

B decay anomalies from non-Abelian local horizontal symmetry

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Recent anomalies in $B \rightarrow K^{(*)} \ell \ell$ meson decays are consistent with exchange of a heavy Z' vector boson. Here we try to connect such new physics to understanding the origin of flavor, by gauging generation number. Phenomenological and theoretical considerations suggest that the smallest viable flavor symmetry [not including any extra $U(1)$ factors] is chiral $SU(3)_L \times SU(3)_R$, which acts only on generation indices and does not distinguish between quarks and leptons. Spontaneous breaking of the symmetry gives rise to the standard model Yukawa matrices, and masses for the 16 Z' -like gauge bosons, one of which is presumed to be light enough to explain the $B \rightarrow K^{(*)} \ell \ell$ anomalies. We perform a bottom-up study of this framework, showing that it is highly constrained by LHC dilepton searches, meson mixing, Z decays and Cabibbo-Kobayashi-Maskawa unitarity. Similar anomalies are predicted for semileptonic decays of B to lighter mesons, with excesses in the ee , $\tau\tau$ channels and deficits in $\mu\mu$, but no deviation in $\nu\nu$. The lightest Z' mass is $\lesssim 6$ TeV if the gauge coupling is $\lesssim 1$.

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I. INTRODUCTION

Particle physicists have long been waiting for some definitive sign of a breakdown in the Standard Model (SM), which generally works so well as to recall Lord Kelvin's famous statement, "There is nothing new to be discovered in physics now. All that remains is more and more precise measurement." But it has also been anticipated that precision measurements, in the context of flavor, could be the most likely harbinger of new physics (NP), since flavor changing neutral currents (FCNCs) are so highly suppressed in the SM [1,2]. The natural progression of such a signal would be a gradual accumulation of tension in some flavor observables. In recent years, tensions have been mounting in semileptonic B decays, which have been measured with increasing accuracy at the LHC [3–8] and B factories [9–13] and point to new sources of lepton-universality violation in nature.

In particular for the $b \rightarrow s \ell \ell$ FCNC transitions, LHCb has found compelling discrepancies in the ratios

$$R_X = \frac{\mathcal{B}(\bar{B} \rightarrow X \mu^+ \mu^-)}{\mathcal{B}(\bar{B} \rightarrow X e^+ e^-)} \quad (1)$$

for decays into $X = K, K^*$, which are predicted to be very close to 1 in the SM [14–16]. The measured values are

$R_K = 0.745 \pm 0.09 \pm 0.036$ [3], 2.6σ below the SM prediction, and $R_{K^*} = 0.660_{-0.070}^{+0.110} \pm 0.024$ (low q^2), $R_{K^*} = 0.685_{-0.069}^{+0.113} \pm 0.047$, where q^2 is the invariant mass of the lepton pair [7]. The significance of the discrepancy in each bin is $2.2 - 2.5\sigma$. Moreover, an angular analysis of $B \rightarrow K^* \mu \mu$ [6] suggests a 3.4σ discrepancy.

The quantity R_X is particularly interesting because hadronic uncertainties in the decay rate cancel to a high degree in the ratio, making this a "clean" observable (see e.g. [16]). Other measurements such as branching fractions and the $B \rightarrow K^* \mu \mu$ angular observables mentioned above are not so theoretically clean, but it is interesting that their inclusion tends to reinforce the evidence from clean observables only [16–25], a further indication that the effect could be real. The best fits are provided by NP contributions involving the effective operators,

$$\mathcal{H}_w \supset -\frac{\alpha_{\text{em}}}{4\pi v^2} \lambda_{bs}^{(t)} [C_{b_L \ell_L}(\mu) (\bar{s}_L \gamma_\mu b_L) (\bar{\ell}_L \gamma^\mu \ell_L) + C_{b_L \ell_R}(\mu) (\bar{s}_L \gamma_\mu b_L) (\bar{\ell}_R \gamma^\mu \ell_R)], \quad (2)$$

where $\lambda_{bs}^{(t)} = V_{tb} V_{ts}^*$ and the SM contributions are $C_{b_L \ell_L}(m_b) = 8.64$, $C_{b_L \ell_R}(m_b) = -0.18$ [26]. Since $|C_{b_L \ell_L}| \gg |C_{b_L \ell_R}|$, it is possible to fit the data well with NP contributions to the left-handed leptonic operators $C_{b_L \ell_L}$ alone.

The $B \rightarrow K^{(*)} \ell \ell$ anomalies have inspired many model-building efforts, with the most popular proposals involving exchange of heavy Z' vector bosons [27–51], leptoquarks [52–74] or loop-induced transitions [56,75–79]. The data can be well fit in simplified models that are designed to

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address only $R_{K^{(*)}}$, but one naturally hopes that the complete picture would shed greater light on one of the biggest puzzles of the SM, the origin of flavor. If flavor symmetry is local and spontaneously broken, then heavy Z' gauge bosons would inevitably arise, possibly having couplings with the right flavor structure for explaining the anomalies [30,31,44,46]. This is the approach we take, with the goal of adopting the smallest non-Abelian flavor symmetry group that seems to be consistent with the observations, while fully accounting for the structure of the SM Yukawa matrices.

The simplest possibility for a generational symmetry as the origin of flavor would be to couple all SM fermions vectorially to a single $SU(3)_H$ generation group. Although global fits to $B \rightarrow K^{(*)}\ell\ell$ decays disfavor purely vectorial currents to the quarks, it was noted in Ref. [44] that chiral currents can arise for the flavor-changing transitions if only left-handed quarks need to be rotated when diagonalizing the quark masses; the full flavor symmetry group must include a $U(1)_{B-L}$ factor to account for neutrino masses in this model.

Here we consider a different possibility, by assuming the larger chiral group $SU(3)_L \times SU(3)_R$ with no $U(1)$ factor. In addition, we attempt to give a detailed account of the origin of the SM fermion masses within the same framework, as explained in Sec. II. It turns out to be highly constrained, with phenomenological requirements restricting the model-building choices at almost every step. We make a number of predictions for collider searches and precision studies that are imminently testable, as explained in Sec. III. Further consequences of the model, focusing on physics above the scale needed to explain $R_{K^{(*)}}$, are discussed in Sec. IV. We summarize the distinctive features of our model and its differences with previous proposals in Sec. V. Appendix A presents the constraints on possible lepton flavor violation that may be present in the model, while Appendix B explains why a simpler related model, with vectorial $SU(3)_H$ flavor group and no $U(1)$ factor, is not viable.

II. MODEL

In order to generate the SM fermion masses and to cancel anomalies, we add a set of fermions $U_{L,R}, D_{L,R}, E_{L,R}, N_{L,R}$ and scalar fields $\Phi_{u,d,l,\nu}, \mathcal{M}, \Phi_6$, that transform as shown in Table I,

$$\begin{aligned} \mathcal{L}_{\text{yuk}} = & \lambda_u \bar{Q}_L \tilde{H} U_R + \lambda'_u \bar{U}_L \Phi_u u_R + \lambda''_u \bar{U}_L \mathcal{M} U_R \\ & + \lambda_d \bar{Q}_L H D_R + \lambda'_d \bar{D}_L \Phi_d d_R + \lambda''_d \bar{D}_L \mathcal{M} D_R \\ & + \lambda_l \bar{L}_L H E_R + \lambda'_l \bar{E}_L \Phi_l l_R + \lambda''_l \bar{E}_L \mathcal{M} E_R \\ & + \lambda_\nu \bar{L}_L \tilde{H} N_R + \lambda'_\nu \bar{N}_L \Phi_\nu \nu_R + \lambda''_\nu \bar{N}_L \mathcal{M} N_R \\ & + \lambda_6 \nu_R^T \Phi_6 \nu_R. \end{aligned} \quad (3)$$

The new fermions play a double role, by canceling the anomalies of $SU(3)_L \times SU(3)_R$, and by generating the SM Yukawa couplings. A Z_2 symmetry under which the right-

TABLE I. Field content and charges of model. The first three lines are the SM fermions, including right-handed neutrinos, while the following contain the new field content.

Field	$U(1)_y$	$SU(2)_L$	$SU(3)_c$	$SU(3)_L$	$SU(3)_R$	Z_2
Q_L	$\frac{1}{6}$	2	3	3	1	1
L_L	$-\frac{1}{2}$	2	1	3	1	1
u_R	$\frac{2}{3}$	1	3	1	3	-1
d_R	$-\frac{1}{3}$	1	3	1	3	-1
e_R	-1	1	1	1	3	-1
ν_R	0	1	1	1	3	-1
$\Phi_u, \Phi_d, \Phi_e, \Phi_\nu$	0	1	1	1	8	-1
\mathcal{M}	0	1	1	3	$\bar{3}$	1
$\Phi_{8,1}, \Phi_{8,2}$	0	1	1	8	1	1
Φ_6	0	1	1	1	6	1
U_R	$\frac{2}{3}$	1	3	3	1	1
D_R	$-\frac{1}{3}$	1	3	3	1	1
E_R	-1	1	1	3	1	1
N_R	0	1	1	3	1	1
U_L	$\frac{2}{3}$	1	3	1	3	1
D_L	$-\frac{1}{3}$	1	3	1	3	1
E_L	-1	1	1	1	3	1
N_L	0	1	1	1	3	1

handed SM fermions and Φ_f are charged prevents direct flavor-universal mass terms such as $\bar{U}_R u_R$. The scalars $\Phi_{u,d,l,\nu}, \mathcal{M}, \Phi_6$ are present to spontaneously break this symmetry and to dynamically generate the flavor structure of the SM, as we now show.

For simplicity, we take \mathcal{M} to get VEVs proportional to the unit matrix

$$\lambda_f'' \langle \mathcal{M} \rangle = M_f \cdot \mathbf{1} \quad (4)$$

while $\langle \Phi_f \rangle$ may be more complicated. We further assume that $M_f \gg \lambda_f' \langle \Phi_f \rangle$, with the possible exception of $f = u$ because of the large top quark mass. On the other hand $\langle \Phi_6 \rangle$ is much greater than the other VEVs, so that the right-handed neutrinos are very heavy.

Integrating out the heavy fields gives rise to the dimension-5 and -7 operators

$$\begin{aligned} \mathcal{L}_{\text{yuk}} = & \frac{1}{\Lambda_u} \bar{Q}_L \tilde{H} \Phi_u u_R + \frac{1}{\Lambda_d} \bar{Q}_L H \Phi_d d_R + \frac{1}{\Lambda_l} \bar{L}_L H \Phi_l l_R \\ & + \frac{1}{\Lambda_\nu^2} (\bar{L} \tilde{H}) \Phi_\nu \frac{1}{\lambda_6 \langle \Phi_6 \rangle} \Phi_\nu^T (\tilde{H} \bar{L})^T + \text{H.c.} \end{aligned} \quad (5)$$

that become the fermion mass matrices. The mass scales in the denominators are given by $\Lambda_f^{-1} = \lambda_f \lambda_f' / M_f$, except possibly for $f = u$ where we use the more exact expression

$$\Lambda_u^{-1} = \frac{\lambda_u \lambda'_u}{\sqrt{M_u^2 + (\lambda'_u \langle \Phi_u \rangle)^2}}. \quad (6)$$

Below, we will assume that $\langle \Phi_u \rangle$ is a diagonal matrix, and only the $\langle \Phi_u \rangle_{33}$ element will be large enough to matter in this more exact expression.

A. Fermion and gauge boson masses

After the Φ_f scalars and the Higgs field get VEVs, the SM fermions get Dirac masses

$$(m_f)_{ij} = \frac{v}{\Lambda_f} \langle \Phi_f \rangle_{ij} \quad (7)$$

(with Λ_u taken to be a matrix as explained above) while the neutrinos get the Majorana mass matrix

$$\tilde{m}_\nu = \frac{v^2}{\Lambda_\nu^2} \langle \Phi_\nu \rangle \frac{1}{\lambda_6 \langle \Phi_6 \rangle} \langle \Phi_\nu \rangle^T \quad (8)$$

with $v \langle \Phi_\nu \rangle / \Lambda_\nu$ playing the role of the Dirac mass matrix in the seesaw formula.

The large VEV $\langle \Phi_6 \rangle$ breaks $SU(3)_R \times SU(3)_L \rightarrow SU(3)_L$ at a high scale, so we henceforth ignore the heavy gauge bosons associated with $SU(3)_R$.¹ The terms that give masses to the $SU(3)_L$ gauge bosons are

$$\mathcal{L}_{\text{gb}} = g_L^2 \sum_{i=1,2} \text{tr}([\Phi_{8,i}^\dagger, A_H][A_H, \Phi_{8,i}]) + g_L^2 \text{tr}(\mathcal{M}^\dagger \mathcal{M} A_L^2), \quad (9)$$

where $A_L^\mu = T^A A_{L,A}^\mu$ with generators of the fundamental representation, and g_L is the $SU(3)_L$ gauge coupling. We will show that the B decay anomalies motivate us to further break $SU(3)_L \rightarrow U(1)_8$, the $U(1)$ subgroup whose gauge boson $Z' = A_L^8$ couples to the diagonal generator T^8 . This is the reason for including the $\Phi_{8,i}$ octet scalars. It suffices to have VEVs of the form $\langle \Phi_{8,1} \rangle = \alpha T^1$, $\langle \Phi_{8,2} \rangle = \beta T^2$, with $\alpha, \beta \gg \text{TeV}$ to give large masses to all components of A_L except $A_{H,8}$ as desired.

The identification of T^8 as a special direction in the space of generators implies a choice of basis for the fermion flavors. We are assuming that in this basis, the mass matrices of the quarks and charged leptons are diagonal, in the limit where Cabibbo-Kobayashi-Maskawa (CKM) mixing is neglected. To include CKM mixing, we will make the simplifying assumption that the up-like mass matrix $(m_u)_{ij}$ is diagonal, and all the mixing comes from $(m_d)_{ij}$. This choice is particularly convenient for revealing that our model enjoys the properties of minimal flavor violation (MFV) [80,81]; all the FCNCs that arise from Z' exchange explicitly have the same CKM structure as in the SM.

We emphasize that the assumption of $(m_u)_{ij}$ being diagonal is not crucial to the more general framework presented here. It would also be consistent to have off-

¹In general the mass eigenstates are mixtures of A_L and A_R , but if $\langle \Phi_6 \rangle \gg \langle \mathcal{M} \rangle$ as assumed, then the lightest 8 of the 16 gauge bosons will be mostly A_L , with a very small admixture of A_R . For simplicity we will henceforth consider A_R to be decoupled and ignore this small mixing.

diagonal contributions to $(m_u)_{ij}$ similar in relative size to those in $(m_d)_{ij}$. For example, suppose that the fermion masses are diagonalized as usual by unitary transformations $f_L \rightarrow V_f^{L\dagger} f_L$, $f_R \rightarrow V_f^{R\dagger} f_R$, such that $V_u^L = V_d^{L\dagger} = \sqrt{V_{\text{CKM}}} \equiv \mathbf{1} + \frac{1}{2}\delta V - \frac{1}{8}\delta V^2 + \dots$, where $\delta V = V_{\text{CKM}} - \mathbf{1}$. Then the predictions we present in the following would be similar to those in the simpler case where $V_u^L = \mathbf{1}$, $V_d^L = V_{\text{CKM}}^\dagger$. The flavor-changing couplings of Z' to down-type quarks would be approximately half as large, this amount being shifted into those of the up-like quark sector. Detailed predictions would change but the overall picture, including MFV structure, would be preserved. We defer the study of such generalizations for possible future work.

B. Currents

Diagonalizing the gauge boson mass matrix determines the mass eigenstates as $\hat{A}_L^a = O_{ab} A_L^b$, where O is an orthogonal matrix. Our model is such that $Z' \cong \hat{A}_L^8$ is the lightest gauge boson, whose exchange is the origin of anomalous $B \rightarrow K^{(*)} \mu^+ \mu^-$ decays. In general, a Z' that has only flavor-diagonal couplings could couple to the linear combination of generators,

$$O_{1A} T^A \cong T^8 + \frac{\epsilon}{\sqrt{3}} T^3 = \frac{1}{2\sqrt{3}} \begin{pmatrix} 1 + \epsilon & 0 & 0 \\ 0 & 1 - \epsilon & 0 \\ 0 & 0 & -2 \end{pmatrix} \quad (10)$$

[which must be traceless since they belong to $SU(3)$]. It turns out that, to avoid large FCNCs affecting $K-\bar{K}$ mixing, ϵ must be negligibly small. Such operators, with complex coefficients, are directly induced by exchange of the $A_L^{1,2}$ gauge bosons coupling to $T^{1,2}$, which constrains their masses to be at the scale $g_L \langle \Phi_{8,i} \rangle \gtrsim 10^4 \text{ TeV}$. Diagonalization of the gauge boson mass matrix then reveals that $\epsilon \sim \mathcal{M}^2 / \Phi_8^2 \lesssim 10^{-8}$, since $m_{Z'} \cong g_L \mathcal{M}$ is at the TeV scale. We therefore ignore ϵ in the following.

The fermion masses are diagonalized as usual by unitary transformations $f_L \rightarrow V_f^{L\dagger} f_L$, $f_R \rightarrow V_f^{R\dagger} f_R$. Then couplings of Z' to fermions in the mass basis are given by the left-handed currents,

$$g_L Z'_\mu \bar{f}_L [V_f^L T^8 V_f^{L\dagger}] \gamma^\mu f_L, \quad (11)$$

where by our simplifying assumption $V_u^L = \mathbf{1}$ and the CKM mixing is entirely due to $V_d^L = V_{\text{CKM}}^\dagger$.

Then as discussed in Sec. II A, flavor mixing in the down-quark sector has a structure resembling the MFV hypothesis. The diagonal couplings to (left-handed) quarks are given by

$$\frac{g_L}{2\sqrt{3}} \begin{cases} 1, & u, d, c, s \\ -2, & b, t \end{cases} \quad (12)$$

while the off-diagonal ones are

$$-\frac{\sqrt{3}}{2} g_L \begin{cases} V_{ts} V_{td}^*, & s \rightarrow d \\ V_{tb} V_{td}^*, & b \rightarrow d \\ V_{tb} V_{ts}^*, & b \rightarrow s \end{cases}. \quad (13)$$

For the left-handed leptons, we require that $V_l T^8 V_l^\dagger$ is nearly diagonal, to avoid tree-level lepton flavor changing neutral currents. Moreover, the diagonal elements must violate flavor universality to explain the $R_{K^{(*)}}$ anomalies. We assume that

$$V_l^L T^8 V_l^{L\dagger} \cong \frac{1}{2\sqrt{3}} \begin{pmatrix} 1 & \epsilon_1 & \epsilon_2 \\ \epsilon_1^* & -2 & \epsilon_3 \\ \epsilon_2^* & \epsilon_3^* & 1 \end{pmatrix} \quad (14)$$

which preserves the eigenvalues of (10) for $\epsilon_i \ll 1$. Hence, V_l is approximately of the form

$$V_l^L \cong \begin{pmatrix} 0 & 1 & 0 \\ 0 & 0 & 1 \\ 1 & 0 & 0 \end{pmatrix} \quad (15)$$

which is just a permutation, and has determinant +1. This is the unique SU(3) transformation that takes the generator T^8 into a diagonal form in which muons couple more strongly than electrons, as indicated by the $R_{K^{(*)}}$ anomaly (and having the right sign, as will be established below), hence (15) is forced upon us. The right-handed rotation V_l^R is still unconstrained, while the transformation V_ν^L is now determined in terms of the Pontecorvo–Maki–Nakagawa–Sakata (PMNS) neutrino mixing matrix U , as $V_\nu^L = UV_\nu^R$. For simplicity we will impose lepton flavor conservation by taking $\epsilon_i = 0$ in the remainder. If one relaxes this assumption, the experimental constraints on lepton flavor violation require that $|\epsilon_1| < 0.007$, $|\epsilon_{2,3}| < 0.7$, as shown in Appendix A.

It is worth emphasizing that, while there is considerable freedom in choosing VEVs of the Φ_f fields to obtain the flavor structure of the quark and lepton Z' currents, there is also an important restriction: the generators are traceless, forcing Z' to couple with similar strength to all quarks and leptons, in addition to the phenomenologically motivated $b \rightarrow s\ell\ell$ coupling. This leads to interesting constraints and predictions as we now explore.

III. CONSTRAINTS AND PREDICTIONS

Having defined the model, there is only one combination of parameters, $g_L/m_{Z'}$, which is left to fit the $R_{K^{(*)}}$

anomalies. Once this is done, a number of predictions for related FCNC semileptonic meson decays, neutral meson oscillations, Z -decays, and violation of unitarity of the CKM matrix follow. In addition a level of dilepton pair production at the LHC is predicted that is close to current constraints. We discuss these issues in the following.

A. Explaining the $R_{K^{(*)}}$ anomalies

The contributions to $b \rightarrow s\ell\ell$ processes from purely semileptonic operators in the SM are contained in Eq. (2), where the Wilson coefficients are independent of the lepton flavor. Global fits to $R_{K^{(*)}}$ point to new lepton flavor nonuniversal contributions to these operators and including other $b \rightarrow s\mu\mu$ data suggests that part of this NP appears in the muonic operators. Contributions to other operators, such as those involving b_R or different Lorentz structures, are disfavored as discussed in Refs. [16,18,82].

From (13) and (14) it follows that our model produces lepton-specific contributions precisely to $\mathcal{O}_{b_l\ell_L}$,

$$\delta C_{b_l\mu_L} = -2\delta C_{b_Le_L} = -2\delta C_{b_l\tau_L} = -\frac{g_L^2}{m_{Z'}^2} \frac{2\pi v^2}{\alpha}. \quad (16)$$

The CKM coefficient $\lambda_{bs}^{(t)}$ of this contribution has factored out with the SM normalization in Eq. (2), which is a consequence of the MFV-like structure of the Z' couplings to the quarks.

In Fig. 1 we show the trajectory of our model as a function of $m_{Z'}/g_L$ in the $(C_{b_l\mu_L}, C_{b_Le_L})$ plane compared to the best fit point to R_K, R_{K^*} and $B_s \rightarrow \mu\mu$ of Ref. [16]. In the lower-left corner we also show the projection of the χ^2 along $m_{Z'}/g_L$. As one can see, our model gives an excellent fit to the data, with a $\chi^2 = 2.8$ for 3 degrees of freedom, which represents a 4.22σ improvement over the SM. The best fit point and 1σ error interval is

$$\frac{m_{Z'}}{g_L} = 5.3_{-0.6}^{+0.9} \text{ TeV}. \quad (17)$$

We have checked that adding the angular observables of $B \rightarrow K^*\mu\mu$ in a global fit slightly narrows the constraint on $C_{b_l\mu_L}$ but does not have a significant impact on the best solution or improvement with respect to the SM.²

Equation (16) predicts excesses for the branching ratios of $B \rightarrow K^{(*)}\tau\tau$, $B \rightarrow K^{(*)}ee$, which are approximately half

²In Ref. [18] a fit was performed for models that are similar to ours, which gives a significantly stronger bound on $C_{b_Le_L}$. This stems from their inclusion of two data points whose respective preferred solutions for the minimum of χ^2 are regions of parameter space with small overlap; these are inclusive $B \rightarrow X_s ee$ and $B^+ \rightarrow K^+ ee$. We do not include these observables in our fit (17); doing so we find that the improvement in χ^2 is comparable and the best fit value is shifted to $m_{Z'}/g_L \approx 5.1$ TeV.

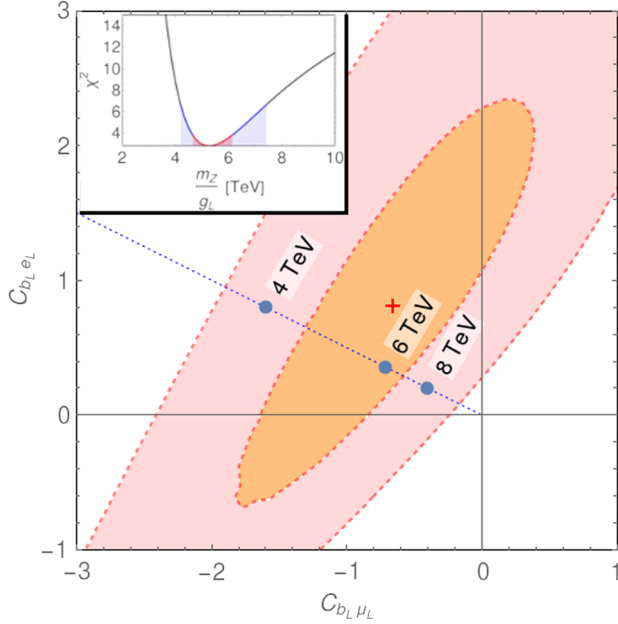


FIG. 1. 1σ and 3σ bounds given by the measurements of R_{K^*} , R_K and the branching fraction of $B_s \rightarrow \mu\mu$ in the $(C_{b_L \mu_L}, C_{b_L e_L})$ plane. The line corresponds to the trajectory of the contributions of the Z' to these Wilson coefficients in our model, where we have indicated benchmark points of $m_{Z'}/g_L$, whereas the embedded plot shows the projected χ^2 . The absolute minimum in the plane, indicated by a red cross corresponds to $\chi^2_{\min} = 2.64$.

the deficit in $B \rightarrow K^{(*)} \mu\mu$. A further consequence of the MFV couplings to the quarks is that similar effects should be measured in $B \rightarrow M l^+ l^-$, where M is a zero-strangeness meson. In particular we predict

$$R_\pi \simeq R_K, \quad R_\rho \simeq R_{K^*}, \quad (18)$$

for the decay channels with pions and ρ mesons.³

B. $d_i \rightarrow d_j \bar{\nu} \nu$ decays

Along with the charged leptons, our Z' couples to neutrinos and hence contributes to rare decays such as $B \rightarrow K^{(*)} \bar{\nu} \nu$ and $K \rightarrow \pi \bar{\nu} \nu$. Interestingly, the former will be searched by Belle II and the latter will be better measured by the NA62 experiment in the coming year. In the SM the $d_i \rightarrow d_j \bar{\nu} \nu$ decays are induced by the low-energy operator

$$\mathcal{H}_w \supset -\frac{\alpha_{\text{em}}}{4\pi v^2} \lambda_{ij}^{(t)} C_{\nu_e} (\bar{d}_j \gamma_\mu d_{iL}) (\bar{\nu}_e \gamma_\mu \nu_{eL}), \quad (19)$$

where $C_{\nu_e} \simeq -12.7$ [84]. The contributions of the Z' are

³Large differences in form factors between the channels could in principle modify this prediction, but such differences are disfavored by approximate SU(3)-flavor symmetry in the light-quark sector of QCD, and by explicit calculations [83].

$$\delta C_{\nu_e} = \delta C_{\nu_\mu} = -\delta C_{\nu_\tau}/2 = \frac{g_L^2}{m_{Z'}^2} \frac{\pi v^2}{\alpha} = 0.37. \quad (20)$$

Even though large deviations are predicted for decays into individual neutrino flavors, what the experiments observe are the “invisible” $B \rightarrow K^{(*)}$ and $K \rightarrow \pi$ rates, in which the absolute contributions from the neutrino flavors are summed over. An important consequence of the tracelessness of the current, Eq. (12), together with the fact that the matrix element contributing to this process is the same for all neutrino flavors, is that the net interference of the Z' and SM contributions vanishes. The NP contribution to the branching fraction is thus given by the quadratic terms,

$$\frac{\text{BR}^{Z'}}{\text{BR}^{\text{SM}}} = \frac{|\delta C_{\nu_e}|^2 + |\delta C_{\nu_\mu}|^2 + |\delta C_{\nu_\tau}|^2}{3|C_{\nu_e}|^2} \simeq 2 \times 10^{-3}, \quad (21)$$

an effect that will be hardly detectable in forthcoming experiments.

C. $\Delta F = 2$ transitions

Neutral-meson mixing receives tree-level contributions from Z' -exchange in our model, yielding

$$\delta \mathcal{H}_w = \frac{3g_L^2}{4m_{Z'}^2} (\lambda_{ij}^{(t)})^2 (\bar{d}_j \gamma^\mu d_{Li}) (\bar{d}_j \gamma_\mu d_{Li}), \quad (22)$$

which has the same operator structure and combination of CKM matrix elements as the box diagram of the top quark in the SM. Parametrizing the deviation from the SM of the ϵ_K parameter in $K-\bar{K}$ mixing by [85]

$$C_{\epsilon_K} = \frac{\text{Im} \langle K^0 | \mathcal{H}_w | \bar{K}^0 \rangle}{\text{Im} \langle K^0 | \mathcal{H}_w^{\text{SM}} | \bar{K}^0 \rangle}, \quad (23)$$

we obtain $C_{\epsilon_K} = 1.14 \pm 0.04$ using Eq. (17), while the current experimental constraint is $C_{\epsilon_K} = 1.04 \pm 0.11$ at 1σ [85]; the latter sets the lower bound $m_{Z'}/g_L \geq 5.1$ TeV, which is quite close to our best fit value (17). In the case of $B_q-\bar{B}_q$ mixing the SM contribution is dominated by the top-loop diagram and its weak phase is aligned with that of the Z' . Thus only the mass differences Δm_{B_q} are constraining, which can be parametrized by

$$C_{B_q} = \left| \frac{\langle B_q | \mathcal{H}_w | \bar{B}_q \rangle}{\langle B_q | \mathcal{H}_w^{\text{SM}} | \bar{B}_q \rangle} \right|. \quad (24)$$

We obtain $C_{B_q} = 1.12 \pm 0.03$ which is within the experimental limits $C_{B_s} = 1.070 \pm 0.088$ and $C_{B_d} = 1.03 \pm 0.11$ [85] and gives the slightly weaker bound $m_{Z'}/g_L \geq 4.8$ TeV. The predictions for C_i and the experimental constraints are summarized in Fig. 2.

There is a potentially dangerous contribution to $K-\bar{K}$ mixing from the loop diagram 3 from exchange of the heavy Φ_d and D particles. In the limit where all the states of

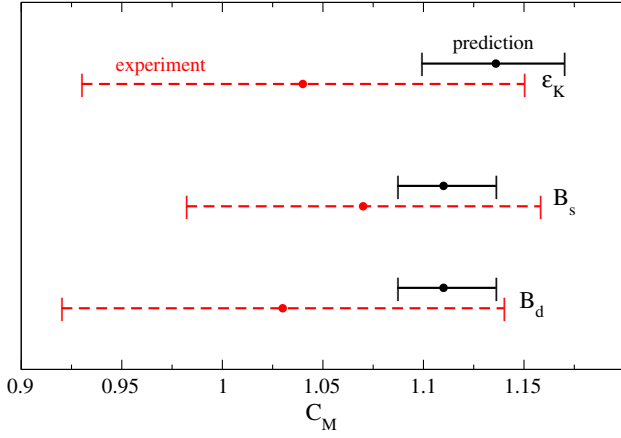


FIG. 2. Predictions and experimentally allowed ranges for the neutral meson mixing parameters (23), (24).

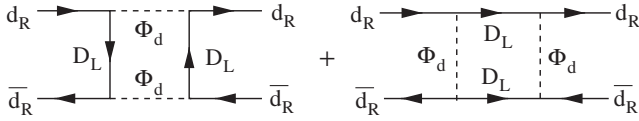


FIG. 3. Loop contribution to neutral meson mixing.

the Φ_d octet are degenerate, the contribution to the amplitude $(\bar{d}_j \gamma^\mu d_{R,i})(\bar{d}_j \gamma^\mu d_{R,i})$ is proportional to the product of SU(3) generators,

$$\sum_{r,s} \sum_{A,B} (T^A)_{ri} (T^A)_{js} (T^B)_{jr} (T^B)_{si} = 0, \quad \text{for } i \neq j. \quad (25)$$

However if there are mass splittings, then FCNCs get generated. For example if Φ_d^1 which couples to T^1 has mass-squared splitting δm_Φ^2 , we find that the operator relevant to $K-\bar{K}$ mixing is

$$\frac{\lambda_d^4 \delta m_\Phi^2}{196\pi^2 m_\Phi^4} (\bar{d}_R \gamma^\mu s_R)^2. \quad (26)$$

Since the coefficient is real, it is constrained at the level of $1/(10^3 \text{ TeV})^2$ [85]. We do not predict the masses m_Φ or splittings δm_Φ^2 here; it would require constructing the full potential of the scalars which is beyond the scope of this work. Nothing ostensibly precludes choosing $\delta m_\Phi^2/m_\Phi^4$ to be sufficiently small.

D. Collider searches for resonant and nonresonant dileptons

A crucial test comes from the search for resonant production of Z' that decays to $\mu^+\mu^-$ and e^+e^- [86]. In our model, production occurs from all flavors of quarks in the proton (but is dominated by the u, d contributions), according to the couplings (12). The branching ratio for decays into muons (electrons) is $B = \frac{1}{12}(\frac{1}{48})$, from

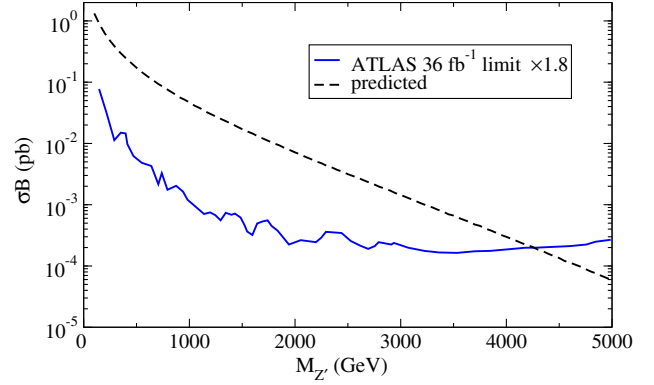


FIG. 4. Production cross section for Z' times branching ratio into muons, versus $m_{Z'}$, as predicted by model and as constrained by ATLAS [86]. The rescaling of the limit by 1.8 corrects for the low branching ratio into ee in our model; see text.

(12) and (14). Using MADGRAPH [87] to predict the resulting production cross section σ at 13 TeV center of mass energy, with QCD correction of $K = 1$, and Eq. (17) to determine g_L , we find the product σB versus $m_{Z'}$ shown in Fig. 4.

The published ATLAS limit applies to models in which equal numbers of electrons and muons are produced. In our model, since primarily muons are produced, and the efficiency for detecting electrons is greater than for muons, the limit is relaxed. In the most interesting mass bin for our purposes, 3–6 TeV, the relative efficiency for electron versus muon detection is $r = 0.45/0.32 = 1.4$. The bound on σB is then relaxed by the factor $(1+r)/(1+r/4) = 1.8$ [88], using $B(ee) = B(\mu\mu)/4$. This leads to the limit $m_{Z'} > 4.3 \text{ TeV}$, which when combined with (17) implies a gauge coupling $g_L \gtrsim 0.7$. Thus another prediction of this model is that the Z' should appear soon in searches for resonant dimuons, if the gauge coupling is not much greater than ~ 1 .

Recently a complementary recasting of dilepton constraints was done in Ref. [89], pointing out that they could also limit the size of effective 4-fermion operators induced by integrating out a heavy Z' , even if its mass is beyond the reach of resonant production at LHC. Coefficients of the operators $(\bar{Q}_1 \gamma^\mu Q_1)(\bar{L}_{e,\mu} \gamma^\mu L_{e,\mu})$ involving first generation left-handed quarks and contributing to $pp \rightarrow e^+e^-$ and $pp \rightarrow \mu^+\mu^-$ are bounded using the resonant dilepton searches. The dimensionless coefficients are identified in our model and constrained as

$$\begin{aligned} C_{Q_1 L_e}^{(1)} &= -\frac{g_L^2 v^2}{6m_{Z'}^2} \\ &= (-1.8 \pm 0.4) \times 10^{-4} \notin [0.0, 1.75] \times 10^{-3} \\ C_{Q_1 L_\mu}^{(1)} &= \frac{2g_L^2 v^2}{3m_{Z'}^2} \\ &= (7.2 \pm 1.9) \times 10^{-4} \in [-5.73, 14.2] \times 10^{-4}, \end{aligned} \quad (27)$$

where the 2σ allowed ranges are given. There is some tension at the level of $\sim 2\sigma$ in the $pp \rightarrow e^+e^-$ channel, where the range is asymmetric because of a deficit of background events in some invariant mass bins. This analysis reinforces the conclusion that dilepton searches could soon reveal evidence for our model, or exclude it.

Interestingly, an independent constraint on the $C_{Q_1 L_e}^{(1)}$ Wilson coefficient arises from parity-violating observables in atomic and electron-proton scattering experiments: $C_{Q_1 L_e}^{(1)} = (1.6 \pm 1.1) \times 10^{-3}$ [90], which is consistent with the bound in Eq. (27). There is no analogous constraint on $C_{Q_1 L_\mu}^{(1)}$, but the muonic coupling of the Z' can be tested using neutrino trident production [91], which in our case leads to the lower limit $m_{Z'}/g_L \gtrsim 700$ GeV.

E. Z and W couplings to fermions

As we discuss in detail below, in order to have a Z' with a mass of ~ 5 TeV, at least some of the exotic fermions U, D, E and N must be in the multi-TeV range. Although such masses are still out of reach for direct searches at the LHC, they can affect low-energy observables like the couplings of weak bosons to the SM fermions. By integrating out all the heavy states in Eq. (3), at the electroweak scale they produce the effective operators

$$\begin{aligned} \mathcal{L}_{\text{eff}} \supset & -\frac{1}{4} \left(\frac{\lambda_d^2}{M_d^2} - \frac{\lambda_u^2}{M_u^2} \right) (H^\dagger i \overleftrightarrow{D}_\mu H) (\bar{Q}_L \gamma^\mu Q_L) \\ & -\frac{1}{4} \left(\frac{\lambda_d^2}{M_d^2} + \frac{\lambda_u^2}{M_u^2} \right) (H^\dagger i \overleftrightarrow{D}_\mu^I H) (\bar{Q}_L \gamma^\mu \tau^I Q_L) \\ & -\frac{1}{4} \left(\frac{\lambda_\ell^2}{M_\ell^2} - \frac{\lambda_\nu^2}{M_\nu^2} \right) (H^\dagger i \overleftrightarrow{D}_\mu H) (\bar{L}_L \gamma^\mu L_L) \\ & -\frac{1}{4} \left(\frac{\lambda_\ell^2}{M_\ell^2} + \frac{\lambda_\nu^2}{M_\nu^2} \right) (H^\dagger i \overleftrightarrow{D}_\mu^I H) (\bar{L}_L \gamma^\mu \tau^I L_L), \end{aligned} \quad (28)$$

where $\overleftrightarrow{D}_\mu = D_\mu - \bar{D}_\mu$ and $\overleftrightarrow{D}_\mu^I = \tau^I D_\mu - \bar{D}_\mu \tau^I$, both of which act trivially in generation space.⁴ These can readily be converted into modifications of the Z and W couplings of the SM left-handed fields,

$$\delta g_{Z,L}^f = -\frac{v^2 \lambda_f^2}{M_f^2}, \quad \delta g_W^{q,\ell} = -\frac{v^2}{2} \left(\frac{\lambda_{d,\ell}^2}{M_{d,\ell}^2} + \frac{\lambda_{u,\nu}^2}{M_{u,\nu}^2} \right). \quad (29)$$

The Z couplings to the fermions have been measured very precisely at LEP. The strongest constraint is on the coupling to the leptons, $\delta g_{Z,L}^\ell = (-0.0952 \pm 0.215) \times 10^{-3}$ which leads to the bound $M_\ell/\lambda_\ell \geq 7.7$ TeV at

⁴We have omitted the contributions to operators of the type $(H^\dagger H)(\bar{Q}_L H d_R), \dots$ which are also generated by integrating out the heavy fermions but that are only weakly constrained by Higgs-coupling measurements.

TABLE II. Upper limit on λ_f/λ_f'' at 95% C.L. from LEP constraints on $Z \rightarrow f\bar{f}$ decays.

$f =$	u	d	ℓ	ν
$\lambda_f/\lambda_f'' <$	2.2	0.87	0.69	1.6

95% C.L. [90]. The invisible width of the Z leads to the bound $\delta g_{Z,L}^\ell = (-1.32 \pm 0.72) \times 10^{-3}$ or $M_\nu/\lambda_\nu > 3.3$ TeV at 95% C.L. Similarly, for the couplings to the up- and down-type quarks we get $M_u/\lambda_u \geq 2.4$ TeV and $M_d/\lambda_d \geq 6.1$ TeV at 95% C.L.

Generally these bounds can be satisfied even if M_f is not very large by taking λ_f sufficiently small. The exact formula for the heavy fermion masses, Eq. (6), implies that $\lambda_f > m_f/v$ [since $\lambda' \Phi / \sqrt{M^2 + (\lambda' \Phi)^2} < 1$], where m_f is the mass of the heaviest SM fermion of type f . Thus for down-type quarks, the Z decay bound can be satisfied even if $M_d = 21$ GeV. However for the up-type quarks we have $\lambda_u \gtrsim 1$, hence $M_u \geq 2.4$ TeV. These constraints can be expressed in terms of the couplings λ_f, λ_f'' by using Eqs. (9) and (17) [see the discussion around Eqs. (32) and (33)], resulting in upper limits shown in Table II.

In the case of the charged currents, the strongest bound stems from the first-row unitarity test of the CKM matrix [92],

$$\begin{aligned} \Delta_{\text{CKM}} &= |\tilde{V}_{ud}|^2 + |\tilde{V}_{us}|^2 + |\tilde{V}_{ub}|^2 - 1 \\ &\simeq v^2 \left(\frac{\lambda_\ell^2}{M_\ell^2} + \frac{\lambda_\nu^2}{M_\nu^2} - \frac{\lambda_d^2}{M_d^2} - \frac{\lambda_u^2}{M_u^2} \right), \end{aligned} \quad (30)$$

where in the second line we have used the corrections in Eq. (28). The experimental bound is $\Delta_{\text{CKM}} = (-4.2 \pm 5.2) \times 10^{-4}$, while the contributions to (30) from the charged leptons can be as large as $(v/7.8 \text{ TeV})^2 \cong 5 \times 10^{-4}$, from the Z-decay constraint. If the other contributions are no larger (even though the Z decay bounds would allow them to be so), the constraint is satisfied without any need for tuning of parameters. This is the case if $\lambda_f/\lambda_f'' \lesssim \sqrt{5 \times 10^{-4} \langle \mathcal{M} \rangle} / v \cong 0.68$, which is consistent with the Z decay limits in Table II.

IV. UV IMPLICATIONS

The discussion so far has been focused on explaining the $R_{K^{(*)}}$ anomalies while satisfying other flavor-changing constraints on the low-energy limit of the theory. Here we return to the higher-energy regime to explore how this relates to the masses of the heavier gauge bosons, and the mechanism of fermion mass generation.

A. Hierarchy of scales

We require the octet scalars $\Phi_{8,i}$ to get VEVs proportional to the generators $T^{1,2}$ in order to give large masses to

all the Z' s that couple to generators other than T^8 . Supposing that $\langle \Phi_8 \rangle = \alpha T^1$, $\langle \Phi_{8,2} \rangle = \beta T^2$ and no other VEVs are present, the gauge boson mass matrix is

$$M_{\text{gb}}^2 = g_L^2 \text{diag} \left(\beta^2, \alpha^2, \gamma^2, \frac{1}{4}\gamma^2, \frac{1}{4}\gamma^2, \frac{1}{4}\gamma^2, \frac{1}{4}\gamma^2, 0 \right), \quad (31)$$

where $\gamma^2 = \alpha^2 + \beta^2$. Separate contributions from two octet fields are required to avoid a second vanishing eigenvalue, that would lead to large FCNCs amongst light quarks. The first two of the states in (31) couple to $T^{1,2}$, which mediate $s \rightarrow d$ transitions. Constraints from $K-\bar{K}$ mixing require that $g_L \alpha, g_L \beta \gtrsim 10^4$ TeV, since their exchange generally produces $(\bar{d}_L \gamma^\mu s_L)^2$ with a coefficient whose imaginary part is not suppressed. (Rotation $T^{1,2} \rightarrow V_{\text{CKM}}^\dagger T^{1,2} V_{\text{CKM}}$ to the quark mass basis does not affect this conclusion.)

Since T^8 commutes with $\langle \Phi_{8,i} \rangle$, only the second term in (7) contributes to the Z' mass. Recalling the simplifying assumption that the \mathcal{M} VEV is proportional to the unit matrix, we have

$$m_{Z'}^2 = g_L^2 \langle \mathcal{M} \rangle^2. \quad (32)$$

Then using (17) the heavy fermion masses are given by

$$M_f = \lambda_f'' \langle \mathcal{M} \rangle = \lambda_f'' \times 5.3 \text{ TeV}. \quad (33)$$

Assuming the couplings $\lambda_f'' \lesssim 1$, this implies that all the heavy fermions are within the reach of the LHC. Current limits on vectorlike quark masses are still close to 1 TeV [93,94].

B. R_D, R_{D^*}

It is interesting to ask whether the present framework could also accommodate the anomalies observed in the decays $B \rightarrow D^{(*)} \tau \nu$. It would require the presence of a heavy W' boson in addition to the Z' . In principle this could be accomplished by extending the gauge symmetry to $\text{SU}(6)_L \times \text{SU}(3)_R$, where $\text{SU}(6)_L$ contains the SM gauge group $\text{SU}(2)_L$. The additional W' gauge bosons then arise from the breaking of $\text{SU}(6)_L \rightarrow \text{SU}(2)_L \times \text{SU}(3)_L$. CKM-like mixing would produce the generation-changing interaction

$$\frac{g_L^2 V_{cb}}{4m_{W'}^2} (\bar{c}_L \gamma^\mu b_L) (\bar{\tau}_L \gamma_\mu \nu_L). \quad (34)$$

Although such an operator can provide a good fit to $R_{D^{(*)}}$, there are two problems in the present framework. First, Eq. (14) also predicts the operator $(\bar{c}_L \gamma^\mu b_L) (\bar{\mu}_L \gamma_\mu \nu_L)$ with coefficient -2 times that in (34), which does not fit the observations [95,96]. Second, as shown in Ref. [97], the required mass for W' to fit $R_{D^{(*)}}$ is too small to satisfy LHC

constraints, given that the W' couples to light quarks through the generator (10).

C. Asymptotic freedom

With the particle content listed in Table I, the g_L and g_R couplings remain asymptotically free. The contributions to the β functions are

$$\frac{16\pi^2}{g_L^3} \beta(g_L) = -11 + \frac{16}{3} + \frac{1}{2} + 1 = -\frac{25}{6} \quad (35)$$

$$\frac{16\pi^2}{g_R^3} \beta(g_R) = -11 + \frac{16}{3} + \frac{1}{2} + 2 + \frac{5}{6} = -\frac{7}{3} \quad (36)$$

from the gauge bosons, fermions, bifundamental scalars, and octets [plus sextet in (36)], respectively. We have assumed that the Φ_f^A component fields are real, so that the matrices Φ_f and hence the SM Yukawa matrices are Hermitian, which is phenomenologically allowed [98].

The beta function (35) is only valid above the scale $\langle \Phi_{8,i} \rangle \sim 10^4$ TeV at which $\text{SU}(3)_L$ is restored. Between this scale and $m_{Z'}$, we should consider the evolution of g_L as the gauge coupling of the $\text{U}(1)$ associated with Z' . Its beta function is given by

$$\beta(g_L) = \frac{g_L^3}{12\pi^2} \times \left(4 + \frac{3}{8} \right), \quad (37)$$

where the respective contributions from the fermions and bosons are shown. Using the initial condition $g_L = 0.7$ at $\mu = m_{Z'} = 4.2$ TeV, which would saturate the current bound from ATLAS dilepton searches, this would lead to a Landau pole at scale $\mu \sim 10^{12}$ TeV. However, asymptotic freedom takes over well before, at 10^4 TeV, so the theory has good UV behavior.

D. Neutrino masses

A further consequence of the structure of the currents is that we are forced to take the transformation V_l^L that diagonalizes the lepton mass matrix to be close to the permutation (15). This means that the lepton masses have to be in an unusual order in the original basis, $\text{diag}(m_\mu, m_\tau, m_e)$. As mentioned above, this fixes the left-handed neutrino rotation in terms of the PMNS matrix U to be

$$V_\nu^L = UV_l^L \cong \begin{pmatrix} -0.15 & 0.82 & 0.54 \\ 0.62 & -0.35 & 0.70 \\ 0.77 & 0.44 & -0.45 \end{pmatrix}. \quad (38)$$

From this, one can infer the form of the seesaw neutrino mass matrix in the original basis, before diagonalization:

$$\begin{aligned} \tilde{m}_\nu &\cong m_{\nu_3} \begin{pmatrix} 0.65 & 0.30 & -0.28 \\ 0.30 & 0.21 & -0.24 \\ -0.28 & -0.24 & 0.28 \end{pmatrix}, & \text{normal} \\ &\cong m_{\nu_3} \begin{pmatrix} 0.98 & 0.12 & 0.09 \\ 0.12 & 0.32 & -0.44 \\ 0.09 & -0.44 & 0.69 \end{pmatrix}, & \text{inverted} \end{aligned}$$

depending upon whether the mass hierarchy is normal or inverted. We assumed that $m_{\nu_1} \ll m_{\nu_{2,3}}$.

V. DISCUSSION

Our model has similarities to that of Ref. [44], in which the gauged flavor symmetry is $SU(3)_H \times U(1)_{B-L}$ acting vectorially on the SM fermions. The horizontal symmetry is the same as we have considered except for the fact that it is not chiral and it includes an extra $U(1)$ factor. This leads to a number of important phenomenological differences between the models. First, right-handed currents are present in Ref. [44] (though they are taken to be flavor diagonal), while they are presumed to be negligible in ours. Second, the flavor generators in [44] are not traceless like in Eqs. (10) and (14). Third, since $B - L$ is opposite for quarks and leptons, the currents for quarks and leptons are different linear combinations of T^8 and 1 in [44], whereas they are the same in our model. The presence of the $U(1)_{B-L}$ factor in [44] leads to a Landau pole at scales $\sim 10^{10}$ GeV, which is not present in our model. Fourth, our model requires no charged lepton flavor violation, whereas it is essential for generating the coupling to muons in [44]. Moreover, we have explored the connection between flavor symmetry breaking and the Yukawa matrices of the SM in our framework.

One consequence of the tracelessness of our generators has already been noted: new contributions to the decays $B \rightarrow K\ell\bar{\nu}$ or $B \rightarrow \pi\ell\bar{\nu}$ are negligible, because the interference with the SM contribution vanishes. Another is that sizable couplings of Z' to all three generations cannot be avoided. In Ref. [44], VEVs for the fundamentals that break $SU(3)_H \times U(1)_{B-L} \rightarrow U(1)_h$ are chosen such that the leptonic generator couples only to the third generator, before mass mixing. By assuming the mixing is small, the branching ratio of $Z' \rightarrow \mu\mu$ (and even more so $Z' \rightarrow ee$) can be suppressed, making it easier to satisfy ATLAS constraints on resonant dilepton production. Our model does not have this option, leading to mild tension in this observable. The traceless generators also imply that no gauge kinetic mixing will arise between the Z' and the SM $U(1)$ hypercharge at one loop. Hence, potentially strong constraints from diboson production [99] are evaded in our model.

Although phenomenologically complete, our analysis does not address how difficult it might be to construct a potential for all the scalar fields that leads to the desired pattern of VEVs, or perhaps a similar one that is

nevertheless viable. This is probably challenging, and might best be postponed pending further experimental evidence in favor of the model. There are a number of new physics signals that should be close to being observable, in addition to direct production of the Z' at LHC. These include a positive contribution to the ϵ_K parameter for $K-\bar{K}$ mixing, a negative contribution to the first-row CKM unitarity test (30), an enhancement of the decay width for $Z \rightarrow \ell\ell$, and vectorlike quarks and leptons at the few-TeV scale.

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APPENDIX A: L-VIOLATING DECAYS

Nothing forbids the entries ϵ_i in the leptonic currents (14), which are constrained by lepton-flavor violating decays such as $\mu \rightarrow 3e$, $\tau \rightarrow 3l$ at the level of 10^{-12} and 10^{-8} in the respective branching ratios. By comparing the NP and SM Wilson coefficients for the exotic decays versus the allowed ones, we find that

$$\begin{aligned} |\epsilon_1| &\lesssim 10^{-6} \frac{2\sqrt{2}G_F}{2g_L^2/(12m_{Z'}^2)} = 0.0067 \\ |\epsilon_{2,3}| &\lesssim 10^{-4} \frac{2\sqrt{2}G_F}{2g_L^2/(12m_{Z'}^2)} = 0.67 \end{aligned} \quad (\text{A1})$$

using Eq. (17).

At one loop, these couplings also give rise to $\mu \rightarrow e\gamma$ and $\tau \rightarrow l\gamma$, through the transition magnetic moment operator $\mu_{ij}(\bar{l}_L^j[\not{q} \cdot \mathbf{A}]l_R^i)$, where q^μ is the photon momentum. We find that

$$\mu_{ij} = \frac{eg_L^2}{384\pi^2 m_{Z'}^2} \begin{cases} -\epsilon_1 m_\mu \ln \frac{m_{Z'}^2}{m_\mu^2}, & \mu \rightarrow e\gamma \\ 2\epsilon_2 m_\tau \ln \frac{m_{Z'}^2}{m_\tau^2}, & \tau \rightarrow e\gamma \\ -\epsilon_3 m_\tau \ln \frac{m_{Z'}^2}{m_\tau^2}, & \tau \rightarrow \mu\gamma. \end{cases} \quad (\text{A2})$$

Using the decay width $\delta\Gamma_{ji} = |\mu_{ij}|^2(m_i^2 - m_j^2)^2/(8\pi m_i)$, and the PDG limits [100] on the radiative decays, we find weaker limits than in (A1):

$$|\epsilon_1| < 0.011, \quad |\epsilon_2| < 4.2, \quad |\epsilon_3| < 5.1, \quad (\text{A3})$$

where we took $m_{Z'} = 6$ TeV to evaluate the logarithms.

The muon anomalous magnetic moment is related to the $\mu \rightarrow e\gamma$ transition moment in (A2) by taking $\epsilon_1 \rightarrow 8$. This

gives a contribution to $(g-2)_\mu/2 = 4 \times 10^{-11}$, smaller than the observed discrepancy by a factor of 75.

Lepton flavor violating decays of vector mesons, for example $J/\psi \rightarrow \mu e$, have branching ratios of order $|\epsilon_i g_L^2 m_{J/\psi}^2 / 32 e^2 m_Z^2|^2 \lesssim 10^{-15} |\epsilon_i|^2$, far below current limits $\sim 10^{-7}$. Pseudoscalar mesons have chirality-suppressed decays to purely leptonic final states. The perturbation to the branching ratio of $B_d \rightarrow \mu\mu$ is predicted to be $\delta B/B \cong \sqrt{2m_B/\pi\Gamma}(V_{td}g_L^2 f_B m_\mu)/(12m_{Z'}^2) \cong 0.08$, which is smaller than the experimental error of 0.4. For the L -violating decays such as $B_{s,d} \rightarrow \mu e$, since there is no interference with the SM the predicted signal is much smaller and gives no useful limits on ϵ_i .

APPENDIX B: VECTORIAL FLAVOR SYMMETRY

One could imagine constructing a similar model to the one we have proposed, but using a vectorial $SU(3)_H$ flavor symmetry instead of $SU(3)_L \times SU(3)_R$. The same interactions as in Eq. (3) could be written, but the fields \mathcal{M}_f would have to be in the 8 representation rather than bifundamental, and a discrete symmetry would be required

to forbid large flavor-universal contributions to the Yukawa matrices involving only SM fields. The flavor-conserving quark and lepton currents would be vectorial, while the FCNCs of the quarks would be left handed as in the model of Ref. [44]. A good fit to $R_{K^{(*)}}$ can still be obtained with vectorial leptonic currents; some authors would argue that this is even preferred [19].

There are several major drawbacks however. First, the sextet field cannot get a large VEV to produce heavy right-handed neutrino masses while leaving a relatively light Z' , making the origin of neutrino masses problematic. [The extra $U(1)_{B-L}$ factor allowed Ref. [44] to overcome this problem.] Second, the tension with dilepton searches is multiplied by having vectorial couplings to the Z' . For resonance searches, the production cross section increases by a factor of 2, while for the nonresonant constraints the number of operators simultaneously contributing to the signal with equal strength is quadrupled, creating a significant tension in all channels but especially electrons. Finally, asymptotic freedom of the gauge coupling is badly spoiled by the large matter content, including ten octet scalars and a heavy copy of the SM fermions, leading to a Landau pole at a relatively low scale.

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