ϵ'/ϵ anomaly and neutron EDM in $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry

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ABSTRACT: The Standard Model prediction for ϵ'/ϵ based on recent lattice QCD results exhibits a tension with the experimental data. We solve this tension through W_R^+ gauge boson exchange in the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with ‘charge symmetry’, whose theoretical motivation is to attribute the chiral structure of the Standard Model to the spontaneous breaking of $SU(2)_R \times U(1)_{B-L}$ gauge group and charge symmetry. We show that $M_{W_R} < 58 \text{ TeV}$ is required to account for the ϵ'/ϵ anomaly in this model. Next, we make a prediction for the neutron EDM in the same model and study a correlation between ϵ'/ϵ and the neutron EDM. We confirm that the model can solve the ϵ'/ϵ anomaly without conflicting the current bound on the neutron EDM, and further reveal that almost all parameter regions in which the ϵ'/ϵ anomaly is explained will be covered by future neutron EDM searches, which leads us to anticipate the discovery of the neutron EDM.

KEYWORDS: Beyond Standard Model, CP violation

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Contents

1	Introduction	1
2	$SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry	3
3	ϵ'/ϵ	6
3.1	Wilson coefficients for $\Delta S = 1$ operators	6
3.2	Hadronic matrix elements	8
3.3	Numerical analysis of ϵ'/ϵ	9
4	Neutron electric dipole moment	11
4.1	Wilson coefficients for operators contributing to the neutron EDM	11
4.2	Hadronic matrix elements	12
4.2.1	Four-quark operators $\mathcal{O}_{1q'q}$, $\mathcal{O}_{2q'q}$	12
4.2.2	Other CP-violating operators $\mathcal{O}_{1q}$, $\mathcal{O}_{2q}$ and $\mathcal{O}_3$	19
4.3	Numerical analysis of neutron EDM versus ϵ'/ϵ	19
5	Summary and discussions	21
A	Operators of $\Delta S = 1$ processes	23
B	Loop functions	24
C	CP-violating operators that contribute to the neutron EDM	24

1 Introduction

The direct CP violation in $K \rightarrow \pi\pi$ decay parametrized by the ϵ' parameter is sensitive to physics beyond the Standard Model (SM) due to the suppressed SM contribution. Recent calculation of the hadronic matrix elements with lattice QCD [1, 3, 4] enables us to evaluate the $K \rightarrow \pi\pi$ decay amplitude without relying on any hadron model. On the basis of the above calculation, the same collaboration has reported that the SM prediction is separated from the experimental value [5–7] by 2.1σ , and other groups [9, 10] have also obtained predictions for ϵ'/ϵ that show a discrepancy of 2.9σ and 2.8σ , respectively. More importantly, the lattice result corroborates the calculation with dual QCD approach [11, 12], which has derived a theoretical upper bound on ϵ'/ϵ that is violated by the experimental data and has thus claimed anomaly in this observable. (However, ref. [13] presents a different calculation that claims the absence of the anomaly.) Some authors have tackled this ϵ'/ϵ anomaly in new physics scenarios, such as a general right-handed current [14], the Littlest

Higgs model with T-parity [15], supersymmetry [16–18], non-standard interaction with Z' and/or Z [19, 20], vector-like quarks [21] and $SU(3)_c \times SU(3)_L \times U(1)_X$ gauge group [22].

The $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge extension of the SM is a well-motivated framework for addressing the ϵ'/ϵ puzzle, because the flavor mixing matrix for right-handed quarks automatically introduces new CP-violating phases, and W_R^+ gauge boson exchange contributes to $\Delta F = 1$ processes at tree level while it contributes to $\Delta F = 2$ processes at loop levels so that other experimental constraints, in particular the constraint from $\text{Re}(\epsilon)$, are readily evaded. Previously, ref. [14] has shown that a general $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with an arbitrary right-handed quark mixing can solve the ϵ'/ϵ discrepancy. However, a major theoretical motivation for the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model lies in its capability of explaining the origin of the chiral nature of the SM, which is achieved by adding either the left-right parity [24] or the ‘charge symmetry’ [27].¹ The left-right parity requires invariance of the theory under the Lorentzian parity transformation plus the exchange of $SU(2)_L$ and $SU(2)_R$ gauge groups, while the charge symmetry requires invariance under the charge conjugation plus the exchange of $SU(2)_L$ and $SU(2)_R$, both of which endow the model with a symmetric structure for the left and right-handed fermions at high energies.

In this paper, we study ϵ'/ϵ in the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry. As a consequence of the charge symmetry, the Yukawa matrices are complex symmetric matrices, which restricts the quark mixing matrix associated with W_R^+ to be the complex conjugate of the SM Cabibbo-Kobayashi-Maskawa (CKM) matrix multiplied by a new CP phase factor for each quark flavor. Given the above restriction, one can evaluate ϵ'/ϵ only in terms of two new CP phases, the mass of W_R^+ and the ratio of two vacuum expectation values (VEVs) of the bifundamental scalar, which leads to a specific prediction for the model parameters.

Our analysis on ϵ'/ϵ proceeds as follows. By integrating out W_R , W_L and the top quark, we obtain the Wilson coefficients for $\Delta S = 1$ operators that contribute to $K \rightarrow \pi\pi$ decay. The anomalous dimension matrix is divided into the same two 18×18 pieces for 36 operators, for which leading order expressions are obtainable from refs. [30, 31]. The hadronic matrix elements for current-current operators are seized from the lattice results [3, 4]. We find that among new physics operators, $(\bar{s}u)_L(\bar{u}d)_R$ and $(\bar{s}u)_R(\bar{u}d)_R$ (each with two ways of color contraction) both dominantly contribute to ϵ'/ϵ . Their contributions are of the same order because the Wilson coefficients of the $(\bar{s}u)_L(\bar{u}d)_R$ operators are suppressed by the hierarchy of two bifundamental scalar VEVs $v_1/v_2 = \tan\beta$, which is about m_b/m_t if there is no fine-tuning in accommodating the top and bottom quark Yukawa couplings, whereas this suppression does not enter into the Wilson coefficients of the $(\bar{s}u)_R(\bar{u}d)_R$ operators. On the other hand, the lattice computation has confirmed that the hadronic matrix elements for the former operators are enhanced compared to the latter. Thus, these operators possibly equally contribute to ϵ'/ϵ . This result is in contrast to the study of ref. [14], which has concentrated solely on the $(\bar{s}u)_L(\bar{u}d)_R$ operators.

¹The ‘charge symmetry’ is inspired by D -parity [28] in the $SO(10)$ grand unification theory. However, the model we consider cannot be embedded in the $SO(10)$ theory, since we assume the charge symmetry breaking scale to be below $O(100)$ TeV.

Once the ϵ'/ϵ anomaly is explained in the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry, correlated predictions for other CP violating observables are of interest. In particular, the neutron electric dipole moment (EDM), an observable sensitive to CP violation in the presence of CPT invariance, receives significant contributions from four-quark operators in $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ models [14, 34–38],² allowing us to discuss future detectability of the neutron EDM in relation to the ϵ'/ϵ anomaly.

Our analysis on the neutron EDM starts by integrating out W_R , W_L and top quark to obtain the Wilson coefficients for CP-violating operators. The leading order expression for the anomalous dimension matrix is found in refs. [41, 42, 44, 45]. Regarding the hadronic matrix elements of CP-violating operators, we reveal that the pion VEV $\langle\pi^0\rangle$ induced by four-quark operators [33] gives the leading contribution to the neutron EDM, which is enhanced by the quark mass ratio $m_s/(m_u + m_d)$ in comparison to the rest. This enhancement is understood as follows: since $W_R^+ W_L^+$ mixing gives rise to CP-odd and isospin-odd interactions, the pion VEV $\langle\pi^0\rangle$, which is isospin-odd, can arise without the factor of $m_d - m_u$, and thus can be directly proportional to $1/(m_u + m_d)$. The pion VEV induces a CP-violating coupling for neutron n , Σ^- baryon, and kaon K^+ without the factor of $m_d - m_u$ because the $\bar{n}\Sigma^- K^+$ vertex is not isospin-even. Consequently, the CP-violating coupling for n, Σ^-, K^+ can appear with the factor of $m_s/(m_u + m_d)$. This coupling contributes to the neutron EDM at the leading chiral order through charged baryon-meson loops. Considering the above-mentioned importance of the pion VEV, we in this paper investigate meson condensation, the resultant CP-violating baryon-meson couplings, and their contributions to the neutron EDM through baryon-meson loops, using chiral perturbation theory.

This paper is organized as follows: in section 2, we review the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry, with emphasis on new sources of CP violation. In section 3, we present the Wilson coefficients for $\Delta S = 1$ operators in the model, their RG evolutions and the hadronic matrix elements for these operators. The numerical result for ϵ'/ϵ is shown at the end of the section. In section 4, we give the Wilson coefficients for CP-violating operators contributing to the neutron EDM. Special care is taken in evaluating meson condensates and their impact on the neutron EDM. The final result is a prediction for the neutron EDM in light of the ϵ'/ϵ anomaly. Section 5 is devoted to summary and discussions.

2 $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry

We consider $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge theory with charge symmetry. The field content is in table 1. Hereafter, the fields are expressed in a way that they transform under a $SU(2)_L \times SU(2)_R$ gauge transformation as

$$\begin{aligned} \Phi &\rightarrow e^{i\theta_L^a \tau^a} \Phi e^{-i\theta_R^b \tau^b}, & \Delta_L &\rightarrow e^{i\theta_L^a \tau^a} \Delta_L e^{-i\theta_L^a \tau^a}, & \Delta_R &\rightarrow (e^{i\theta_R^b \tau^b})^* \Delta_R (e^{-i\theta_R^b \tau^b})^*, \\ Q_L^i &\rightarrow e^{i\theta_L^a \tau^a} Q_L^i, & Q_R^{ci} &\rightarrow (e^{i\theta_R^b \tau^b})^* Q_R^{ci}, & L_L^i &\rightarrow e^{i\theta_L^a \tau^a} L_L^i, & L_R^{ci} &\rightarrow (e^{i\theta_R^b \tau^b})^* L_R^{ci}, \end{aligned} \quad (2.1)$$

²See also refs. [39, 40].

Field	Lorentz SO(1, 3)	SU(3) _C	SU(2) _L	SU(2) _R	U(1) _{B-L}
Q_L^i	(2 , 1)	3	2	1	1/3
Q_R^{ci}	(2 , 1)	$\bar{\mathbf{3}}$	1	2	-1/3
L_L^i	(2 , 1)	1	2	1	-1
L_R^{ci}	(2 , 1)	1	1	2	1
Φ	1	1	2	2	0
Δ_L	1	1	3	1	2
Δ_R	1	1	1	3	-2

Table 1. Field content and charge assignments. i labels the three generations.

with θ_L^a and θ_R^a being gauge parameters for SU(2)_L and SU(2)_R, respectively. We demand the theory to be invariant under the following ‘charge symmetry’ transformation:

$$\begin{aligned} &\text{charge conjugation of all gauge fields,} \\ &\text{and SU(2)}_L \leftrightarrow \text{SU(2)}_R, \quad Q_L^i \leftrightarrow Q_R^{ci}, \quad L_L^i \leftrightarrow L_R^{ci}, \quad \Phi \leftrightarrow \Phi^T, \quad \Delta_L \leftrightarrow \Delta_R. \end{aligned} \quad (2.2)$$

The part of the Lagrangian describing SU(2)_L × SU(2)_R × U(1)_{B-L} and Yukawa interactions of quarks is given by

$$\begin{aligned} -\mathcal{L} \supset & Q_L^{i\dagger} \bar{\sigma}_\mu \left(\frac{1}{2} g_L \sigma^a W_L^{a\mu} + \frac{1}{3} g_X X^\mu \right) Q_L^i + Q_R^{ci\dagger} \bar{\sigma}_\mu \left(-\frac{1}{2} g_R (\sigma^a)^T W_R^{a\mu} - \frac{1}{3} g_X X^\mu \right) Q_R^{ci} \\ & + (Y_q)_{ij} Q_L^{i\dagger} \Phi \epsilon_s (Q_R^{cj})^* + (\tilde{Y}_q)_{ij} Q_L^{i\dagger} (\epsilon_g^T \Phi^* \epsilon_g) \epsilon_s (Q_R^{cj})^* + \text{H.c.}, \end{aligned} \quad (2.3)$$

where g_L , g_R and g_X are the gauge coupling constants for SU(2)_L, SU(2)_R and U(1)_{B-L} gauge groups, respectively, and Y_q and $\tilde{Y}_q$ are the quark Yukawa couplings. ϵ_s denotes the antisymmetric tensor for Lorentz spinors and ϵ_g denotes that for the fundamental representation of SU(2)_L or SU(2)_R. Invariance under the charge symmetry transformation eq. (2.2) leads to the following tree-level relations:

$$g_L = g_R, \quad (Y_q)_{ij} = (Y_q)_{ji}, \quad (\tilde{Y}_q)_{ij} = (\tilde{Y}_q)_{ji}. \quad (2.4)$$

The SU(2)_R triplet scalar Δ_R develops a VEV, v_R , to break SU(2)_R × U(1)_{B-L} → U(1)_Y, and the bi-fundamental scalar Φ takes a VEV configuration,

$$\langle \Phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v \sin \beta & 0 \\ 0 & v \cos \beta e^{i\alpha} \end{pmatrix}, \quad (2.5)$$

to break SU(2)_L × U(1)_Y → U(1)_{em}, where α is the spontaneous CP phase. The VEV of Δ_L is hereafter neglected, as it is severely constrained from ρ -parameter. The resultant mass matrices for W_L^a , W_R^a and X gauge bosons read

$$\begin{aligned} -\mathcal{L} \supset & \begin{pmatrix} W_L^- & W_R^- \end{pmatrix} \begin{pmatrix} g_L^2 v^2/4 & -g_L g_R \sin(2\beta) e^{-i\alpha} v^2/4 \\ -g_L g_R \sin(2\beta) e^{i\alpha} v^2/4 & g_R^2 (v_R^2 + v^2/4) \end{pmatrix} \begin{pmatrix} W_L^+ \\ W_R^+ \end{pmatrix} \\ & + \frac{1}{2} \begin{pmatrix} W_L^3 & W_R^3 & X \end{pmatrix} \begin{pmatrix} g_L^2 v^2/4 & -g_L g_R v^2/2 & 0 \\ -g_L g_R v^2/2 & g_R^2 (2v_R^2 + v^2/4) & -2g_L g_X v_R^2 \\ 0 & -2g_L g_X v_R^2 & 2g_X^2 v_R^2 \end{pmatrix} \begin{pmatrix} W_L^3 \\ W_R^3 \\ X \end{pmatrix}. \end{aligned} \quad (2.6)$$

The mass matrix for the charged gauge bosons is diagonalized as

$$-\mathcal{L} \supset (W^- \ W'^-) \begin{pmatrix} M_W^2 & 0 \\ 0 & M_{W'}^2 \end{pmatrix} \begin{pmatrix} W^+ \\ W'^+ \end{pmatrix}, \quad \begin{pmatrix} W_L^+ \\ W_R^+ \end{pmatrix} = \begin{pmatrix} \cos \zeta & -e^{-i\alpha} \sin \zeta \\ e^{i\alpha} \sin \zeta & \cos \zeta \end{pmatrix} \begin{pmatrix} W^+ \\ W'^+ \end{pmatrix},$$

$$\sin(2\zeta) = \frac{2g_L g_R \sin(2\beta) v^2}{(g_L^2 + g_R^2) v^2 + 4g_R^2 v_R^2 - 8M_W^2}. \quad (2.7)$$

For $v_R \gg v$ and $g_L = g_R$, we have an important relation for ζ ,

$$\sin \zeta \simeq \sin(2\beta) \frac{M_W^2}{M_{W'}^2}, \quad (2.8)$$

which indicates that when we assume $\tan \beta \simeq m_b/m_t$ so that the top and bottom Yukawa couplings are naturally derived, the W_L - W_R mixing angle ζ is smaller than $M_W^2/M_{W'}^2$ by the factor $2m_b/m_t \sim 0.05$.

The quark mass matrices are given by³

$$-\mathcal{L} \supset (M_u)_{ij} U_L^{i\dagger} U_R^j + (M_d)_{ij} D_L^{i\dagger} D_R^j + \text{H.c.},$$

$$M_u = \frac{v}{\sqrt{2}} \left(\sin \beta Y_q + \cos \beta e^{-i\alpha} \tilde{Y}_q \right), \quad M_d = \frac{v}{\sqrt{2}} \left(\cos \beta e^{i\alpha} Y_q + \sin \beta \tilde{Y}_q \right), \quad (2.9)$$

which we diagonalize as $M_u = V_{uL}^\dagger \text{diag}(m_u, m_c, m_t) V_{uR}$ and $M_d = V_{dL}^\dagger \text{diag}(m_d, m_s, m_b) V_{dR}$, with $V_{uL}, V_{uR}, V_{dL}, V_{dR}$ being unitary matrices. However, since Y_q and $\tilde{Y}_q$ are complex symmetric matrices, so are M_u and M_d , and one can most generally write

$$V_{uR} = \begin{pmatrix} e^{i\phi_u} & 0 & 0 \\ 0 & e^{i\phi_c} & 0 \\ 0 & 0 & e^{i\phi_t} \end{pmatrix} V_{uL}^*, \quad V_{dR} = \begin{pmatrix} e^{i\psi_d} & 0 & 0 \\ 0 & e^{i\psi_s} & 0 \\ 0 & 0 & e^{i\psi_b} \end{pmatrix} V_{dL}^*. \quad (2.10)$$

Hence, the SM CKM matrix, $V_L = V_{uL} V_{dL}^\dagger$, and the corresponding flavor mixing matrix for right-handed quarks, $V_R = V_{uR} V_{dR}^\dagger$, are related as

$$V_R = V_{uR} V_{dR}^\dagger = \begin{pmatrix} e^{i\phi_u} & 0 & 0 \\ 0 & e^{i\phi_c} & 0 \\ 0 & 0 & e^{i\phi_t} \end{pmatrix} V_{uL}^* V_{dL}^T \begin{pmatrix} e^{-i\psi_d} & 0 & 0 \\ 0 & e^{-i\psi_s} & 0 \\ 0 & 0 & e^{-i\psi_b} \end{pmatrix}$$

$$= \begin{pmatrix} e^{i\phi_u} & 0 & 0 \\ 0 & e^{i\phi_c} & 0 \\ 0 & 0 & e^{i\phi_t} \end{pmatrix} V_L^* \begin{pmatrix} e^{-i\psi_d} & 0 & 0 \\ 0 & e^{-i\psi_s} & 0 \\ 0 & 0 & e^{-i\psi_b} \end{pmatrix}. \quad (2.11)$$

³ $U_R^i \equiv \epsilon_s (U_R^{ci})^*, \ D_R^i \equiv \epsilon_s (D_R^{ci})^*.$

Eventually, the part of the Lagrangian eq. (2.3) describing flavor-changing W, W' interactions is recast, in the unitary gauge, into the form,

$$\begin{aligned}
 -\mathcal{L} &\supset \frac{g_L}{\sqrt{2}} (V_L)_{ij} U_L^{i\dagger} W_L^{+\mu} \bar{\sigma}_\mu D_L^j + \frac{g_R}{\sqrt{2}} (V_L^*)_{ij} e^{i(\phi_i - \psi_j)} U_R^{i\dagger} W_R^{+\mu} \sigma_\mu D_R^j + \text{H.c.} \\
 &= \frac{1}{\sqrt{2}} \bar{U}^i W^{+\mu} \gamma_\mu \left\{ g_L (V_L)_{ij} \cos \zeta P_L + g_R (V_L^*)_{ij} e^{i(\phi_i - \psi_j + \alpha)} \sin \zeta P_R \right\} D^j \\
 &\quad + \frac{1}{\sqrt{2}} \bar{U}^i W'^{+\mu} \gamma_\mu \left\{ -g_L (V_L)_{ij} e^{-i\alpha} \sin \zeta P_L + g_R (V_L^*)_{ij} e^{i(\phi_i - \psi_j)} \cos \zeta P_R \right\} D^j + \text{H.c.},
 \end{aligned} \tag{2.12}$$

where U^i and D^i denote the Dirac fields of the up and down-type quarks, respectively.

In this paper, we adopt the following convention for the quark phases and $\phi_u, \phi_c, \phi_t, \psi_d, \psi_s, \psi_b$: first, we redefine the phases of five quarks to render the CKM matrix in the standard form,

$$V_L = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix}. \tag{2.13}$$

Next, we redefine $\phi_c, \phi_t, \psi_d, \psi_s, \psi_b$ to set

$$\phi_u = 0. \tag{2.14}$$

Phase convention fixed in this way, all sources of CP violation are parametrized by $\text{Im}[(V_L)_{cd}]$, $\text{Im}[(V_L)_{cs}]$, $\text{Im}[(V_L)_{td}]$, $\text{Im}[(V_L)_{ts}]$, $\text{Im}[(V_L)_{cd}]$, the newly-defined $\phi_c, \phi_t, \psi_d, \psi_s, \psi_b$, and α .

3 ϵ'/ϵ

3.1 Wilson coefficients for $\Delta S = 1$ operators

We match the $\text{SU}(3)_C \times \text{SU}(2)_L \times \text{SU}(2)_R \times \text{U}(1)_{B-L}$ gauge theory with charge symmetry to the effective QCD $\times$ QED theory in which W, W' bosons and the top quark are integrated out. In the effective theory, the $\Delta S = 1$ Hamiltonian is parametrized as

$$\begin{aligned}
 \mathcal{H}_{\Delta S=1} &= \frac{G_F}{\sqrt{2}} \left\{ \sum_{i=1,2,\dots,10,1c,2c} (C_i O_i + C'_i O'_i) + \sum_{j=1,2,1c,2c} (C_j^{RL} O_j^{RL} + C_j^{LR} O_j^{LR}) \right. \\
 &\quad \left. + \sum_{k=g,\gamma} (C_k O_k + C'_k O'_k) \right\} + \text{H.c.},
 \end{aligned} \tag{3.1}$$

where operators O 's are defined in appendix A. We determine the Wilson coefficients as follows: we approximate $g_R = g_L$ by ignoring difference in RG evolutions of g_L and g_R at scales below $M_{W'}$. Also, for each Wilson coefficient, if multiple terms have an identical phase, we only consider the one in the leading order of $M_W^2/M_{W'}^2$ or $\sin \zeta$. By integrating

out W' , one obtains the following leading-order matching conditions at a scale $\mu \sim M_{W'}$ (note our convention with $\phi_u = 0$):

$$C'_2 = \frac{M_W^2}{M_{W'}^2} (V_L)_{us} (V_L^*)_{ud} e^{i(\psi_s - \psi_d)} \cos^2 \zeta, \quad (3.2)$$

$$C'_{2c} = \frac{M_W^2}{M_{W'}^2} (V_L)_{cs} (V_L^*)_{cd} e^{i(\psi_s - \psi_d)} \cos^2 \zeta, \quad (3.3)$$

$$C'_4 = C'_6 = \frac{M_W^2}{M_{W'}^2} \frac{\alpha_s}{4\pi} \cos^2 \zeta \sum_{i=u,c,t} (V_L)_{is} (V_L^*)_{id} e^{i(\psi_s - \psi_d)} \frac{1}{2} F_1(y_i), \quad C'_3 = C'_5 = -\frac{1}{3} C'_4, \quad (3.4)$$

$$C'_7 = C'_9 = \frac{M_W^2}{M_{W'}^2} \frac{\alpha}{4\pi} \cos^2 \zeta \sum_{i=u,c,t} (V_L)_{is} (V_L^*)_{id} e^{i(\psi_s - \psi_d)} \frac{2}{3} E_{1d}(y_i), \quad (3.5)$$

$$\delta C_g = \frac{M_W^2}{M_{W'}^2} \cos^2 \zeta \frac{m_d}{m_s} \sum_{i=u,c,t} (V_L)_{is} (V_L^*)_{id} F_2(y_i), \quad (3.6)$$

$$\delta C_\gamma = \frac{M_W^2}{M_{W'}^2} \cos^2 \zeta \frac{m_d}{m_s} \sum_{i=u,c,t} (V_L)_{is} (V_L^*)_{id} E_{2d}(y_i), \quad (3.7)$$

$$\delta C'_g = \frac{M_W^2}{M_{W'}^2} \cos^2 \zeta \sum_{i=u,c,t} (V_L)_{is} (V_L^*)_{id} F_2(y_i), \quad (3.8)$$

$$\delta C'_\gamma = \frac{M_W^2}{M_{W'}^2} \cos^2 \zeta \sum_{i=u,c,t} (V_L)_{is} (V_L^*)_{id} E_{2d}(y_i), \quad (3.9)$$

with $y_i \equiv m_i^2 / M_{W'}^2$.

By further integrating out W and the top quark, one gains the following leading-order matching conditions at a scale $\mu \sim M_W$ (note our convention with $\phi_u = 0$):

$$C_2 = (V_L^*)_{us} (V_L)_{ud} \cos^2 \zeta, \quad (3.10)$$

$$C_{2c} = (V_L^*)_{cs} (V_L)_{cd} \cos^2 \zeta, \quad (3.11)$$

$$C_2^{RL} = (V_L)_{us} (V_L)_{ud} e^{i(\psi_s - \alpha)} \sin \zeta \cos \zeta, \quad C_2^{LR} = (V_L^*)_{us} (V_L^*)_{ud} e^{i(-\psi_d + \alpha)} \sin \zeta \cos \zeta, \quad (3.12)$$

$$C_{2c}^{RL} = (V_L)_{cs} (V_L)_{cd} e^{i(-\phi_c + \psi_s - \alpha)} \sin \zeta \cos \zeta, \quad C_{2c}^{LR} = (V_L^*)_{cs} (V_L^*)_{cd} e^{i(\phi_c - \psi_d + \alpha)} \sin \zeta \cos \zeta, \quad (3.13)$$

$$C_4 = C_6 = \frac{\alpha_s}{4\pi} \cos^2 \zeta \sum_{i=u,c,t} (V_L^*)_{is} (V_L)_{id} \frac{1}{2} F_1(x_i), \quad C_3 = C_5 = -\frac{1}{3} C_4, \quad (3.14)$$

$$C_7 = C_9 = \frac{\alpha}{4\pi} \cos^2 \zeta \sum_{i=u,c,t} (V_L^*)_{is} (V_L)_{id} \frac{2}{3} E_{1d}(x_i), \quad (3.15)$$

$$\Delta C_g = \sum_{i=u,c,t} \left\{ \cos^2 \zeta (V_L^*)_{is} (V_L)_{id} F_2(x_i) + \sin \zeta \cos \zeta \frac{m_i}{m_s} (V_L)_{is} (V_L)_{id} e^{i(-\phi_i + \psi_s - \alpha)} F_3(x_i) \right\}, \quad (3.16)$$

$$\Delta C_\gamma = \sum_{i=u,c,t} \left\{ \cos^2 \zeta (V_L^*)_{is} (V_L)_{id} E_{2d}(x_i) + \sin \zeta \cos \zeta \frac{m_i}{m_s} (V_L)_{is} (V_L)_{id} e^{i(-\phi_i + \psi_s - \alpha)} E_{3d}(x_i) \right\}, \quad (3.17)$$

$$\Delta C'_g = \sum_{i=u,c,t} \left\{ \cos^2 \zeta \frac{m_d}{m_s} (V_L^*)_{is} (V_L)_{id} F_2(x_i) + \sin \zeta \cos \zeta \frac{m_i}{m_s} (V_L^*)_{is} (V_L^*)_{id} e^{i(\phi_i - \psi_d + \alpha)} F_3(x_i) \right\}, \quad (3.18)$$

$$\Delta C'_\gamma = \sum_{i=u,c,t} \left\{ \cos^2 \zeta \frac{m_d}{m_s} (V_L^*)_{is} (V_L)_{id} E_{2d}(x_i) + \sin \zeta \cos \zeta \frac{m_i}{m_s} (V_L^*)_{is} (V_L^*)_{id} e^{i(\phi_i - \psi_d + \alpha)} E_{3d}(x_i) \right\}, \quad (3.19)$$

$$\text{with } x_i \equiv m_i^2 / M_W^2,$$

where loop functions F_1, F_2, F_3 and E_{1d}, E_{2d}, E_{3d} are defined in appendix B. We are aware that the dipole operators receive two contributions with different phases when W' is integrated out and when W is. The two are expressed as $\delta C_g, \delta C_\gamma, \delta C'_g, \delta C'_\gamma$ and $\Delta C_g, \Delta C_\gamma, \Delta C'_g, \Delta C'_\gamma$, respectively.

We take into account RG evolutions of the Wilson coefficients at order $O(\alpha_s)$. The fact that four sets of operators, $(\{O_i\}), (\{O'_i\})$ ($i = 1, 2, \dots, 10, 1c, 2c$), $(\{O_j^{RL}\}), (\{O_j^{LR}\})$ ($j = 1, 2, 1c, 2c$), do not mix with each other facilitates the computation. For $(\{C_i\}), (\{C_j^{RL}\})$ and $(\{C_j^{LR}\})$, we assume that their initial conditions at scale $\mu = M_W$ are given by eqs. (3.10)–(3.15) and solve the RG equations from $\mu = M_W$ to the scale for which the lattice results are reported. For $(\{C'_i\})$, we assume that their initial conditions at $\mu = M_{W'}$ are provided by eqs. (3.2)–(3.9) and solve the RG equations from $\mu = M_{W'}$ to the scale of lattice results. Finally, we compute RG evolutions of the coefficients of the dipole operators (C_g, C_γ) and (C'_g, C'_γ) , which receive contributions from $(\{C_i, C_j^{RL}\})$ and $(\{C'_i, C_j^{LR}\})$, respectively. The $O(\alpha_s)$ RG equations for $(\{C_i\})$ and $(\{C_i^{RL}\})$ are found in ref. [31], and those for $(\{C_i^{LR}\})$ and (C_g, C_γ) are in ref. [30].

3.2 Hadronic matrix elements

We employ the lattice calculations of hadronic matrix elements $\langle (\pi\pi)_I | O_i | K^0 \rangle$ for $i = 1, 2, \dots, 10$ for $I = 0, 2$ reported by RBC/UKQCD in refs. [3, 4].

Since lattice calculations for the matrix elements of O_1^{LR} and O_2^{LR} are missing, we estimate them from the RBC/UKQCD results using isospin symmetry. In the limit of exact isospin symmetry, we find, for $\Delta I = 3/2$ amplitudes,

$$\begin{aligned} \langle (\pi\pi)_{I=2} | O_7 | K^0 \rangle &= \langle (\pi\pi)_{I=2} | (\bar{s}d)_L \left(\bar{u}u - \frac{1}{2} \bar{d}d - \frac{1}{2} \bar{s}s \right)_R | K^0 \rangle \\ &= \frac{3}{4} \langle (\pi\pi)_{I=2} | (\bar{s}d)_L (\bar{u}u - \bar{d}d)_R | K^0 \rangle \end{aligned} \quad (3.20)$$

$$= \frac{3}{4} \langle (\pi\pi)_{I=2} | \sqrt{2} (\bar{s}u)_L \sqrt{2} (\bar{u}d)_R | K^0 \rangle \quad (3.21)$$

$$= \frac{3}{2} \langle (\pi\pi)_{I=2} | O_2^{LR} | K^0 \rangle, \quad (3.22)$$

where we have discarded $\Delta I = 1/2$ part when obtaining the second line, and when deriving the third line, we have inserted Clebsch-Gordan coefficients for constructing the $\Delta I = 3/2$

operator from a $\Delta I = 1/2$ one and a $\Delta I = 1$ one. For $\Delta I = 1/2$ amplitudes, we find

$$\langle (\pi\pi)_{I=0} | \left(\frac{4}{3}O_7 + \frac{2}{3}O_5 \right) | K^0 \rangle = \langle (\pi\pi)_{I=0} | \{ (\bar{s}d)_L(\bar{u}u + \bar{d}d)_R + (\bar{s}d)_L(\bar{u}u - \bar{d}d)_R \} | K^0 \rangle \quad (3.23)$$

$$= \langle (\pi\pi)_{I=0} | \left\{ \sqrt{\frac{3}{2}}(\bar{s}u)_L\sqrt{2}(\bar{u}d)_R - \sqrt{\frac{1}{3}}\sqrt{\frac{3}{2}}(\bar{s}u)_L\sqrt{2}(\bar{u}d)_R \right\} | K^0 \rangle \quad (3.24)$$

$$= (\sqrt{3} - 1) \langle (\pi\pi)_{I=0} | O_2^{LR} | K^0 \rangle, \quad (3.25)$$

where in the first line, we have separated $(\bar{u}u)_R$ into $\Delta I = 0$ and $\Delta I = 1$ parts, and in the second line, we have inserted Clebsch-Gordan coefficients for constructing the $\Delta I = 1/2$ operator from a $\Delta I = 1/2$ one and a $\Delta I = 0$ one or a $\Delta I = 1$ one. The same relations hold between the matrix elements for O_1^{LR} and O_8, O_6 .

The hadronic matrix elements for the chromo-dipole operators O_g, O'_g are extracted from the calculation based on dual QCD approach [46]. Note that the above calculation is corroborated by the fact that it is consistent with a lattice calculation of the K - $\pi\pi$ one by chiral perturbation theory.

3.3 Numerical analysis of ϵ'/ϵ

The definition for the decay amplitudes of $K^0 \rightarrow \pi\pi$ is

$$A_0 e^{i\delta_0} = \langle (\pi\pi)_{I=0} | \mathcal{H}_{\Delta S=1} | K^0 \rangle, \quad A_2 e^{i\delta_2} = \langle (\pi\pi)_{I=2} | \mathcal{H}_{\Delta S=1} | K^0 \rangle, \quad (3.26)$$

where $\delta_{0,2}$ represent the strong phases. In terms of the above amplitudes, one writes the direct CP violation parameter divided by the indirect one as

$$\text{Re} \left(\frac{\epsilon'}{\epsilon} \right) = \text{Re} \left(\frac{i\omega e^{i(\delta_2 - \delta_0)}}{\sqrt{2}} \right) \left(\frac{\text{Im}A_2}{\text{Re}A_2} - \frac{\text{Im}A_0}{\text{Re}A_0} \right), \quad (3.27)$$

where $\omega = \text{Re}A_2/\text{Re}A_0$ is a suppression factor due to the $\Delta I = 1/2$ rule. For the strong phases, we use the values of refs. [3, 4], $\delta_2 = 23.8 \pm 5.0$ degree and $\delta_0 = -11.6 \pm 2.8$ degree. For the real parts of the decay amplitudes, we employ the experimental data [8], $\text{Re}A_2 = 1.479 \times 10^{-8}$ GeV and $\text{Re}A_0 = 33.20 \times 10^{-8}$ GeV, which leads to $\omega = 4.454 \times 10^{-2}$. In our analysis, we separate the SM and new physics contributions as

$$\text{Re} \left(\frac{\epsilon'}{\epsilon} \right) = \text{Re} \left(\frac{\epsilon'}{\epsilon} \right)_{\text{SM}} + \text{Re} \left(\frac{\epsilon'}{\epsilon} \right)_{\text{NP}}. \quad (3.28)$$

For the SM part, we quote the calculation in the literature $\text{Re}(\epsilon'/\epsilon)_{\text{SM}} = (1.38 \pm 6.90) \times 10^{-4}$ [4]. It is the new physics part,

$$\text{Re} \left(\frac{\epsilon'}{\epsilon} \right)_{\text{NP}} = \text{Re} \left(\frac{i\omega e^{i(\delta_2 - \delta_0)}}{\sqrt{2}} \right) \left(\frac{\text{Im}A_2^{\text{NP}}}{\text{Re}A_2} - \frac{\text{Im}A_0^{\text{NP}}}{\text{Re}A_0} \right), \quad (3.29)$$

that we compute in this paper. In doing so, we approximate $\cos^2 \zeta = 1$ in the Wilson coefficients eqs. (3.10)–(3.19), so that the SM contribution is separated from the new physics one at the operator level.

Figure 1. Numerical result of $\text{Re}(\epsilon'/\epsilon)$. The blue band represents the 1σ range of experimental data given by PDG [8], while model predictions with specific choice of phases are shown by the lines.

Figure 2. Numerical result of $\text{Re}(\epsilon'/\epsilon)$. The blue band represents the experimental data given by PDG [8], while each red dot corresponds to the model prediction with a randomly generated set of $(\alpha - \psi_d, \alpha - \psi_s)$.

In the analysis, we fix the ratio of the bifundamental scalar VEVs at its natural value as $\tan\beta = m_b/m_t$. We have found numerically that for $M_{W'} > 1$ TeV, the chromo-dipole contribution to $\text{Re}(\epsilon'/\epsilon)$ does not exceed $\mathcal{O}(10^{-4})$ and is hence safely neglected.⁴ Consequently, only two combinations of new CP phases, $\alpha - \psi_d$ and $\alpha - \psi_s$, and the W' mass determine the new physics contribution.

First, we choose specific values for the new CP phases in the calculation of ϵ'/ϵ to illustrate the model prediction. In figure 1, the prediction for $\text{Re}(\epsilon'/\epsilon)$ is presented with specific choices of $\alpha - \psi_d$ and $\alpha - \psi_s$.

Next, we randomly vary $\alpha - \psi_d$ and $\alpha - \psi_s$ in the range $[0, 2\pi]$, since they are free parameters. In figure 2, we show the region of $\text{Re}(\epsilon'/\epsilon)$ obtained by varying $\alpha - \psi_d$ and $\alpha - \psi_s$. One observes that $M_{W'} < 58$ TeV is necessary for 1σ explanation of the anomaly.

⁴When we use a calculation based on the chiral quark model in ref. [48] to evaluate the hadronic matrix elements of the chromo-dipole operators, we are again lead to the result that the chromo-dipole contribution to $\text{Re}(\epsilon'/\epsilon)$ is below $\mathcal{O}(10^{-4})$ for $M_{W'} > 1$ TeV.

We have confirmed that among the terms of $\Delta S = 1$ Hamiltonian eq. (3.1), $\sum_{i=1,2} C'_i O'_i$ and $\sum_{j=1,2} (C_j^{RL} O_j^{RL} + C_j^{LR} O_j^{LR})$ are the leading sources of the new physics contribution.

4 Neutron electric dipole moment

4.1 Wilson coefficients for operators contributing to the neutron EDM

In the effective QCD×QED theory in which W, W' and the top quark are integrated out, the part of the CP-violating Hamiltonian that contributes to the neutron EDM is parametrized as

$$\mathcal{H}_{nEDM} = \frac{G_F}{\sqrt{2}} \left\{ \sum_q \sum_{i=1,2,4,5} \mathcal{C}_{iq} \mathcal{O}_{iq} + \mathcal{C}_3 \mathcal{O}_3 + \sum_{q \neq q'} \sum_{i=1}^4 \mathcal{C}_{iq'q} \mathcal{O}_{iq'q} \right\}, \quad (4.1)$$

where operators $\mathcal{O}$'s are defined in appendix C.

We determine the Wilson coefficients $\mathcal{C}$'s as follows: again, for each coefficient, if multiple terms have an identical phase, we exclusively consider the one in the leading order of $M_W^2/M_{W'}^2$, or $\sin \zeta$. By integrating out W and the top quark, one obtains the following leading-order matching conditions at $\mu \sim M_W$ (note our convention with $\phi_u = 0$):

$$\mathcal{C}_{2du} = -\mathcal{C}_{2ud} = 4 \sin \zeta \cos \zeta \operatorname{Im} \left[(V_L)_{ud} (V_L)_{ud} e^{i(\psi_d - \alpha)} \right], \quad (d \rightarrow s), \quad (4.2)$$

$$\mathcal{C}_{2dc} = -\mathcal{C}_{2cd} = 4 \sin \zeta \cos \zeta \operatorname{Im} \left[(V_L)_{cd} (V_L)_{cd} e^{i(-\phi_c + \psi_d - \alpha)} \right], \quad (d \rightarrow s), \quad (4.3)$$

$$\mathcal{C}_{1u} = \frac{e}{4\pi^2} \sin \zeta \cos \zeta \sum_{i=d,s,b} \frac{m_i}{m_u} \operatorname{Im} \left[(V_L)_{ui} (V_L)_{ui} e^{i(\psi_i - \alpha)} \right] E_{3u}(x_i), \quad (4.4)$$

$$\mathcal{C}_{1c} = \frac{e}{4\pi^2} \sin \zeta \cos \zeta \sum_{i=d,s,b} \frac{m_i}{m_u} \operatorname{Im} \left[(V_L)_{ui} (V_L)_{ui} e^{i(-\phi_c + \psi_i - \alpha)} \right] E_{3u}(x_i), \quad (4.5)$$

$$\mathcal{C}_{1d} = \frac{e}{4\pi^2} \sin \zeta \cos \zeta \sum_{i=u,c,t} \frac{m_i}{m_d} \operatorname{Im} \left[(V_L)_{id} (V_L)_{id} e^{i(-\phi_i + \psi_d - \alpha)} \right] E_{3d}(x_i), \quad (d \rightarrow s), \quad (4.6)$$

$$\mathcal{C}_{2u} = -\frac{g_s}{4\pi^2} \sin \zeta \cos \zeta \sum_{i=d,s,b} \frac{m_i}{m_u} \operatorname{Im} \left[(V_L)_{ui} (V_L)_{ui} e^{i(\psi_i - \alpha)} \right] F_3(x_i), \quad (4.7)$$

$$\mathcal{C}_{2c} = -\frac{g_s}{4\pi^2} \sin \zeta \cos \zeta \sum_{i=d,s,b} \frac{m_i}{m_u} \operatorname{Im} \left[(V_L)_{ui} (V_L)_{ui} e^{i(-\phi_c + \psi_i - \alpha)} \right] F_3(x_i), \quad (4.8)$$

$$\mathcal{C}_{2d} = \frac{g_s}{4\pi^2} \sin \zeta \cos \zeta \sum_{i=u,c,t} \frac{m_i}{m_d} \operatorname{Im} \left[(V_L)_{id} (V_L)_{id} e^{i(-\phi_i + \psi_d - \alpha)} \right] F_3(x_i), \quad (d \rightarrow s), \quad (4.9)$$

$$\mathcal{C}_3 \simeq \frac{4g_s^3}{(16\pi^2)^2} \sin \zeta \cos \zeta \frac{m_t}{m_b} F_3(x_t) \operatorname{Im} \left[(V_L)_{tb} (V_L)_{tb} e^{i(-\phi_t + \psi_b - \alpha)} \right], \quad (4.10)$$

with $x_i \equiv m_i^2/M_W^2$,

where loop functions F_1, F_2, F_3 and $E_{1d}, E_{2d}, E_{3d}, E_{3u}$ are defined in appendix B. In eq. (4.10) (which corresponds to the Weinberg operator [49]), we present the dominant part proportional to m_t . Terms obtained by integrating out W' possess the same phases as eqs. (4.2)–(4.10) and are simply suppressed by $M_W^2/M_{W'}^2$, compared to eqs. (4.2)–(4.10). They are therefore neglected in our analysis.

The RG equations at order $O(\alpha_s)$ for the Wilson coefficients are obtainable in refs. [41, 42, 44, 45]. We assume that the initial conditions at $\mu = M_W$ for the RG equations are given by eqs. (4.2)–(4.10), and solve the equations from $\mu = M_W$ to $\mu = 1$ GeV. At the 1 GeV scale, we evaluate the hadronic matrix elements.

4.2 Hadronic matrix elements

4.2.1 Four-quark operators $\mathcal{O}_{1q'q}$, $\mathcal{O}_{2q'q}$

In the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry, the Wilson coefficients for the four-quark operators $\mathcal{O}_{1q'q} = (\bar{q}'q)(\bar{q}i\gamma_5 q)$ and $\mathcal{O}_{2q'q} = (\bar{q}'_\alpha q'_\beta)(\bar{q}_\beta i\gamma_5 q_\alpha)$ ($q \neq q'$; $q, q' = u, d, s$) are particularly large. Therefore, we scrutinize how these operators contribute to the neutron EDM. Operators $\mathcal{O}_{1q'q}$ contribute in the following three ways:

- The first one is through meson condensation [33]; $\mathcal{O}_{1q'q}$ operators give rise to tadpole terms for pseudoscalar mesons and induce their VEVs. These VEVs generate CP-violating interactions for baryons and mesons, which contribute to the neutron EDM through baryon-meson loop diagrams.
- The second one is through hadronic matrix elements of $\mathcal{O}_{1q'q}$ with baryons and mesons, $\langle BM | \mathcal{O}_{1q'q} | B \rangle$ (B denotes a baryon and M a meson), which contribute to the neutron EDM through baryon-meson loop diagrams.
- The third one is directly through the hadronic matrix element of $\mathcal{O}_{1q'q}$ with neutrons and photon.

On the other hand, $\mathcal{O}_{2q'q}$ operators do not yield meson condensation, but do contribute to the neutron EDM in the latter two ways. Later, it will be shown that the contribution from the pion VEV $\langle \pi^0 \rangle$, which belongs to the first category, is enhanced by the factor $m_s/(m_u + m_d)$ compared to the latter two. We therefore investigate how $\mathcal{O}_{1q'q}$ operators bring about meson condensation, thereby contributing to the neutron EDM.

We are aware that if Peccei-Quinn mechanism [50] exists, it affects the meson condensation and also induces an effective non-zero $\bar{\theta}$ term due to incomplete cancellation between the genuine $\bar{\theta}$ term and the axion VEV. Alternatively, it is logically possible to assume $\bar{\theta} = 0$ without Peccei-Quinn mechanism, by considering an unknown mechanism or through fine-tuning, in which case we do not need to take into account the effect of Peccei-Quinn mechanism or that of non-zero $\bar{\theta}$. In this paper, we consider both cases where (i) one has $\bar{\theta} = 0$ without Peccei-Quinn mechanism, and (ii) Peccei-Quinn mechanism is at work.

We start from the case with $\bar{\theta} = 0$ without Peccei-Quinn mechanism. The meson condensation contribution is evaluated by the following steps:

(1). First, we implement $\sum \mathcal{C}_{1q'q} \mathcal{O}_{1q'q}$ part of the Hamiltonian eq. (4.1) into the meson chiral Lagrangian. To this end, we rewrite

$$\sum_{q \neq q'; q, q' = u, d, s} \mathcal{C}_{1q'q} \mathcal{O}_{1q'q} = \sum_{i, j, k, l = u, d, s} \left\{ i \mathcal{C}_{ijkl}^{LRLR} (\bar{q}_{iL} q_{jR}) (\bar{q}_{kL} q_{lR}) + i \mathcal{C}_{ijkl}^{RLLR} (\bar{q}_{iR} q_{jL}) (\bar{q}_{kL} q_{lR}) \right\} - (L \leftrightarrow R), \quad (4.11)$$

$$\text{with } \mathcal{C}_{ijkl}^{LRLR} = \mathcal{C}_{ijkl}^{RLLR} \equiv \sum_{q \neq q'} \mathcal{C}_{1q'q} \delta_{iq'} \delta_{jq'} \delta_{kq} \delta_{lq}. \quad (4.12)$$

It then becomes clear that the theory would be invariant (except for $U(1)_A$ anomaly) if coefficients $\mathcal{C}_{ijkl}^{LRLR}$ and $\mathcal{C}_{ijkl}^{RLLR}$ transformed under $U(3)_L \times U(3)_R$ rotations, $L \times R$, as

$$\begin{aligned} \mathcal{C}_{ijkl}^{LRLR} &\rightarrow \sum_{m, n, o, p} (L)_{im} (L)_{ko} \mathcal{C}_{mnop}^{LRLR} (R^\dagger)_{nj} (R^\dagger)_{pl}, \\ \mathcal{C}_{ijkl}^{RLLR} &\rightarrow \sum_{m, n, o, p} (R)_{im} (L)_{ko} \mathcal{C}_{mnop}^{RLLR} (L^\dagger)_{nj} (R^\dagger)_{pl}. \end{aligned}$$

From the above transformation property and the parity invariance of QCD, the meson chiral Lagrangian at order $O(p^2)$ plus the leading CP-violating terms is found to be (remind that $\bar{\theta} = 0$ has been assumed)

$$\begin{aligned} \mathcal{L}_{\text{mesons}} &= \frac{F_\pi^2}{4} \text{tr} \left[(D_\mu U)^\dagger D^\mu U + \chi (U + U^\dagger) \right] + \frac{F_0^2 - F_\pi^2}{12} \text{tr} \left[U D_\mu U^\dagger \right] \text{tr} \left[U^\dagger D^\mu U \right] \\ &+ a_0 \text{tr} \left[\log U - \log U^\dagger \right]^2 \\ &+ \frac{G_F}{\sqrt{2}} \sum_{i, j, k, l = u, d, s} \left\{ i \mathcal{C}_{ijkl}^{LRLR} \left(c_1 [U]_{ji} [U]_{lk} - c_1 [U^\dagger]_{ji} [U^\dagger]_{lk} + c_2 [U]_{li} [U]_{jk} - c_2 [U^\dagger]_{li} [U^\dagger]_{jk} \right) \right. \\ &\left. + i \mathcal{C}_{ijkl}^{RLLR} \left(c_3 [U^\dagger]_{ji} [U]_{lk} - c_3 [U]_{ji} [U^\dagger]_{lk} \right) \right\}, \end{aligned} \quad (4.13)$$

where $\mathcal{C}_{ijkl}^{LRLR}$, $\mathcal{C}_{ijkl}^{RLLR}$ have been defined in eq. (4.12). Here, U is a nonlinear representation of the nine Nambu-Goldstone bosons that transforms under $U(3)_L \times U(3)_R$ rotations $L \times R$ as $U \rightarrow R U L^\dagger$, and χ includes the quark mass term, which are given by

$$U = \exp \left[\frac{2i}{\sqrt{6} F_0} \eta_0 I_3 + \frac{2i}{F_\pi} \Pi \right], \quad \Pi \equiv \begin{pmatrix} \frac{1}{2} \pi^0 + \frac{1}{2\sqrt{3}} \eta_8 & \frac{1}{\sqrt{2}} \pi^+ & \frac{1}{\sqrt{2}} K^+ \\ \frac{1}{\sqrt{2}} \pi^- & -\frac{1}{2} \pi^0 + \frac{1}{2\sqrt{3}} \eta_8 & \frac{1}{\sqrt{2}} K^0 \\ \frac{1}{\sqrt{2}} K^- & \frac{1}{\sqrt{2}} \bar{K}^0 & -\frac{1}{\sqrt{3}} \eta_8 \end{pmatrix}, \quad (4.14)$$

$$I_3 \equiv \text{diag}(1, 1, 1),$$

$$\chi = 2B_0 \text{diag}(m_u, m_d, m_s). \quad (4.15)$$

$[U]_{ij}$ denotes the (i, j) component of matrix U . F_π is the pion decay constant in the chiral limit and F_0 is the decay constant for η_0 , which we approximate as $F_0 \simeq F_\pi$. B_0 satisfies $B_0 \simeq m_\pi^2 / (m_u + m_d)$. The term with $\log U$ represents instanton effects, whose expression is exact in the large N_c limit [51], and a_0 satisfies $48a_0/F_0^2 \simeq m_\eta^2 + m_{\eta'}^2 - 2m_K^2$. c_1 , c_2 and c_3

are unknown low energy constants (LECs), which can be estimated by naïve dimensional analysis [53] as

$$c_1 \sim c_2 \sim c_3 \sim \frac{(4\pi F_\pi)^6}{(4\pi)^4}. \quad (4.16)$$

(2). The CP-violating part of the Lagrangian eq. (4.13) contains tadpole terms for mesons, which lead to non-zero meson VEVs. Assuming that electric charge and strangeness are not broken spontaneously, we obtain the following potential for neutral mesons π^0 , η_8 and η_0 :

$$\begin{aligned} V(\pi^0, \eta_8, \eta_0) = & F_\pi^2 B_0 \left\{ m_u \cos \left(\frac{\pi^0}{F_\pi} + \frac{\eta_8}{\sqrt{3}F_\pi} + \frac{2\eta_0}{\sqrt{6}F_0} \right) + m_d \cos \left(-\frac{\pi^0}{F_\pi} + \frac{\eta_8}{\sqrt{3}F_\pi} + \frac{2\eta_0}{\sqrt{6}F_0} \right) \right. \\ & + m_s \cos \left(-\frac{2\eta_8}{\sqrt{3}F_\pi} + \frac{2\eta_0}{\sqrt{6}F_0} \right) \Big\} - 24 \frac{a_0}{F_0^2} (\eta_0)^2 \\ & - \frac{G_F}{\sqrt{2}} 2c_1 \left\{ (\mathcal{C}_{1ud} + \mathcal{C}_{1du}) \sin \left(\frac{2\eta_8}{\sqrt{3}F_\pi} + \frac{4\eta_0}{\sqrt{6}F_0} \right) \right. \\ & + (\mathcal{C}_{1us} + \mathcal{C}_{1su}) \sin \left(\frac{\pi^0}{F_\pi} - \frac{\eta_8}{\sqrt{3}F_\pi} + \frac{4\eta_0}{\sqrt{6}F_0} \right) \\ & + (\mathcal{C}_{1ds} + \mathcal{C}_{1sd}) \sin \left(-\frac{\pi^0}{F_\pi} - \frac{\eta_8}{\sqrt{3}F_\pi} + \frac{4\eta_0}{\sqrt{6}F_0} \right) \Big\} \\ & - \frac{G_F}{\sqrt{2}} 2c_3 \left\{ (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) \sin \left(-\frac{2\pi^0}{F_\pi} \right) + (\mathcal{C}_{1us} - \mathcal{C}_{1su}) \sin \left(-\frac{\pi^0}{F_\pi} - \frac{\sqrt{3}\eta_8}{F_\pi} \right) \right. \\ & \left. + (\mathcal{C}_{1ds} - \mathcal{C}_{1sd}) \sin \left(\frac{\pi^0}{F_\pi} - \frac{\sqrt{3}\eta_8}{F_\pi} \right) \right\}. \end{aligned} \quad (4.17)$$

The above potential is minimized with non-zero meson VEVs, $\langle \pi^0 \rangle$, $\langle \eta_8 \rangle$ and $\langle \eta_0 \rangle$. Insofar as we are concerned with vertices with one meson, the physical modes of π^0 , η_8 and η_0 fields can be approximated as

$$\pi_{\text{phys}}^0 \simeq \pi^0 - \langle \pi^0 \rangle, \quad \eta_{8\text{phys}} \simeq \eta_8 - \langle \eta_8 \rangle, \quad \eta_{0\text{phys}} \simeq \eta_0 - \langle \eta_0 \rangle. \quad (4.18)$$

In the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry, there hold relations $\mathcal{C}_{1ud} \simeq -\mathcal{C}_{1du}$ and $|\mathcal{C}_{1ud}| \gg |\mathcal{C}_{1sq}|, |\mathcal{C}_{1qs}|$ ($q = u, d$). When $(\mathcal{C}_{1ud} + \mathcal{C}_{1du})$ and $\mathcal{C}_{1sq}, \mathcal{C}_{1qs}$ are neglected accordingly, one finds, for small VEVs,

$$\begin{aligned} \frac{\langle \pi^0 \rangle}{F_\pi} & \simeq \frac{G_F}{\sqrt{2}} (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) \frac{c_3}{B_0 F_\pi^2} \frac{B_0 F_\pi^2 (m_u + m_d) m_s + 8a_0 (m_u + m_d + 4m_s)}{B_0 F_\pi^2 m_u m_d m_s + 8a_0 (m_u m_d + m_d m_s + m_s m_u)}, \\ \frac{\langle \eta_8 \rangle}{F_\pi} & \simeq \frac{G_F}{\sqrt{2}} (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) \frac{c_3}{\sqrt{3} B_0 F_\pi^2} (m_d - m_u) \frac{B_0 F_\pi^2 m_s + 24a_0}{B_0 F_\pi^2 m_u m_d m_s + 8a_0 (m_u m_d + m_d m_s + m_s m_u)}, \\ \frac{\langle \eta_0 \rangle}{F_0} & \simeq \frac{G_F}{\sqrt{2}} (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) \frac{\sqrt{2} c_3}{\sqrt{3} B_0 F_\pi^2} (m_d - m_u) \frac{B_0 F_\pi^2 m_s}{B_0 F_\pi^2 m_u m_d m_s + 8a_0 (m_u m_d + m_d m_s + m_s m_u)}. \end{aligned} \quad (4.19)$$

Note that $\langle \eta_8 \rangle$ and $\langle \eta_0 \rangle$ are proportional to $m_d - m_u$. This is because these VEVs are isospin singlets and hence must be constructed from the product of isospin-odd coefficient $\mathcal{C}_{1ud} - \mathcal{C}_{1du}$ and isospin-odd mass term $m_d - m_u$. In contrast, $\langle \pi^0 \rangle$ does not contain $m_d - m_u$ because this VEV is isospin-violating. Since $20a_0 \sim B_0 F_\pi^2 m_s$ holds empirically, we find from eq. (4.19) that $\langle \pi^0 \rangle$ is much larger than $\langle \eta_8 \rangle$ and $\langle \eta_0 \rangle$ by the factor $m_s/(m_d - m_u)$.

(3). Meson condensation breaks CP symmetry (and $U(3)_L \times U(3)_R$ symmetry) and induces CP-violating interactions for baryons and mesons. To study these interactions, we write the baryon chiral Lagrangian at order $O(p^2)$ as (terms irrelevant in the current discussion are omitted)

$$\begin{aligned} \mathcal{L}_{\text{baryons}} = & \text{tr} [\bar{B} i \gamma^\mu (\partial_\mu B + [\Gamma_\mu, B]) - M_B \bar{B} B] \\ & - \frac{D}{2} \text{tr} [\bar{B} \gamma^\mu \gamma_5 \{\xi_\mu, B\}] - \frac{F}{2} \text{tr} [\bar{B} \gamma^\mu \gamma_5 [\xi_\mu, B]] - \frac{\lambda}{2} \text{tr} [\xi_\mu] \text{tr} [\bar{B} \gamma^\mu \gamma_5 B] \\ & + b_D \text{tr} [\bar{B} \{\chi_+, B\}] + b_F \text{tr} [\bar{B} [\chi_+, B]] + b_0 \text{tr} [\chi_+] \text{tr} [\bar{B} B] + \dots, \end{aligned} \quad (4.20)$$

where B represents baryons and ξ_L, ξ_R include mesons as

$$B = \begin{pmatrix} \frac{1}{\sqrt{2}} \Sigma^0 + \frac{1}{\sqrt{6}} \Lambda^0 & \Sigma^+ & p \\ \Sigma^- & -\frac{1}{\sqrt{2}} \Sigma^0 + \frac{1}{\sqrt{6}} \Lambda^0 & n \\ \Xi^- & \Xi^0 & -\frac{2}{\sqrt{6}} \Lambda^0 \end{pmatrix}, \quad (4.21)$$

$$U = \xi_R \xi_L^\dagger, \quad (4.22)$$

$$\xi_R = \xi_L^\dagger. \quad (4.23)$$

Γ_μ is a covariant derivative for baryons, ξ_μ is a combination of meson fields, and χ_+ contains quark masses, which are defined as

$$\Gamma_\mu \equiv \frac{1}{2} \xi_R^\dagger (\partial_\mu - i r_\mu) \xi_R + \frac{1}{2} \xi_L^\dagger (\partial_\mu - i l_\mu) \xi_L, \quad (4.24)$$

$$\xi_\mu \equiv i \xi_R^\dagger (\partial_\mu - i r_\mu) \xi_R - i \xi_L^\dagger (\partial_\mu - i l_\mu) \xi_L, \quad (4.25)$$

$$\chi_+ \equiv 2B_0 \xi_L^\dagger \text{diag}(m_u, m_d, m_s) \xi_R + 2B_0 \xi_R^\dagger \text{diag}(m_u, m_d, m_s) \xi_L. \quad (4.26)$$

M_B is the baryon mass in the chiral limit. We insert meson VEVs $\langle \pi^0 \rangle, \langle \eta_8 \rangle, \langle \eta_0 \rangle$ into the baryon chiral Lagrangian eq. (4.20) and extract CP-violating interaction terms involving neutron n . We thus obtain

$$\begin{aligned} \mathcal{L}_{\text{baryons}} \supset & \bar{g}_{nn\pi} \bar{n} n \pi_{\text{phys}}^0 + \bar{g}_{nn8} \bar{n} n \eta_{8\text{phys}} + \bar{g}_{nn0} \bar{n} n \eta_{0\text{phys}} + \bar{g}_{np\pi} (\bar{p} n \pi^+ + \bar{n} p \pi^-) \\ & + \bar{g}_{n\Sigma^0 K^0} (\bar{\Sigma}^0 n \bar{K}^0 + \bar{n} \Sigma^0 K^0) + \bar{g}_{n\Sigma^- K^+} (\bar{\Sigma}^+ n K^- + \bar{n} \Sigma^- K^+) \\ & + \bar{g}_{n\Lambda K} (\bar{\Lambda} n \bar{K}^0 + \bar{n} \Lambda K^0), \end{aligned} \quad (4.27)$$

where the coupling constants are given by

$$\bar{g}_{nn\pi} = \frac{B_0}{F_\pi} \left[4 \{ -b_0(m_u + m_d) - (b_D + b_F)m_d \} \frac{\langle \pi^0 \rangle}{F_\pi} + \frac{4}{\sqrt{3}} \{ b_0(m_d - m_u) + (b_D + b_F)m_d \} \left(\frac{\langle \eta_8 \rangle}{F_\pi} + \sqrt{2} \frac{\langle \eta_0 \rangle}{F_0} \right) \right], \quad (4.28)$$

$$\bar{g}_{nn8} = \frac{B_0}{F_\pi} \left[\frac{4}{\sqrt{3}} \{ b_0(m_d - m_u) + (b_D + b_F)m_d \} \frac{\langle \pi^0 \rangle}{F_\pi} - \frac{4}{3} \{ b_0(m_u + m_d + 4m_s) + b_D(m_d + 4m_s) + b_F(m_d - 4m_s) \} \frac{\langle \eta_8 \rangle}{F_\pi} - \frac{4\sqrt{2}}{3} \{ b_0(m_u + m_d - 2m_s) + b_D(m_d - 2m_s) + b_F(m_d + 2m_s) \} \frac{\langle \eta_0 \rangle}{F_0} \right], \quad (4.29)$$

$$\bar{g}_{nn0} = \frac{B_0}{F_0} \left[\frac{4\sqrt{2}}{\sqrt{3}} \{ b_0(m_d - m_u) + (b_D + b_F)m_d \} \frac{\langle \pi^0 \rangle}{F_\pi} - \frac{4\sqrt{2}}{3} \{ b_0(m_u + m_d - 2m_s) + b_D(m_d - 2m_s) + b_F(m_d + 2m_s) \} \frac{\langle \eta_8 \rangle}{F_\pi} - \frac{8}{3} \{ b_0(m_u + m_d + m_s) + b_D(m_d + m_s) + b_F(m_d - m_s) \} \frac{\langle \eta_0 \rangle}{F_0} \right], \quad (4.30)$$

$$\bar{g}_{np\pi} = \frac{B_0}{F_\pi} (b_D + b_F) \left[\sqrt{2}(m_d - m_u) \frac{\langle \pi^0 \rangle}{F_\pi} - \frac{2\sqrt{2}}{\sqrt{3}} (m_u + m_d) \left(\frac{\langle \eta_8 \rangle}{F_\pi} + \sqrt{2} \frac{\langle \eta_0 \rangle}{F_0} \right) \right], \quad (4.31)$$

$$\bar{g}_{n\Sigma^0 K^0} = \frac{B_0}{F_\pi} (b_D - b_F) \left[-\frac{1}{2} (3m_d + m_s) \frac{\langle \pi^0 \rangle}{F_\pi} + \frac{1}{2\sqrt{3}} (m_d - 5m_s) \frac{\langle \eta_8 \rangle}{F_\pi} + \frac{2\sqrt{2}}{\sqrt{3}} (m_d + m_s) \frac{\langle \eta_0 \rangle}{F_0} \right], \quad (4.32)$$

$$\bar{g}_{n\Sigma^- K^+} = \frac{B_0}{F_\pi} (b_D - b_F) \left[-\frac{1}{\sqrt{2}} (3m_u + m_s) \frac{\langle \pi^0 \rangle}{F_\pi} + \frac{1}{\sqrt{6}} (-m_u + 5m_s) \frac{\langle \eta_8 \rangle}{F_\pi} - \frac{4}{\sqrt{3}} (m_u + m_s) \frac{\langle \eta_0 \rangle}{F_0} \right] \quad (4.33)$$

$$\bar{g}_{n\Lambda K} = \frac{B_0}{F_\pi} (b_D + 3b_F) \left[-\frac{1}{2\sqrt{3}} (3m_d + m_s) \frac{\langle \pi^0 \rangle}{F_\pi} + \frac{1}{6} (m_d - 5m_s) \frac{\langle \eta_8 \rangle}{F_\pi} + \frac{2\sqrt{2}}{3} (m_d + m_s) \frac{\langle \eta_0 \rangle}{F_0} \right]. \quad (4.34)$$

Note in particular that $\langle \pi^0 \rangle$ enters into the expression for $\bar{g}_{n\Sigma^- K^+}$ eq. (4.33) without the factor of $m_d - m_u$, which is allowed because the coupling $\bar{g}_{n\Sigma^- K^+}$ violates isospin. It follows that $\bar{g}_{n\Sigma^- K^+}$ is enhanced by the factor $m_s/(m_u + m_d)$, as it contains a term $m_s \langle \pi^0 \rangle$.

We compare the above meson-VEV-induced CP-violating couplings with those arising from direct hadronic matrix elements of $\mathcal{O}_{1q'q}$ and $\mathcal{O}_{2q'q}$. The latter are estimated by naïve

dimensional analysis [53] as⁵ (B and M represent any baryon and meson, respectively)

$$\bar{g}_{BBM}|_{\text{direct}} \sim \frac{G_F}{\sqrt{2}} \sum_{i=1,2} \sum_{q,q'} |\mathcal{C}_{iq'q}| \frac{1}{F_\pi} \frac{(4\pi F_\pi)^3}{(4\pi)^2}. \quad (4.35)$$

On the other hand, $\bar{g}_{n\Sigma^-K^+}$ eq. (4.33), for example, is estimated to be

$$\bar{g}_{n\Sigma^-K^+} \simeq -\frac{B_0}{F_\pi} (b_D - b_F) \frac{m_s \langle \pi^0 \rangle}{\sqrt{2} F_\pi} \sim -\frac{G_F}{\sqrt{2}} (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) (b_D - b_F) \frac{4m_s}{\sqrt{2}(m_u + m_d)} \frac{(4\pi F_\pi)^6}{(4\pi)^4 F_\pi^3}, \quad (4.36)$$

where eq. (4.19) and the naïve dimensional analysis on c_3 eq. (4.16) are in use. Noting that $(b_D - b_F)(4\pi F_\pi) \sim 1$ holds numerically, we observe that the meson VEV contribution eq. (4.36) dominates over the direct hadronic matrix element one eq. (4.35) by the factor $m_s/(m_u + m_d)$. This fact allows us to neglect the latter contribution in the rest of the analysis.

(4). The neutron EDM receives contributions from baryon-meson loop diagrams involving a CP-violating coupling of eqs. (4.27), a CP-conserving baryon-meson axial-vector coupling and a photon coupling. We refer to the loop calculation of ref. [58] performed with infrared regularization [60, 61], from which the neutron EDM, d_n , is obtained as

$$d_n|_{\text{loop}} = \frac{e}{8\pi^2 F_\pi} \left\{ \frac{\bar{g}_{np\pi}}{\sqrt{2}} (D + F) \left(\frac{1}{\bar{\epsilon}} - 1 - \log \frac{m_\pi^2}{\mu^2} + \frac{\pi m_\pi}{2m_N} \right) - \frac{\bar{g}_{n\Sigma^-K^+}}{\sqrt{2}} (D - F) \left(\frac{1}{\bar{\epsilon}} - 1 - \log \frac{m_K^2}{\mu^2} + \frac{\pi m_K}{2m_N} - \frac{\pi(m_\Sigma - m_N)}{m_K} \right) \right\}. \quad (4.37)$$

Here, the divergent part $1/\bar{\epsilon} \equiv 1/\epsilon - \gamma_E + \log(4\pi)$ and the scale μ stem from dimensional regularization in $4 - 2\epsilon$ dimension with mass parameter μ . In fact, the baryon chiral Lagrangian contains a LEC which cancels the above divergence and whose finite part contributes to the neutron EDM. The impact of the finite part of the LEC is assessed by naïve dimensional analysis [53] as

$$d_n|_{\text{LEC}} \sim \frac{G_F}{\sqrt{2}} \sum_{i=1,2} \sum_{q,q'} |\mathcal{C}_{iq'q}| e \frac{4\pi F_\pi}{(4\pi)^2}. \quad (4.38)$$

On the other hand, from eqs. (4.19) and (4.33) and the estimate on c_3 eq. (4.16), the finite part of the loop contribution eq. (4.37) is estimated to be

$$\begin{aligned} d_n|_{\text{loop}} &\sim \frac{e}{8\pi^2 F_\pi} \frac{\bar{g}_{n\Sigma^-K^+}}{\sqrt{2}} (D - F) \\ &\sim -\frac{e}{8\pi^2 F_\pi} \frac{G_F}{\sqrt{2}} (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) (b_D - b_F) \frac{2m_s}{m_u + m_d} \frac{(4\pi F_\pi)^6}{(4\pi)^4 F_\pi^3} (D - F). \end{aligned} \quad (4.39)$$

Since $(b_D - b_F)(4\pi F_\pi) \sim 1$ and $D - F \sim 1$, we find that the loop contribution eq. (4.39) dominates over the LEC one eq. (4.38) by the factor $m_s/(m_u + m_d)$. It is thus justifiable

⁵There are also studies in which the direct hadronic matrix elements are estimated with vacuum saturation approximation [54–57] and with hadron models [35, 36].

to estimate d_n by simply extracting the finite part of the loop contribution. We further set $\mu = m_N$, since m_N is a natural cutoff scale, and arrive at

$$d_n \sim \frac{e}{8\pi^2 F_\pi} \left\{ \frac{\bar{g}_{np\pi}}{\sqrt{2}} (D + F) \left(-1 - \log \frac{m_\pi^2}{m_N^2} + \frac{\pi m_\pi}{2m_N} \right) - \frac{\bar{g}_{n\Sigma^- K^+}}{\sqrt{2}} (D - F) \left(-1 - \log \frac{m_K^2}{m_N^2} + \frac{\pi m_K}{2m_N} - \frac{\pi(m_\Sigma - m_N)}{m_K} \right) \right\}. \quad (4.40)$$

Next, we study the case with Peccei-Quinn mechanism. We incorporate the axion field, a , into the meson Lagrangian eq. (4.13) by performing $U(3)_A$ chiral rotations to remove the gluon theta term and transform the quark fields as

$$\begin{aligned} u_L &\rightarrow e^{-i\alpha_u/2} u_L, & u_R &\rightarrow e^{i\alpha_u/2} u_R, & d_L &\rightarrow e^{-i\alpha_d/2} d_L, & d_R &\rightarrow e^{i\alpha_d/2} d_R, \\ s_L &\rightarrow e^{-i\alpha_s/2} s_L, & s_R &\rightarrow e^{i\alpha_s/2} s_R, \end{aligned} \quad (4.41)$$

where $\alpha_u, \alpha_d, \alpha_s$ include the axion field a as

$$\begin{aligned} \alpha_u &= \frac{m_d m_s}{m_u m_d + m_d m_s + m_s m_u} \left(\frac{a}{f_a} + \bar{\theta} \right), \\ \alpha_d &= \frac{m_s m_u}{m_u m_d + m_d m_s + m_s m_u} \left(\frac{a}{f_a} + \bar{\theta} \right), \\ \alpha_s &= \frac{m_u m_d}{m_u m_d + m_d m_s + m_s m_u} \left(\frac{a}{f_a} + \bar{\theta} \right), \end{aligned}$$

with f_a denoting the axion decay constant and $\bar{\theta}$ being the genuine theta term. (With the above choice of $\alpha_u, \alpha_d, \alpha_s$, the axion does not mix with π^0 or η_8 .) As a result, the axion field is associated with the quark masses and the coefficients $\mathcal{C}_{1q'q}$, and can thus be implemented in the meson chiral Lagrangian through these terms. Accordingly, the meson potential eq. (4.17) is modified to the potential of π^0, η_8, η_0 and axion a ,

$$\begin{aligned} V(\pi^0, \eta_8, \eta_0, a) &= F_\pi^2 B_0 \left\{ m_u \cos \left(\frac{\pi^0}{F_\pi} + \frac{\eta_8}{\sqrt{3}F_\pi} + \frac{2\eta_0}{\sqrt{6}F_0} + \alpha_u \right) \right. \\ &\quad + m_d \cos \left(-\frac{\pi^0}{F_\pi} + \frac{\eta_8}{\sqrt{3}F_\pi} + \frac{2\eta_0}{\sqrt{6}F_0} + \alpha_d \right) \\ &\quad + m_s \cos \left(-\frac{2\eta_8}{\sqrt{3}F_\pi} + \frac{2\eta_0}{\sqrt{6}F_0} + \alpha_s \right) \Big\} - 24 \frac{a_0}{F_0^2} (\eta_0)^2 \\ &\quad - 2c_1 \left\{ (\mathcal{C}_{1ud} + \mathcal{C}_{1du}) \sin \left(\frac{2\eta_8}{\sqrt{3}F_\pi} + \frac{4\eta_0}{\sqrt{6}F_0} + \alpha_u + \alpha_d \right) \right. \\ &\quad + (\mathcal{C}_{1us} + \mathcal{C}_{1su}) \sin \left(\frac{\pi^0}{F_\pi} - \frac{\eta_8}{\sqrt{3}F_\pi} + \frac{4\eta_0}{\sqrt{6}F_0} + \alpha_u + \alpha_s \right) \\ &\quad + (\mathcal{C}_{1ds} + \mathcal{C}_{1sd}) \sin \left(-\frac{\pi^0}{F_\pi} - \frac{\eta_8}{\sqrt{3}F_\pi} + \frac{4\eta_0}{\sqrt{6}F_0} + \alpha_d + \alpha_s \right) \Big\} \\ &\quad - 2c_3 \left\{ (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) \sin \left(-\frac{2\pi^0}{F_\pi} - \alpha_u + \alpha_d \right) \right. \\ &\quad + (\mathcal{C}_{1us} - \mathcal{C}_{1su}) \sin \left(-\frac{\pi^0}{F_\pi} - \frac{\sqrt{3}\eta_8}{F_\pi} - \alpha_u + \alpha_s \right) \\ &\quad + (\mathcal{C}_{1ds} - \mathcal{C}_{1sd}) \sin \left(\frac{\pi^0}{F_\pi} - \frac{\sqrt{3}\eta_8}{F_\pi} - \alpha_d + \alpha_s \right) \Big\}. \end{aligned} \quad (4.42)$$

where it should be reminded that $\alpha_u, \alpha_d, \alpha_s$ are functions of a . The minimization condition for eq. (4.42) yields meson VEVs $\langle \pi^0 \rangle$, $\langle \eta_8 \rangle$, $\langle \eta_0 \rangle$ and an axion VEV $\langle a \rangle$. When only the term $(\mathcal{C}_{1ud} - \mathcal{C}_{1du})$ is non-zero, these VEVs are given by

$$\begin{aligned} \frac{\langle \pi^0 \rangle}{F_\pi} &\simeq \frac{G_F}{\sqrt{2}} (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) \frac{c_3}{B_0 F_\pi^2} \frac{m_u + m_d + 4m_s}{m_u m_d + m_d m_s + m_s m_u}, \\ \frac{\langle \eta_8 \rangle}{F_\pi} &\simeq \frac{G_F}{\sqrt{2}} (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) \frac{\sqrt{3} c_3}{B_0 F_\pi^2} (m_d - m_u) \frac{1}{m_u m_d + m_d m_s + m_s m_u}, \quad \frac{\langle \eta_0 \rangle}{F_0} \simeq 0, \\ \frac{\langle a \rangle}{f_a} + \bar{\theta} &\simeq \frac{G_F}{\sqrt{2}} (\mathcal{C}_{1ud} - \mathcal{C}_{1du}) \frac{2c_3}{B_0 F_\pi^2} (m_d - m_u) \frac{1}{m_u m_d}. \end{aligned} \quad (4.43)$$

The VEVs of π^0 and η_8 remain of the same order as the case without Peccei-Quinn mechanism, and hence they contribute to the neutron EDM in an analogous way. The axion VEV no longer cancels the genuine $\bar{\theta}$ term and the leftover induces an effective $\bar{\theta}$ term; we estimate its contribution by employing the result of ref. [62] as

$$d_n|_{\text{ind } \bar{\theta}} = -(2.7 \pm 1.2) \times 10^{-16} \left(\frac{\langle a \rangle}{f_a} + \bar{\theta} \right) e \text{ cm}. \quad (4.44)$$

The final result is the sum of the meson VEV contribution estimated analogously to eq. (4.40), plus eq. (4.44).

4.2.2 Other CP-violating operators $\mathcal{O}_{1q}$, $\mathcal{O}_{2q}$ and $\mathcal{O}_3$

The contributions of the dipole operators in eq. (C.1) and the Weinberg operator in eq. (C.2) to the neutron EDM can be obtained with the QCD sum rule. The former is calculated in ref. [63] while the latter is in ref. [64], resulting in the following relations:

$$d_n|_{\text{quark}} = 0.47d_d - 0.12d_u + e(0.18d_d^c - 0.18d_u^c - 0.008d_s^c), \quad (4.45)$$

$$d_n|_{\text{quark}}^{\text{PQ}} = 0.47d_d - 0.12d_u + e(0.35d_d^c + 0.17d_u^c), \quad (4.46)$$

$$d_n|_{\text{Weinberg}} = \frac{G_F}{\sqrt{2}} e g_s \mathcal{C}_3 \times (10 - 30) \text{ MeV}, \quad (4.47)$$

where r.h.s. must be evaluated at 1 GeV. In eqs. (4.45), (4.46), d_q and d_q^c ($q = u, d, s$), so-called quark EDM and quark chromo-EDM, are defined as,

$$d_q(\mu) = -\frac{G_F}{\sqrt{2}} e e_q \mathcal{C}_{1q}(\mu) m_q(\mu), \quad (4.48)$$

$$d_q^c(\mu) = -\frac{G_F}{\sqrt{2}} \mathcal{C}_{2q}(\mu) m_q(\mu). \quad (4.49)$$

Equations (4.45) and (4.46) represent the quark EDM contributions without and with Peccei-Quinn mechanism, respectively. For the case without Peccei-Quinn mechanism, we have taken $\bar{\theta} = 0$.

4.3 Numerical analysis of neutron EDM versus ϵ'/ϵ

For numerical analysis of d_n , we employ the following values: the chiral-limit pion decay constant F_π is obtained from a lattice calculation as $F_\pi = 86.8 \text{ MeV}$ [65]. D, F have

been measured to be $D = 0.804$ and $F = 0.463$. For b_D, b_F , we quote the result of refs. [66, 67] with a NLO calculation in Lorentz covariant baryon chiral perturbation theory with decuplet contributions, which reads $b_D = 0.161 \text{ GeV}^{-1}$ and $b_F = -0.502 \text{ GeV}^{-1}$. Since the same calculation formalism, combined with experimental data $\sigma_{\pi N} \simeq 59(7) \text{ MeV}$, predicts a small value of the strange quark contribution to the nucleon mass σ_s [66], we infer that these values of b_D, b_F are most robust. For the quark masses, we adopt lattice results in ref. [68], $m_{ud}(2 \text{ GeV}) = 3.373 \text{ MeV}$ and $m_s(2 \text{ GeV}) = 92.0 \text{ MeV}$, and further evaluate QCD five-loop RG evolutions to obtain the masses at 1 GeV in $\overline{MS}$ scheme, which are used in our analysis. Also, we exploit an estimate $m_u/m_d = 0.46$ [68].

The main source of uncertainty in our analysis is the unknown LEC c_3 in the meson chiral Lagrangian eq. (4.13). The other unknown LEC c_1 is ineffective, because the Wilson coefficients satisfy $|\mathcal{C}_{1ud} - \mathcal{C}_{1u}| \gg |\mathcal{C}_{1ud} + \mathcal{C}_{1u}|, |\mathcal{C}_{1sq}|, |\mathcal{C}_{qs}|$. Our calculations of loop-induced d_n eq. (4.40) and axion-induced d_n eq. (4.40) are hence proportional to c_3 and subject to $O(1)$ uncertainty originating from its naïve dimensional analysis eq. (4.16). The fact that our results depend only on one LEC c_3 is good news, because it excludes the possibility of accidental cancellation between contributions with different LECs. Another source of uncertainty is the renormalization scale μ in the loop calculation eq. (4.37), but this is subdominant compared to the uncertainty of c_3 .

In the analysis, the ratio of the bifundamental scalar VEVs is again fixed as $\tan \beta = m_b/m_t$. The values of the new CP phases $\phi_c, \phi_t, \psi_d, \psi_s, \psi_b, \alpha$ are randomly generated. We find that the contribution of the Weinberg operator is suppressed by roughly $10^{-7} - 10^{-9}$ compared with that of the four-quark operators, and thus we neglect it in the analysis.

First, we show the numerical result for the neutron EDM without the constraint from ϵ'/ϵ in figure 3. One observes that the contribution of the four-quark operators is dominant over that of the quark EDMs.

As stated previously, an effective $\bar{\theta}$ term is induced in the presence of Peccei-Quinn mechanism. In figure 4, we additionally show the numerical prediction based on eq. (4.44). One finds that the induced $\bar{\theta}$ gives subleading contribution to the neutron EDM.

Next, the correlated prediction for $|d_n|$ and $\text{Re}(\epsilon'/\epsilon)$ is presented in figures 5 and 6 in the cases without and with Peccei-Quinn mechanism, respectively. Here, small contributions from the quark EDMs are neglected. The cases with and without Peccei-Quinn mechanism yield almost identical results because the induced $\bar{\theta}$ has a subdominant effect, as seen in figure 4. We observe that $M_{W'} = 20 \text{ TeV}$ and 50 TeV can be consistent with the data on $\text{Re}(\epsilon'/\epsilon)$ at 1σ level, whereas the case with $M_{W'} = 70 \text{ TeV}$ cannot explain it. However, the case with $M_{W'} = 20 \text{ TeV}$ has already been excluded by the current bound on the neutron EDM, and only $M_{W'} = 50 \text{ TeV}$ can be compatible with the neutron EDM bound and the data on $\text{Re}(\epsilon'/\epsilon)$. Figures 5 and 6 further inform us that almost all parameter points that account for the $\text{Re}(\epsilon'/\epsilon)$ data will be covered by future neutron EDM searches [71]. Therefore, unless the tree-level $\bar{\theta}$ in the case without Peccei-Quinn mechanism miraculously cancels the contribution of the model, we anticipate the discovery of the neutron EDM in the near future.

Figure 3. Prediction for the neutron EDM in the case with $\bar{\theta} = 0$ without Peccei-Quinn mechanism. Only the contributions of four-quark operators and quark-level EDMs including both quark EDM and chormo-EDM are shown. A dashed line represents the current bound on the neutron EDM [69], while a dashed dotted line stands for the future bound [71].

Figure 4. Comparison between the contribution of the induced $\bar{\theta}$ to the neutron EDM and others in the presence of Peccei-Quinn mechanism. A gray dashed line represents the current bound of EDM [69] while a black dashed dotted line stands for the future bound [71].

5 Summary and discussions

We have addressed the ϵ'/ϵ anomaly in the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge extension of the SM with charge symmetry. Since the charge symmetry gives strong restrictions on the mixing matrix for right-handed quarks, ϵ'/ϵ can be evaluated only in terms of two new CP phases $\alpha - \psi_d$ and $\alpha - \psi_s$, the mass of W' gauge boson (mostly composed of W_R), and the bifundamental scalar VEV ratio $\tan \beta$. By fixing $\tan \beta$ at its natural value m_b/m_t , and by randomly varying $\alpha - \psi_d$ and $\alpha - \psi_s$, we have shown that $M_{W'} < 58$ TeV must be satisfied to account for the experimental value of ϵ'/ϵ at 1σ level.

Figure 5. Correlation plot for the direct CP violation in $K \rightarrow \pi\pi$ decay and the neutron EDM in the case with $\bar{\theta} = 0$ without Peccei-Quinn mechanism. A gray dashed line and a black dashed dotted line represent the current [69] and the future [71] bounds on the neutron EDM, while a cyan band stands for the 1σ range of the direct CP violation in $K \rightarrow \pi\pi$ decay obtained from PDG [8].

Figure 6. The same figure as figure 5 with Peccei-Quinn mechanism.

Next, we have made a prediction for the neutron EDM d_n when the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ model with charge symmetry solves the ϵ'/ϵ anomaly. We have investigated the contribution of meson condensates induced by four-quark operators, and revealed that the π^0 VEV dominantly contributes to the neutron EDM, whose impact is enhanced by $m_s/(m_u + m_d)$ compared to other contributions. This enhancement is attributable to the isospin violating coupling of W' gauge boson, which allows the π^0 VEV to arise without

the factor of $m_d - m_u$. Additionally, we have found that the induced $\bar{\theta}$ term in the presence of Peccei-Quinn mechanism yields only a subleading effect on d_n . On the basis of the above observations, we have shown that the ϵ'/ϵ anomaly can be explained without conflicting the current experimental bound on d_n , and that the parameter space where the ϵ'/ϵ data are accounted for will be almost entirely covered by future experiments [71].

We comment on the constraint from $\text{Re}(\epsilon)$ on the model. Since W' gauge boson contributes to $\Delta F = 2$ processes only at loop levels, for $M_{W'} > 20 \text{ TeV}$, its contribution to $\text{Re}(\epsilon)$ is safely below the experimental bound [73]. However, the heavy neutral scalar particles coming from the bifundamental scalar induce $\Delta F = 2$ processes at tree level. Since their mass is of the same order as or below $M_{W'}$ if there is no fine-tuning in the scalar potential, these particles may lead to a tension with the data on $\text{Re}(\epsilon)$ [73] (constraint from $\text{Re}(\epsilon)$ on general left-right models is found in ref. [74], and that on the model with left-right parity is in ref. [75]) (for early studies on the $\text{Re}(\epsilon)$ constraint, see, e.g., ref. [76]).

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A Operators of $\Delta S = 1$ processes

$$O_1 = (\bar{s}_\alpha u_\beta)_L (\bar{u}_\beta d_\alpha)_L, \quad O_2 = (\bar{s}u)_L (\bar{u}d)_L, \quad (\text{A.1})$$

$$O_{1c} = (\bar{s}_\alpha c_\beta)_L (\bar{c}_\beta d_\alpha)_L, \quad O_{2c} = (\bar{s}c)_L (\bar{c}d)_L, \quad (\text{A.2})$$

$$O_1^{RL} = (\bar{s}_\alpha u_\beta)_R (\bar{u}_\beta d_\alpha)_L, \quad O_2^{RL} = (\bar{s}u)_R (\bar{u}d)_L, \quad (\text{A.3})$$

$$O_{1c}^{RL} = (\bar{s}_\alpha c_\beta)_R (\bar{c}_\beta d_\alpha)_L, \quad O_{2c}^{RL} = (\bar{s}c)_R (\bar{c}d)_L, \quad (\text{A.4})$$

$$O_3 = \sum_{q=u,d,s} (\bar{s}d)_L (\bar{q}q)_L, \quad O_4 = \sum_{q=u,d,s} (\bar{s}_\alpha d_\beta)_L (\bar{q}_\beta q_\alpha)_L, \quad (\text{A.5})$$

$$O_5 = \sum_{q=u,d,s} (\bar{s}d)_L (\bar{q}q)_R, \quad O_6 = \sum_{q=u,d,s} (\bar{s}_\alpha d_\beta)_L (\bar{q}_\beta q_\alpha)_R, \quad (\text{A.6})$$

$$O_7 = \frac{3}{2} \sum_{q=u,d,s} (\bar{s}d)_L e_q (\bar{q}q)_R, \quad O_8 = \frac{3}{2} \sum_{q=u,d,s} (\bar{s}_\alpha d_\beta)_L e_q (\bar{q}_\beta q_\alpha)_R, \quad (\text{A.7})$$

$$O_9 = \frac{3}{2} \sum_{q=u,d,s} (\bar{s}d)_L e_q (\bar{q}q)_L, \quad O_{10} = \frac{3}{2} \sum_{q=u,d,s} (\bar{s}_\alpha d_\beta)_L e_q (\bar{q}_\beta q_\alpha)_L, \quad (\text{A.8})$$

$$O_g = \frac{g_s}{8\pi^2} m_s \bar{s} \sigma_{\mu\nu} G^{a\mu\nu} T^a P_L d, \quad O_\gamma = \frac{e}{8\pi^2} m_s \bar{s} \sigma_{\mu\nu} F^{\mu\nu} P_L d, \quad (\text{A.9})$$

where $(\bar{q}q')_L \equiv \bar{q}\gamma_\mu(1 - \gamma_5)q'$ and $(\bar{q}q')_R \equiv \bar{q}\gamma_\mu(1 + \gamma_5)q'$, α, β are color indices, and color summation is taken in each quark bilinear unless α, β are displayed. $e_u = 2/3$ and $e_d = e_s = -1/3$. The operators O'_i, O_j^{LR} are obtained by interchanging $L \leftrightarrow R$ in the corresponding operators.

B Loop functions

The loop functions in the main text are defined as follows:

$$\begin{aligned}
 F_1(x) &= \frac{x(-18 + 11x + x^2)}{12(1-x)^3} + \frac{x^2(-15 + 16x - 4x^2)}{6(1-x)^4} \log x + \frac{2}{3} \log x + \frac{2}{3}, \\
 F_2(x) &= \frac{x(2 + 5x - x^2)}{4(1-x)^3} + \frac{3x^2 \log x}{2(1-x)^4}, \\
 F_3(x) &= \frac{4 + x + x^2}{2(1-x)^2} + \frac{3x \log x}{(1-x)^3}, \\
 E_{1d}(x) &= \frac{25x^2 - 19x^3}{36(1-x)^3} + \frac{x^2(6 + 2x - 5x^2)}{18(1-x)^4} \log x + \frac{4}{9} \log x + \frac{4}{9}, \\
 E_{2d}(x) &= \frac{x(7 - 5x - 8x^2)}{12(1-x)^3} + \frac{x^2(2 - 3x)}{2(1-x)^4} \log x, \\
 E_{3d}(x) &= \frac{20 - 31x + 5x^2}{6(1-x)^2} + \frac{x(2 - 3x)}{(1-x)^3} \log x, \\
 E_{3u}(x) &= \frac{8 - 16x + 2x^2}{3(1-x)^2} + \frac{x(1 - 3x)}{(1-x)^3} \log x.
 \end{aligned}$$

C CP-violating operators that contribute to the neutron EDM

$$\mathcal{O}_{1q} = -\frac{e}{2} e_q m_q \bar{q} \sigma_{\mu\nu} i \gamma_5 q F^{\mu\nu}, \quad \mathcal{O}_{2q} = -\frac{g_s}{2} m_q \bar{q} \sigma_{\mu\nu} i \gamma_5 T^a q G^{a\mu\nu}, \quad (\text{C.1})$$

$$\mathcal{O}_3 = -\frac{1}{6} f^{abc} \epsilon^{\mu\nu\rho\sigma} \eta^{\tau\nu} G_{\tau\mu}^a G_{\nu\rho}^b G_{\sigma\nu}^a, \quad (\text{C.2})$$

$$\mathcal{O}_{4q} = \bar{q} q \bar{q} i \gamma_5 q, \quad \mathcal{O}_{5q} = \bar{q} \sigma_{\mu\nu} q \bar{q} \sigma^{\mu\nu} i \gamma_5 q, \quad (\text{C.3})$$

$$\mathcal{O}_{1q'q} = \bar{q}' q' \bar{q} i \gamma_5 q, \quad \mathcal{O}_{2q'q} = \bar{q}'_\alpha q'_\beta \bar{q}_\beta i \gamma_5 q_\alpha, \quad (\text{C.4})$$

$$\mathcal{O}_{3q'q} = \bar{q}' \sigma^{\mu\nu} q' \bar{q} \sigma_{\mu\nu} i \gamma_5 q, \quad \mathcal{O}_{4q'q} = \bar{q}'_\alpha \sigma^{\mu\nu} q'_\beta \bar{q}_\beta \sigma_{\mu\nu} i \gamma_5 q_\alpha, \quad (\text{C.5})$$

where $q', q = u, d, s$ and $q' \neq q$. α, β are color indices, and color summation is taken in each quark bilinear unless α, β are displayed.

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