

Leptoquark-vectorlike quark model for the CDF m_W , $(g-2)_\mu$, $R_{K^{(*)}}$ anomalies, and neutrino masses

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Very recently, a substantial 7σ deviation of the W -boson mass from the Standard Model prediction has been reported by the CDF Collaboration. Furthermore, the Muon $g-2$ Experiment recently confirmed the long-standing tension in $(g-2)_\mu$. Besides, the updated result from the LHCb Collaboration found evidence for the breaking of lepton universality in beauty-quark decays, which shows a 3.1σ discrepancy and is consistent with their previous measurements. Motivated by several of these drawbacks of the Standard Model, in this work, we propose a model consisting of two scalar leptoquarks and a vectorlike quark to simultaneously address the W -boson mass shift, the $(g-2)_\mu$, and anomalies in the neutral current transitions of the B -meson decays. The proposed model also sheds light on the origin of neutrino mass and can be fully tested at the future colliders.

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I. INTRODUCTION

The W -boson mass, M_W , is a precisely measured quantity, and even a slight deviation from the predicted value would hint toward physics beyond the Standard Model (SM). The SM predicts $M_W^{\text{SM}} = (80.357 \pm 0.004)$ GeV, which agrees with the most up-to-date PDG value $M_W^{\text{PDG}} = (80.379 \pm 0.012)$ GeV at the 2σ confidence level [1]. Very recently, the CDF Collaboration reported a new precision measurement of M_W using their full 8.8 fb^{-1} dataset that yields [2]

$$M_W^{\text{CDF-2022}} = (80.4335 \pm 0.0094) \text{ GeV}, \quad (1)$$

which deviates from the SM prediction by 7σ , clearly indicating the presence of new physics (NP) [3–71].

The muon's anomalous magnetic moment (AMM) $\Delta a_\mu = (g-2)_\mu/2$, is another quantity that has recently been measured with unprecedented accuracy in the Muon $g-2$ Experiment [72]. The result from this experiment is in complete agreement with the previously measured value at BNL [73]. When these two results are combined, it shows a large 4.2σ discrepancy compared to the SM prediction [74] (for original works, see Refs. [75–94]):

$$\Delta a_\mu = (2.51 \pm 0.59) \times 10^{-9}, \quad (2)$$

hinting towards physics beyond the SM (BSM); for a recent review, see, e.g., [95].

In addition, lepton flavor universality (LFU) violating B -meson decays has been persistently observed in a series of experiments [96–100]. The most noteworthy deviation is observed in neutral-current transitions associated with the $R_K - R_{K^*}$ ratios, which are defined as

$$R_K = \frac{\text{Br}(B \rightarrow K\mu^+\mu^-)}{\text{Br}(B \rightarrow Ke^+e^-)}, \quad R_{K^*} = \frac{\text{Br}(B \rightarrow K^*\mu^+\mu^-)}{\text{Br}(B \rightarrow K^*e^+e^-)}. \quad (3)$$

LFU in the SM predicts these ratios to be unity with uncertainties less than 1%. However, the most precise measurement by LHCb [100] finds a deficit with a significance of 3.1σ for R_K ratio. There are several other related observables for which LHCb also found deficits with respect to the SM prediction, which are of order $\mathcal{O}(1.5 - 3.5)\sigma$; for a comprehensive list, see, e.g., [101]. However, if only the theoretically clean observables: R_K, R_{K^*} ratios and $\text{BR}(B_s \rightarrow \mu^+\mu^-)$ are taken into account, the data are found to be in 4.2σ tension with the SM [102]; for a recent review, see [103]. On the other hand, when both theoretically clean and dirty observables are considered, global analyses show preferences compared to the SM hypothesis with pulls more than 7σ (see [102,104–110] for theoretical assumptions and data included in these fits).

On top of these downsides mentioned above, neutrinos remain massless in the SM. On the contrary, several experiments discovered nonzero masses of the neutrinos

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via observations of neutrino oscillations [111–117]. This work proposes a simultaneous explanation of the W -boson mass shift, the tension in the $(g-2)_\mu$, and the anomalies in the neutral current transitions in the B -meson decays, as well as neutrino oscillation data. The proposed model employs two scalar leptoquarks (LQs [118,119]): $\tilde{R}_2 \sim (3, 2, 1/6)$ and $S_3 \sim (\bar{3}, 3, 1/3)$ and a vectorlike quark (VLQ) $\psi \sim (3, 2, -5/6)$. Nonzero mixing between \tilde{R}_2 and S_3 LQs leads to loop corrections to W -boson self-energy explaining the CDF anomaly. Utilizing this same mixing, the $(g-2)_\mu$ receives a large NP contribution via the mass flip of the VLQ inside the loop. The S_3 LQ, with its interactions with the SM fermions, addresses the discrepancies in the rare decays of B mesons based on the neutral current $b \rightarrow s\ell\ell$ transitions. Furthermore, nonzero mixing between \tilde{R}_2 and S_3 LQs is also responsible for generating neutrino mass at the one-loop order, and the model put forward in this work can be tested in the ongoing and future experiments.

II. PROPOSAL

In this work, we propose a new leptoquark-vectorlike quark model that contains three BSM particles: (i) an iso-doublet LQ, $\tilde{R}_2(3, 2, 1/6)$, (ii) an iso-triplet LQ, $S_3(\bar{3}, 3, 1/3)$, and (iii) an iso-doublet vectorlike quark, $\psi(3, 2, -5/6)$. Here, the quantum numbers are shown under the SM gauge group $SU(3) \times SU(2) \times U(1)$. Furthermore, we assign a baryon number of $1/3$ ($-1/3$) to ψ , \tilde{R}_2 (S_3). The corresponding component fields of these particles are defined in the following way:

$$\tau.S_3 = \begin{pmatrix} S^{1/3} & \sqrt{2}S^{4/3} \\ \sqrt{2}S^{-2/3} & -S^{1/3} \end{pmatrix}, \quad (4)$$

$$\tilde{R}_2 = \begin{pmatrix} R^{2/3} \\ R^{-1/3} \end{pmatrix}, \quad \psi = \begin{pmatrix} \psi^{-1/3} \\ \psi^{-4/3} \end{pmatrix}. \quad (5)$$

Yukawa sector: The relevant part of the Yukawa Lagrangian is given by

$$\begin{aligned} \mathcal{L}_Y \supset & y_i^S \bar{Q}_L^c \epsilon(\tau.S_3) L_L + \hat{y}_L^\psi \bar{\psi}_L \tilde{R}_2 \ell_R + y^R \bar{d}_R L_L \epsilon \tilde{R}_2 \\ & + \hat{y}_R^\psi \bar{L}_L(\tau.S_3) \psi_R + m_\psi \bar{\psi}_L \psi_R + \text{H.c.} \end{aligned} \quad (6)$$

One more term is allowed by the gauge symmetries: $y' \bar{\psi}_L d_R \epsilon H^*$. Once the EW symmetry is broken, it generates a mixing between $\psi^{-1/3}$ and $d^{-1/3}$ via $\bar{\psi}_L d_R v / \sqrt{2}$, for which we assume the Yukawa coupling to be negligibly small. Therefore, the mass generation of the SM fermions remains unaltered. Note that, baryon number assignments as described above forbid two terms: $\mathcal{L}_Y \not\supset \bar{\psi}_R^c u_R \tilde{R}_2$ and $\mathcal{L}_Y \not\supset \bar{Q}_L^c \epsilon(\tau.S_3) Q_L$.

First we focus on a muon-philic scenario, then the Yukawa couplings in Eq. (6) take the following form:

$$y^S = \begin{pmatrix} 0 & 0 & 0 \\ 0 & y_{s\mu}^S & 0 \\ 0 & y_{b\mu}^S & 0 \end{pmatrix}, \quad \hat{y}_L^\psi = \begin{pmatrix} 0 \\ y_L^\psi \\ 0 \end{pmatrix}, \quad \hat{y}_R^\psi = \begin{pmatrix} 0 \\ y_R^\psi \\ 0 \end{pmatrix}. \quad (7)$$

These forms of $\hat{y}_{L,R}^\psi$ are required to avoid excessive cLFV and provide large NP contribution to $(g-2)_\mu$. The texture of y^S is chosen to explain the anomalies in the neutral current transitions, to be discussed later in the text. For the simplicity of our work, we take all model parameters to be real.

Scalar sector: The relevant terms in the scalar potential take the following form:

$$V \supset m_R^2 \tilde{R}_2^\dagger \tilde{R}_2 + m_S^2 S_3^\dagger S_3 + \{\mu H^\dagger(\tau.S_3)\tilde{R}_2 + \text{H.c.}\}. \quad (8)$$

The cubic coupling μ leads to mixing between \tilde{R}_2 and S_3 components, which is crucial in addressing both the W -boson mass and $(g-2)_\mu$ anomalies within the proposed model. Remarkably, the existence of this cubic term also allows neutrinos to have nonzero masses at the one-loop order. In this theory, μ is one of the most important parameters, and in the limit $\mu \rightarrow 0$, one gets $\Delta m_W \rightarrow 0$, $\Delta a_\mu \rightarrow 0$, as well as $m_\nu \rightarrow 0$. Note that, to reduce the number of parameters and for the simplicity of our study, scalar quartic couplings are assumed to be somewhat smaller and are not included in Eq. (8). Consequently, all mass splittings are only a function of the trilinear coupling μ .

From the above potential, the mass matrices in the $\{S^Q, R^Q\}$ basis are given by

$$\begin{aligned} M_{2/3}^2 &= \begin{pmatrix} m_S^2 & \mu v \\ \mu v & m_R^2 \end{pmatrix}, \\ M_{1/3}^2 &= \begin{pmatrix} m_S^2 & -\mu v / \sqrt{2} \\ -\mu v / \sqrt{2} & m_R^2 \end{pmatrix}, \end{aligned} \quad (9)$$

with $v = 246$ GeV. We denote the weak and mass eigenstates with X and \hat{X} , respectively, which are related by

$$S^{\pm Q} = c_x \hat{S}^{\pm Q} - s_x \hat{R}^{\pm Q}, \quad (10)$$

$$R^{\pm Q} = s_x \hat{S}^{\pm Q} + c_x \hat{R}^{\pm Q}, \quad (11)$$

with $x = \theta, \phi$ for $Q = 2/3, 1/3$. Masses and mixing for \hat{X}^Q states then take the form

$$m_{\hat{S}_Q, \hat{R}_Q}^2 = \frac{1}{2} \{m_S^2 + m_R^2 \pm [(m_S^2 - m_R^2)^2 + a_Q \mu^2 v^2]^{1/2}\}, \quad (12)$$

$$\sin 2x = \frac{b_Q \mu v}{m_{\hat{S}_Q}^2 - m_{\hat{R}_Q}^2}, \quad (13)$$

where, $a_Q(b_Q) = 4(2)$ and $2(-\sqrt{2})$ for $Q = 2/3$ and $1/3$, respectively. In the above analysis, we have adopted the convention of $m_S > m_R$.

W-boson mass shift: The effects of NP phenomena on the electroweak (EW) gauge sector are parametrized in terms of oblique parameters [120,121] S , T , and U . Then, the shift in the W -boson mass from the NP can be calculated as a function of these oblique parameters [122],

$$m_W^2 = m_{W,\text{SM}}^2 \left\{ 1 + \frac{\alpha_{em}(c_w^2 T - \frac{1}{2} S + \frac{c_w^2 - s_w^2}{4s_w^2} U)}{c_w^2 - s_w^2} \right\}. \quad (14)$$

When the new CDF data are taken into account in a global electroweak precision fit, the oblique parameters would deviate from their previous (PDG) SM predictions, which several studies have already analyzed [15,23,26], and updated the 2σ allowed ranges for S , T , and U

parameters in light of the CDF result. By incorporating these new sets of values of oblique parameters in our numerical analysis, we find that for our model with TeV scale LQs, the mass splitting of the mixed LQ states must be of order $\Delta m_{LQ} \sim \mathcal{O}(100)$ GeV to be compatible with the result reported by the CDF Collaboration. It is noteworthy to mention that while the CDF II data alone show a 7σ deviation; taking the world average that includes previous measurements (that are compatible with the SM) from the Tevatron and LEP experiments as well as LHC would reduce the tension somewhat (for quantitative analysis, see, for example, [15,23,26]).

In our model, NP contributions to these parameters originate from the mass splittings among the component fields as a result of mixing between the same charged states from \hat{R}_2 and S_3 . We obtain the following one-loop correction to the T parameter for our model,

$$\begin{aligned} \Delta T = \frac{N_c}{16\pi s_w^2 m_W^2} & \left\{ (s_\phi s_\theta - \sqrt{2} c_\phi c_\theta)^2 \hat{F}[m_{\hat{S}^{-2/3}}, m_{\hat{S}^{-1/3}}] + (s_\phi c_\theta + \sqrt{2} c_\phi s_\theta)^2 \hat{F}[m_{\hat{R}^{-2/3}}, m_{\hat{S}^{-1/3}}] \right. \\ & + (c_\phi s_\theta + \sqrt{2} s_\phi c_\theta)^2 \hat{F}[m_{\hat{S}^{-2/3}}, m_{\hat{R}^{-1/3}}] + (c_\phi c_\theta - \sqrt{2} s_\phi s_\theta)^2 \hat{F}[m_{\hat{R}^{-2/3}}, m_{\hat{R}^{-1/3}}] \\ & + (\sqrt{2} c_\phi)^2 \hat{F}[m_{\hat{S}^{1/3}}, m_{\hat{S}^{-4/3}}] + (-\sqrt{2} s_\phi)^2 \hat{F}[m_{\hat{R}^{1/3}}, m_{\hat{S}^{-4/3}}] \\ & \left. - \frac{c_\phi^2 s_\phi^2}{2} \hat{F}[m_{\hat{S}^{1/3}}, m_{\hat{R}^{-1/3}}] - \frac{c_\theta^2 s_\theta^2}{2} \hat{F}[m_{\hat{S}^{2/3}}, m_{\hat{R}^{-2/3}}] \right\}, \quad (15) \end{aligned}$$

with,

$$\hat{F}(m_1, m_2) = m_1^2 + m_2^2 - \frac{2m_1^2 m_2^2}{m_1^2 - m_2^2} \log\left(\frac{m_1^2}{m_2^2}\right). \quad (16)$$

In contrast, we find that the NP contribution to ΔS is small compared to ΔT , which in our model takes the following form:

$$\begin{aligned} \Delta S = \frac{N_c}{\pi m_Z^2} & \left\{ -\frac{1}{3} \mathcal{B}_{22}[m_Z^2, m_{\hat{S}^{4/3}}^2, m_{\hat{S}^{4/3}}^2] + \frac{1}{48} (3 + c_{2\theta})(1 + 3c_{2\theta}) \mathcal{B}_{22}[m_Z^2, m_{\hat{S}^{2/3}}^2, m_{\hat{S}^{2/3}}^2] \right. \\ & + \frac{1}{96} (9 - 20c_{2\theta} + 3c_{4\theta}) \mathcal{B}_{22}[m_Z^2, m_{\hat{R}^{2/3}}^2, m_{\hat{R}^{2/3}}^2] - \frac{1}{24} (1 + 3c_{2\phi}) s_\phi^2 \mathcal{B}_{22}[m_Z^2, m_{\hat{S}^{1/3}}^2, m_{\hat{S}^{1/3}}^2] \\ & + \frac{1}{24} (-1 + 3c_{2\phi}) c_\phi^2 \mathcal{B}_{22}[m_Z^2, m_{\hat{R}^{1/3}}^2, m_{\hat{R}^{1/3}}^2] + \frac{1}{8} s_{2\theta}^2 \mathcal{B}_{22}[m_Z^2, m_{\hat{S}^{2/3}}^2, m_{\hat{R}^{2/3}}^2] \\ & \left. + \frac{1}{8} s_{2\phi}^2 \mathcal{B}_{22}[m_Z^2, m_{\hat{S}^{1/3}}^2, m_{\hat{R}^{1/3}}^2] \right\}; \quad (17) \end{aligned}$$

here, the expression for the loop function $\mathcal{B}_{22}(q^2, m_1^2, m_2^2)$ is defined as [123] (with $x_k = m_k^2/q^2$):

$$\begin{aligned} \mathcal{B}_{22}(q^2, m_1^2, m_2^2) = \frac{q^2}{24} & \left\{ 2 \ln q^2 + \ln(x_1 x_2) + [(x_1 - x_2)^3 - 3(x_1^2 - x_2^2) + 3(x_1 - x_2)] \ln \frac{x_1}{x_2} \right. \\ & \left. - \left[2(x_1 - x_2)^2 - 8(x_1 + x_2) + \frac{10}{3} \right] - [(x_1 - x_2)^2 - 2(x_1 + x_2) + 1] h(x_1, x_2) - 6H(x_1, x_2) \right\}, \quad (18) \end{aligned}$$

with

$$H(x_1, x_2) = \frac{x_1 + x_2}{2} - \frac{x_1 x_2}{x_1 - x_2} \ln \frac{x_1}{x_2}, \quad (19)$$

and

$$h(x_1, x_2) = \begin{cases} -2\sqrt{\Delta} \left[\tan^{-1} \frac{x_1 - x_2 + 1}{\sqrt{\Delta}} - \tan^{-1} \frac{x_1 - x_2 - 1}{\sqrt{\Delta}} \right], \\ 0, \\ \sqrt{-\Delta} \ln \frac{x_1 + x_2 - 1 + \sqrt{-\Delta}}{x_1 + x_2 - 1 - \sqrt{-\Delta}}, \end{cases} \quad (20)$$

$$\Delta = 2(x_1 + x_2) - (x_1 - x_2)^2 - 1, \quad (21)$$

for $\Delta > 0$, $\Delta = 0$, and $\Delta < 0$, respectively. For previous works on leptoquarks effects in EW oblique parameters, see, for example, Refs. [124–128].

Muon AMM: In this theory, the AMM of the muon receives NP contributions as shown in Fig. 1, which can be expressed as [129,130],

$$\Delta a_\mu = -\frac{m_\mu N_c}{4\pi^2} \sum_k \left\{ \Gamma_k^L \Gamma_k^R \frac{m_\psi}{m_{\phi_k}^2} F_k \left(\frac{m_\psi^2}{m_{\phi_k}^2} \right) + [|\Gamma_k^L|^2 + |\Gamma_k^R|^2] \frac{m_\mu}{m_{\phi_k}^2} G_k \left(\frac{m_\psi^2}{m_{\phi_k}^2} \right) \right\}; \quad (22)$$

here the sum is taken over $k = \{\hat{S}^{2/3}, \hat{R}^{2/3}, \hat{S}^{1/3}, \hat{R}^{1/3}\}$ for which we define $\{\Gamma_k^R, \Gamma_k^L\} = \{y_L^\psi s_\theta, \sqrt{2}y_R^\psi c_\theta\}$, $\{y_L^\psi c_\theta, -\sqrt{2}y_R^\psi s_\theta\}$, $\{y_L^\psi s_\phi, -y_R^\psi c_\phi\}$, and $\{y_L^\psi c_\phi, y_R^\psi s_\phi\}$, respectively. While $\hat{X}^{2/3}$ is propagating in the loop $Q_\psi = -1/3$, and for $\hat{X}^{-1/3}$ it is $Q_\psi = -4/3$. And the functions $F(x)$, $G(x)$ are given by

$$F(x) = f(x) + Q_\psi g(x), \quad G(x) = \tilde{f}(x) + Q_\psi \tilde{g}(x), \quad (23)$$

$$f(x) = \frac{x^2 - 1 - 2x \log x}{4(x-1)^3}, \quad g(x) = \frac{x - 1 - \log x}{2(x-1)^2}, \quad (24)$$

$$\tilde{f}(x) = \frac{2x^3 + 3x^2 - 6x + 1 - 6x^2 \log x}{24(x-1)^4}, \quad \tilde{g}(x) = \frac{1}{2}f(x). \quad (25)$$

By considering only the dominant chirally enhanced terms, $(g-2)_\mu$ takes the following simpler form:

$$\Delta a_\mu = \frac{3m_\mu m_\psi y_L^\psi y_R^\psi}{8\pi^2} \left\{ s_{2\phi} \left[\frac{F\left(\frac{m_\psi^2}{m_{\hat{S}^{1/3}}^2}\right) - F\left(\frac{m_\psi^2}{m_{\hat{R}^{1/3}}^2}\right)}{m_{\hat{S}^{1/3}}^2} - \frac{F\left(\frac{m_\psi^2}{m_{\hat{S}^{2/3}}^2}\right) - F\left(\frac{m_\psi^2}{m_{\hat{R}^{2/3}}^2}\right)}{m_{\hat{S}^{2/3}}^2} \right] - \sqrt{2}s_{2\theta} \left[\frac{F\left(\frac{m_\psi^2}{m_{\hat{S}^{2/3}}^2}\right) - F\left(\frac{m_\psi^2}{m_{\hat{R}^{2/3}}^2}\right)}{m_{\hat{S}^{2/3}}^2} - \frac{F\left(\frac{m_\psi^2}{m_{\hat{S}^{1/3}}^2}\right) - F\left(\frac{m_\psi^2}{m_{\hat{R}^{1/3}}^2}\right)}{m_{\hat{S}^{1/3}}^2} \right] \right\}. \quad (26)$$

If the external photon leg is removed from the Feynman diagram Fig. 1, then the corresponding diagram contributes to

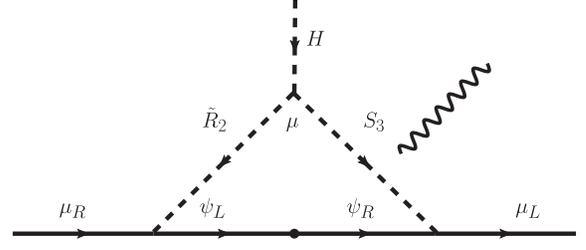


FIG. 1. Leading order NP contribution to muon AMM (in the weak basis). Photon can be attached to either the fermion line or the scalar line.

the muon mass. Such a chirally enhanced contribution can, in principle, generate a large mass correction δm_ℓ to the lepton [95,131,132]. Theoretically, this can be absorbed by adjusting the pole mass of the lepton, m_ℓ^0 , which, however, introduces a large degree of fine-tuning. To avoid a fine-tuned solution, in this work, we adopt the criterion $\delta m_\ell \lesssim m_\ell^0$ [133].

As aforementioned, for the AMM of the muon, we adopt the theoretical estimate Eq. (2) quoted in the 2020 white paper [74]. The SM prediction given in [74] is based on the estimate of the leading-order hadronic vacuum polarization (HVP) contribution (a_μ^{hvp}), evaluated from a data-driven approach (using a dispersion integral involving hadronic cross-section data). This method has an error of 0.6%, subject to experimental uncertainties associated with measured cross-section data. On the other hand, lattice QCD calculations for a_μ^{hvp} , generally face numerous technical challenges. Despite that, a recent lattice computation, namely, by the BMW Collaboration [134], quotes an error of only 0.8%. If this computation of a_μ^{hvp} is considered, then the combined measurement from E989 and E821, when compared to the SM prediction, reduces to 1.5σ from 4.2σ . Moreover, two more groups, CLS/Mainz group [135] and ETMC [136] just recently released their lattice computations which show consistency with the BMW result. However, the results of [135] are somewhat in tension with the previous lattice computations for light quarks by the RBC/UKQCD Collaboration [137] and ETMC [138]. Instead of considering the HVP from the data-driven method, if the result from [135] is adopted, the tension between the SM prediction for $(g-2)_\mu$ and experiment would be reduced to 2.9σ . Concerning the other new lattice result of [136], it agrees with the BMW (CLS/Mainz) group at the level of 1.0σ (1.3σ). However, if these new lattice results hold, they point towards a large $\sim 4.2\sigma$ discrepancy with the low-energy $e^+e^- \rightarrow$ hadrons cross-section data with respect to Standard Model predictions (see also [139–142]).

$R_K - R_{K^}$ anomalies*: The effective Hamiltonian responsible for processes of the form $B \rightarrow K^{(*)} \ell^+ \ell^-$ can be described by

$$\mathcal{H}_{\text{eff}}^{dd\ell\ell} = -\frac{4G_F}{\sqrt{2}} V_{tj} V_{ti}^* \left(\sum_{X=9,10} C_X^{ij,\ell\ell'} \mathcal{O}_X^{ij,\ell\ell'} \right) + \text{H.c.}, \quad (27)$$

where the effective operators are defined as

$$\mathcal{O}_9^{ij,\ell\ell'} = \frac{\alpha}{4\pi} (\bar{d}_i \gamma^\mu P_L d_j) (\bar{\ell} \gamma_\mu \ell'), \quad (28)$$

$$\mathcal{O}_{10}^{ij,\ell\ell'} = \frac{\alpha}{4\pi} (\bar{d}_i \gamma^\mu P_L d_j) (\bar{\ell} \gamma_\mu \gamma_5 \ell'). \quad (29)$$

Now, the part of the Lagrangian relevant for the $R_{K^{(*)}}$ observable is the first term given in Eq (6). We choose to work with the ‘‘down-type diagonal’’ flavor ansatz for which the CKM matrix enters in the interactions associated with the up-type quarks. Then S_3 couplings to SM fermions contain a term of the form

$$\mathcal{L}_{S_3} \supset (-\sqrt{2} y^S)_{ij} \bar{d}_L^c i \mu_{Lj} S^{4/3} + \text{H.c.}, \quad (30)$$

which leads to a $b \rightarrow s \mu^+ \mu^-$ transition as shown in Fig 2. After integrating out the heavy leptoquark and combining the Yukawa part of the Lagrangian associated to S_3 as given above, the relevant Wilson coefficients contained in Eq. (27) at the LQ mass scale generating such neutral current processes take the form

$$\Delta C_9^{\mu\mu} = -\Delta C_{10}^{\mu\mu} = \frac{v^2}{V_{tb} V_{ts}^*} \frac{\pi}{\alpha_{em}} \frac{y_{b\mu}^S (y_{s\mu}^S)^*}{M_{4/3}^2}. \quad (31)$$

A global fit to the data that includes all $b \rightarrow s \mu \mu$ observables, the $R_{K^{(*)}}$ ratios, and $B_s \rightarrow \mu^- \mu^+$ branching ratio prefers $\Delta C_9^{\mu\mu} = -\Delta C_{10}^{\mu\mu} = -0.39 \pm 0.07$ [104] (see also [102,143]).

Neutrino mass: Nonzero mixing between the \tilde{R}_2 and S_3 LQs and BSM Yukawa interactions of the SM fermions with these LQs give rise to neutrino oscillations [144,145] in this theory (for LQ effects in leptonic processes, see, e.g., [146,147]). Feynman diagram that leads to nonzero neutrino mass is shown in Fig. 3, and the neutrino mass formula takes the following form [144]:

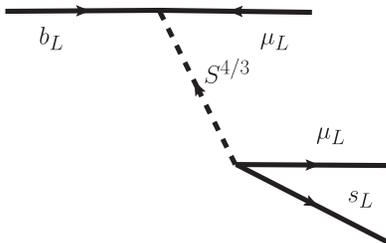


FIG. 2. Feynman diagram leading to $b \rightarrow s \mu^+ \mu^-$ transition.

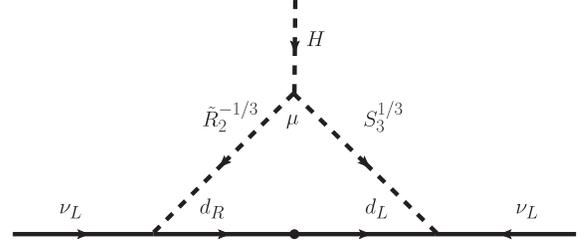


FIG. 3. Feynman diagram leading to nonzero neutrino masses (in the weak basis).

$$\begin{aligned} \mathcal{M}_{ij}^\nu &= \frac{\sin 2\phi}{32\pi^2} \sum_{k=d,s,b} m_k [(y^R)_{ki} (y^S)_{kj} + (y^R)_{kj} (y^S)_{ki}] \\ &\times \left[\frac{m_{\tilde{S}_{1/3}}^2 \ln \frac{m_k^2}{m_{\tilde{S}_{1/3}}^2}}{m_{\tilde{S}_{1/3}}^2 - m_k^2} - \frac{m_{\tilde{R}_{1/3}}^2 \ln \frac{m_k^2}{m_{\tilde{R}_{1/3}}^2}}{m_{\tilde{R}_{1/3}}^2 - m_k^2} \right]. \end{aligned} \quad (32)$$

Since down-type quark masses are much smaller than the LQ masses, the above formula can be further simplified,

$$\mathcal{M}^\nu \approx \frac{\sin 2\phi}{16\pi^2} \log \left(\frac{m_{\tilde{R}_{1/3}}}{m_{\tilde{S}_{1/3}}} \right) \{ (y^R)^T m_D y^S + (y^S)^T m_D y^R \}. \quad (33)$$

Collider constraints: At LHC, \tilde{R}_2 and S_3 LQs can be pair produced [148,149] via gluon-fusion $pp \rightarrow \text{LQLQ}^\dagger$. Once produced, each of these LQs would decay to SM fermions. Several searches for LQ pairs have been made at ATLAS and CMS for different final states with or without neutrinos. The strongest constraints for our scenario come from decay of these LQs to a third generation quark and a second generation charged lepton, namely, $b\mu$ and $t\mu$. For 100% branching ratio to the $pp \rightarrow b\bar{b}\mu^+\mu^-$ ($pp \rightarrow t\bar{t}\mu^+\mu^-$) channel, LHC provides a lower bound of $m_{\text{LQ}} \gtrsim 1.7(1.5)$ TeV [150,151]. For similar processes with third generation charged lepton, the corresponding bounds are $m_{\text{LQ}} \gtrsim 1$ TeV ($pp \rightarrow b\bar{b}\tau^+\tau^-$) and $m_{\text{LQ}} \gtrsim 1.4$ TeV ($pp \rightarrow t\bar{t}\tau^+\tau^-$), respectively [152,153].

The single production of LQ becomes only relevant for larger Yukawa couplings to the first- and second-generation quarks, which is not the case in our scenario. For a similar reasons, nonresonant dilepton searches at the LHC do not provide strong constraints for the parameter space we are interested in.

VLQs can also be pair produced at the LHC through gluon fusion. If LQs are lighter than VLQs, then each VLQ would mostly decay to a muon and a LQ leading to $pp \rightarrow t\bar{t}\mu^+\mu^-\ell^+\ell^-$, $b\bar{b}\mu^+\mu^-\ell^+\ell^-$. Processes of these types have been previously considered in Ref. [154]. We take the LHC bounds on VLQ from [155,156] that typically correspond to $m_\psi \gtrsim 1.3\text{--}1.4$ TeV.

On the other hand, if LQs are heavier, then pair produced VLQs would still decay to SM fermion final states as mentioned above, via effective 4-fermion operator of the

form $y^\psi y^{\text{LQ}}/M_{\text{LQ}}^2(\bar{\psi}\ell)(\ell q)$; for details see Ref. [154], where y^{LQ} and y^ψ represent generic Yukawa couplings of the LQs and VLQ. For TeV scale VLQ with order one couplings, the decay is prompt even for LQs masses as heavy as 100 TeV. However, if the corresponding LQ couplings with the SM fermions are very small, then the VLQ can become long-lived and form R hadrons [157]. For such a scenario, by comparing with the limits on production cross sections for gluinos and squarks, a lower bound of $m_\psi \gtrsim 1.5$ TeV on the VLQ mass is obtained in Ref. [158].

Future collider prospects: Here we point out that future experiments such as the multi-TeV muon collider (MuC) [159–162] and 100 TeV future circular hadron collider (FCC-hh) [163–165] will probe the entire parameter space of the theory relevant for the B -meson anomaly. The most efficient way to probe this scenario is via the predominant LQ interactions with the muons since the new physics cannot appear at an arbitrarily high energy scale ($2 \rightarrow 2$ fermion scattering amplitudes associated to $b \rightarrow s\ell\ell$ transitions saturate the unitarity bound below 80 TeV [166]).

From pair production, FCC-hh will rule out LQ masses almost up to 10 TeV [167]. On the other hand, Drell-Yan (DY) $pp \rightarrow \mu\bar{\mu}$ from nonresonant t -channel contribution, a large portion of the parameter space in the Yukawa-mass plane will be ruled out [167] leaving part of the parameter space unconstrained (corresponding to $y_{b\mu} \approx y_{s\mu}$ that minimizes the contribution $pp \rightarrow \mu\bar{\mu}$).

For muon colliders, the inverted Drell-Yan (IDY) channel $\mu\bar{\mu} \rightarrow jj$ will constrain the Yukawa-mass plane, which is similar to the FCC-hh scenario (with DY processes). Remarkably, MuC could directly observe an s -channel resonance in the $\mu q \rightarrow \mu j$ (due to the quark content inside the muon) for masses up to $\sim s_0^{1/2}$, which would be the most promising on-shell process at muon colliders. When the LQ pair production, IDY, and $\mu\mu \rightarrow \mu j$ processes are combined, MuC10 will probe the entire parameter space [167].

Results and discussion.—One of the most important parameters in this model is the scalar cubic coupling μ , which mixes the two LQs. As described above, for $\mu \rightarrow 0$, $(g-2)_\mu, \Delta m_W, m_\nu \rightarrow 0$. First, we explicitly demonstrate the required range of μ to correctly reproduce the electroweak oblique parameters $(\Delta T, \Delta S)$ consistent with the recent CDF II measurement. This is portrayed in Fig. 4 by randomly varying the LQ mass parameters in the ranges $m_R \in (1, 5)$ TeV, $m_S - m_R \in (1, 500)$ GeV. In making this plot, μ is restricted to vary in the range $\mu \lesssim 0.7$ TeV; this upper bound is chosen here for the sake of clarity and illustration. The 1σ and 2σ ranges that favor CDF II data in the $(\Delta T, \Delta S)$ plane are taken from Ref. [15] that performed a global fit to the electroweak data. This figure shows that $\mu \gtrsim 350$ GeV is essential to resolve the anomaly. Note that in our scenario with only the μ parameter responsible for splitting the masses of the LQ states, the CDF anomaly can only be addressed at the 2σ confidence limit (C.L.). It is because the ΔS parameter

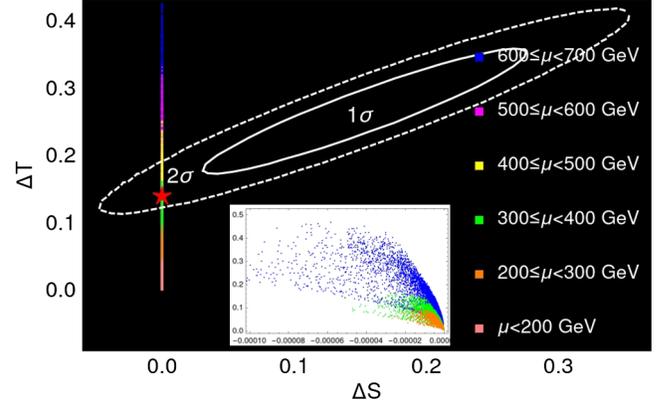


FIG. 4. Expected values of the μ parameter to incorporate CDF mass shift of the W boson; see text for details. The 1σ and 2σ regions are obtained from Ref. [15] that performs a global electroweak fit, including the CDF II data. For clarity, part of the parameter space is zoomed in for three different ranges of μ parameter. The red star represents a particular benchmark point discussed later in the text.

given in Eq. (17) turns out to be always tiny; however, addressing the CDF anomaly within 1σ C.L. demands $\Delta S \in [0.03, 0.27]$, whereas it is $\Delta S \in [-0.048, 0.35]$ within 2σ C.L. In Fig. 4 the red star represents a particular benchmark scenario discussed later in the text.

The nontrivial functional dependence of $(g-2)_\mu$ on μ is given in Eq. (26), which we graphically illustrate in Fig. 5. Here (for Figs. 5, 6, and 8), we randomly scan over the relevant parameters in the ranges: $m_{S,R} \in [1, 15]$ TeV, $\mu \in [0.1, 5]$ TeV, $m_\psi \in [1, 10]$ TeV, and $-y_L^\psi \cdot y_L^\psi \in [10^{-5}, 1]$. The points in green (yellow) correspond to solutions that simultaneously satisfy the $(g-2)_\mu$ anomaly at the 1σ (2σ) and the CDF anomaly at the 2σ C.L. The red points that are consistent with the CDF anomaly at the 2σ C.L., however, fail to reproduce the muon AMM within its 2σ values. As can be seen from Fig. 5, reproducing the correct value of

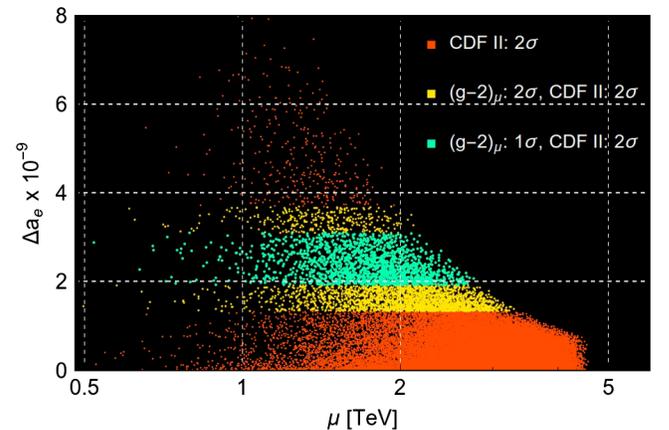


FIG. 5. Dependence of $(g-2)_\mu$ on the scalar cubic coupling μ . See text for details.

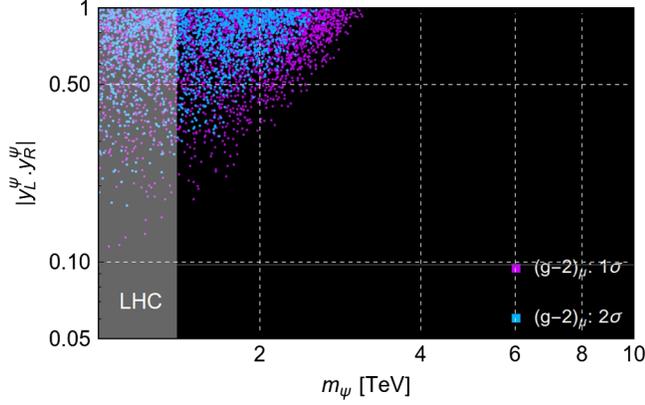


FIG. 6. Correlations between the product of the Yukawa couplings $|y_L^\psi \cdot y_R^\psi|$ and the mass m_ψ of the VLQ. See text for details.

Δa_μ requires μ in between $\mathcal{O}(0.5)$ to $\mathcal{O}(3)$ TeV. As shown in [168], even though the current LHC measurements [169,170] of $h \rightarrow \mu\mu$ allow a large trilinear coupling, future colliders such as the FCC may be able to measure this coupling and constrain the theory parameter space.

Moreover, in Fig. 6, we depict the correlations between the product of the Yukawa couplings $|y_L^\psi \cdot y_R^\psi|$ and the mass of the VLQ, m_ψ . This figure exhibits the required range of the Yukawa couplings as a function of m_ψ . Correctly replicating the measured value of the muon AMM given in Eq. (2) calls for a somewhat small mass $m_\psi \lesssim 3$ TeV for the VLQ. Due to these low-mass vectorlike quarks, our model can be distinctively searched for at the LHC. Since the LQs as well as the VLQ predominantly couple to the muons, see Eq. (7), pair-produced VLQs will lead to a unique signature $pp \rightarrow 4\mu + t\bar{t}$, either via on-shell or off-shell leptoquarks (see discussion above), as shown in Fig. 7. However, a dedicated collider study is beyond the scope of this work. Branching ratios to other up-type quark final states are expected to be suppressed by CKM elements, and diagrams with neutrinos in the final states are not shown.

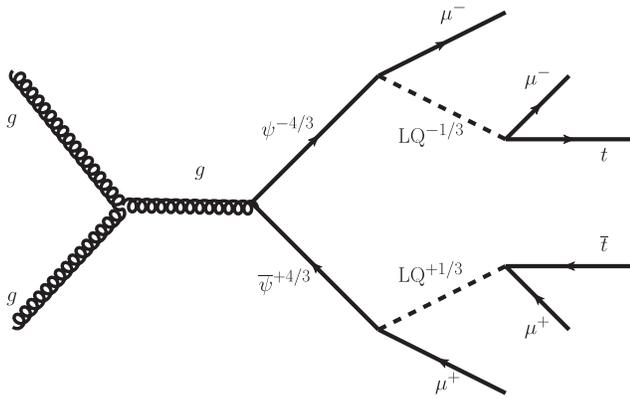


FIG. 7. Unique signature of the proposed model.

Even though both $(g-2)_\mu$ and W -boson mass shift heavily depend on the μ parameter, they are largely insensitive to the scale of the LQ masses. From a naive estimation, $\Delta a_\mu \sim 3m_\mu m_\psi \theta_{LQ}^2 y_{LQ}^2 / (16\pi^2 m_{LQ}^2) \ln[m_\psi^2 / m_{LQ}^2]$, where a sum over all LQ contributions must be taken with appropriate signs (the relative signs play important role). Then, for TeV scale VLQ, with order unity Yukawa couplings and maximal mixing, LQ masses as heavy as $m_{LQ} \sim \mathcal{O}(10)$ TeV still provide correct order Δa_μ , as depicted in Fig. 8.

Finally, we exemplify how this model can address the anomaly in the neutral current transitions of the B -meson decays as well as incorporate neutrino oscillation data. This, however, requires a careful fit to neutrino observables via minimization of a χ^2 function; a random scan over the parameters is not sufficient. This is why we demonstrate the viability of our model with a specific benchmark for which we choose $m_S = 2$, $m_R = 1.8$ TeV and mass of the VLQ is $m_\psi = 1.5$ and $\mu = 0.62$ TeV. This specific benchmark point along with $y_L^\psi y_R^\psi = -0.3$ corresponds to $\Delta a_\mu = 2.5 \times 10^{-9}$ [consistent with Eq. (2) at the 1σ C.L.] and $[\Delta T, \Delta S] = [0.14038, -8 \times 10^{-6}]$ (this point represents the red star shown in Fig. 4). Note, however, that this fitting procedure is highly nontrivial since the same Yukawa couplings addressing the $R_K - R_{K^*}$ anomalies also enter in neutrino observables. In fact, with the texture for y^S with only two nonzero elements as given in Eq. (7), neutrino masses and mixings utilizing the formula Eq. (33) cannot be accommodated. To generate viable neutrino masses and mixings, a few more entries must be introduced in y^S , which would lead to both charged lepton flavor violation as well as flavor violations in the quark sector [171–173] (see also [145,174–183]).

Our detailed numerical analysis shows that a nonzero 23-block in y^S is insufficient to satisfy neutrino oscillation

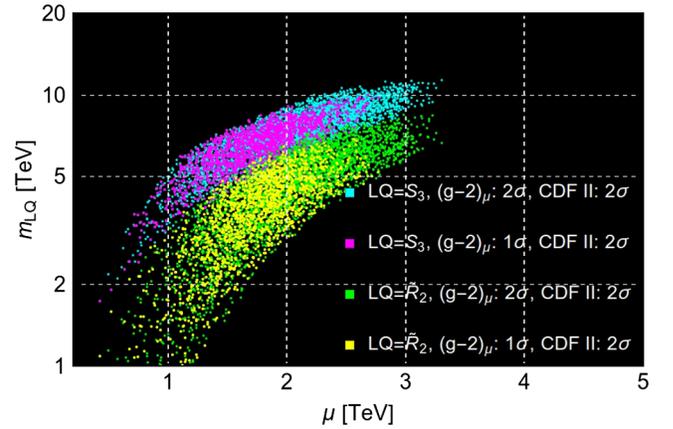


FIG. 8. Correlations between the μ parameter and the masses of the LQs; see text for details. For an example, here we have chosen $Q = 2/3$; the other case with $Q = -1/3$ shows indistinguishable behavior.

data. Therefore, we also introduce a nonzero 31 entry, which is constrained from $\mu - e$ transition (through CKM rotations). The most crucial cLFV for the given texture is $\tau \rightarrow \mu\gamma$, via which this model may be probed at the upcoming experiment [184]. In addition, various other flavor violating processes are considered in our numerical analysis that includes $Z \rightarrow \ell\ell$, $\ell \rightarrow \ell'\ell''\ell'''$, τ decays into mesons, and several meson decay observables as well as meson-antimeson oscillations; for details, see, e.g., [171,172]. The observables that provide the most stringent constraints on the model parameters are summarized in the Appendix. As aforementioned, the 23 block in y^S plays an important role in fitting neutrino observables; some of its entries address the $R_{K^{(*)}}$ anomalies and are required to be sizable. By randomly varying these couplings, interrelations between $BR[\tau \rightarrow \mu\gamma]$ and $B \rightarrow K^{(*)}\nu\bar{\nu}$ observables are depicted in Fig. 9 along with their respective experimental bounds. Clearly, the prospective measurement of the $B \rightarrow K^{(*)}\nu\bar{\nu}$ signal at the Belle II experiment shows a promising avenue to test our model.

In addition to the above-mentioned parameters, from a combined fit, we obtain the following Yukawa couplings addressing $R_K - R_{K^*}$ anomalies as well as neutrino oscillation data:

$$y^S = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0.01216 & -0.01495 \\ 0.07729 & 0.19480 & 0.07439 \end{pmatrix}, \quad (34)$$

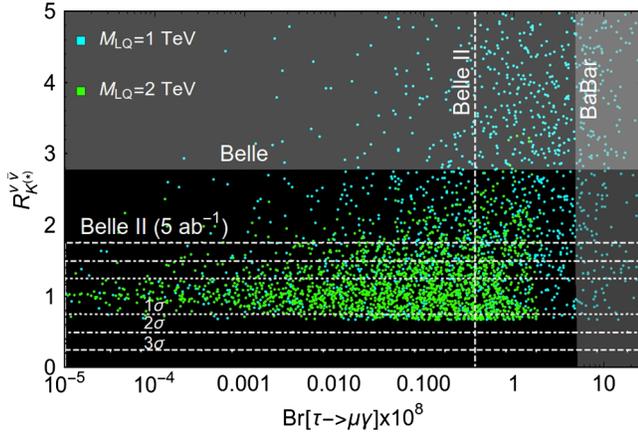


FIG. 9. The results of random scans showing the correlations between $BR[\tau \rightarrow \mu\gamma]$ and $R_{K^{(*)}}^{\nu\bar{\nu}}$. Current (future) bound on $BR[\tau \rightarrow \mu\gamma]$ from *BABAR* Collaboration [185] (Super B Factory [184]) is shown by shaded light-gray area (vertical dashed white line). For the observable $R_{K^{(*)}}^{\nu\bar{\nu}}$, current experimental limit from Belle collaboration [186] is presented by shaded light-gray area. The regions bounded by dotted, dot-dashed, and dashed white (horizontal) lines depict the projected reach (1 ± 0.25) for $R_{K^{(*)}}^{\nu\bar{\nu}}$ at Belle II [187] for 5 ab^{-1} of data with 1σ , 2σ , and 3σ C.L., respectively, assuming the best-fit value is SM-like.

$$y^R = 10^{-7} \begin{pmatrix} 1.2583 & 0.29906 & -0.72677 \\ 0.21836 & -0.60107 & -0.03121 \\ 0.12504 & -0.88446 & -1.49160 \end{pmatrix}. \quad (35)$$

This fit corresponds to the following neutrino observables:

$$(m_1, m_2, m_3) = (1.59 \times 10^{-2}, 8.60, 50.27) \text{ meV}, \quad (36)$$

$$(\sin^2 \theta_{12}, \sin^2 \theta_{23}, \sin^2 \theta_{13}) = (0.309, 0.574, 0.02224), \quad (37)$$

which are in excellent agreement with experimental data and satisfy all flavor constraints.

UV completion: Before concluding, we briefly discuss the possible ultraviolet (UV) complete model of the proposed scenario. For demonstration, we choose $SU(5)$ GUT, which is the minimal simple group containing the entire SM gauge group, i.e., $SU(5) \supset SU(3)_c \times SU(2)_L \times U(1)_Y$. As is well known, the minimal $SU(5)$ GUT, namely, the Georgi-Glashow model [188] is incompatible with the observed charged fermion masses. To overcome this drawback, Georgi and Jarlskog proposed [189] to include a 45_H . Interestingly, 45_H Higgs contains the S_3 LQ; $45 \supset (3, 3, -1/3)$. On the other hand, \tilde{R}_2 LQ can be embedded in a new Higgs in the 10_H dimensional representation; $10 \supset (3, 2, 1/6)$. Finally, the BSM fermion ψ can emerge from 24_F dimensional fermionic representation since $24 \supset (3, 2, -5/6) + (\bar{3}, 2, 5/6)$.

Within this framework, the Yukawa couplings giving rise to $(g-2)_\mu$ originate from two independent interactions, $10_F 24_F 10_H^*$ and $\bar{5}_F 24_F 45_H$. Recall that in $SU(5)$ GUT, the SM fermions are embedded in $\bar{5}_F + 10_F$ representations, and in the Georgi-Glashow model, the scalar 5_H contains the SM Higgs. Finally, the cubic coupling of Fig. 1 emanates from the interaction term of the form $10_H^* 5_H 45_H$. The Yukawa coupling responsible for addressing the B -meson anomaly in the neutral current transition arises from $10_F \bar{5}_F 45_H^*$ interaction. In addition to the last two interaction terms as aforementioned, neutrino mass generation utilizes another Yukawa coupling with the \tilde{R}_2 LQ that appears from $\bar{5}_F \bar{5}_F 10_H$ term in the Lagrangian.

One requires additional fine-tuning on top of the usual doublet-triplet splitting to keep the required multiplets (S_3, \tilde{R}_2, ψ) light compared to their partners, which we assume to reside at the GUT scale. Further tuning of parameters would be required to keep more states light to achieve successful gauge coupling unification, which we do not attempt to address in this work. For the scalars there are enough numbers of free parameters in the scalar potential to achieve this; for details, see, for example, Ref. [190]. For the BSM fermion sector, there exist only two free parameters ($\mathcal{L} \supset m_F 24_F^2 + \lambda 24_F^2 24_H$) that determine the mass of all the multiplets. To be precise, $m_{(1,1,0)} = m_F - X$, $m_{(1,3,0)} = m_F - 3X$, $m_{(8,1,0)} = m_F + 2X$, and $m_{(3,2,-5/6)} = m_F - X/2$, where $X = \lambda v_{\text{GUT}} / \sqrt{30}$ and

the vev of the Higgs in the adjoint representation is $\langle 24_H \rangle = \text{diag}(2, 2, 2, -3, -3)v_{\text{GUT}}/\sqrt{30}$ that breaks the GUT symmetry to the SM gauge group. It is straightforward to see that in the limit of $X \rightarrow 2m_F$ (this leads to an additional fine-tuning on top of the well-known doublet-triplet splitting in generic GUTs), only $\psi + \text{c.c.}$ remain light, whereas the rest of the components of 24_F live close to the GUT scale if the natural choice $m_F \sim v_{\text{GUT}}$ is assumed.

Finally, we comment on proton decay constraints. The \tilde{R}_2 LQ does not contain diquark couplings, hence does not lead to proton decay. However, S_3 LQ, in general, has both lepton-quark as well as quark-quark couplings leading to proton decay. The latter coupling originates from terms of the form $10_F 10_F 45_H$; as long as this coupling is forbidden (or highly suppressed) and S_3 LQ does not mix with any other fields that contain diquark couplings, S_3 LQ can be kept light without any conflict with the stringent experimental limits on proton decay. Both of these requirements can be satisfied within this framework, leading to additional fine-tuning. Furthermore, the texture zeros in the Yukawa couplings that we consider may be arranged by imposing flavor symmetries, which, however, is beyond the scope of this work.

III. CONCLUSION

In this work, we proposed a simple new physics scenario to simultaneously address several puzzles that cannot be accounted for by the Standard Model alone. The model is comprised of two scalar leptoquarks and a vectorlike quark. One of the most crucial parameters in this theory is the mixing parameter between the two leptoquarks. This mixing generates neutrino masses via quantum corrections at one loop, provides additional contributions to the W -boson mass consistent with recent CDF measurement, and plays a nontrivial role in incorporating the long-standing tension in the muon anomalous magnetic moment. The vectorlike quark, assisted with both the leptoquarks, gives rise to the required sizable new physics contributions to the $(g-2)_\mu$ via chirally enhanced terms proportional to its mass. Furthermore, the iso-triplet leptoquark is responsible for accounting for the deviations observed persistently in the $R_{K^{(*)}}$ ratios. By performing a numerical analysis, we have illustrated how to consistently resolve all these mysteries mentioned above by keeping flavor violations under control. Moreover, the model is within reach of the current and future upgrades of the LHC and has the potential to be fully probed by the future colliders such as the future circular hadron collider and multi-TeV muon collider.

APPENDIX: CONSTRAINTS ON MODEL PARAMETERS

In this appendix, we provide relevant expressions of all the NP contributions to various flavor violating processes. Since the Yukawa couplings of the \tilde{R}_2 LQ are small as

required for neutrino mass generation, here we focus on the constraints associated with the S_3 LQ. Since we have chosen flavor conserving Yukawa couplings of the vectorlike quark, ψ does not lead to flavor violation.

1. LFV: $\ell \rightarrow \ell'\gamma$

In our model, the dominant LFV process arise from $\ell \rightarrow \ell'\gamma$. The effective Lagrangian leading to such radiative decays of the charged leptons is given by [191]

$$\mathcal{L}_{\ell \rightarrow \ell'\gamma} = \frac{e}{2} \bar{\ell}' i \sigma^{\mu\nu} F_{\mu\nu} (\sigma_L^{\ell\ell'} P_L + \sigma_R^{\ell\ell'} P_R) \ell. \quad (\text{A1})$$

Then the branching ratios associated to these process are calculated by the following formula [191]:

$$\text{Br}(\ell \rightarrow \ell'\gamma) = \frac{\tau_\ell \alpha m_\ell^3}{4} (|\sigma_L^{\ell\ell'}|^2 + |\sigma_R^{\ell\ell'}|^2), \quad (\text{A2})$$

where τ_ℓ is the lifetime of the initial state lepton. The expressions of $\sigma_{L,R}$ are as follows [119,191]:

$$\sigma_{L,S_3}^{if} = \frac{iN_c}{16\pi^2 M_3^2} m_i \left\{ (Vy^S)_{qf}^* (Vy^S)_{qi} \frac{-1}{12} + (y^S)_{qf}^* y_{qi}^S \frac{1}{3} \right\}, \quad (\text{A3})$$

$$\sigma_{R,S_3}^{if} = \frac{iN_c}{16\pi^2 M_3^2} m_f \left\{ (Vy^S)_{qf}^* (Vy^S)_{qi} \frac{-1}{12} + (y^S)_{qf}^* y_{qi}^S \frac{1}{3} \right\}. \quad (\text{A4})$$

Here V is the Cabibbo-Kobayashi-Maskawa (CKM) mixing matrix. The current experimental limits on these processes are [192,193]

$$\text{Br}(\mu \rightarrow e\gamma) < 4.2 \times 10^{-13}, \quad (\text{A5})$$

$$\text{Br}(\tau \rightarrow e\gamma) < 3.3 \times 10^{-8}, \quad (\text{A6})$$

$$\text{Br}(\tau \rightarrow \mu\gamma) < 4.4 \times 10^{-8}. \quad (\text{A7})$$

Among these, $\tau \rightarrow \mu\gamma$ provides the most stringent constraint in our scenario, and the respective future sensitivity is of order $\text{BR}[\tau \rightarrow \mu\gamma] \sim 10^{-9}$ [184].

2. Z decays: $Z \rightarrow \ell\ell'$

Leptonic decays of the Z boson receive contributions from the LQs that constraint the Yukawa couplings. These processes are explained with the following effective Lagrangian:

$$\delta\mathcal{L}_{\text{eff}}^{Z \rightarrow \ell\ell'} = \frac{g}{\cos\theta_W} \sum_{f,i,j} \bar{f} \gamma^\mu (g_{f_L}^{ij} P_L + g_{f_R}^{ij} P_R) f_j Z_\mu. \quad (\text{A8})$$

Here g is the $SU(2)_L$ gauge coupling; g^{ij} are dimensionless couplings measured with great accuracy at the LEP [194] that provide stringent constraints on the associated Yukawa couplings for a fixed LQ mass. NP contributions to these dimensionless couplings can be expressed as follows [195]:

$$\begin{aligned} \text{Re}[\delta g_{L,R}^{\ell}]^{ij} &= \frac{3w_{ij}^u (w_{ii}^u)^*}{16\pi^2} \left[(g_{u_{L,R}} - g_{u_{R,L}}) \frac{x_t(x_t - 1 - \log x_t)}{(x_t - 1)^2} \right] \\ &+ \frac{x_Z}{16\pi^2} \sum_{q=u,c} w_{qj}^u (w_{qi}^u)^* \left[g_{u_{L,R}} \left(\log x_Z - \frac{1}{6} \right) + \frac{g_{\ell_{L,R}}}{6} \right] \\ &+ \frac{x_Z}{16\pi^2} \sum_{q=d,s,b} w_{qj}^d (w_{qi}^d)^* \left[g_{d_{L,R}} \left(\log x_Z - \frac{1}{6} \right) + \frac{g_{\ell_{L,R}}}{6} \right]. \end{aligned} \quad (\text{A9})$$

For δg_L (δg_R), $w_{ij}^u = -(V^* y^S)_{ij}$, $w_{ij}^d = -\sqrt{2} y_{ij}^S$ ($w_{ij}^u = 0$, $w_{ij}^d = 0$) for S_3 LQ. The LEP Collaboration provides the following limits on these NP contributions [196]:

$$\text{Re}[\delta g_L^{ee}] \leq 3.0 \times 10^{-4}, \quad (\text{A10})$$

$$\text{Re}[\delta g_L^{\mu\mu}] \leq 1.1 \times 10^{-3}, \quad (\text{A11})$$

$$\text{Re}[\delta g_L^{\tau\tau}] \leq 5.8 \times 10^{-4}. \quad (\text{A12})$$

3. $\mu \rightarrow e$ conversion

The S_3 LQ mediates $\mu \rightarrow e$ transition in nuclei at the tree level, and the rate of which can be calculated from the following formula [197]:

$$CR(\mu - e) = \frac{\Gamma^{\mu-e}}{\Gamma_{\text{capture}}(Z)}, \quad (\text{A13})$$

$$\Gamma^{\mu-e} = 2G_F^2 |(2V^{(p)} + g_{LV}^{(u)} V^{(n)}) g_{LV}^{(u)}|^2, \quad (\text{A14})$$

$$g_{LV}^{(u)} = \frac{-2v^2}{m_{S_3}^2} (V^* y^S)_{u\ell'} (V^* y^S)_{u\ell'}^*. \quad (\text{A15})$$

Here, $\Gamma_{\text{capture}}(Z)$ is the total capture rate for a nucleus with atomic number Z , which is $13.07 \times 10^6 \text{ s}^{-1}$ for gold, and the corresponding nuclear form factors in units of $m_\mu^{5/2}$ are given by $V^{(p)} = 0.0974$, $V^{(n)} = 0.146$ [197]. The current sensitivity implies [198]

$$CR(\mu - e) < 7 \times 10^{-13}, \quad (\text{A16})$$

whereas the future projected sensitivity is expected to make almost 4 orders of magnitude improvement over the current limit, i.e., $CR(\mu \rightarrow e) < 10^{-16}$ [199–205].

4. $B \rightarrow K^{(*)} \nu \bar{\nu}$ decays

Contributions to the left-handed currents in the $b \rightarrow s \ell \ell$ process unavoidably imply contributions to $B \rightarrow K^{(*)} \nu \bar{\nu}$ decays which are well constrained by experiments. The S_3 LQs can induce $B \rightarrow K^{(*)} \nu \bar{\nu}$ decay at the tree-level via $d_k \rightarrow d_j \nu \bar{\nu}$ processes. The Wilson coefficient responsible for such decays associated with $b \rightarrow s$ transition takes the form

$$C_L^{fi} = \frac{\pi v^2}{2V_{tb} V_{ts}^* \alpha} \frac{y_{bi}^S (y_{sf}^S)^*}{M_3^2}. \quad (\text{A17})$$

Then the branching ratio for $B \rightarrow K^{(*)} \nu \bar{\nu}$ can be expressed as [206]

$$R_{K^{(*)}}^{\nu \bar{\nu}} = \frac{1}{3|C_L^{\text{SM}}|^2} \sum_{i,f=1}^3 |\delta^{fi} C_L^{\text{SM}} + C_L^{fi}|^2, \quad (\text{A18})$$

where $C_L^{\text{SM}} = -1.47/\sin^2 \theta_W$ is the SM contribution. The Belle Collaboration limits these ratios to be $R_{K^*}^{\nu \bar{\nu}} < 3.9$ and $R_{K^0}^{\nu \bar{\nu}} < 2.7$ [186].

5. $B_s^0 - \bar{B}_s^0$ oscillation

S_3 contributes to meson–antimeson mixing, and this NP contribution to $B_s^0 - \bar{B}_s^0$ mixing can be described by the following effective Lagrangian [207]:

$$\mathcal{L}_{\text{eff}}^{\Delta B=2} = -(C_1^{\text{SM}} + C_1^{\text{NP}}) (\bar{b}_L \gamma_\mu s_L)^2. \quad (\text{A19})$$

Here the SM part is $C_1^{\text{SM}} = 2.35/(4\pi^2) (V_{tb} V_{ts}^* G_F m_W)^2$ [208] and the NP contribution at the heavy scale (Λ) is given by [119,207,209,210],

$$C_1^{\text{NP}} = \frac{1}{128\pi^2} \frac{5}{M_3^2} \left(\sum_{\ell} y_{b\ell}^S (y_{s\ell}^S)^* \right)^2. \quad (\text{A20})$$

Here we neglect the evolution of C_1^{NP} from high scale to the m_w scale, which is only relevant for precision calculation. Then the mass difference is given by

$$\Delta m_{B_s}^{\text{SM+NP}} = \Delta m_{B_s}^{\text{SM}} \left| 1 + \frac{C_1^{\text{NP}}}{C_1^{\text{SM}}} \right|, \quad (\text{A21})$$

where the SM prediction is $\Delta m_{B_s}^{\text{SM}} = (18.3 \pm 2.7) \times 10^{12} \text{ s}^{-1}$ [211,212]. This mass difference has been measured in the experiments [194,213] with great accuracy, leading to strong constraints on the NP contribution [214]:

$$|C_1^{\text{NP}}| < 2.01 \times 10^{-5} \text{ TeV}^{-2}. \quad (\text{A22})$$

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