Implications of hidden gauged $U(1)$ model for B anomalies

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(Received 21 December 2017; published 5 June 2018)

We propose a hidden gauged $U(1)_H$ Z' model to explain deviations from the standard model (SM) values in lepton flavor universality known as R_K and R_D anomalies. The Z' only interacts with the SM fermions via their mixing with vectorlike doublet fermions after the $U(1)_H$ symmetry breaking, which leads to $b \rightarrow s\mu\mu$ transition through the Z' at tree level. Moreover, introducing an additional mediator, inert-Higgs doublet, yields $b \rightarrow c\tau\nu$ process via charged scalar contribution at tree level. Using FLAVIO package, we scrutinize adequate sizes of the relevant Wilson coefficients to these two processes by taking various flavor observables into account. It is found that significant mixing between the vectorlike and the second generation leptons is needed for the R_K anomaly. A possible explanation of the R_D anomaly can also be simultaneously addressed in a motivated situation, where a single scalar operator plays a dominant role, by the successful model parameters for the R_K anomaly.

DOI: 10.1103/PhysRevD.97.115003

I. INTRODUCTION

While there is null result from direct searches for new physics at the Large Hadron Collider, recent LHCb measurements of lepton flavor universality in B physics have shown deviations from the standard model (SM) predictions. The measurement of $R_{K^{*0}} = \mathcal{B}(B^0 \rightarrow K^{*0}\mu^+\mu^-)/\mathcal{B}(B^0 \rightarrow K^{*0}e^+e^-)$ in two different kinematic regions by LHCb collaboration gives [1]:

$$R_{K^{*0}} = \begin{cases} 0.66^{+0.11}_{-0.07} \pm 0.03 & (2m_\mu)^2 < q^2 < 1.1 \text{ GeV}^2 \\ 0.69^{+0.11}_{-0.07} \pm 0.05 & 1.1 \text{ GeV}^2 < q^2 < 6 \text{ GeV}^2. \end{cases} \quad (1)$$

These values are compatible with the SM predictions [2–11] within 2.1 – 2.3σ and 2.4 – 2.5σ , respectively. The ratio of $R_K = \mathcal{B}(B^+ \rightarrow K^+\mu^+\mu^-)/\mathcal{B}(B^+ \rightarrow K^+e^+e^-)$ is also given by LHCb collaboration [12] which shows 2.6σ deviation from the SM predictions [2]:

$$R_K = 0.745^{+0.090}_{-0.074} \pm 0.036 \quad 1 \text{ GeV}^2 < q^2 < 6 \text{ GeV}^2. \quad (2)$$

Furthermore, charged current decays of $\bar{B} \rightarrow D^{(*)}l^-\bar{\nu}_l$, which have been measured by the BABAR [13,14], Belle [15–18], and LHCb [19], also indicate discrepancies in the ratios of $R_{D^{(*)}} = \mathcal{B}(\bar{B} \rightarrow D^{(*)}\tau^-\bar{\nu}_\tau)/\mathcal{B}(\bar{B} \rightarrow D^{(*)}l^-\bar{\nu}_l)$ ($l = e, \mu$). The SM predictions are estimated in [20–22] for $R(D^*)$ and [23–25] for $R(D)$, and the current experimental average values are roughly 4σ deviations from the SM values. In recent years, these anomalies have been paid more attention, and possibilities for the explanation in various extensions of the SM are discussed: Z' models [26–35], leptoquark [33,36–40] for the R_K anomaly, and the charged Higgs [41–48], W' models [28,49,50] and leptoquark [38,40,51–55] for the R_D anomaly.¹

We propose a hidden gauged $U(1)_H$ model to explain the R_K anomaly, in which a Z' gauge boson interacts with the SM particles only through mediator particles. The mediators are vectorlike doublet fermions, whose masses are assumed to be around $O(1)$ TeV. The Z' couplings to the SM particles appear with symmetry breaking of the hidden gauged $U(1)_H$, which results in $b \rightarrow s\mu\mu$ transition through the tree-level Z' exchange. Furthermore, introducing an another mediator, we can have a relevant $b \rightarrow c$ transition to the R_D anomaly. The role of the additional mediator is played by an inert-Higgs doublet, and the charged Higgs component of the inert Higgs can induce the $b \rightarrow c$

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¹For model-independent analyses, see [11,56–61].

transition at tree level. Since some essential couplings for the $b \rightarrow s\mu\mu$ are also involved in $b \rightarrow c\tau^-\bar{\nu}$, the former process can affect a possibility of explaining the R_D anomaly.

Typically, the $b \rightarrow s$ transition is tightly constrained by various observables such as $B_s - \bar{B}_s$, $B_s \rightarrow \mu\mu$, etc. While in the lepton sector, the Z' coupling to the muon contributes to neutrino trident production as discussed in [62,63]. At the same time, we carefully examine that there is no significant flavor changing neutral current Z couplings, although the mixing between the mediator and the SM fermions could be significantly large. In addition to these constraints, we also address direct searches of the Z' and the charged Higgs at collider experiments. To avoid flavor constraints from the neutral Higgses, we take degenerate masses on the inert Higgs spectra. We utilize FLAVIO package [9] to perform a comprehensive analysis on various flavor observables, and obtain the global fit to the Wilson coefficients on the R_K and R_D anomalies. Applying these results to our hidden gauged model, we scrutinize the favored parameter space on explaining these anomalies.

The structure of this paper is as follows. We introduce the hidden gauged extension of the SM in Sec. II. In Sec. III, the mixing between the vectorlike and SM fermions and the Z' couplings are presented. Using the FLAVIO, we perform a global fit on the relevant Wilson coefficients to the R_K and R_D anomalies in Sec. IV. Subsequently, various experimental searches and constraints are discussed in Sec. VI. In Sec. V, we apply the obtained results in previous sections to the model parameters. Finally, Sec. VII is devoted to conclusions.

II. THE MODEL

We consider a hidden sector extension of the SM under the hidden gauged $U(1)_H$ symmetry [64]. Although all the SM particles are not charged under the $U(1)_H$ symmetry, the $U(1)_H$ gauge boson can couple to them through mediators, such as new vectorlike fermions or scalars. The mediators have both the $U(1)_H$ and the SM gauge charges. The schematic framework of this model is described in Fig. 1.

For the R_K anomaly, essential mediators are the vectorlike fermions. These fermions could be $SU(2)_L$ singlet, doublet, or multiplets, with one or more generations. Here, we focus on one minimal fermion assignments: all the vectorlike fermions are $SU(2)_L$ doublets, with only one generation.² To be free of gauge anomaly, the new vectorlike fermions need to possess appropriate quantum numbers. We assign the SM quantum numbers on the vectorlike

²Another minimal choice is to introduce only vectorlike singlet fermions with one generation. Since our purpose is to illustrate how the $U(1)_H$ works for the R_K anomaly, the fermion assignment is not so important. We expect that the analysis is quite similar to different fermion assignments.

FIG. 1. The schematic description of the hidden gauged $U(1)_H$ model.

fermions such that they mix with the SM fermions at tree level.

Moreover, if we introduce additional mediators such as a charged scalar, the R_D anomaly could also be explained as a bonus. In this work, we add a second scalar doublet H' as an inert type [64,65] to keep our setup simple.

The new particle contents and their quantum numbers are shown in Table I. Under this charge assignments, the model is free of gauge anomaly. The $U(1)_H$ symmetry spontaneously breaks when a singlet scalar Φ , which is only charged under the $U(1)_H$, obtains a nonzero vacuum expectation value (VEV) v_Φ . They are cast into the form

$$\Phi = \frac{s + ia}{\sqrt{2}}, \quad \langle \Phi \rangle = \frac{v_\Phi}{\sqrt{2}}, \quad (3)$$

and the Lagrangian is written as

$$\mathcal{L}_{\text{kin}} = D_\mu \Phi^\dagger D_\mu \Phi - V(H, \Phi, H'), \quad (4)$$

with the covariant derivative $D_\mu \Phi = (\partial_\mu + ig'Z'_\mu)\Phi$. Here, g' is the gauge coupling of the $U(1)_H$, and the gauge boson mass is given by $m_{Z'} = g'v_\Phi$. In the above potential, H describes the SM $SU(2)$ doublet, and the two Higgs doublets are parametrized by

$$H = \begin{pmatrix} G^+ \\ \frac{v+h+iG^0}{\sqrt{2}} \end{pmatrix}, \quad H' = \begin{pmatrix} H'^+ \\ \frac{H'+iA'}{\sqrt{2}} \end{pmatrix}. \quad (5)$$

TABLE I. The new particle contents and their quantum numbers in the gauge group $SU(2)_L \times U(1)_Y \times U(1)_H$. Here $Q_{L,R}$ and $L_{L,R}$ denote the vectorlike representation under the gauge group. All the SM particles are not charged under $U(1)_H$. The H' doublet is optional: without H' added, only R_K anomaly is addressed; with H' , we could also explain the R_D anomaly. The vectorlike fermion χ is also optional, which is Z_2 -odd under $U(1)_H \rightarrow Z_2$ symmetry breaking, and thus it is the dark matter candidate.

	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_H$
$Q_{L,R} = (\tilde{U}_{L,R}, \tilde{D}_{L,R})^T$	3	2	$+\frac{1}{6}$	+1
$L_{L,R} = (\tilde{N}_{L,R}, \tilde{L}_{L,R})^T$	1	2	$-\frac{1}{2}$	+1
Φ	1	1	0	+1
H' (optional)	1	2	$\frac{1}{2}$	+1
χ (optional)	1	1	0	-1

The potential $V(H, \Phi, H')$ is not relevant to our flavor study, and for simplicity, it is taken as

$$V(H, \Phi, H') = -\mu_\Phi^2 |\Phi|^2 + \lambda_\Phi |\Phi|^4 + \mu_{H'}^2 |H'|^2 + \lambda_{H'} |H'|^4 + \lambda_{H\Phi} |H|^2 |\Phi|^2 + \lambda_{H'\Phi} |H'|^2 |\Phi|^2,$$

which respects a $Z_2(H') \times Z_2(\Phi)$ symmetry. The $\mu_{H'}$ parameter is positive, and thus H' does not get the VEV.

The vectorlike fermions are defined by

$$Q_{L,R} = \begin{pmatrix} \tilde{U}_{L,R} \\ \tilde{D}_{L,R} \end{pmatrix}, \quad L_{L,R} = \begin{pmatrix} \tilde{N}_{L,R} \\ \tilde{L}_{L,R} \end{pmatrix}, \quad (6)$$

and their Dirac masses are

$$-\mathcal{L}_{\text{mass}} = m_Q \bar{Q}_L Q_R + m_L \bar{L}_L L_R + \text{H.c.} \quad (7)$$

All the new fermions are charged under $U(1)_H$, as shown in Table I.³ The right-handed vectorlike fermions can couple to the left-handed SM fermions through the Yukawa interactions of Φ

$$\begin{aligned} -\mathcal{L}_{\text{Yukawa}} &= \Phi^\dagger \sum_i (\bar{q}_{Li} Y_{q_i Q_R} Q_R + \bar{\ell}_{Li} Y_{\ell_i L_R} L_R) + \text{H.c.}, \\ &\supset \frac{v_\Phi}{\sqrt{2}} (Y_{d_i \tilde{D}_R} \bar{d}_{Li} \tilde{D}_R + Y_{u_i \tilde{U}_R} \bar{u}_{Li} \tilde{U}_R \\ &\quad + Y_{\ell_i \tilde{L}_R} \bar{\ell}_{Li} \tilde{L}_R + Y_{\nu_i \tilde{N}_R} \bar{\nu}_{Li} \tilde{N}_R + \text{H.c.}), \end{aligned} \quad (8)$$

where the $SU(2)$ symmetry imposes a relation⁴

$$Y_{d_i \tilde{D}_R} = V_{\text{CKM}}^\dagger Y_{u_i \tilde{U}_R}, \quad Y_{\ell_i \tilde{L}_R} = V_{\text{PMNS}}^\dagger Y_{\nu_i \tilde{N}_R} \quad (9)$$

with a generation index $i = 1, 2, 3$, Cabibbo-Kobayashi-Maskawa (CKM) matrix V_{CKM} and Pontecorvo-Maki-Nakagawa-Sakata (PMNS) matrix V_{PMNS} . Similar situations are studied in [62,66] for the quark sectors and [67] for the lepton sectors. The interactions in Eq. (8) describe the mixing between the vectorlike and SM fermions, which yields $\bar{s}bZ'$ and $\bar{\mu}\mu Z'$ couplings after diagonalization of mass matrices. These Z' couplings are originated from three Yukawa components, $Y_{b\tilde{D}_R}$, $Y_{s\tilde{D}_R}$, and $Y_{\mu\tilde{L}_R}$.

Apart from the right-handed vectorlike fermions, the left-handed vectorlike fermions have couplings to the right-handed SM fermions through the inert doublet

³These new fermions contribute to the kinetic mixing between $U(1)_Y$ and $U(1)_H$ at the one loop order. The mixing parameter is estimated to be roughly $\frac{eq}{(4\pi)^2} \ln \frac{m_Q}{m_L}$. Here for simplicity, we take small or zero mass splitting between the vectorlike quarks and leptons, and thus the kinetic mixing is very tiny and negligible.

⁴If the vectorlike fermions are the $SU(2)$ singlet, in general, they introduce more free Yukawa couplings than the doublet case here.

$$\begin{aligned} -\mathcal{L}'_{\text{Yukawa}} &= y'_{Q_L u_i} \bar{Q}_L \tilde{H}' u_R + y'_{Q_L d_i} \bar{Q}_L H' d_R \\ &\quad + y'_{L_L e_i} \bar{L}_L H' e_R + \text{H.c.}, \end{aligned} \quad (10)$$

with $\tilde{H}' = i\sigma_2 H'^*$. Unlike the Yukawa couplings in Eq. (8), the above interactions do not contribute to the mass matrices of the fermions. $b \rightarrow c\tau\nu$ transition is induced by the charged scalar in the inert doublet, in which the vectorlike fermions play a role in yielding $\bar{c}bH^+$ and $\bar{\nu}\tau H^+$ couplings.

Since we introduce a hidden gauge sector, it is likely to have hidden particles, which are not charged under the SM gauge group at all. If so, the lightest uncolored one would be the dark matter candidate. Here we assume it is the WIMP dark matter. When we study the phenomenology of the hidden sector particles, typically only the lightest dark matter candidate plays an important role. For simplicity, we introduce a vectorlike fermion χ , which only carries the $U(1)_H$ charge

$$\mathcal{L} = \bar{\chi} i\gamma^\mu D_\mu \chi, \quad D_\mu = \partial_\mu - ig' Z'_\mu. \quad (11)$$

To make it a dark matter candidate, we assign a Z_2 charge for the χ . Under $U(1)_H$ breaking down to the Z_2 symmetry, the χ is Z_2 -odd while all other mediators and the SM particles are Z_2 even.

III. MASS MATRICES AND COUPLINGS

The interactions in Eq. (8) imply off-diagonal parts in fermion mass matrices, which are described by

$$\begin{aligned} &(\bar{f}_{i,L} \quad \bar{\tilde{F}}_L) \begin{pmatrix} \frac{v}{\sqrt{2}} y_i & \frac{v_\Phi}{\sqrt{2}} Y_{f_i, \tilde{F}_R} \\ 0 & M_F \end{pmatrix} \begin{pmatrix} f_{i,R} \\ \tilde{F}_R \end{pmatrix} + \text{H.c.} \\ &\equiv \bar{f}_{La} \mathcal{M}^f f_{Ra} + \text{H.c.}, \end{aligned} \quad (12)$$

with $(f_i, \tilde{F})_{L,R}^T = (u_i, \tilde{U})_{L,R}^T, (d_i, \tilde{D})_{L,R}^T, (\ell_i, \tilde{L})_{L,R}^T, i = 1, \dots, 3$ and $a = 1, \dots, 4$. y_i is the SM Higgs Yukawa coupling, and its VEV is $v \simeq 246$ GeV. The matrix $\mathcal{M}^f$ is diagonalized by unitary matrices U_L^f and U_R^f

$$U_L^{f\dagger} \mathcal{M}^f U_R^f = \text{diag}(m_{f_1}, m_{f_2}, m_{f_3}, m_{f_4}), \quad (13)$$

which gives relationships between the original gauge eigenstate f_a and new mass eigenstate f'_a as $f'_a = (U_R^f)^\dagger f_R$ and $f'_L = (U_L^f)^\dagger f_L$. This diagonalization can affect original gauge interactions such as the W and Z boson couplings. However, it should be noted that, since the left-handed vectorlike fermions have the same SM charges as those of the SM fermions, the Z boson couplings to the left-handed fermions do not change. We will see this situation later.

In the charged-lepton sector, the mixing between the second generation lepton μ and vectorlike lepton $\tilde{L}$ is essential for the $b \rightarrow s\mu\mu$ transition. Assuming there is only

one mixing, namely, $Y_{\mu\tilde{L}_R} \neq 0$, the matrices U_L^ℓ and U_R^ℓ can be described by⁵

$$U_L^{\ell\dagger} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & c_{\alpha_L} & 0 & -s_{\alpha_L} \\ 0 & 0 & 1 & 0 \\ 0 & s_{\alpha_L} & 0 & c_{\alpha_L} \end{pmatrix},$$

$$U_R^\ell = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & c_{\alpha_R} & 0 & s_{\alpha_R} \\ 0 & 0 & 1 & 0 \\ 0 & -s_{\alpha_R} & 0 & c_{\alpha_R} \end{pmatrix} \quad (14)$$

where $s_\alpha \equiv \sin \alpha$ and $c_\alpha \equiv \cos \alpha$. The mixing angle α_L and α_R are approximately given by

$$s_{\alpha_L} \simeq \frac{v_\Phi Y_{\mu\tilde{L}_R}/\sqrt{2}}{\sqrt{v_\Phi^2 Y_{\mu\tilde{L}_R}^2/2 + M_L^2}},$$

$$s_{\alpha_R} \simeq \frac{m_\mu}{\sqrt{v_\Phi^2 Y_{\mu\tilde{L}_R}^2/2 + M_L^2}} s_{\alpha_L}. \quad (15)$$

Since we suppose that $v_\Phi \sim O(100)$ GeV, $Y_{\mu\tilde{L}} \sim O(0.1-1)$ and $M_L \sim O(1)$ TeV, it turned out that the angle $\alpha_L \gg \alpha_R$.

For the down-quark sector, the Yukawa couplings $Y_{b\tilde{D}_R}$ and $Y_{s\tilde{D}_R}$ are necessary for the R_K anomaly, which implies the nonzero values of $Y_{t\tilde{U}_R}$ and $Y_{c\tilde{U}_R}$ with the same order of magnitude. Although, apart from the charged-lepton sector, the diagonalization is somewhat complicated due to the two mixing parameters, we obtain similar situation

$$|1 - U_L^{u\dagger}|_{ij} \gg |1 - U_R^u|_{ij}, \quad |1 - U_L^{d\dagger}|_{ij} \gg |1 - U_R^d|_{ij}. \quad (16)$$

Again, the deviations of the right-handed mixing matrices from identity are suppressed by order $\frac{m_q}{m_Q}$.

For the neutrino parts, its Lagrangian contains

$$\mathcal{L}_{\text{neutrino}} \supset -M_L \tilde{N}_L \tilde{N}_R - \frac{v_\Phi}{\sqrt{2}} Y_{\nu_i \tilde{N}_R} \tilde{\nu}_{Li} \tilde{N}_R + \text{H.c.} \quad (17)$$

As explained in the previous section, the second term is related to that of the charged lepton due to the $SU(2)$. As long as we keep only nonzero $Y_{\mu\tilde{L}_R}$, the coupling is given by $Y_{\nu_i \tilde{N}_R} = (V_{\text{PMNS}})_{i2} Y_{\mu\tilde{L}_R}$. For sake of simplicity, we assume that the V_{PMNS} is a unit matrix, which results in only nonzero $Y_{\nu_\mu \tilde{N}_R}$ with more direct relationship of $Y_{\nu_\mu \tilde{N}_R} = Y_{\mu\tilde{L}_R}$. As in the same way of the charged-lepton sector, if we rotate the neutrino fields with

$$\begin{pmatrix} \nu'_{Li} \\ \nu'_{L4} \end{pmatrix} = U_L^{\nu\dagger} \begin{pmatrix} \nu_{Li} \\ \tilde{N}_L \end{pmatrix},$$

$$U_L^{\nu\dagger} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & c_{\beta_L} & 0 & s_{\beta_L} \\ 0 & 0 & 1 & 0 \\ 0 & -s_{\beta_L} & 0 & c_{\beta_L} \end{pmatrix} \quad (18)$$

where

$$s_{\beta_L} = \frac{v_\Phi Y_{\mu\tilde{L}_R}/\sqrt{2}}{\sqrt{M_L^2 + v_\Phi^2 Y_{\mu\tilde{L}_R}^2/2}} \simeq s_{\alpha_L}, \quad (19)$$

the Lagrangian in Eq. (17) becomes

$$\mathcal{L}_{\text{neutrino}} = -\sqrt{M_L^2 + v_\Phi^2 Y_{\mu\tilde{L}_R}^2/2} \tilde{\nu}'_{L4} \tilde{N}_R + \text{H.c.} \quad (20)$$

We find that $U_L^\nu \simeq U_L^\ell$ due to the approximate $SU(2)$ symmetry, which deviates once the lepton masses are taken into account.

After diagonalizing the mass matrices, the Z boson couplings to the lepton sector are

$$\mathcal{L}_Z \supset \frac{e}{s_W c_W} Z_\mu \left[\bar{e}'_L \gamma^\mu U_L^{\ell\dagger} (T_L^3 + s_W^2) U_L^\ell e'_L \right. \\ \left. + \bar{e}'_R \gamma^\mu U_R^{\ell\dagger} (T_R^3 + s_W^2) U_R^\ell e'_R \right. \\ \left. - \bar{\nu}'_L \gamma^\mu U_L^{\nu\dagger} T_L^3 U_L^\nu \nu'_L + \frac{1}{2} \tilde{N}_R \gamma^\mu \tilde{N}_R \right], \quad (21)$$

where T_L^3 and T_R^3 are defined as $T_L^3 = -\frac{1}{2} \text{diag}(1, 1, 1, 1)$ and $T_R^3 = -\frac{1}{2} \text{diag}(0, 0, 0, 1)$. It is seen that the left-handed Z couplings are diagonal, while the flavor-violating Z couplings to the right-handed charged leptons exist. However, the right-handed mixing angle is tiny, which results in the small flavor-violating couplings. Similarly, in the quark sectors, while the Z couplings to the left-handed quarks do not receive any corrections from the field definitions, the right-handed couplings have the flavor-violating parts but they are suppressed.

The W couplings to the leptons are given by

$$\mathcal{L}_W \supset \frac{e}{\sqrt{2} s_W} W_\mu^- [\bar{e}'_L \gamma^\mu U_L^{\ell\dagger} \text{diag}(1, 1, 1, 1) U_L^\ell \nu'_L \\ + \bar{e}'_{Ra} \gamma^\mu (U_R^{\ell\dagger})_{a4} \tilde{N}_R] + \text{H.c.}, \quad (22)$$

Because of the approximate $SU(2)$ symmetry, the left-handed couplings have $U_L^{\ell\dagger} \text{diag}(1, 1, 1, 1) U_L^\ell \simeq 1$. The right-handed interactions are also present with a suppression factor $U_R^{\ell\dagger}$. On the other hand, the quark sector has a somewhat different situation due to the CKM structure. The model parameters $Y_{d_i \tilde{D}_R}$, v_Φ and m_Q , which decide

⁵Here, for simplicity, complex phases are not written down.

FIG. 2. Examples of the left-handed Z' couplings induced by mixing between the vectorlike and the SM fermions with mass insertion notation. Here the double line represents the vectorlike fermions.

$U_{L,R}^u$ and $U_{L,R}^d$, should be chosen to realize the measured CKM values. In our analysis, we use moderate values of the model parameters which satisfy the CKM values within error bars.

One of our main interests is the Z' couplings, which are given by

$$\mathcal{L}_{Z'} \supset -g' Z'_\mu [\bar{e}'_L \gamma^\mu U_L^{\ell\ddagger} \text{diag}(0, 0, 0, 1) U_L^\ell e'_L + \bar{d}'_L \gamma^\mu U_L^{d\ddagger} \text{diag}(0, 0, 0, 1) U_L^d d'_L]. \quad (23)$$

These interactions can be described by integrating out the heavy fermions

$$(U_L^{\ell\ddagger} \text{diag}(0, 0, 0, 1) U_L^\ell)_{ij} \simeq \frac{v_\Phi^2}{2m_{\tilde{L}}^2} Y_{i\tilde{L}_R} Y_{j\tilde{L}_R}^*, \quad (24)$$

$$(U_L^{d\ddagger} \text{diag}(0, 0, 0, 1) U_L^d)_{ij} \simeq \frac{v_\Phi^2}{2m_{\tilde{Q}}^2} Y_{d\tilde{L}_R} Y_{j\tilde{L}_R}^*. \quad (25)$$

and this situation is schematically drawn in Fig. 2. If we focus on only relevant Yukawa couplings $Y_{b\tilde{D}_R}$, $Y_{s\tilde{D}_R}$, and $Y_{\mu\tilde{L}_R}$ to the R_K anomaly, they yield

$$\mathcal{L}_{Z'} \supset -Z'_\mu (g_{sb}^L \bar{s}_L \gamma^\mu b_L + g_{\mu\mu}^L \bar{\mu}_L \gamma^\mu \mu_L + g_{bs}^L \bar{b}_L \gamma^\mu s_L + g_{ss}^L \bar{s}_L \gamma^\mu s_L + g_{bb}^L \bar{b}_L \gamma^\mu b_L), \quad (26)$$

where

$$g_{d_i d_j}^L = g' \frac{v_\Phi^2}{2m_{\tilde{Q}}^2} Y_{d_i \tilde{D}_R} Y_{d_j \tilde{D}_R}^*, \quad g_{\mu\mu}^L = g' \frac{v_\Phi^2}{2m_{\tilde{L}}^2} Y_{\mu \tilde{L}_R} Y_{\mu \tilde{L}_R}^*, \quad (27)$$

proportional to the mixing angles, and we omit the prime in the fermion fields. The first two terms in Eq. (26) lead to the $b \rightarrow s\mu\mu$ transition, and the three Yukawa couplings also produce the Z' couplings to the up quark and neutrino sectors, such as $\bar{t}tZ'$ and $\bar{\nu}_\mu \nu_\mu Z'$, due to the $SU(2)$ symmetry. Note that the right-handed Z' couplings are also present, however, they are induced by the inert scalar loop and are numerically suppressed. Therefore, our current study does not take them into account.

The other intriguing interactions are the Yukawa couplings between the inert-Higgs doublet and the SM fermions, in particular, the charged Higgs couplings $\bar{c}bH^+$ and $\bar{\nu}_\tau H^+$. Using mass-insertion method as in the Z' coupling, we obtain the Yukawa couplings

$$\mathcal{L}_{H'} \supset -(\mathcal{Y}_{c_R b_L} \bar{c}_R b_L + \mathcal{Y}_{c_L b_R} \bar{c}_L b_R) H'^+ - \mathcal{Y}_{\nu_\mu \tau_R} \bar{\nu}_{L\mu} \tau_R H'^+ + \text{H.c.}, \quad (28)$$

with

$$\begin{aligned} \mathcal{Y}_{c_R b_L} &= -\frac{v_\Phi}{\sqrt{2}m_Q} y_{Q_L c}^* Y_{b\tilde{D}_R}^*, \\ \mathcal{Y}_{c_L b_R} &= \frac{v_\Phi}{\sqrt{2}m_Q} Y_{c\tilde{D}_R} y_{Q_L b}', \\ \mathcal{Y}_{\nu_\mu \tau_R} &= \frac{v_\Phi}{\sqrt{2}m_L} Y_{\nu_\mu \tilde{N}_R} y_{L_L \tau}'. \end{aligned} \quad (29)$$

These interactions are originated from the mixing between the vectorlike and SM fermions as seen in Fig. 3. Their dependences on $v_\Phi/m_{Q,L}$ are different from those in Eq. (27). The neutral components in the inert Higgs also induce the Yukawa couplings as in Eq. (28). In our study, we assume that their masses are degenerate. This assumption produces a simple situation where some scalar operators disappear due to the degeneracy as mentioned later.

As we will see in the next section, while the explanation of R_K anomaly needs nonzero g_{sb}^L and $g_{\mu\mu}^L$, the R_D anomaly requires nonzero $\mathcal{Y}_{cb}$ and $\mathcal{Y}_{\nu\tau}$. From the phenomenological point of view, we consider a minimal setup in which these relevant couplings are exclusively focused on. Simultaneously, it implies nonzero $Y_{b\tilde{D}_R}$, $Y_{s\tilde{D}_R}$, $Y_{\mu\tilde{L}_R}$ and $y'_{Q_L b}$, $y'_{Q_L c}$, $y'_{L_L \tau}$. From now on, we denote them as

$$Y_b \equiv Y_{b\tilde{D}_R}, \quad Y_s \equiv Y_{s\tilde{D}_R}, \quad Y_\mu \equiv Y_{\mu\tilde{L}_R}, \quad (30)$$

and

$$y'_\tau \equiv y'_{L_L \tau}, \quad y'_b \equiv y'_{Q_L b}, \quad y'_c \equiv y'_{Q_L c}. \quad (31)$$

FIG. 3. Examples of the charged scalar couplings through mixing between the vectorlike and the SM fermions with mass insertion notation. Here the double line represents the vectorlike fermions.

We also assume that the Yukawa couplings are real, and a flavor hierarchy $|Y_b| \gg |Y_s|$.

IV. B DECAY ANOMALIES

The relevant effective Hamiltonians to $b \rightarrow s\mu^+\mu^-$ and $b \rightarrow c\tau^-\bar{\nu}$ are given by [68–70]

$$\mathcal{H}_{\text{eff}}(b \rightarrow s\mu^+\mu^-) = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \frac{e^2}{16\pi^2} [C_9(\bar{s}\gamma^\mu P_L b)(\bar{\mu}\gamma_\mu \mu) + C_{10}(\bar{s}\gamma^\mu P_L b)(\bar{\mu}\gamma_\mu \gamma_5 \mu)], \quad (32)$$

$$\mathcal{H}_{\text{eff}}(b \rightarrow c\tau^-\bar{\nu}) = \frac{4G_F}{\sqrt{2}} V_{cb} [C_S(\bar{c}P_R b)(\bar{\tau}P_L \nu) + C'_S(\bar{c}P_L b)(\bar{\tau}P_L \nu)]. \quad (33)$$

The Wilson coefficients of C_9 and C_{10} are induced by the Z' interactions in Fig. 2. Integrating out the Z' boson, the Wilson coefficients are given by

$$C_9 = -C_{10} = -\frac{1}{2m_{Z'}^2 C_{\text{SM}}^{bs}} g_{sb}^L g_{\mu\mu}^L = -\frac{v_\Phi^2}{8C_{\text{SM}}^{bs} m_Q^2 m_L^2} Y_s Y_b Y_\mu^2, \quad (34)$$

with $C_{\text{SM}}^{bs} = \frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \frac{e^2}{16\pi^2}$. The final expression in Eq. (34) is obtained by using the couplings defined in Eq. (27). Although $m_{Z'}$ dependence appears in the middle of the Eq. (34), it is canceled out by those in g_{sb}^L and $g_{\mu\mu}^L$. Thus, v_Φ is finally left only in the numerator. The nonzero Yukawa, Y_b , Y_s , and Y_μ , lead to $Y_q = (V_{\text{CKM}})_{qs} Y_s + (V_{\text{CKM}})_{qb} Y_b$ for $q = u, c, t$ and $Y_{\nu_i} = (V_{\text{PMNS}})_{\nu_i\mu} Y_\mu$ for $i = e, \mu, \tau$.

For the R_D anomaly, neglecting the neutrino flavor textures which is irrelevant to our study, we simply assume that the PMNS matrix is a unit matrix, which results in $Y_{\nu_\mu} = Y_\mu$. The Yukawa coupling allows the flavor-violating charged Higgs coupling $\bar{\tau} P_L \nu_\mu H^-$. As seen in [33], such a new flavor-violating process would also contribute to the R_D anomaly. Therefore, the charged scalar interactions in Fig. 3 produce

$$C'_S = -\frac{1}{C_{\text{SM}}^{bc} m_{H'}^2} \mathcal{Y}_{c_R b_L} \mathcal{Y}_{\nu_\mu \tau_R}^* = \frac{v_\Phi^2}{2C_{\text{SM}}^{bc} m_{H'}^2 m_Q m_L} y'_c y'_\tau Y_b Y_\mu, \quad (35)$$

$$C_S = -\frac{1}{C_{\text{SM}}^{bc} m_{H'}^2} \mathcal{Y}_{c_L b_R} \mathcal{Y}_{\nu_\mu \tau_R}^* = -\frac{v_\Phi^2}{2C_{\text{SM}}^{bc} m_{H'}^2 m_Q m_L} y'_b y'_\tau Y_c Y_\mu, \quad (36)$$

with $C_{\text{SM}}^{bc} = \frac{4G_F}{\sqrt{2}} V_{cb}$. The third terms are written by Eq. (28). It is seen that the model parameters appearing in C_9 and C_{10} in the Eq. (34) are correlated to C'_S and C_S except $m_{H'}$ and

$y'_{b,c,\tau}$. Given that $Y_c \sim Y_s$ due to $SU(2)$ symmetry, we expect that C'_S plays a relevant role in the R_D anomaly. As pointed out in [47], a constraint from $\text{BR}(B_c^- \rightarrow \tau \bar{\nu})$ renders explanations for the measured value of R_D with the scalar operators impossible. Therefore, in what follows, our analysis focuses on only R_D taking $C_S = 0$.

In order to determine the possible sizes of these Wilson coefficients, we use FLAVIO 0.21 [9]:

- (i) C_9 and C_{10} The observables and the corresponding experimental measurements listed in the Appendix in Ref. [11] are used, and Table III in the last page summarizes them. We take into account all known correlations among observables and approximate the uncertainties as Gaussian. From Eq. (34), our model holds the relation $C_9 = -C_{10}$, therefore essentially only one Wilson coefficient needs to be fitted. Assuming all the UV parameters are real, we define the following real parameters in the global fit:

$$C_{9,\text{new}} = C_9 \cdot \exp(i \arg(C_{\text{SM}}^{bs})) = -C_{10,\text{new}} = -C_{10} \cdot \exp(i \arg(C_{\text{SM}}^{bs})) \quad (37)$$

The Bayesian method is employed in the global fit with following procedures: we first obtain the likelihood function with a single argument $C_{9,\text{new}}$ from FLAVIO FASTFIT class, and assume a uniform prior probability of $C_{9,\text{new}}$ ranges from -3 , to 3 . We then use pymultinest [71] to implement a Monte Carlo sampling, and finally obtain the posterior probability shown in Fig. 4 with corresponding 1σ and 2σ uncertainties, which is plotted using SUPERPLOT [72]. In the figure, the best-fit point is indicated

FIG. 4. The likelihood as function of the $C_{9,\text{new}}$ from the Bayesian fit. The blue and green lines on the top of curve represent the 1σ and 2σ regions. The star at the bottom is the best-fit point.

FIG. 5. The one-dimensional χ^2 fit on C'_S for the measured value of R_D with $C_S = 0$. The brown, green, and cyan lines correspond to 1σ , 2σ , and 3σ regions. The red and light red regions are excluded by the current LHC limit, $\text{BR}(B_c^- \rightarrow \tau^- \bar{\nu}) < 30\%$ [47], and recasted LEP limit, $\text{BR}(B_c^- \rightarrow \tau^- \bar{\nu}) < 10\%$ [73], respectively.

by a star, and 1σ and 2σ regions are represented by blue and green lines on the top of curve. The obtained values are in agreement with the current observed branching ratio for $B^{0,\pm} \rightarrow K^{*0,\pm} \mu\mu$, $B_s \rightarrow \phi\mu\mu$, $B \rightarrow X_s\mu\mu$, and $B_s \rightarrow \mu\mu$.

(ii) C'_S

For the R_D anomaly, our numerical analysis utilizes the experimental results in Ref. [14,15]. Considering the R_D as a function of C'_S , we derive the 1 , 2 , and 3σ region by demanding the $\Delta\chi^2 = \chi^2 - \chi^2_{\min} = 1, 4, 9$. Figure 5 shows the value of $\Delta\chi^2$ depending on C'_S . The 1 , 2 , and 3σ regions correspond to brown, green, and cyan lines, and they are given by $(-1.047, -0.689) \cup (0.689, 1.047)$, $(-1.185, -0.405) \cup (0.405, 1.185)$, and $(-1.309, 1.309)$. The red and light red regions represent the constraints $\text{BR}(B_c^- \rightarrow \tau^- \bar{\nu}) < 30\%$ and $\text{BR}(B_c^- \rightarrow \tau^- \bar{\nu}) < 10\%$, respectively. Taking the limit $\text{BR}(B_c^- \rightarrow \tau^- \bar{\nu}) < 30\%$ interpreted by the lifetime of the B_c^- meson [47], it gives the allowed region with $|C'_S| < -0.8454$. Moreover, if the limit $\text{BR}(B_c^- \rightarrow \tau^- \bar{\nu}) < 10\%$ [73] which recasts LEP data, it is seen that the current setup can marginally explain the $R(D)$ anomaly within 2σ region.

Finally, we summarize the obtained numerical values in Table II, where the limit $\text{BR}(B_c^- \rightarrow \tau^- \bar{\nu}) < 30\%$ is imposed on C'_S .⁶

⁶The 2σ range of C'_S implies that $R_D \sim 0.25$ which is consistent with the SM prediction.

V. EXPERIMENTAL SEARCHES AND CONSTRAINTS

The hidden gauge boson Z' encounters direct constraints from collider searches. Depending on the Z' mass, $m_{Z'} = g'v_\Phi$, there are different limits on the signal rate of the Z' production. Here, we discuss the following cases: light Z' case ($m_{Z'} < m_Z$) and heavy Z' case ($m_{Z'} > m_Z$).

- (i) When the Z' is lighter than the Z mass, some viable parameter regions exist. Since the Z' has no coupling to the electron, LEP searches cannot provide direct constraint on the light Z' . Furthermore, the Tevatron [85,86] and LHC [87,88] searches for Z' to dilepton final state only apply to the case of $m_{Z'} > 100$ GeV. The relevant limit to the light Z' case comes from the LHC searches at $pp \rightarrow Z \rightarrow 4\mu$. Its typical SM process is through the off-shell Z mediated 2μ decay, while the light Z' could be on-shell when $m_{Z'} < m_Z - 2m_\mu$. A detailed analysis on how to recast the current LHC search limit has been done in Ref. [63]. Mapping their analyses to our model, we obtain the constraint shown in Fig. 6. In the figure, the orange line indicates the g' coupling as a function of $m_{Z'}$ with $v_\Phi = 700$ GeV. The excluded region corresponds to the blue region. For $v_\Phi = 700$ GeV, we find that the region of $m_{Z'} \lesssim 10$ GeV and $50 \text{ GeV} \lesssim m_{Z'}$ still have some spaces for $10^{-3} < g' < 0.14$. However, one should note that for $m_{Z'} \lesssim 10$ GeV, we are not able to integrate out the Z' particle when deriving the Wilson coefficients C_9 and C_{10} in Eq. (34).

- (ii) In the case of the heavy Z' , the LHC searches in the dimuon final states put the tightest constraints on its mass. The current limit on $m_{Z'}$ is around 3–4 TeV [87,88]. In order to apply this limit to our case, we should first take into account the suppressed coupling to the first generation quarks by the small mixing with vectorlike fermions, which results in small production cross section. Moreover, although the LHC searches assume that the decay branching ratio to the muon or electron lepton is 100%, our Z' can also decay into the SM quarks, and other light $U(1)_H$ charged particles. Therefore, the branching ratio could be much smaller than the assumed value. We estimate $pp \rightarrow Z'$ production cross section with MADGRAPH [89], and demonstrate the constraint on $\sigma(pp \rightarrow Z' \rightarrow \mu\mu)$ in Fig. 7 with benchmark values: $v_\Phi = 700$ GeV, $Y_b = -1$ and $Y_s = 0.0184$. In addition to the muons, another decay processes of the bottom, top, neutrino, charged scalar, scalar (s, H), and pseudoscalar (A) are also present. A naive estimation implies that the Z' decay branching ratio to the muons is roughly $1/10.5$. The cross section times branching ratio is shown in the red curve. According to the figure, the Z' mass can be as low as 1.5 TeV. Furthermore, the branching ratio can be

TABLE II. The desired regions of the Wilson coefficients for the measured R_K and R_D values. The limit of $\text{BR}(B_c^- \rightarrow \tau^- \bar{\nu}) < 30\%$ is taken for C'_S .

	$C_9 = -C_{10}$	$C_{9,\text{new}} = -C_{10,\text{new}}$	C'_S with $C_S = 0$
1σ	$[-0.742, -0.643]$	$[0.515, 0.72]$	$[-0.8454, -0.689] \cup [0.689, 0.8454]$
2σ	$[-0.767, -0.618]$	$[0.433, 0.843]$	$[-0.8454, -0.405] \cup [0.405, 0.8454]$

reduced if additional decay channels come from unknown hidden gauge sector particles, which are irrelevant to our flavor study. However, if we allow opening additional decay channels, the constraint

could be significantly relaxed. For example, taking the branching ratio to be $\text{BR}(Z' \rightarrow \mu\mu) = 1/20$, which is shown in the green curve, the Z' mass can be lower than 750 GeV.

TABLE III. List of observable used in global fit of $R_{K^{(*)}}$ anomaly.

Observable	$[q_{\min}^2, q_{\max}^2]$ [GeV ²]	Experiments
<i>Branching ratios</i>		
$\frac{d}{dq^2} \text{BR}(B \rightarrow X_s \mu\mu)$	[1, 6], [14.2, 25]	BABAR [74]
$\frac{d}{dq^2} \text{BR}(B^+ \rightarrow K^{*+} \mu\mu)$	[0.0, 2.0], [2.0, 4.3]	CDF [75]
$\frac{d}{dq^2} \text{BR}(B^\pm \rightarrow K^\pm \mu\mu)$	[2, 4], [4, 6], [15, 19]	LHCb [76]
$\frac{d}{dq^2} \text{BR}(B^\pm \rightarrow K^\pm \mu\mu)$	[0.0, 2.0], [2.0, 4.3]	CDF [75]
$\frac{d}{dq^2} \text{BR}(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2], [2, 3], [3, 4], [4, 5], [15, 22]	LHCb [76]
$\frac{d}{dq^2} \text{BR}(B^0 \rightarrow K^{*0} \mu\mu)$	[0.0, 2.0], [2.0, 4.3]	CDF [75]
$\frac{d}{dq^2} \text{BR}(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [77]
$\frac{d}{dq^2} \text{BR}(B^0 \rightarrow K^{*0} \mu\mu)$	[1, 2], [2, 4.3]	CMS [78,79]
$\frac{d}{dq^2} \text{BR}(B^0 \rightarrow K^{*0} \mu\mu)$	[0.0, 2.0], [2.0, 4.3]	CDF [75]
$\frac{d}{dq^2} \text{BR}(B_s \rightarrow \phi \mu\mu)$	[2.5, 4], [4, 6], [15, 22]	LHCb [76]
$\frac{d}{dq^2} \text{BR}(B_s \rightarrow \phi \mu\mu)$	[1, 6]	CDF [75]
R_{K^*}	[1, 6], [15, 19]	LHCb [80]
R_K	[0.045, 1.1], [1.1, 6.0]	LHCb [1]
$\text{BR}(B_s \rightarrow \mu\mu), \text{BR}(B^0 \rightarrow \mu\mu)$	[1.0, 6.0]	LHCb [12]
	...	LHCb [81]
<i>Angular observables</i>		
$\langle A_{FB} \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[0.0, 2.0], [2.0, 4.3]	CDF [75]
$\langle A_{FB} \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [82]
$\langle F_L \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1, 2], [2, 4.3], [4.3, 6]	CMS [78,79,83]
$\langle F_L \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[0.0, 2.0], [2.0, 4.3]	CDF [75]
$\langle F_L \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [82]
$\langle F_L \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1, 2], [2, 4.3], [4.3, 6]	CMS [78,79,83]
$\langle S_3 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[0.04, 2.0], [2.0, 4.0], [4.0, 6.0]	ATLAS [84]
$\langle S_3 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [82]
$\langle S_4 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1, 2], [2, 4.3], [4.3, 6]	CMS [78,79,83]
$\langle S_4 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [82]
$\langle S_5 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1, 2], [2, 4.3], [4.3, 6]	CMS [78,79,83]
$\langle S_5 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [82]
$\langle P_1 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[0.04, 2.0], [2.0, 4.0], [4.0, 6.0]	ATLAS [84]
$\langle P_1 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [82]
$\langle P_4 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1, 2], [2, 4.3], [4.3, 6]	CMS [78,79,83]
$\langle P_4 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[0.04, 2.0], [2.0, 4.0], [4.0, 6.0]	ATLAS [84]
$\langle P_5 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [82]
$\langle P_5 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[0.04, 2.0], [2.0, 4.0], [4.0, 6.0]	ATLAS [84]
$\langle P_5 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1.1, 2.5], [2.5, 4], [4, 6], [15, 19]	LHCb [82]
$\langle \bar{F}_L \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[1, 2], [2, 4.3], [4.3, 6]	CMS [78,79,83]
$\langle \bar{F}_L \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[2, 5], [15, 19]	LHCb [82]
$\langle \bar{S}_3 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[2, 5], [15, 19]	LHCb [82]
$\langle \bar{S}_4 \rangle(B^0 \rightarrow K^{*0} \mu\mu)$	[2, 5], [15, 19]	LHCb [82]

FIG. 6. Constraint on low mass Z' from $Z \rightarrow 4\mu$ search at LHC. The blue region is excluded by the current LHC searches. The orange line represents our benchmark model with $v_\phi = 700$ GeV.

Next, we briefly comment on the constraints on the mediators. In [90], the vectorlike quark search is discussed with its decay into Wb , and the current limit is $m_Q > 1.1$ TeV. On the other hand, studies about the multilepton final state [91] show that the limit on the vectorlike lepton is $m_L > 500$ GeV. In our numerical analysis, we take the benchmark point $m_Q = m_L = 1.4$ TeV, which is consistent with these limits. Regarding the inert-Higgs doublet, the charged scalar receives the tightest bounds. The LHC searches for $H^\pm$ are all associated with the top quark production through $tbH^\pm$ coupling. Therefore, we can evade any constraints from the LHC by setting $y'_t = y'_b = 0$ such that the $tbH^\pm$ coupling vanishes, and it follows that $C_S = 0$. Actually, we investigate the scenario that y'_t and y'_b are nonzero and calculate the production cross section for the process $pp \rightarrow \bar{t}bH^+$ with MADGRAPH for a benchmark point: $y'_t = y'_b = 1$, $m_Q = 1.4$ TeV, $v_\phi = 700$ GeV, $Y_b = -1$, and $Y_s = 0.0184$. It is found that $m_{H^\pm} > 300$ GeV is consistent with the ATLAS experiments [92,93] provided that $\text{BR}(H^+ \rightarrow \tau\nu) < 80\%$. Again, if taking into account the $H^\pm$ decay channel, $H^{'+} \rightarrow \bar{t}\bar{b}$, $H^{'+} \rightarrow c\bar{b}$, $H^{'+} \rightarrow c\bar{s}$, and $H^{'+} \rightarrow \tau\nu$, we find that the requirement $\text{BR}(H^+ \rightarrow \tau\nu) < 80\%$ is easily satisfied.

In the following, we consider various constraints from flavor physics and low energy observables:

(i) $B_s - \bar{B}_s$ mixing

The $\bar{s}bZ'$ coupling induces a relevant operator to the $B_s - \bar{B}_s$ mixing

$$\mathcal{H}_{\text{eff}} = -C_V^{LL}(\bar{s}_L\gamma^\mu b_L)(\bar{s}_L\gamma_\mu b_L) \quad (38)$$

with $C_V^{LL} = -(g_{sb}^L)^2/m_{Z'}^2$. We calculate the allowed region of this operator using the measurements of ΔM_s for $B_s - \bar{B}_s$ mixing in Ref. [94]. In our analysis, the Bayesian fit is also employed. The likelihood function is approximated by Gaussian probability distribution function with mean as

FIG. 7. Constraint on Z' with $v_\phi = 700$ GeV, $Y_b = -1$, and $Y_s = 0.0184$. The blue line is the 95% C.L. exclusion limit by ATLAS experiments [87]. The red and green curves correspond to the branching ratio of $\text{BR}(Z' \rightarrow \mu\mu) = 1/10.5$ and $\text{BR}(Z' \rightarrow \mu\mu) = 1/20$ (including new hidden particles), respectively.

difference of the theoretical prediction and experimental central value, and the variance as the quadratic sum of the theoretical uncertainty and the experimental uncertainty. The prior probability distribution is taken as uniform distribution between -3×10^{-10} to 3×10^{-10} . The fitting results are shown in Fig. 8, and it is found that $|C_V^{LL}| < 1.29 \times 10^{-11}$ at 1σ . Note that, although scalar operators $\mathcal{O}_S^{LL}$ and $\mathcal{O}_S^{RR}$ listed in [95] are caused by neutral and CP -odd scalars in the inert-Higgs doublets, they vanish as long as the two scalar masses are degenerate. Also, we do not consider an operator $\mathcal{O}_S^{LR}$ since its Wilson coefficient is proportional to y'_s that is irrelevant to our current study.

FIG. 8. The normalized likelihood as function of the C_V^{LL} coefficient using the Bayesian fit.

FIG. 9. The value of $\text{Br}(B_s \rightarrow \mu\mu)$ as a function of $C_{9,\text{new}} = -C_{10,\text{new}}$. The orange band is the 1σ region given by the latest measurement at LHCb [81]. The dark blue and light blue are the 1 and 2σ region for explaining the R_K anomaly.

(ii) $B_s \rightarrow \mu\mu$

The nonzero $C_{10,\text{new}}$ can also contribute to the $B_s \rightarrow \mu\mu$ process. The current measurement of $\text{Br}(B_s \rightarrow \mu\mu)$ is given by LHCb [81], which is $\text{Br}(B_s \rightarrow \mu\mu) = 3.0 \pm 0.6^{+0.3}_{-0.2} \times 10^{-9}$. In Fig. 9, the red line is the theoretical value of $\text{Br}(B_s \rightarrow \mu\mu)$ as a function of $C_{9,\text{new}} (= -C_{10,\text{new}})$. The 1σ region of the LHCb experimental value is shown by the orange band. The 1 and 2σ regions for the global fit of R_K anomaly correspond to the dark and light blue bands. It is seen that the red line on the blue bands is within the 1σ region of $\text{Br}(B_s \rightarrow \mu\mu)$.

(iii) Neutrino trident production

Neutrino trident production gives an upper limit on $g_{\mu\mu}^L/m_{Z'} = v_\Phi Y_{\mu L_R}^2/2m_L^2$. Its cross section is given by [63,96,97]

$$\frac{\sigma_{\text{SM+NP}}}{\sigma_{\text{SM}}} = \frac{1}{1 + (1 + 4s_W^2)^2} \left[\left\{ 1 + v^2 \left(\frac{g_{\mu\mu}^L}{m_{Z'}} \right)^2 \right\}^2 + \left\{ 1 + 4s_W^2 + v^2 \left(\frac{g_{\mu\mu}^L}{m_{Z'}} \right)^2 \right\}^2 \right]. \quad (39)$$

The experimental measurement [98] shows $\sigma_{\text{exp}}/\sigma_{\text{SM}} = 0.82 \pm 0.28$. For $m_L = 1.4$ TeV and $v_\Phi = 700$ GeV, the 1σ region gives an upper bound on the Yukawa coupling $|Y_{\mu L_R}| \lesssim 2.1$.

(iv) $B \rightarrow K\nu\bar{\nu}$

In the same way as $b \rightarrow s\mu\mu$ process, $b \rightarrow s\nu\bar{\nu}$ transition is also induced as shown in Fig. 10. The effective Hamiltonian is [33]

$$\mathcal{H}_{\text{eff}}(b \rightarrow s\nu_i\bar{\nu}_j) = -\frac{\alpha G_F}{\sqrt{2}\pi} V_{tb} V_{ts}^* C_L^{ij} (\bar{c} \gamma^\mu P_L b) \times (\bar{\nu}_i \gamma_\mu (1 - \gamma_5) \nu_j). \quad (40)$$

FIG. 10. The $b \rightarrow s\nu\bar{\nu}$ transition induced by the Z' exchanges. Here the double line represents the vectorlike fermions.

Our current setup produces only the 2nd generation neutrino at final states, therefore the Wilson coefficient is described by

$$C_L^{22} = -\frac{1}{\sqrt{2}m_{Z'}^2} \frac{\pi}{\alpha G_F V_{ts} V_{ts}^*} g_{sb}^L g_{\nu_\mu \nu_\mu}^L, \quad (41)$$

with $g_{\nu_\mu \nu_\mu}^L = g_{\mu\mu}^L$. Following the bound on the Wilson coefficient obtained in [33], we find that $-13\text{Re}[C_L^{22}] + |C_L^{22}|^2 \leq 473$.

(v) $b \rightarrow s\gamma$

The Yukawa couplings in Eq. (28) contribute to $b \rightarrow s\gamma$ process through the charged scalar loop. In Fig. 11, two possible diagrams are drawn. It is seen that while the left diagram contains the product $y'_b y'_c$ related to C_S and C'_S , the diagram on the right is proportional to $y'_c y'_s$. Therefore, we include only the left diagram which contains two relevant parameters to $b \rightarrow c$ transition. The effective operators are given by [99]

$$\mathcal{H}_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* (C_7 \mathcal{O}_7 + C_8 \mathcal{O}_8) \quad (42)$$

with

$$\mathcal{O}_7 = \frac{e}{16\pi^2} m_b \bar{s}_L \sigma^{\mu\nu} b_R F_{\mu\nu} \quad (43)$$

$$\mathcal{O}_8 = \frac{e}{16\pi^2} m_b \bar{s}_L T^a \sigma^{\mu\nu} b_R G_{\mu\nu}^A, \quad (44)$$

The Wilson coefficients in our model are

FIG. 11. $b \rightarrow s\gamma$ process through the charged scalar loop. The double lines represent the vectorlike fermions $\tilde{U}$ or $\tilde{D}$. The left diagram is proportional to $y'_b y'_c$, while the right diagram contains $y'_c y'_s$.

$$C_7 = \frac{v^2}{V_{tb}V_{ts}^*} \mathcal{Y}_{c_L b_R} \mathcal{Y}_{c_R s_L}^* \left(\frac{C_{7,XY}^0(x_c)}{m_b m_c} + \frac{C_{7,YY}^0(x_c)}{m_c^2} \right), \quad (45)$$

$$C_8 = \frac{v^2}{V_{tb}V_{ts}^*} \mathcal{Y}_{c_L b_R} \mathcal{Y}_{c_R s_L}^* \left(\frac{C_{8,XY}^0(x_c)}{m_b m_c} + \frac{C_{8,YY}^0(x_c)}{m_c^2} \right), \quad (46)$$

where $x_c = m_c^2/m_{H'}^2$ and the loop functions $C_{7,8}^0$ are listed in [99]. The Wilson coefficients C_7 and C_8 are proportional to a product of $Y_s Y_c y'_b y'_c$. While Y_s and Y_c are predicted by the explanation of the R_K anomaly, y'_b and y'_c are related to C'_S and C_S in Eqs. (35) and (36). According to a global fit of the Wilson coefficients in the $b \rightarrow s$ transition in [3], contributions from new physics to $C_7(\mu_b = 5 \text{ GeV})$ are allowed between $-0.03 \leq C_7(\mu_b) \leq 0.032$ for 1σ . Taking that the charged scalar mass is 300 GeV, we find that the loop functions are roughly $O(10^{-5}-10^{-6})$. In addition, the Wilson coefficients receive suppression proportional to $Y_s Y_c \sim 10^{-4}$. Later, we will check this constraint numerically. It should be noted that, as in the $B_s - \bar{B}_s$ mixing, the neutral scalars do not contribute to the process due to cancellations between H' and A' contributions for the degenerate mass spectra in the inert Higgs.

In addition to the above constraints, lepton-flavor-violating and -conserving processes, such as $\tau \rightarrow 3\mu$, $\tau \rightarrow \mu\gamma$ and muon $g-2$ can also be induced. However, these processes also need irrelevant parameters to the B anomalies,⁷ therefore we do not deal with them here.

Finally, we briefly discuss the hidden particles which are only charged under the $U(1)_H$ group. Although these hidden particles are not relevant to the B physics observables, it provides additional signatures. We assume that there is a WIMP dark matter candidate in the hidden sector, and the dark matter particle χ should annihilate into SM particles during freeze-out at early universe. In the non-relativistic approximation, the annihilation cross section times relative velocity $\langle\sigma v\rangle$ can be decomposed as

$$\langle\sigma v\rangle = a + bv^2 + \mathcal{O}(v^4), \quad (47)$$

where a and b are the s -wave and p -wave cross sections. We only take the dominant contributions into account. When $m_\chi < m_{Z'}$, the dominant annihilation channel is through the s -channel process $\bar{\chi}\chi \rightarrow f\bar{f}$, where $f = \mu, \nu_\mu, b, t$:

⁷For example, one assumption that $y'_\mu = 0$ results in disappearing 1-loop contributions to the muon $g-2$ from new scalars.

$$\langle\sigma v\rangle_{s\text{-channel}}(\bar{\chi}\chi \rightarrow f\bar{f}) = \frac{g^2 g_{Z'ff}^2}{2\pi} \frac{2m_\chi^2 + m_f^2}{4m_\chi^2 - m_{Z'}^2} \sqrt{1 - \frac{m_f^2}{m_\chi^2}}, \quad (48)$$

where $g_{Z'ff}$ denotes the Z' coupling to the fermion f . When $m_\chi > m_{Z'}$, additional t -channel $\bar{\chi}\chi \rightarrow Z'Z' \rightarrow f\bar{f}f\bar{f}$ appears. The annihilation cross section is

$$\langle\sigma v\rangle_{t\text{-channel}}(\bar{\chi}\chi \rightarrow Z'Z') = \frac{g'^4}{4\pi} \frac{m_\chi^2}{(2m_\chi^2 - m_{Z'}^2)^2} \left(1 - \frac{m_{Z'}^2}{m_\chi^2}\right)^{3/2}. \quad (49)$$

The thermal relic density is written as

$$\Omega_{\text{DM}} h^2 = \frac{m_\chi n_\chi}{\rho_{\text{crit}}/h^2} = \frac{s_0 h^2}{\sqrt{\frac{\pi}{45}} \rho_{\text{crit}}} \frac{1}{M_{\text{pl}} \sqrt{g_{\text{eff}} \mathcal{I}(x_f)}}, \quad (50)$$

where s_0 is the entropy density of the present universe, ρ_{crit} the critical density, n_χ the dark matter number, h the hubble parameter, g_{eff} the effective degree of freedom during freeze-out, and $\mathcal{I}(x_f) = \frac{a+3b/x_f}{x_f}$, with $x_f = m_\chi/T_f$ and T_f the temperature during freeze-out. There are also constraints from the direct detection measurements. However, we expect the limit is not so tight for the light Z' , because the size of the gauge coupling g' is typically smaller than 0.1. Furthermore, in our model the dark matter only dominantly couples to the third generation quarks inside nucleons. This kinds of scenarios have been studied in Ref. [100]. From their study, we find that for the light Z' , if $g' \sim 0.05$, the direct detection constraints could be escaped. We will perform detailed study on the dark matter sector in future.

VI. RESULTS

We first consider the effective operators for the R_K anomaly in terms of $X_q \equiv v_\Phi Y_q / \sqrt{2} m_Q$ ($q = b, s$) and $X_\mu \equiv v_\Phi Y_\mu / \sqrt{2} m_L$. With these expressions, the Eq. (37) is described by

$$C_{9,\text{new}} (= -C_{10,\text{new}}) = -\frac{1}{2v_\Phi^2 |C_{\text{SM}}^{bs}|} X_b X_s X_\mu^2. \quad (51)$$

This expression consists of only relevant model parameters through the combinations of X_q and X_μ that are related to the mixing angles in the quark and lepton sectors. It should be noted that the value of $C_{9,\text{new}}$ in Table II is positive, therefore either X_b or X_s should be negative to compensate for the minus sign in the above expression. The left plot in Fig. 12 shows 1σ and 2σ regions with green and light green bands in $(-X_b X_s, X_\mu^2)$ plane with the fixed value of

FIG. 12. [Left] 1σ and 2σ regions of $C_{9,\text{new}} (= -C_{10,\text{new}})$ are drawn with green and light green bands, respectively. v_Φ is fixed at 700 GeV. The best fitted value for the R_K anomaly is indicated by the green line, and the red line represents $-X_b X_s = 2.3 \times 10^{-3}$ as a benchmark points. The blue region is excluded by the $B_s - \bar{B}_s$ mixing. [Right] The size of $C_{9,\text{new}}$ as a function of Y_μ with the fixed value $-X_b X_s = 2.3 \times 10^{-3}$. We take $m_Q = m_L = 1.4$ TeV.

$v_\Phi = 700$ GeV. The best fitted value is drawn by a green line, and the blue region is excluded by $B_s - \bar{B}_s$ mixing. We also consider the constraint from the $B \rightarrow K \nu \nu$ process, however, the present parameter region can evade the restriction. It is seen that $X_b X_s$ is tightly constrained by the $B_s - \bar{B}_s$ mixing and the allowed size is roughly $O(10^{-3})$ at most. And, it follows that a somewhat large value of $X_\mu^2 \sim O(0.1)$ is necessary to achieve the required Wilson coefficients for the R_K . Such a large value can be obtained by taking a large Y_μ or a small m_L . Here, we fix the mass m_L at 1.4 TeV and see how large Yukawa coupling, Y_μ , is needed at one benchmark region of $-X_b X_s = 2.3 \times 10^{-3}$. The reason we choose this representative value is to make sure that the value of Y_μ is not too large, which may spoil our vev insertion approximation, and simultaneously avoid the restriction from the $B_s - \bar{B}_s$ mixing constraint. The region is indicated by a red line in the left panel in Fig. 12, and the benchmark value results in $Y_b Y_s = -0.0184$ if $m_Q = m_L = 1.4$ TeV. Since our assumption is $|Y_b| \gg |Y_s|$, we simply take $Y_b = -1$ and $Y_s = 0.0184$. In the right plot of Fig. 12, the size of $C_{9,\text{new}}$ is shown as a function of Y_μ . Here, instead of the mass-insertion method, we use g_{sb}^L and $g_{\mu\mu}^L$ obtained after diagonalization of the mass matrices. All bands and line in the figure correspond to those in the left one. The 2σ region requires $1.2 \leq Y_\mu \leq 1.9$, and the best fitted value stays around $Y_\mu \sim 1.5$. The values of $Y_b = -1, Y_s = 0.0184$ and $Y_\mu = 1.5$ yield $g_{sb}^L/g' = -2.17 \times 10^{-3}$ and $g_{\mu\mu}^L/g' = 0.22$, which leads to $C_{9,\text{new}} = 0.6$. The current size of $g_{\mu\mu}^L/m_{Z'}$ gives the cross section of the neutrino

trident production in Eq. (39), $\sigma_{\text{SM+NP}}/\sigma_{\text{SM}} = 1.0$, which is within the 1σ error. Later, we also use these Yukawa couplings as the benchmark values for the discussion of the R_D anomaly.

Figure 13 shows the experimental data and theoretical predictions on the B decay branching ratios in different energy bins for $\text{dBR}(B^0 \rightarrow K^0 \mu^+ \mu^-)/dq^2$ (upper left), $\text{dBR}(B^0 \rightarrow K^{*0} \mu^+ \mu^-)/dq^2$ (upper right), $\text{dBR}(B^\pm \rightarrow K^\pm \mu^+ \mu^-)/dq^2$ (lower left), $\text{dBR}(B^+ \rightarrow K^{*+} \mu^+ \mu^-)/dq^2$ (lower right). The experimental results obtained by CDF [75], LHCb [1,77], CMS [78,79] are plotted. The purple bands correspond to the SM predictions with uncertainties obtained by varying the input parameters with the Gaussian distribution. The green bands represent the best fitted value of $C_{9,\text{new}}$, which is indicated by the green line in Fig. 12, with the uncertainties obtained by the same way as the purple bands. The width of the purple and green bands are not the same, and their ratios are roughly proportional to $|C_{\text{SM}}^{bs}/(C_{\text{SM}}^{bs} + C_i)|^2$ ($i = 9, 10$). It is found that our model provides a better fit on the data than the SM.

Fig 14 represents the value of $R_{K^{*0}}$ in low q^2 bins. As in Fig. 13, the purple and green bands are the theoretical predictions of the SM and the current setup with their uncertainties. The dots indicate the experimental measured values in Ref. [1]. It is seen that, in the lower q^2 bin of $[0.045, 1.1]$, the best fitted value is away from the experimental data with the 1σ uncertainty, though it is still better than the SM prediction.

Finally, we discuss the possibility of explaining the R_D anomaly using the parameter space which could explain the

FIG. 13. The B decay branching ratios $\text{dBR}(B^0 \rightarrow K^0 \mu^+ \mu^-)/dq^2$ (upper left), $\text{dBR}(B^0 \rightarrow K^{*0} \mu^+ \mu^-)/dq^2$ (upper right), $\text{dBR}(B^\pm \rightarrow K^\pm \mu^+ \mu^-)/dq^2$ (lower left), $\text{dBR}(B^+ \rightarrow K^{*+} \mu^+ \mu^-)/dq^2$ (lower right), $R_{\mu\ell}(B^+ \rightarrow K^{*+} \ell^+ \ell^-)$ (bottom) for different energy bins. The experimental data are expressed by dots with uncertainties. The purple and green bands represent the theoretical predictions of the SM and our model with the best fitted value of $C_{9,\text{new}}$, respectively. The theoretical uncertainties are calculated by varying the input parameters with the Gaussian uncertainties.

R_K anomaly. Once Y_b and Y_s are determined, Y_c is also fixed due to the $SU(2)$ symmetry. Our benchmark values with $Y_b = -1$ and $Y_s = 0.018$ yield $Y_c = -0.02$ where the Y_b part dominantly gives a contribution to it since $(V_{\text{CKM}})_{cb} Y_b > Y_s$. As expected, the smallness lead to irrelevance of C_S unless other model parameters compensate for it.

The correlation between C'_S and $C_{9,\text{new}}$ is described by

$$C'_S = -\frac{4C_{\text{SM}}^{bs} y'_c y'_\tau m_Q m_L}{C_{\text{SM}}^{bc} Y_s Y_\mu m_{H'}^2} C_{9,\text{new}} \equiv -\frac{4C_{\text{SM}}^{bs}}{C_{\text{SM}}^{bc}} A C_{9,\text{new}}. \quad (52)$$

Given that both C'_S and C_9 are $O(0.1)$, the factor of C_9 should be $O(1)$. Since $4C_{\text{SM}}^{bs}/C_{\text{SM}}^{bc} \sim 2.4 \times 10^{-3}$, the

coefficient A needs to be the same order of magnitude for the compensation. In Fig. 15, we see the correlation changing the coefficient A . The green bands represent the 1 and 2 sigma regions of C_9 , and the pink bands are those of C'_S . Here, we assume that A is positive and impose only $\text{Br}(B_c \rightarrow \tau^- \bar{\nu}) < 30\%$ on C'_S . The black lines correspond to $A = 200, 400, 600$, and 800 from top to bottom. The overlapped region between two 1σ bands implies that $400 \lesssim A \lesssim 600$. Using the benchmark values obtained in Fig. 12, we find that

$$A = \frac{y'_c y'_\tau}{0.018 \times 1.5} \left(\frac{1.4 \times 10^3}{m_{H'}} \right)^2. \quad (53)$$

Therefore, the desirable size of A can be acquired if $y'_c y'_\tau \sim O(1)$ and $m_{H'} \sim O(100)$ GeV.

FIG. 14. $R_{K^{*0}}$ for the low q^2 bins. The dots and bands are the same definitions as in Fig. 13.

In Fig. 16, we show the 1σ and 2σ regions of C'_S in (y'_c, y'_τ) plane with the fixed charged scalar mass, $m_{H^\pm} = 300$ GeV. The model parameters that correlate with the R_K anomaly are fixed at $Y_b = -1$, $Y_s = 0.018$, $Y_\mu = 1.5$, $v_\phi = 700$ GeV and $m_Q = m_L = 1.4$ TeV. This setup yields that $A \sim 800$, therefore, we obtain that $y'_c y'_\tau \sim O(0.1)$. As long as we focus on only C'_S , the size of y'_b is not predicted. Therefore, our current setup does not affect $b \rightarrow s\gamma$ process. However, it is found that, even if we take $y'_c = 1$ and $y'_b = 1$, $C_7(\mu_b)$ roughly becomes 8.0×10^{-4} which is consistent with the current observed value described in Sec. V.

FIG. 15. The correlation between $C_{9,\text{new}}$ and C'_S . The green and pink bands are 1 and 2 sigma regions listed in Table II. The black lines correspond to $A = 200, 400, 600$, and 800 from top to bottom.

FIG. 16. 1 and 2σ regions for C'_S . It is taken that $m_{H^\pm} = 300$ GeV, and the benchmark parameters for the R_K anomaly: $v_\phi = 700$ GeV, $m_Q = m_L = 1.4$ TeV, $Y_b = -1$, $Y_s = 0.018$ and $Y_\mu = 1.5$.

VII. CONCLUSION

We have studied a hidden gauged $U(1)_H$ extension of the SM including vectorlike fermion (and scalar) mediators and investigated possibilities of explanations for the existing R_K (and R_D) anomalies in B physics. The vectorlike fermions have the same SM charges as the SM doublet fermions, and they mix with the left-handed SM fermions. The new scalar is an inert doublet, which possesses the Yukawa interaction between the left-handed vectorlike fermions and the right-handed SM fermions. While the vectorlike fermions play a role in inducing $b \rightarrow s\mu\mu$ transition through the Z' exchange, the inert doublet yields significant charged scalar couplings for $b \rightarrow c\tau\nu$ transition.

In order to determine desired sizes of the Wilson coefficients, we utilized the FLAVIO package and performed a global Bayesian fit including several observables, $B^{0,\pm} \rightarrow K^{*0,\pm}\mu\mu$, $B_s \rightarrow \phi\mu\mu$, $B_s \rightarrow X_s\mu\mu$ and $B_s \rightarrow \mu\mu$. At the same time, we also considered various flavor constraints, $B_s - \bar{B}_s$ mixing, $B \rightarrow K\nu\bar{\nu}$, neutrino trident productions, $b \rightarrow s\gamma$, and $B_c^- \rightarrow \tau^-\bar{\nu}$. It is found that the absolute values of $C_9 (= -C_{10})$ and C'_S are roughly required to be $O(0.5)$ for explaining the measured values of R_K and R_D . However, the favored range of C'_S are almost excluded by $B_c^- \rightarrow \tau^-\bar{\nu}$ if we impose the severe limit on $\text{BR}(B_c^- \rightarrow \tau^-\bar{\nu}) < 10\%$.

The desirable sizes of C_9 and C_{10} can be achieved by somewhat large Y_μ , which describes the mixing parameter between the vectorlike and SM leptons, since those in the quark sector Y_b and Y_s are highly constrained by the $B_s - \bar{B}_s$ mixing. Our benchmark point with $m_Q = m_L = 1.4$ TeV, $v_\phi = 700$ GeV, and $Y_b Y_s = -0.018$ needs $Y_\mu \sim 1.5$ for the

successful scenario of the R_K anomaly. This benchmark leads to suppressed C_S due to tiny Y_c , and negative C'_S is obtained by taking $Y_b < 0$. The possible region of C'_S requires $O(0.1)$ couplings of y'_c and y'_τ .

Since this model contains new particles, there are promising collider and astrophysical signatures. Current LHC searches on $pp \rightarrow \mu\bar{\mu}$ and $pp \rightarrow 4\mu$ put strong constraints on the Z' masses, although the light region of $M_{Z'} < 100$ GeV is still allowed. Furthermore, since the $U(1)$ gauge symmetry is hidden, it is likely that other hidden particles also exist, such as dark matter candidate, in the hidden sector. Although the hidden sector particles do not contribute to the B physics observables, they are able to contribute to the dark matter relic density and possible indirect detections, and would relax

the constraints on the heavy Z' . On the other hand, the charged Higgs is not tightly constrained due to its Yukawa texture. We anticipate that the future collider searches will explore larger mass regions for the Z' and charged Higgs, and these collider signatures and potential dark matter signatures could be able to validate this model.

ACKNOWLEDGMENTS

We are grateful to our referees for valuable comments. H. L. thanks David Straub for helpful discussions about FLAVIO and Aniket Joglekar for valuable discussion on the diagonalization of fermion mass matrix. J. H. Y. and K. F. are supported by DOE Grant No. DE-SC0011095.

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