R d^m \), with \( \alpha, \beta \) color indices. These ‘left-right’ operators violate CP as long as the corresponding couplings have an imaginary part. The first two \( \Delta S = 0 \) operators in Eq. 2 will induce EDMs, while the second pair violate strangeness by one unit and contribute to \( \epsilon' \). Note that no \( \Delta F = 2 \) operators are generated at tree level. The matching at the \( W \) boson mass scale gives,

\[
C_{1LR}^{ijlm}(m_W) = \frac{4G_F}{\sqrt{2}} V_{lm}^* \xi_{ij}, \quad C_{2LR}^{ijlm}(m_W) = 0,
\]

(3)

while the \( \mathcal{O}_{2LR} \) operators are induced through QCD renormalization.

After evolving the above Lagrangian to \( \mu \approx 3 \, \text{GeV} \), the contributions to \( \epsilon' \) and EDMs still require the matrix elements of the left-right operators, which we obtain from ChiPT. Before moving on to the chiral realization of the left-right operators, however, we slightly rewrite the relevant parts of the effective Lagrangian,

\[
\mathcal{L} = \mathcal{L}_{m_q=0}^{\text{QCD}} - \bar{q} \mathcal{M} q + \bar{q} \left[ m_3 t_3 + m_6 t_6 + m_8 t_8 \right] i \gamma_5 q + \mathcal{L}_{LR},
\]

(4)

where \( t^a \) are the \( SU(3) \) generators, \( q \) is a triplet of quark fields \( q = (u, d, s) \), and \( \mathcal{M} = \text{diag}(m_u, m_d, m_s) \). In addition, we assumed that the strong CP problem is solved by a Peccei-Quinn mechanism. Finally, the left-right operators induce couplings which couple the neutral mesons, \( \pi^0, K^0, \eta \), to the vacuum. To avoid such couplings when constructing the Chiral Lagrangian we perform a \( SU(3)_L \times SU(3)_R \) rotation to eliminate these terms, this introduces \( m_{3,6,8} \) which are specified in the next section.

2.2 Chiral Lagrangian

To construct the chiral Lagrangian it is useful to note that the left-right operators can schematically be written as, \( (\bar{q}^\gamma \mu t^a e^\beta) \mathcal{P}_{LR} \). Such operators belong to the \( 8L \times 8R \) representation of \( SU(3)_L \times SU(3)_R \) and so transform as \( t^a \rightarrow L^a L^1, t^b \rightarrow R^b R^1 \), where \( L, R \in SU(3)_{LR} \).

Using these transformation properties and the well-known chiral realization of the usual QCD Lagrangian, the leading-order mesonic Lagrangian is given by

\[
\mathcal{L}_\pi = \frac{F_0^2}{4} \text{Tr} \left( \partial_\mu U \partial^\mu U^\dagger \right) + \frac{F_0^2}{4} \text{Tr} \left( U \chi^\dagger + U^\dagger \chi \right)
\]
up to isospin corrections, which gives $Q_I$ for $\delta a_1 - i\delta a_2)(\delta b_1 + i\delta b_2) + C_{i LR}^{us \, us}(\delta a_1 - i\delta a_2)(\delta b_4 + i\delta b_5) + h.c.$

$$+ C_{i LR}^{ud \, us}(\delta a_4 - i\delta a_5)(\delta b_1 + i\delta b_2) + C_{i LR}^{us \, ud}(\delta a_1 - i\delta a_2)(\delta b_4 + i\delta b_5) + h.c. \right] ,$$

(5)

here $F_0$ is the pion decay constant in the Chiral limit, and $U$ is the usual matrix of pseudo Nambu Goldstone bosons in the notation of $17$. Furthermore,

$$\chi = 2B (M + i (m_3 l_3 + m_6 t_6 + m_8 t_8)) \right] ,$$

(6)

with

$$m_3 = - \sum_{i=1,2} r_i \text{Im} \left( C_{i LR}^{ud \, ud} + \frac{1}{2} C_{i LR}^{us \, us} \right), \quad m_6 = \frac{1}{2} \sum_{i=1,2} r_i \text{Im} \left( C_{i LR}^{ud \, us} + C_{i LR}^{us \, ud} \right),$$

$$m_8 = - \frac{\sqrt{3}}{2} \sum_{i=1,2} r_i \text{Im} C_{i LR}^{us \, us} ,$$

(7)

where $r_i = \frac{F_0^2}{2} A_{i \, LR}$.

The second and third lines in the Lagrangian in Eq. 5 involve terms that will induce EDMs and $\epsilon'$, respectively. These contributions depend on the low energy constants (LECs), $A_{i \, LR}$, which can be related to the known matrix elements of the SM electroweak penguin operators, $Q_7$ and $Q_8 \, 22,23$. These SM operators also transform as $8_L \times 8_R$ and therefore induce similar terms as $O_{i \, LR}$ in the Chiral Lagrangian, with the same LECs, $A_{i \, LR}$. As a result, we can express their matrix elements for $K_L \rightarrow \pi \pi$ in terms of $A_{i \, LR}$. At leading order (LO) in ChiPT, together with recent lattice results $24$, this gives

$$A_{1 \, LR}(3 \, \text{GeV}) = \frac{1}{\sqrt{3}F_0} \langle (\pi\pi)_{I=2}|Q_7|K^0 \rangle + \mathcal{O} \left( m_K^2 \right) \simeq (2.2 \pm 0.13) \, \text{GeV}^2 ,$$

$$A_{2 \, LR}(3 \, \text{GeV}) = \frac{1}{\sqrt{3}F_0} \langle (\pi\pi)_{I=2}|Q_8|K^0 \rangle + \mathcal{O} \left( m_K^2 \right) \simeq (10.1 \pm 0.6) \, \text{GeV}^2 .$$

(8)

3 Contribution to CP violation in the kaon sector

3.1 Direct CP violation

The measure of CP direct violation in $K_L \rightarrow \pi \pi$ decays is given by,

$$\text{Re} \left( \frac{\epsilon'}{\epsilon} \right) = \text{Re} \left( \frac{i\omega e^{i\phi_2 - \phi_0}}{\sqrt{2}e} \right) \left[ \frac{\text{Im} A_2}{\text{Re} A_2} - \frac{\text{Im} A_0}{\text{Re} A_0} \right] ,$$

(9)

where $A_{0,2} e^{i\delta_{0,2}}$ are the amplitudes for final states with total isospin $I = 0, 2$, $\omega = \text{Re} A_2/\text{Re} A_0$, and $\epsilon$ denotes the CP violation in $K^- - K$ mixing. The left-right operators contribute to the imaginary parts of the amplitudes, $\text{Im} A_{0,2}$, while the remaining quantities in the above expression are well known experimentally. The Chiral Lagrangian of the previous section, Eq. 5, together with the determination of the LECs, Eq. 8, now allows us to calculate the contributions to these amplitudes. One would expect such a calculation to be subject to $\mathcal{O}(m_K^2/A_s^2)$ corrections due to the fact that Eq. 8 is a LO ChiPT prediction. Luckily, this is not the case for the $I = 2$ amplitude. The reason is that, after an isospin decomposition, the $I = 3/2$ parts of the $O_{1(2)} \, LR$ and $Q_{7(8)}$ coincide. Thus, the right-handed contributions to the $I = 2$ amplitudes can be determined up to isospin corrections, which gives

$$\text{Im} A_2(\xi) = \frac{1}{6\sqrt{2}} \text{Im} \left[ \left( C_{1 \, LR}^{ud \, us} - C_{1 \, LR}^{us \, ud} \right) \langle (\pi\pi)_{I=2}|Q_7|K^0 \rangle \right. + \left. \left( C_{2 \, LR}^{ud \, us} - C_{2 \, LR}^{us \, ud} \right) \langle (\pi\pi)_{I=2}|Q_8|K^0 \rangle \right] .$$

(10)
contributions, including lattice uncertainties, give this contribution can be estimated reliably. Using the conventions of Naive-dimensional-analysis estimates for the additional direct pieces suggest they are roughly an order of magnitude smaller than the indirect contributions. As such, we follow 

$$\Delta = \frac{m_K^2}{\Lambda^2}$$

where

$$\Delta I = 1/2$$

and estimate this contribution to $\epsilon'$, as it is suppressed by the lattice determinations of the matrix elements.

### 3.2 CP violation in mixing

Apart from the direct CP violation in kaon decays, the right-handed interactions can also induce CP violation in mixing, $\epsilon_K$. Although $\xi_{ud}$ and $\xi_{us}$ do not induce tree-level $\Delta S = 2$ operators, they do contribute to $\epsilon_K$ through short- and long-distance effects. The former arise through box diagrams involving the $\xi$ couplings. However, due to the chirality of the vertices, the box diagrams linear in $\xi$ require one internal and one external quark mass insertion, i.e. they are suppressed by $\frac{m_{ud}}{m_W^2}$. The short-distance contributions are therefore negligible.

The long-distance contribution arises from the combination of a $\Delta S = 1$ left-right interaction with a $\Delta S = 1$ SM charged current. The Chiral realizations of these operators lead to diagrams where $K^0$ mixes into a pion or eta meson, which then mixes into a $\bar{K}^0$. This involves the LECs of the left-right operators, $\mathcal{A}_{iLR}$, as well as those for the weak charged current. Here we follow $^{17}$ and estimate this contribution to $\epsilon_K$ by the tree-level diagrams (which are non-zero at NLO) and assign a 50% uncertainty due to unknown NLO counterterms.

### 4 Contribution to hadronic EDMs

The contributions of the left-right operators to hadronic EDMs can be calculated by first matching to an extension of chiral effective field theory that contains CP-violating hadronic interactions $^{27,28}$. Chiral power counting then predicts $^{27,28}$ that contributions of the four-quark operators to nuclear EDMs are dominated by long-range pion-exchange between nucleons $^b$. The leading pion-nucleon couplings, $\bar{g}_{0,1}$, are induced by the left-right operators in several ways. Firstly, there is a direct contribution whose LEC involves matrix elements of the form, $\langle N\pi|O_{iLR}|N\rangle$, which are currently unknown. A second contribution arises due to the rotation performed to align the vacuum, mentioned in section 2. The relevant meson-baryon Lagrangian then takes the following form,

$$\mathcal{L}_{\pi N} = b_0 \text{Tr} (\bar{B}B) \text{Tr} \chi + b_D \text{Tr} (\bar{B}\{\chi,B\}) + b_F \text{Tr} (\bar{B}[\chi,B]) + \mathcal{L}_{\text{direct}}, \quad (11)$$

where $B$ represents the octet of baryon fields, notation is as in $^{17}$, and $\chi = u^\dagger u^\dagger + u^\dagger u$.

Here the direct contributions to $\bar{g}_{0,1}$, with unknown LECs, are due to $\mathcal{L}_{\text{direct}}$. The contributions induced by vacuum alignment arise through $\chi$ and depend on $\mathcal{A}_{iLR}$, and $b_{0,D,F}$. Since $b_{0,D,F}$ can be related to the baryon mass splittings $^{29,30,31,32}$, and $\mathcal{A}_{iLR}$ are known from Eq. 8 this contribution can be estimated reliably. Using the conventions of $^{17}$ for $\bar{g}_{0,1}$, the indirect contributions, including lattice uncertainties, give

$$\bar{g}_0 \frac{2F_\pi}{F} = - (0.16 \pm 0.03) \times 10^{-5} \text{Im}(V_{us}^* \xi_{us}),$$

$$\bar{g}_1 \frac{2F_\pi}{F} = - (2.9 \pm 0.33) \times 10^{-5} \text{Im}(V_{us}^* \xi_{us}) - (5.7 \pm 0.67) \times 10^{-5} \text{Im}(V_{ud}^* \xi_{ud}). \quad (12)$$

Naive-dimensional-analysis estimates for the additional direct pieces suggest they are roughly an order of magnitude smaller than the indirect contributions. As such, we follow $^{17}$ and use the indirect piece as the central values for $\bar{g}_{0,1}$ and conservatively assign a 50% uncertainty due to the direct piece.

$^b$Note that chiral power counting has not been tested for systems as large as $^{199}$Hg or $^{225}$Ra.
4.1 The neutron EDM

Although $\bar{g}_{0,1}$ give the dominant contributions to nuclear EDMs, for the neutron EDM additional counterterms appear at the same order. One has\textsuperscript{33,34}

$$d_n = \bar{d}_n(\mu) + \frac{e g A \bar{g}_1}{(4\pi F)^2} \left( \frac{\bar{g}_0}{g_1} \left( \log \frac{m^2}{\mu^2} - \frac{\pi m}{2m_N} \right) + \frac{1}{4} \left( \kappa_1 - \kappa_0 \right) \frac{m^2}{m^2_N} \log \frac{m^2}{\mu^2} \right), \quad (13)$$

where $g_A \simeq 1.27$ is the nucleon axial charge, and $\kappa_1 = 3.7$ and $\kappa_0 = -0.12$ are related to the nucleon magnetic moments. $\bar{d}_n(\mu)$ is a counterterm, which is again unknown. We estimate its size by the $\mu$ dependence of the loop contributions, which we obtain by varying $\mu$ from $m_K$ to $m_N$ in Eq. 13. The resulting sizes are in agreement with naive-dimensional-analysis estimates. As a result, we take Eq. 13 as the central value with $\bar{d}_n(\mu) = 0$. We estimate the uncertainties by the combination of the errors on $\bar{g}_{0,1}$ discussed above together with the variation due to the $\mu$ dependence\textsuperscript{17}.

4.2 Nuclear EDMs

As already mentioned, for nuclear EDMs, the dominant contributions should be captured by $\bar{g}_{0,1}$, implying that no further unknown LECs enter the expressions in this case. Nuclear calculations, within large uncertainties, predict\textsuperscript{35,36,37,38,39,40,28,41,42}

$$d_D = -(0.18 \pm 0.02) \frac{\bar{g}_1}{2F} \text{ e fm},$$
$$d_{Hg} = (2.8 \pm 0.6) \cdot 10^{-4} \cdot \left( 0.13^{+0.05}_{-0.07} \frac{\bar{g}_0}{2F} + 0.25^{+0.89}_{-0.63} \frac{\bar{g}_1}{2F} \right) \text{ e fm},$$
$$d_{Ra} = (7.7 \pm 0.8) \cdot 10^{-4} \cdot \left( -19^{+6.4}_{-5.7} \frac{\bar{g}_0}{2F} + 76^{+227}_{-25} \frac{\bar{g}_1}{2F} \right) \text{ e fm}. \quad (14)$$

We set constraints using Eqs. 14 and 13 together with the experimental measurements\textsuperscript{33,44,45,46}.

5 Discussion

Using the expressions in section 3 we show $\epsilon'$ as a function of $\xi_{ud}$ and $\xi_{us}$ in the upper-left and -right panels of Fig. 1, respectively. Here the green band indicates the experimental value, while the solid and dashed blue lines are theory predictions using the SM values of Ref.\textsuperscript{4} and\textsuperscript{25}, respectively. These panels show that the tension can be alleviated when the couplings have sizes of $O(10^{-7} - 10^{-6})$. Coefficients of this size naively point towards a scale of $\Lambda = O(100 \text{ TeV})$, although, in specific models this scale can be lowered by small model parameters. The same panels also show the constraints from $\epsilon_K$ and $d_n$, and future $d_{D,Ra}$ sensitivities\textsuperscript{47,48}. In principle, the stringent experimental limit on the mercury EDM also leads to strong constraints if one neglects theoretical uncertainties. However, here we follow the R-fit procedure\textsuperscript{49} to obtain limits. In this case the large nuclear and hadronic uncertainties allow for cancellations which result in a vanishing $d_{Hg}$ EDM.

As can be seen from the Figure, the current $\epsilon_K$ and $d_n$ limits do not rule out the region of interest. For both $\xi_{ud}$ and $\xi_{us}$ a future neutron, deuteron, or radium measurement would be able to probe the regions in which the $\epsilon'$ tension is resolved. To illustrate this more clearly, we show the allowed values of the neutron and radium EDMs in the lower-left and -right panels of Fig. 1. Here the black lines indicate the neutron and radium EDMs in the case that $\xi_{ud,us}$ have the right size to solve the $\epsilon'$ discrepancy. The red points show the values once one accounts for hadronic and nuclear uncertainties. For both $\xi_{ud}$ and $\xi_{us}$, one can see that the projected experimental sensitivities for $d_n$ and $d_{Ra}$ would allow one to rule out the entire parameter space. In summary, a right-handed explanation of the tension in $\epsilon'$ points to a scale of new physics of the order of $O(100 \text{ TeV})$, which is not within reach of direct searches. Nonetheless, future EDM experiments would be able to falsify this scenario.
Figure 1 – The top-left (-right) panel shows the value of Re $\epsilon'/\epsilon$ as a function of Im $\xi_{ud}$ (Im $\xi_{us}$). The solid and dashed blue bands indicate the theory prediction (see text), while the experimental value is shown in green (all at 1$\sigma$). The vertical lines indicate the current/future sensitivities of $\epsilon_K$ and $d_{n,D,Ra}$ experiments, derived using the R-fit procedure. The bottom-left (-right) panel shows the sizes of $d_{Ra}$ and $d_n$, assuming a value for Im $\xi_{ud}$ (Im $\xi_{us}$) that solves the $\epsilon'/\epsilon$ discrepancy. The red points are generated by taking random values of the nuclear and hadronic matrix elements within their allowed ranges. The black lines result from taking the central values of these matrix elements.

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