

Implications for new physics from a novel puzzle in $\bar{B}_{(s)}^0 \rightarrow D_{(s)}^{(*)+} \{\pi^-, K^-\}$ decays

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Recently, the standard model predictions for the B -meson hadronic decays, $\bar{B}^0 \rightarrow D^{(*)+} K^-$ and $\bar{B}_s^0 \rightarrow D_s^{(*)+} \pi^-$, have been updated based on the QCD factorization approach. This improvement sheds light on a novel puzzle in the B -meson hadronic decays: there are mild but universal tensions between data and the predicted branching ratios. Assuming the higher-order QCD corrections are not huge enough to solve the tension, we examine several new physics interpretations of this puzzle. We find that the tension can be partially explained by a left-handed W' model, which can be compatible with other flavor observables and collider bounds.

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I. INTRODUCTION

To test the standard model (SM) and search for physics beyond the SM, precision measurements of meson decays, especially B -meson decays, have been considerably investigated over the past 30 years. In the meantime, the experimental uncertainty has been surprisingly reduced by experimentalists. On the other hand, theorists have played an equally important role: several approaches that can evaluate the QCD corrections have been invented, and the SM predictions have been sharpened.

Very recently, SM predictions for several B -meson hadronic decays are improved by Ref. [1]:

$$\mathcal{B}(\bar{B}^0 \rightarrow D^+ K^-)_{\text{SM}}^{\text{exp}} = \begin{cases} (1.86 \pm 0.20) \times 10^{-4}, \\ (3.26 \pm 0.15) \times 10^{-4}, \end{cases} \quad (1)$$

$$\mathcal{B}(\bar{B}^0 \rightarrow D^{*+} K^-)_{\text{SM}}^{\text{exp}} = \begin{cases} (2.12 \pm 0.15) \times 10^{-4}, \\ (3.27^{+0.39}_{-0.34}) \times 10^{-4}, \end{cases} \quad (2)$$

$$\mathcal{B}(\bar{B}_s^0 \rightarrow D_s^+ \pi^-)_{\text{SM}}^{\text{exp}} = \begin{cases} (3.00 \pm 0.23) \times 10^{-3}, \\ (4.42 \pm 0.21) \times 10^{-3}, \end{cases} \quad (3)$$

$$\mathcal{B}(\bar{B}_s^0 \rightarrow D_s^{*+} \pi^-)_{\text{SM}}^{\text{exp}} = \begin{cases} (2.0 \pm 0.5) \times 10^{-3}, \\ (4.3^{+0.9}_{-0.8}) \times 10^{-3}, \end{cases} \quad (4)$$

where the upper numbers are the PDG averages of the experimental data [2], while the lower ones are the SM expectation values [1]. These SM predictions are obtained by the QCD factorization (QCDF) [3–5] at leading power in Λ_{QCD}/m_b , where the Wilson coefficients at next-to-next-to-leading logarithmic accuracy are used [6]. Compared to the previous estimations [7], the theoretical uncertainties are significantly reduced thanks to recent developments in the $\bar{B}_{(s)}^0 \rightarrow D_{(s)}^{(*)}$ form factors including order $\mathcal{O}(1/m_c^2)$ corrections within the framework of the heavy-quark expansion [8–11].

These hadronic channels are theoretically clean due to the absence of penguin and annihilation topologies. Furthermore, resultant amplitudes are dominated by the color-favored tree topology.

Above SM predictions deviate from the data at 5.6σ ($D^+ K^-$), 3.1σ ($D^{*+} K^-$), 4.6σ ($D_s^+ \pi^-$), and 2.4σ levels ($D_s^{*+} \pi^-$), respectively. Surprisingly, all deviations are in the same direction and similar size. Note that $\mathcal{B}(\bar{B}^0 \rightarrow D^+ \pi^-)_{\text{SM}} = (3.93^{+0.43}_{-0.42}) \times 10^{-3}$ and $\mathcal{B}(\bar{B}^0 \rightarrow D^{*+} \pi^-)_{\text{SM}} = (3.45^{+0.53}_{-0.50}) \times 10^{-3}$, which are evaluated in Ref. [7], also deviate from the data, $\mathcal{B}(\bar{B}^0 \rightarrow D^+ \pi^-)^{\text{exp}} = (2.52 \pm 0.13) \times 10^{-3}$ and $\mathcal{B}(\bar{B}^0 \rightarrow D^{*+} \pi^-)^{\text{exp}} = (2.74 \pm 0.13) \times 10^{-3}$ [2] at the 3.2σ and 1.4σ levels, respectively.

Within the SM, there are two possibilities that these tensions are alleviated. The first possibility is an input value of $|V_{cb}|$. For $|V_{cb}|$, the authors of Ref. [1] use an average of

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the inclusive and exclusive determinations in the B -meson semileptonic decays: $|V_{cb}| = (41.1 \pm 0.5) \times 10^{-3}$ [9,10]. If one adopts the exclusive $|V_{cb}|$, $|V_{cb}| = (39.25 \pm 0.56) \times 10^{-3}$ [12], amplitudes of the above processes are uniformly reduced by 4.5%. Note that the exclusive $|V_{cb}|$, however, produces an additional 4.2σ level tension in ε_K [13]. See also [11] for a more recent determination of the exclusive $|V_{cb}|$ using the full angular distribution data.

Another possibility is higher-order QCD corrections. The next-to-leading power and next-to-next-to-leading power corrections to the QCDF amplitudes are also estimated by the same authors [1], and the sizes of those corrections to the amplitudes are evaluated as $\mathcal{O}(1)\%$.

The above puzzled situation could be resolved by introducing new physics contributions to $b \rightarrow c\bar{u}q$ transitions, where $q = d$ and s . Furthermore, it is shown that all ratios between these branching fractions are consistent with data [1]. It clearly implies that the new physics effects should be universal in $b \rightarrow c\bar{u}q$ transitions. Therefore, the following questions are interesting: whether such a new physics is still allowed by the other flavor constraints and by the hadron collider constraints, and how much the tensions can be alleviated by a valid new physics model. Below we will refer to this puzzle as $b \rightarrow c\bar{u}q$ anomaly. In this paper, we examine several new physics scenarios to explain the $b \rightarrow c\bar{u}q$ anomaly.

II. FRAMEWORK

We consider the following effective Lagrangian to investigate new physics contributions to $b \rightarrow c\bar{u}q$ processes:

$$\mathcal{L} = -\frac{4G_F}{\sqrt{2}} \sum_q V_{cb} V_{uq}^* \sum_{i=1,2} C_i^q(\mu) \mathcal{Q}_i^q(\mu), \quad (5)$$

with the left-handed current-current operators in the CMM basis [14,15],

$$\mathcal{Q}_1^q = (\bar{c}_L \gamma^\mu T^a b_L)(\bar{q}_L \gamma_\mu T^a u_L), \quad (6)$$

$$\mathcal{Q}_2^q = (\bar{c}_L \gamma^\mu b_L)(\bar{q}_L \gamma_\mu u_L), \quad (7)$$

where $q = d, s$. T^a is the $SU(3)_C$ generator, and V is the Cabibbo-Kobayashi-Maskawa matrix [16,17]. In our analysis, we refrain from adding operators that are absent in the SM, e.g., $(\bar{c}_L b_R)(\bar{q}_L u_R)$. We will discuss this possibility in the last section.

New physics contributions to the Wilson coefficients, $C_1^{q,\text{NP}}$ and $C_2^{q,\text{NP}}$, become involved at the new physics scale Λ . These values are modified by the renormalization-group (RG) evolution from Λ down to the hadronic scale m_b . The leading-order (LO) QCD RG evolution is summarized in the Supplemental Material [18]. For instance, when $\Lambda = 1$ TeV, we obtain an evolution matrix as

$$\begin{pmatrix} C_1^{\text{NP}}(m_b) \\ C_2^{\text{NP}}(m_b) \end{pmatrix} = \begin{pmatrix} 1.36 & -0.87 \\ -0.19 & 1.07 \end{pmatrix} \begin{pmatrix} C_1^{\text{NP}}(1 \text{ TeV}) \\ C_2^{\text{NP}}(1 \text{ TeV}) \end{pmatrix}. \quad (8)$$

It is found that a universal destructive shift in the SM contributions is favored in the $b \rightarrow c\bar{u}q$ anomaly [1]. The preferred size is $\sim -17\%$, which corresponds to $C_2^{d,\text{NP}} = C_2^{s,\text{NP}} = C_2^{\text{NP}}$ and

$$\frac{C_2^{\text{NP}}(m_b)}{C_2^{\text{SM}}(m_b)} = -0.17 \pm 0.03. \quad (9)$$

It is checked that such a new physics contribution is compatible with data of the total decay rates of the B -mesons, τ_{B_s}/τ_{B_d} , and a_d^{fs} [1,19–21]. Another potentially strong constraint comes from the kaon hadronic decays ($s \rightarrow u\bar{u}d$). The CP -conserving parts of the isospin amplitudes, $A_I = \langle (\pi\pi)_I | \mathcal{H}_{\text{eff}}^{|\Delta S|=1} | K \rangle$ for $I = 0, 2$, have been measured very precisely through all $K \rightarrow \pi\pi$ data [22,23]

$$\text{Re}A_0^{\text{exp}} = (3.3201 \pm 0.0018) \times 10^{-7} \text{ GeV}, \quad (10)$$

$$\text{Re}A_2^{\text{exp}} = (1.4787 \pm 0.0031) \times 10^{-8} \text{ GeV}. \quad (11)$$

On the other hand, these theoretical predictions are

$$\text{Re}A_0^{\text{SM}} = (2.99 \pm 0.67) \times 10^{-7} \text{ GeV}, \quad (12)$$

$$\text{Re}A_2^{\text{SM}} = (1.50 \pm 0.15) \times 10^{-8} \text{ GeV}, \quad (13)$$

where the hadronic matrix elements are calculated by the lattice QCD simulations [23–27]. Although the A_2 amplitude is more sensitive to new physics than A_0 , we find that a $\pm 20\%$ new physics contribution to the $s \rightarrow u\bar{u}d$ amplitude could be compatible with the data.

III. MINIMAL FLAVOR VIOLATION

First, we study the most simple possibility for new physics scenario: minimal flavor violation (MFV) hypothesis [28,29]. The detailed calculations for this section can be found in the Supplemental Material [18].

We examine a dimension-six operator, $\mathcal{L} = 1/(2\Lambda^2) (\bar{Q}_L \gamma^\mu Q_L)^2$, whose flavor off-diagonal components are controlled by the quark Yukawa. In the quark mass-diagonal basis, this operator produces $C_2^{q,\text{MFV}}(\Lambda) \sim -1/(2\sqrt{2}G_F\Lambda^2)$. Then, the $b \rightarrow c\bar{u}q$ anomaly in Eq. (9) suggests $\Lambda \lesssim 0.49$ TeV.

Among the various flavor and collider constraints, a nonresonant dijet angular distribution search in the LHC gives the most stringent constraint on this scenario. The result is reported by the ATLAS collaboration at $\sqrt{s} = 13$ TeV with $\int dt \mathcal{L} = 37 \text{ fb}^{-1}$ [30]. We interpret the result

and obtain a 95% C.L. exclusion limit as $\Lambda < 3.7$ TeV, which excludes the suggested $\Lambda \sim 0.49$ TeV. From this collider constraint, we obtain a bound

$$\frac{C_2^{\text{MFV}}(m_b)}{C_2^{\text{SM}}(m_b)} \gtrsim -0.002. \quad (14)$$

Hence, this scenario never explains the $b \rightarrow c\bar{u}q$ anomaly.

IV. SU(2) \times SU(2) \times U(1) MODEL

Next, we consider a new physics model that can produce a more convoluted flavor structure. An extended electro-weak gauge group $\text{SU}(2)_1 \times \text{SU}(2)_2 \times \text{U}(1)_Y$ with heavy vectorlike fermions produces heavy gauge bosons, W'^{\pm} and Z' , interacting with the left-handed SM fermions with a nontrivial flavor structure [31–35]. These flavor structures are controlled by the number of generations of the vectorlike fermions (n_{VF}) and mixings between the SM fermions and vectorlike fermions.

The heavy gauge boson interactions are [35]

$$\begin{aligned} \mathcal{L} = & + \frac{g_{ij}}{2} Z'_\mu \bar{d}_L^i \gamma^\mu d_L^j - \frac{(VgV^\dagger)_{ij}}{2} Z'_\mu \bar{u}_L^i \gamma^\mu u_L^j \\ & - \frac{(Vg)_{ij}}{\sqrt{2}} W'^+_\mu \bar{u}_L^i \gamma^\mu d_L^j + \text{H.c.}, \end{aligned} \quad (15)$$

where u_L, d_L are the mass eigenstates, and a coupling g_{ij} is defined in the d_L basis. In the following, we will take $M_{W'} = M_{Z'} = M_V$ for simplicity. By integrating out W'^{\pm} , new physics contribution $C_2^{q,W'}$ is obtained as

$$C_2^{q,W'}(M_V) = \frac{1}{4\sqrt{2}G_F M_V^2} \frac{(Vg)_{23}(Vg)_{1q}^*}{V_{cb}V_{uq}^*}. \quad (16)$$

In order to generate an uniform shift in both $b \rightarrow c\bar{u}d$ and $b \rightarrow c\bar{u}s$, a SM-like flavor structure in $(Vg)_{1q}$ is required, and hence g_{11} should be nonzero. When only g_{11} is a nonzero entry in g_{ij} , a dangerous $\bar{c}uZ'$ flavor-changing neutral current is generated and it is severely constrained by the D -meson mixing as $|g_{11}|/M_V < \mathcal{O}(10^{-2})$ (TeV) $^{-1}$ [36]. To evade this bound, we follow the $\text{U}(2)^3$ flavor symmetry [37,38] and take $g_{11} = g_{22}$ in g_{ij} in the following analyses. Then the bound from the D -meson mixing is significantly relaxed as $|g_{11}|/M_V \lesssim 16$ (TeV) $^{-1}$.

Another flavor constraint comes from the $K \rightarrow \pi\pi$ data. By permitting a $\pm 20\%$ new physics contribution to the Wilson coefficient of $(\bar{u}_L \gamma^\mu d_L)(\bar{s}_L \gamma_\mu u_L)$ [see Eq. (13)], we obtain

$$|g_{11}|/M_V \lesssim 3.6 \text{ (TeV)}^{-1}. \quad (17)$$

Note that many types of diagrams contribute to $K \rightarrow \pi\pi$ decays, and nonperturbative QCD plays an essential role there. Therefore, this bound is a just reference value.

In addition to g_{11} , another nonzero entry of g_{33} or g_{23} is necessary to produce $C_2^{q,W'}$. Therefore, we consider the following flavor texture:

$$g_{ij} = \begin{pmatrix} g_{11} & 0 & 0 \\ 0 & g_{11} & g_{23} \\ 0 & g_{23} & g_{33} \end{pmatrix}, \quad (18)$$

and will discuss several scenarios in detail. We assume g_{ij} is real for simplicity. Note that when g_{11} is $\mathcal{O}(1)$, production cross sections of the heavy gauge bosons become considerably large in the hadron collider, and hence we will mostly discuss the LHC constraints in each subsection. To evade surveying a dedicated collider constraint for low-mass region where the constraint would be more stringent, the mass range $M_V > 1$ TeV is considered in our analysis.

A. Scenario 1: g_{11} and g_{33}

In this subsection, we take $g_{23} = 0$ and consider a scenario of $g_{ij} = \text{diag}(g_{11}, g_{11}, g_{33})$. Such a flavor structure can be obtained from $n_{\text{VF}} = 1$. In this case, $(Vg)_{23}$ in Eq. (16) comes from $V_{cb}g_{33}$. Since one has a factor of V_{cb} just as the SM, $\sqrt{|g_{11}g_{33}|}/M_V$ should be larger than $\mathcal{O}(1)$ TeV $^{-1}$ to generate new physics contributions to $b \rightarrow c\bar{u}q$ processes (see previous section). Furthermore, a relative sign between g_{11} and g_{33} must be negative to produce the destructive interference with the SM in the $b \rightarrow c\bar{u}q$ decays. A requirement of the $b \rightarrow c\bar{u}q$ anomaly within 2σ level leads to

$$2.6 \text{ (TeV)}^{-1} \lesssim \sqrt{|g_{11}g_{33}|}/M_V \lesssim 3.8 \text{ (TeV)}^{-1}. \quad (19)$$

Therefore, large couplings are necessary in this scenario.

First, let us examine the constraint from the B_s -meson mass difference (ΔM_s). In this scenario, the dominant contribution comes from a W - W' box diagram. We observed that the GIM mechanism still works in this flavor structure, and obtain a simple formula for the W - W' box contribution to ΔM_s ,

$$\frac{\Delta M_s^{W'}}{\Delta M_s^{\text{SM}}} \simeq \eta^2 \frac{2g_{11}g_{33}f'(x_t, x_V)}{g_W^2 f(x_t)}, \quad (20)$$

with $\eta = \alpha_s(M_V)/\alpha_s(m_W)$, $x_t = m_t^2/m_W^2$ and $x_V = M_V^2/m_W^2$, and g_W is the weak coupling. The loop functions are defined in the Supplemental Material [18]. We also have the same shift in B_d -meson mixing, but it is less constrained because of its large theoretical uncertainty. By imposing that the new physics contribution is within 2σ uncertainty of ΔM_s^{SM} [39,40], we obtain

$$\sqrt{|g_{11}g_{33}|}/M_V \lesssim 1.7 \text{ (TeV)}^{-1}. \quad (21)$$

Although ΔM_s bound is incompatible with the $b \rightarrow c\bar{u}q$ anomaly in Eq. (19), we want to know how much this scenario can alleviate the puzzle.

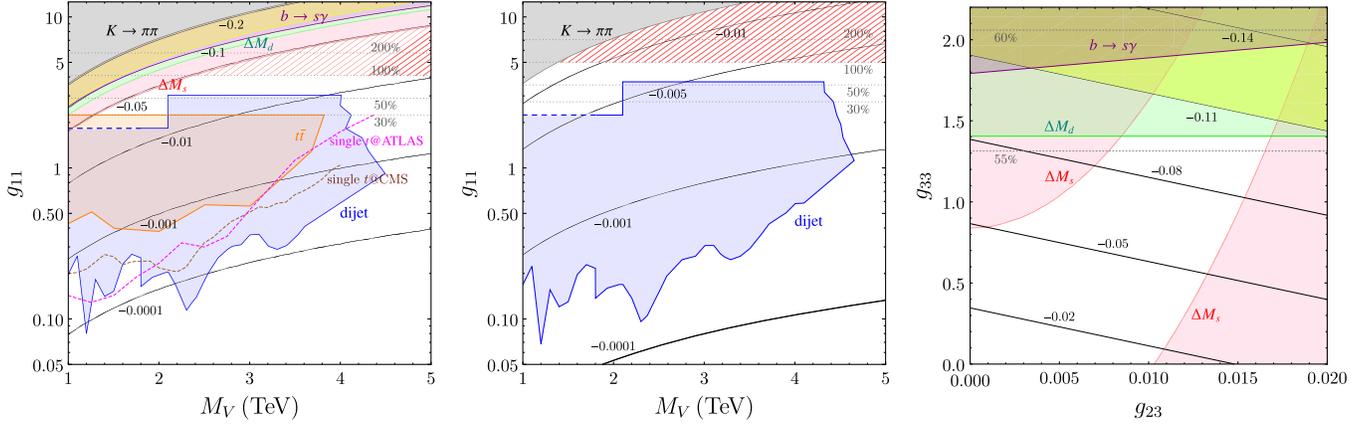


FIG. 1. Contours of $C_2^{\text{NP}}(m_b)/C_2^{\text{SM}}(m_b)$ are presented. The puzzle can be explained at 2σ level in the yellow bands. The blue and orange shaded regions are excluded by the dijet [41–43] and $t\bar{t}$ searches [44,45] at 95% C.L., respectively. The regions above the dashed lines are excluded by the single t searches in the NWA (see text) [48,49]. Furthermore, the gray, red, green, and purple shaded regions are constrained by $K \rightarrow \pi\pi$, ΔM_s , ΔM_d , and $b \rightarrow s\gamma$, respectively. The dotted line indicates Γ_V/m_V and the red-hatched regions represent $\Gamma_V/m_V > 100\%$. Left: scenario 1. We take $g_{33} = -g_{11}$. Middle: scenario 2. We take $g_{23} = -0.01(M_V/\text{TeV})$. Right: scenario 3. We take $M_V = 1 \text{ TeV}$ and $g_{11} = -3.6$.

Next, we consider constraints from resonant productions of the heavy gauge bosons at the LHC. When g_{11} and g_{33} entries are nonzero, Z' is produced via $pp \rightarrow q\bar{q} \rightarrow Z'$ and also $pp \rightarrow b\bar{b} \rightarrow Z'$, while W'^{\pm} is produced through $pp \rightarrow q\bar{q}' \rightarrow W'^{\pm}$ processes. When $M_V \gg m_t$, the decay width of those particle is approximately given as,

$$\Gamma_{V=W',Z'} \simeq \frac{2|g_{11}|^2 + |g_{33}|^2}{16\pi} m_V. \quad (22)$$

We find that relevant collider bounds come from dijet and $t\bar{t}$ searches. The former provides the relevant bound for $|g_{11}| \gg |g_{33}|$, while the latter for $|g_{11}| \lesssim |g_{33}|$.

Currently, both ATLAS and CMS collaborations reported upper limits on the heavy dijet resonance cross section using the data of $\sim 140 \text{ fb}^{-1}$ [41,42]. Since $\mathcal{O}(1)$ couplings are necessary to relax the tension, the decay width can be not small. Therefore, we adopt width-dependent limits on the cross section times the dijet branching ratio. The broader the width is, the weaker the limits become because a characteristic resonance peak is diluted. The search is robust up to $\Gamma_V/M_V = 20\%$ for 1.8–2.1 TeV, and up to $\Gamma_V/m_V = 55\%$ for the heavier region [42]. For the mass range of 1–1.8 TeV, we use an upper limit in Ref. [43], where the narrow width approximation (NWA) is used. As for the heavy $t\bar{t}$ resonance search, CMS reported the width-dependent limit using the data of 36 fb^{-1} up to $\Gamma_V/M_V = 30\%$ [44], while ATLAS reported the result using the data of 139 fb^{-1} in the NWA [45].

We obtained the production cross section of Z' and W'^{\pm} by rescaling the result in Refs. [42,46], where $\sigma(pp \rightarrow q\bar{q}' \rightarrow W'^+) + \sigma(pp \rightarrow q\bar{q}' \rightarrow W'^-) \simeq 2\sigma(pp \rightarrow q\bar{q} \rightarrow Z')$ is used [47]. The excluded regions from the dijet and $t\bar{t}$ searches are shown as the blue and orange shaded regions in Fig. 1(left), respectively.

We also show constraints from the single t searches by using the data of $\sim 36 \text{ fb}^{-1}$ of CMS [48] and ATLAS [49]: the regions above the dashed lines in Fig. 1(left) are excluded. Note that both analyses assume the narrow resonance, and no study exists for broad resonances.

Taking a conservative position, regions above the plateaus of the shaded areas can not be excluded, where the corresponding Γ_V/M_V exceeds the maximum width shown in each experimental result: $\Gamma_V/M_V > 30\%$ in the $t\bar{t}$ search, and $\Gamma_V/M_V > 55\%$ for 2.1–5 TeV and $\Gamma_V/M_V > 20\%$ for 1.8–2.1 TeV in the dijet search. The horizontal blue dashed lines are extrapolations obtained by assuming the analysis of Ref. [43] is valid up to $\Gamma_V/M_V = 20\%$, and should be taken with more care. We note that limits from the dijet angular distribution data, which are not considered here, would also depend on the width-mass ratio and only contact interaction models are investigated [30,50]. Further dedicated analysis would be necessary to exclude such a broad width region.

The red-hatched regions represent $\Gamma_V > m_V$, where a particle picture is no longer valid and one could not discuss any conclusive prediction.

Note that both our study and above experimental analyses have considered only the s -channel productions of W' and Z' , although there are several t -channel contributions. Since the t -channel process does not show a resonant nature, and there is a huge QCD t -channel background in the dijet production, we suppose that inclusion of the t -channel processes in the signal could not amplify the signal-to-noise ratio in the resonance searches. Such t -channel contributions, which are insensitive to the width, would be potentially accessible in the angular distribution search.

As long as we allow the broad width scenario, we find that the bound from ΔM_s in Eq. (21) determines the maximal deviation of C_2^W/C_2^{SM} , which is independent of

the ratio of g_{11} and g_{33} . For these reasons, we conclude $C_2^{W'}/C_2^{\text{SM}} \gtrsim -0.05$ when $g_{23} = 0$.

B. Scenario 2: g_{11} and g_{23}

For the second scenario, we set $g_{33} = 0$ and consider g_{11} and g_{23} in Eq. (18). Such a flavor structure can be obtained when $n_{\text{VF}} = 2$. In this scenario, the $b \rightarrow c\bar{u}q$ anomaly requires

$$0.54 (\text{TeV})^{-1} \lesssim \sqrt{|g_{11}g_{23}|}/M_V \lesssim 0.78 (\text{TeV})^{-1}. \quad (23)$$

Although the size of the required coupling product is much smaller than the previous scenario, a severe bound on g_{23} comes from ΔM_s , where there is a tree-level Z' exchange diagram. We obtain

$$\frac{\Delta M_s^{Z'}}{\Delta M_s^{\text{SM}}} \simeq \eta^2 \frac{16\pi^2 g_{23}^2}{g_W^4 (V_{ts} V_{tb}^*)^2 x_V x_t f(x_t)}, \quad (24)$$

and find that g_{23} always gives a positive shift in ΔM_s . The constraint from ΔM_s is [39,40]

$$|g_{23}|/M_V \lesssim 0.01 (\text{TeV})^{-1}. \quad (25)$$

Therefore, $g_{11} \gtrsim 30 (M_V/\text{TeV}) \gg 4\pi$ is required by Eqs. (23) and (25), which implies that the $b \rightarrow c\bar{u}q$ anomaly can not be explained by this scenario.

In this scenario, $|g_{23}| \ll |g_{11}|$ should be satisfied. This simplifies the collider constraints because the production cross section is controlled only by $|g_{11}|$, and the heavy gauge bosons decay into jets with $\mathcal{B} \simeq 1$. The constraints are shown in Fig. 1(middle). We find $C_2^{W'}/C_2^{\text{SM}} \gtrsim -0.01$, where $g_{23} = -0.01(M_V/\text{TeV})$ is taken.

C. Scenario 3: g_{11} , g_{23} , and g_{33}

To see maximum value of $|C_2^{W'}/C_2^{\text{SM}}|$ in this model, we combine the first and second scenarios: all g_{11} , g_{23} , and g_{33} are non-zero entries. The point of this scenario is that the severe bound from ΔM_s can be turned off by

$$\frac{\Delta M_s^{W'}}{\Delta M_s^{\text{SM}}} + \frac{\Delta M_s^{Z'}}{\Delta M_s^{\text{SM}}} \sim 0, \quad (26)$$

where the W' contribution is destructive and the Z' one is constructive in ΔM_s (see previous subsections). We find, however, that even if the ΔM_s bound is turned off, $g_{11}g_{33}$ is still constrained from the ΔM_d as

$$\sqrt{|g_{11}g_{33}|}/M_V \lesssim 2.3 (\text{TeV})^{-1}. \quad (27)$$

This bound restricts the possible W' contribution to the $b \rightarrow c\bar{u}q$ processes. Also, we have checked a constraint from $b \rightarrow s\gamma$ data. We conclude that the $b \rightarrow s\gamma$ bound is less sensitive than ΔM_d , see the Supplemental Material [18].

Since $|g_{23}| \ll |g_{11}|, |g_{33}|$ still holds in this scenario, the collider constraints are almost the same as the scenario 1. We focus on a parameter region that the all LHC constrains are evaded by the broad width of the heavy gauge bosons. In Fig. 1(right), $C_2^{W'}/C_2^{\text{SM}}$ is shown on g_{23} - g_{33} plane by fixing $M_V = 1$ TeV and $g_{11} = -3.6$ corresponding to the maximum value allowed by the $K \rightarrow \pi\pi$ data in Eq. (17). Eventually, we obtain

$$\frac{C_2^{W'}(m_b)}{C_2^{\text{SM}}(m_b)} \gtrsim -0.10. \quad (28)$$

V. DISCUSSION

Motivated by a recent improvement of the SM predictions on $\bar{B}^0 \rightarrow D^{(*)+}K^-$ and $\bar{B}_s^0 \rightarrow D_s^{(*)+}\pi^-$, we investigated the size of possible several new physics contributions to these processes. In spite of severe bounds from the other flavor observables and the LHC searches, we conclude that a -10% shift in the $b \rightarrow c\bar{u}q$ amplitude is possible by the left-handed W' model. Such a new physics contribution can reduce the tension in the $b \rightarrow c\bar{u}q$ processes.

Since $g_{22} = g_{11}$ is a necessary condition, this model also produces new physics contributions to $b \rightarrow c\bar{c}s$ processes with the same size [51,52]. Although they, e.g., $B^+ \rightarrow J/\psi K^+$, have been measured precisely, the SM predictions suffer from large nonfactorizable corrections [53–55]. We, therefore, expect that the $b \rightarrow c\bar{c}s$ processes are less sensitive than $b \rightarrow c\bar{u}q$.

It is unclear whether the new physics scalar operator can explain the $b \rightarrow c\bar{u}q$ anomaly, but it is an interesting direction to consider it. For instance, within a general two Higgs doublet model, a charged Higgs interaction is [56]

$$\mathcal{L} = -H^+ \bar{u}^i (V\rho_d P_R - \rho_u^\dagger V P_L)_{ij} d^j + \text{H.c.}, \quad (29)$$

where $(V\rho_d)_{23}$ is stringently constrained by ΔM_s via a heavy neutral Higgs exchange, while $(\rho_u^\dagger V)_{23}$ is less constrained by the flavor and collider observables [57,58]. Therefore, a potentially large contribution to the $b \rightarrow c\bar{u}q$ processes would be expected.

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