Connecting $b \to s\ell\bar{\ell}$ anomalies to enhanced rare nonleptonic $\bar{B}_s^0$ decays in Z' model

Gaber Faisel[a] **and Jusak Tandean**[b,c]

[a]*Department of Physics, Faculty of Arts and Sciences, Süleyman Demirel University,*
D685, Isparta 32260, Turkey

[b]*Department of Physics, National Taiwan University,*
No. 1, Sec. 4, Roosevelt Rd., Taipei 10617, Taiwan

[c]*Physics Division, National Center for Theoretical Sciences,*
No. 101, Sec. 2, Kuang Fu Rd., Hsinchu 30013, Taiwan

E-mail: gaberfaisel@sdu.edu.tr, jtandean@yahoo.com

ABSTRACT: The present data on a number of observables in $b \to s\mu^+\mu^-$ processes manifest some tensions with the standard model (SM). Assuming that these anomalies have a new physics origin, we consider the possibility that a Z' boson is responsible for them. We further assume that its interactions with quarks also affect rare nonleptonic decays of the $\bar{B}_s^0$ meson which are purely isospin-violating and tend to be dominated by electroweak-penguin contributions, namely $\bar{B}_s^0 \to (\eta, \eta', \phi)(\pi^0, \rho^0)$. Most of these decays are not yet observed, and their rates are expected to be relatively small in the SM. Taking into account constraints from various measurements, including the evidence for $\bar{B}_s^0 \to \phi\rho^0$ recently seen by LHCb, we find that the Z' effects on $\bar{B}_s^0 \to (\eta, \phi)\pi^0$ can make their rates bigger than the SM predictions by up to an order of magnitude. For $\bar{B}_s^0 \to \eta'\pi^0, (\eta, \eta')\rho^0$, the enhancement factors are at most a few. Since the Z' contributions to the different channels depend on different combinations of its couplings, observations of more of these decays in future experiments, along with improved $b \to s\mu^+\mu^-$ data, will probe this Z' scenario more thoroughly.

Contents

1 Introduction

The latest measurements of various $b \to s\mu^+\mu^-$ processes have turned up some intriguing discrepancies from the expectations of the standard model (SM) of particle physics. Specifically, the LHCb Collaboration in its angular analysis of the decay $B^0 \to K^{*0}\mu^+\mu^-$ found tensions with the SM at the 3.4σ level [1]. This was later confirmed in the Belle experiment on the same process, but with lower statistical confidence [2]. Furthermore, LHCb findings [3, 4] on the ratio R_K of the branching fractions of $B^+ \to K^+\mu^+\mu^-$ and $B^+ \to K^+e^+e^-$ decays and on the corresponding ratio R_{K^*} for $B^0 \to K^{*0}\mu^+\mu^-$ and $B^0 \to K^{*0}e^+e^-$ decays are all below their SM predictions [5–7] by 2.1σ to 2.6σ. In addition, the current data [8–10] on the branching fractions of $B \to K^{(*)}\mu^+\mu^-$ and $B_s \to \phi\mu^+\mu^-$ favor values less than their SM estimates.

Although the statistical significance of the aforesaid deviations from SM expectations is still too low for a definite conclusion, they may be early clues about interactions beyond the SM in $b \to s$ transitions. Recent model-independent theoretical analyses have in fact demonstrated that new physics (NP) could account for these anomalies [11–25]. In view of the possibility that these tentative hints of NP will be confirmed by upcoming experiments, it is of interest to explore the potential implications for other $b \to s$ processes.

Among them are the nonleptonic decays $\bar{B}_s \to \eta\pi^0$, $\bar{B}_s \to \eta'\pi^0$, $\bar{B}_s \to \phi\pi^0$, $\bar{B}_s \to \eta\rho^0$, $\bar{B}_s \to \eta'\rho^0$, and $\bar{B}_s \to \phi\rho^0$. Each of these transitions has a final state with total isospin $I = 1$ and thus fully breaks isospin symmetry, implying that their amplitudes receive no contributions from QCD-penguin operators and arise instead from charmless tree and electroweak-penguin (EWP) operators [26, 27]. The product of Cabibbo-Kobayashi-Maskawa (CKM) matrix elements in the tree contributions is suppressed compared to that in the EWP ones, and the suppression factor is $|V_{us}V_{ub}|/|V_{ts}V_{tb}| \sim 0.02$. Consequently, although the Wilson coefficients of the tree operators are much bigger than those of the EWP operators, the latter turn out to dominate the majority of these channels, and the resulting decay rates are relatively low [26–29]. Most of them are not yet observed, the exception being $\bar{B}_s \to \phi\rho^0$. Evidence for it was detected by LHCb last year [30] with

a branching fraction $\mathcal{B}\big(\bar{B}_s \to \phi\rho^0\big) = (0.27 \pm 0.08)\times 10^{-6}$ [10], which agrees with some of its estimates in the SM within sizable errors [31–36].

The smallness of the rates of $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$ in the SM implies that they may serve as probes of physics beyond it. This has been considered to varying extents in the contexts of different models [34–40]. In this paper, we treat these rare nonleptonic $\bar{B}_s$ decays along similar lines and suppose that the NP influencing them also causes the aforementioned $b \to s\mu^+\mu^-$ anomalies. We adopt in particular a scenario where an electrically neutral and uncolored spin-one particle, the Z' boson, is responsible for the new interactions in these two sets of $b \to s$ transitions. We assume that it couples nonuniversally to SM fermions and does not mix with SM gauge bosons, but it is not necessarily a gauge boson and could even be composite.

Although the possibility of NP effects on $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$ in the Z' context has been entertained before [36–39], our current paper contains new considerations and results which were not available in the previous literature. Firstly, while the past studies separately examined only subsets of these six modes,[1] here we treat all of them at the same time, noting that among B_s decays into two charmless mesons the six are the only ones which are strangeness changing and purely isospin-violating. This allows us to gain a more complete picture than before concerning the Z' contributions, which reveals clearly how they in general modify the different channels in different ways. A second novel aspect of our analysis is that, as stated in the preceding paragraph, we explore a scenario in which the same Z' not only modifies these nonleptonic $\bar{B}_s$ decays, but also gives rise to the $b \to s\ell\bar{\ell}$ anomalies. It turns out that assuming this link between the two sets of processes leads to an important consequence for the Z' interactions, namely that the left-handed bsZ' coupling must be roughly ten times stronger than the right-handed one if both of them exist, as will be detailed later on. This particular finding was absent from the earlier studies [36–39], which did not deal with such a potential link, as most of them appeared before the arrival of the anomalies. A third significant novelty in our analysis is that we take into account the foregoing evidence of $\bar{B}_s \to \phi\rho^0$ recently seen by LHCb [30]. As this new measurement, albeit still with a sizable uncertainty, is compatible with its SM expectations, we will show that the implied room for the Z' influence on the channels with the ρ^0, not only $\bar{B}_s \to \phi\rho^0$ but also $\bar{B}_s \to (\eta, \eta')\rho^0$, is now limited. In contrast, previously $\bar{B}_s \to \phi\rho^0$ and $\bar{B}_s \to \eta\rho^0$ were allowed in refs. [36] and [39], respectively, to have rates exceeding their SM predictions by an order of magnitude. Last but not least, we will nevertheless also demonstrate that the viable Z' parameter space still accommodates the possibility that the pionic channels $\bar{B}_s \to \phi\pi^0$ and $\bar{B}_s \to \eta\pi^0$ can have rates which are about a factor of ten higher than their SM values. Needless to say, this should add to the motivation for intensified efforts in upcoming experiments at LHCb and Belle II to investigate all these decays. The acquired data on them would provide especially useful complementary information about the NP responsible for the $b \to s\ell\bar{\ell}$ anomalies should the latter be established by future measurements to be signals of physics beyond the SM.

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[1]Of the six modes, only $\bar{B}_s \to \phi\pi^0$ was discussed in [37, 38], $\bar{B}_s \to \phi(\pi^0, \rho^0)$ in [36], and $\bar{B}_s \to (\eta, \eta')(\pi^0, \rho^0)$ in [39].

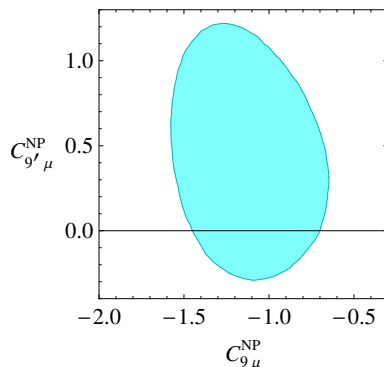**Figure 1**. Allowed 2σ (cyan) region of $C_{9'\mu}^{\rm NP}$ versus $C_{9\mu}^{\rm NP}$ from the global analysis of $b \to s\mu^+\mu^-$ data performed in ref. [18].

The rest of the paper is organized as follows. In section 2, we address the Z' contributions to $b \to s\mu^+\mu^-$ and apply constraints from the relevant empirical information, including that on B_s-$\bar{B}_s$ mixing. In section 3, we examine the impact of the Z' interactions with SM quarks on $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$. To evaluate their amplitudes, we employ the soft-collinear effective theory [41–48]. We show that in the Z' presence the rates of $\bar{B}_s \to \eta'\pi^0, (\eta, \eta')\rho^0$ can increase by as much as factors of a few with respect to the predictions in the SM, while the rates of $\bar{B}_s \to (\eta, \phi)\pi^0$ can exceed their SM values by up to an order of magnitude. We make our conclusions in section 4.

2 Z' interactions in $b \to s\mu\bar{\mu}$ and B_s-$\bar{B}_s$ mixing

Global analyses [18, 19] have found that some of the best fits to the most recent anomalous $b \to s\mu^+\mu^-$ measurements result from effective interactions given by

$$\mathcal{L}_{\rm eff} \supset \frac{\alpha_{\rm e}\lambda_t G_{\rm F}}{\sqrt{2}\,\pi}\big(C_{9\mu}\,\overline{s}\,\gamma^\kappa P_L b + C_{9'\mu}\,\overline{s}\,\gamma^\kappa P_R b\big)\overline{\mu}\,\gamma_\kappa\mu \ + \ {\rm H.c.}\,,$$

$$\lambda_t \ = \ V_{ts}^* V_{tb}\,, \qquad P_{L,R} \ = \ \frac{1}{2}(1 \mp \gamma_5)\,, \tag{2.1}$$

where $C_{9\mu} = C_{9\ell}^{\rm SM} + C_{9\mu}^{\rm NP}$ and $C_{9'\mu} = C_{9'\mu}^{\rm NP}$ are the Wilson coefficients, $\alpha_{\rm e} = 1/133$ is the fine structure constant at the b-quark mass (m_b) scale, and $G_{\rm F}$ is the Fermi constant. The same SM part $C_{9\ell}^{\rm SM}$ occurs in the electron and tau channels $b \to s(e^+e^-, \tau^+\tau^-)$, but they are not affected by the NP. In figure 1, for later use, we display the 2σ (cyan) region of $C_{9'\mu}^{\rm NP}$ versus $C_{9\mu}^{\rm NP}$ permitted by the data, from the global fit carried out in ref. [18].

In the literature, many models possessing some kind of Z' particle with different sets of fermionic couplings have been studied in relation to the $b \to s\mu^+\mu^-$ anomalies [49–88]. In the Z' scenario considered here, the interactions responsible for $C_{9\mu,9'\mu}^{\rm NP}$ in eq. (2.1) are described by

$$\mathcal{L}_{Z'} \supset -\big[\overline{s}\,\gamma^\kappa\big(\Delta_L^{sb} P_L + \Delta_R^{sb} P_R\big)b\,Z'_\kappa \ + \ {\rm H.c.}\big] \ - \ \Delta_V^{\mu\mu}\,\overline{\mu}\,\gamma^\kappa\mu\,Z'_\kappa\,, \tag{2.2}$$

where the constants $\Delta_{L,R}^{sb}$ are generally complex and $\Delta_V^{\mu\mu}$ is real due to the Hermiticity of $\mathcal{L}_{Z'}$. Any other possible Z' couplings to leptons are taken to be negligible. To simplify the
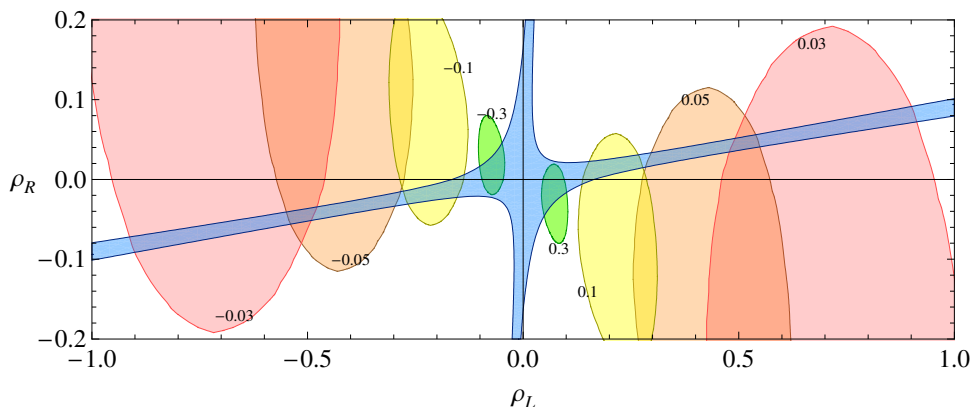
Figure 2. Regions of ρ_R versus ρ_L for $m_{Z'} = 1\,\text{TeV}$ which are consistent with the $C_{9'\mu}^{\text{NP}}$-$C_{9\mu}^{\text{NP}}$ constraint depicted in figure 1 for $\Delta_V^{\mu\mu} = \pm 0.03$ (red), ± 0.05 (orange), ± 0.1 (yellow), and ± 0.3 (green). The blue area fulfills the condition in eq. (2.7) from B_s-$\bar{B}_s$ mixing data.

analysis, hereafter we focus on the special case in which

$$\Delta_L^{sb} = \rho_L V_{ts}^* V_{tb}, \qquad \Delta_R^{sb} = \rho_R V_{ts}^* V_{tb}, \qquad (2.3)$$

where $\rho_{L,R}$ are real numbers, and so they do not supply any new CP-violation phase. Accordingly, for a heavy Z' with mass $m_{Z'}$ we obtain

$$C_{9\mu}^{\text{NP}} = \frac{-\sqrt{2}\,\pi\,\rho_L \Delta_V^{\mu\mu}}{\alpha_{\text{e}} G_{\text{F}}\, m_{Z'}^2}, \qquad C_{9'\mu}^{\text{NP}} = \frac{-\sqrt{2}\,\pi\,\rho_R \Delta_V^{\mu\mu}}{\alpha_{\text{e}} G_{\text{F}}\, m_{Z'}^2}. \qquad (2.4)$$

In figure 2, we illustrate the ranges of ρ_L and ρ_R corresponding to the allowed $C_{9'\mu}^{\text{NP}}$-$C_{9\mu}^{\text{NP}}$ (cyan) region in figure 1 for $m_{Z'} = 1\,\text{TeV}$ and some sample choices of $\Delta_V^{\mu\mu}$, namely ± 0.03 (red), ± 0.05 (orange), ± 0.1 (yellow), and ± 0.3 (green). We note that these $\Delta_V^{\mu\mu}$ values contribute positively to the SM muon anomalous magnetic moment, but with $m_{Z'} = 1\,\text{TeV}$ are too small to explain the disagreement with its measurement [89].

The Z' couplings in eq. (2.2) also induce tree-level effects on $\Delta M_s = 2|M_{12}^s|$, which pertains to B_s-$\bar{B}_s$ mixing and has been measured to be $\Delta M_s^{\text{exp}} = (17.757 \pm 0.021)/\text{ps}$ [10]. We can express the sum of the SM and Z' contributions as [90]

$$M_{12}^s = M_{12}^{s,\text{SM}} \left(1 + 4\tilde{r}\, \frac{\rho_L^2 + \rho_R^2 + \kappa_{LR}\, \rho_L \rho_R}{g_{\text{SM}}^2\, S_0\, m_{Z'}^2} \right), \qquad (2.5)$$

where [90] $\tilde{r} = 0.985$ for $m_{Z'} = 1\,\text{TeV}$ is a QCD factor, $g_{\text{SM}}^2 = 1.7814 \times 10^{-7}\,\text{GeV}^{-2}$, the SM loop function $S_0 \simeq 2.35$ for a top-quark mass $m_t \simeq 165\,\text{GeV}$, and

$$\kappa_{LR} = \frac{6\left(C_1^{LR} \langle Q_1^{LR} \rangle + C_2^{LR} \langle Q_2^{LR} \rangle \right)}{\eta_B \hat{B}_{B_s} f_{B_s}^2 m_{B_s} \tilde{r}}, \qquad (2.6)$$

with [90] $C_1^{LR} = 1 - \alpha_{\text{s}} \left[1/6 + 2\log(m_{Z'}/\mu') \right]/(4\pi)$ and $C_2^{LR} = \alpha_{\text{s}} \left[-1 - 12\log(m_{Z'}/\mu') \right]/(4\pi)$ containing the strong coupling constant α_{s}, all evaluated at a scale $\mu' \sim m_{Z'}$, $\langle Q_1^{LR} \rangle =$

$-0.37\,\mathrm{GeV}^3$, $\langle Q_2^{LR}\rangle = 0.51\,\mathrm{GeV}^3$, $\eta_B = 0.55 \pm 0.01$, and [91, 92] $f_{B_s}\hat{B}_{B_s}^{1/2} = (262.2 \pm 9.7)\,\mathrm{MeV}$. With the central values of these parameters and m_{B_s} from ref. [10], we get $\kappa_{LR} = -11.2$ for $m_{Z'} = 1\,\mathrm{TeV}$.

To apply restrictions on $\rho_{L,R}$ from the B_s-$\bar{B}_s$ mixing data, we impose

$$0.899 \leq \frac{\Delta M_s}{\Delta M_s^{\mathrm{SM}}} = \left|\frac{M_{12}^s}{M_{12}^{s,\mathrm{SM}}}\right| \leq 1.252\,, \tag{2.7}$$

which is the 95% confidence level (CL) range from the latest UTfit global analysis [92]. Since some of the numbers quoted in the last paragraph have uncertainties up to a few percent, we let κ_{LR} vary by up to 10% from its central value when scanning the parameter space for $\rho_{L,R}$ values which conform to eq. (2.7). For $m_{Z'} = 1\,\mathrm{TeV}$, we incorporate the scan result into figure 2, represented by the blue area. Thus, in this figure each overlap of the blue area with one of the other colored ones of a particular $\Delta_V^{\mu\mu}$ value corresponds to the parameter space that can explain the $b \to s\mu^+\mu^-$ anomalies and simultaneously satisfies eq. (2.7). With smaller choices of $|\Delta_V^{\mu\mu}|$, such overlaps could be found at larger $|\rho_{L,R}|$ values. This graph also reveals that in the absence of the right-handed coupling, $\rho_R = 0$, the allowed range of ρ_L would be rather narrow, indicating the importance of nonvanishing ρ_R for gaining bigger viable parameter space [51].

3 Z' contributions to rare nonleptonic $\bar{B}_s$ decays

Given that $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$ change both strangeness and isospin, in the SM their amplitudes proceed from $b \to s$ four-quark operators $O_{1,2}^u$ and $O_{7,8,9,10}$ which are derived from charmless tree and electroweak-penguin diagrams, respectively. In contrast, the QCD-penguin operators $O_{3,4,5,6}$, which preserve isospin symmetry, do not affect these processes.[2] In many models beyond the SM, new interactions may modify the Wilson coefficients C_i of O_i and/or give rise to extra operators $\tilde{O}_i$ which are the chirality-flipped counterparts of O_i. A flavor-violating Z' boson may contribute to some of them, depending on the details of its properties.

In our scenario of interest, besides its couplings in eq. (2.2), the Z' has flavor-conserving interactions with the u and d quarks via

$$\mathcal{L}_{Z'} \supset -\big[\bar{u}\,\gamma^\kappa\big(\Delta_L^{uu}P_L + \Delta_R^{uu}P_R\big)u + \bar{d}\,\gamma^\kappa\big(\Delta_L^{dd}P_L + \Delta_R^{dd}P_R\big)d\big]Z'_\kappa\,, \tag{3.1}$$

the constants $\Delta_{L,R}^{uu,dd}$ being real, but does not couple flavor-diagonally to other quarks. From eqs. (2.2) and (3.1), we can derive tree-level Z'-mediated diagrams contributing to nonleptonic $b \to s$ reactions. For a heavy Z', these diagrams yield

$$\mathcal{L}_{4\text{-quark}}^{Z'} \supset \frac{-\lambda_t}{m_{Z'}^2}\,\bar{s}\,\gamma^\kappa\big(\rho_L P_L + \rho_R P_R\big)b \sum_{q=u,d} \bar{q}\,\gamma_\kappa\big(\Delta_L^{qq}P_L + \Delta_R^{qq}P_R\big)q \tag{3.2}$$

after applying eq. (2.3). It is straightforward to realize that these additional terms bring about modifications to the coefficients of the QCD- and electroweak-penguin operators

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[2]The expressions for O_i, $i = 1, 2, \cdots, 10$, can be found in, e.g., [47].

$O_{3,5,7,9}$ in the SM and also generate their chirality-flipped partners $\tilde{O}_{3,5,7,9}$ [93]. We can express them in the effective Lagrangian for $b \to s$ transitions as

$$\mathcal{L}_{\text{eff}} \supset \sqrt{8}\,\lambda_t G_{\text{F}} \sum_{q=u,d} \left\{ \overline{s}\,\gamma^\kappa P_L b \left[\left(C_3 + \frac{3}{2}C_9 e_q \right) \overline{q}\,\gamma_\kappa P_L q + \left(C_5 + \frac{3}{2}C_7 e_q \right) \overline{q}\,\gamma_\kappa P_R q \right] \right.$$
$$\left. + \overline{s}\,\gamma^\kappa P_R b \left[\left(\tilde{C}_3 + \frac{3}{2}\tilde{C}_9 e_q \right) \overline{q}\,\gamma_\kappa P_R q + \left(\tilde{C}_5 + \frac{3}{2}\tilde{C}_7 e_q \right) \overline{q}\,\gamma_\kappa P_L q \right] \right\}, \quad (3.3)$$

where $C_j = C_j^{\text{SM}} + C_j^{Z'}$ and $\tilde{C}_j = \tilde{C}_j^{Z'}$ for $j = 3, 5, 7, 9$ are the Wilson coefficients. Thus, from eq. (3.2) we have [36, 38, 93]

$$C_{3,5}^{Z'} = \frac{\rho_L\left(-\Delta_{L,R}^{uu} - 2\Delta_{L,R}^{dd}\right)}{6\sqrt{2}\,G_{\text{F}}\,m_{Z'}^2} = \frac{\rho_L\left(-\delta_{L,R} - 3\Delta_{L,R}^{dd}\right)}{6\sqrt{2}\,G_{\text{F}}\,m_{Z'}^2},$$

$$\tilde{C}_{3,5}^{Z'} = \frac{\rho_R\left(-\Delta_{R,L}^{uu} - 2\Delta_{R,L}^{dd}\right)}{6\sqrt{2}\,G_{\text{F}}\,m_{Z'}^2} = \frac{\rho_R\left(-\delta_{R,L} - 3\Delta_{R,L}^{dd}\right)}{6\sqrt{2}\,G_{\text{F}}\,m_{Z'}^2},$$

$$C_{7,9}^{Z'} = \frac{\rho_L\left(-\Delta_{R,L}^{uu} + \Delta_{R,L}^{dd}\right)}{3\sqrt{2}\,G_{\text{F}}\,m_{Z'}^2} = \frac{-\rho_L\,\delta_{R,L}}{3\sqrt{2}\,G_{\text{F}}\,m_{Z'}^2},$$

$$\tilde{C}_{7,9}^{Z'} = \frac{\rho_R\left(-\Delta_{L,R}^{uu} + \Delta_{L,R}^{dd}\right)}{3\sqrt{2}\,G_{\text{F}}\,m_{Z'}^2} = \frac{-\rho_R\,\delta_{L,R}}{3\sqrt{2}\,G_{\text{F}}\,m_{Z'}^2}, \quad (3.4)$$

where

$$\delta_L = \Delta_L^{uu} - \Delta_L^{dd}, \qquad \delta_R = \Delta_R^{uu} - \Delta_R^{dd}. \quad (3.5)$$

As $O_{3,5}$ and $\tilde{O}_{3,5}$ do not break isospin, only $C_{7,9}^{Z'}$ and $\tilde{C}_{7,9}^{Z'}$ contribute to $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$.

To estimate the Z' impact on these decays, we make use of the soft-collinear effective theory (SCET) [41–48], similarly to what was done in ref. [39] in the case of a leptophobic-Z' model. For any one of them, the SCET amplitude at leading order in $\alpha_s(m_b)$ can be written as [48]

$$\mathcal{A}_{\bar{B}_s \to M_1 M_2} = \frac{f_{M_1} G_{\text{F}} m_{B_s}^2}{\sqrt{2}} \left[\int_0^1 d\nu \left(\zeta_J^{BM_2} T_{1J}(\nu) + \zeta_{Jg}^{BM_2} T_{1Jg}(\nu) \right) \phi_{M_1}(\nu) + \zeta^{BM_2} T_1 + \zeta_g^{BM_2} T_{1g} \right]$$
$$+ (1 \leftrightarrow 2), \quad (3.6)$$

where f_M is the decay constant of meson M, the ζ's are nonperturbative hadronic parameters which can be fixed from experiment, the T's are hard kernels which are functions of the Wilson coefficients C_i and $\tilde{C}_i$, and $\phi_M(\nu)$ is the light-cone distribution amplitude of M which is normalized as $\int_0^1 d\nu\,\phi_M(\nu) = 1$. The so-called charming-penguin term, which in this case conserves isospin, is absent from $\mathcal{A}_{\bar{B}_s \to M_1 M_2}$. The hard kernels for the decays of concern are available from the literature [33, 47, 48] and have been listed in table 1, where the flavor states $\eta_q \sim \left(u\bar{u} + d\bar{d}\right)/\sqrt{2}$ and $\eta_s \sim s\bar{s}$ are related to the physical meson states η and η' by $\eta = \eta_q \cos\theta - \eta_s \sin\theta$ and $\eta' = \eta_q \sin\theta + \eta_s \cos\theta$ with mixing angle $\theta = 39.3°$ [47, 48, 95, 96].

In the presence of NP which also generates the extra operators $\tilde{O}_i$, the quantities $c_{2,3}$ and $b_{2,3}$ in table 1 depend not only on C_i and $\tilde{C}_i$, but also on the final mesons M_1 and

Decay mode	T_1	T_2	T_{1g}	T_{2g}
$\bar{B}_s \to \eta_s \pi^0$	0	$\frac{1}{\sqrt{2}}(c_2 - c_3)$	0	$\frac{1}{\sqrt{2}}(c_2 - c_3)$
$\bar{B}_s \to \eta_s \rho^0$	0	$\frac{1}{\sqrt{2}}(c_2 + c_3)$	0	$\frac{1}{\sqrt{2}}(c_2 + c_3)$
$\bar{B}_s \to \eta_q \pi^0$	0	0	0	$c_2 - c_3$
$\bar{B}_s \to \eta_q \rho^0$	0	0	0	$c_2 + c_3$
$\bar{B}_s \to \phi \pi^0$	0	$\frac{1}{\sqrt{2}}(c_2 - c_3)$	0	0
$\bar{B}_s \to \phi \rho^0$	0	$\frac{1}{\sqrt{2}}(c_2 + c_3)$	0	0

Table 1. Hard kernels $T_{1,2,1g,2g}$ for $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$ decays. The hard kernels $T_{rJ,rJg}(\nu)$ for $r = 1, 2$ are obtainable from $T_{r,rg}$, respectively, through the replacement $c_k \to b_k$, where b_k has dependence on ν.

M_2, as well as on the CKM factors λ_t and $\lambda_u = V_{us}^* V_{ub}$. The dependence on M_1 and M_2 arises from the fact that, with regard to the nonzero kernels in this table, for each 4-quark operator the $\bar{B}_s \to M_1$ and vacuum $\to M_2$ matrix elements and their contraction in the amplitude can lead to an overall negative or positive sign for the contribution of the operator, the sign being determined by the chirality combination of the operator and by whether the final mesons are pseudoscalars (PP), vectors (VV), PV, or VP. Thus, for $\bar{B}_s \to (\eta_q, \eta_s)\pi^0$ and $\bar{B}_s \to \phi\rho^0$ we have[3]

$$c_2 = \lambda_u \left(C_2 - \tilde{C}_2 + \frac{C_1 - \tilde{C}_1}{N_c} \right) - \frac{3\lambda_t}{2} \left(C_9 - \tilde{C}_9 + \frac{C_{10} - \tilde{C}_{10}}{N_c} \right),$$

$$c_3 = -\frac{3\lambda_t}{2} \left(C_7 - \tilde{C}_7 + \frac{C_8 - \tilde{C}_8}{N_c} \right),$$

$$b_2 = \lambda_u \left[C_2 - \tilde{C}_2 + \left(1 - \frac{m_b}{\omega_3}\right)\frac{C_1 - \tilde{C}_1}{N_c} \right] - \frac{3\lambda_t}{2} \left[C_9 - \tilde{C}_9 + \left(1 - \frac{m_b}{\omega_3}\right)\frac{C_{10} - \tilde{C}_{10}}{N_c} \right],$$

$$b_3 = -\frac{3\lambda_t}{2} \left[C_7 - \tilde{C}_7 + \left(1 - \frac{m_b}{\omega_2}\right)\frac{C_8 - \tilde{C}_8}{N_c} \right], \tag{3.7}$$

where $N_c = 3$ is the color number and $b_{2,3}$, which are contained in $T_{2J}(\nu)$ and $T_{2Jg}(\nu)$, are also functions of ν via [48] $\omega_2 = \nu m_{B_s}$ and $\omega_3 = (\nu-1)m_{B_s}$. However, for $\bar{B}_s \to (\eta_q, \eta_s)\rho^0$ and $\bar{B}_s \to \phi\pi^0$ we need to make the sign change $-\tilde{C}_i \to +\tilde{C}_i$ in $c_{2,3}$ and $b_{2,3}$.

The expressions in eq. (3.7) generalize the SM ones provided previously in refs. [47, 48]. They also supplied the values of the SM coefficients C_i^{SM} at the m_b scale, $C_{1,2}^{\text{SM}} = (1.11, -0.253)$ and $C_{7,8,9,10}^{\text{SM}} = (0.09, 0.24, -10.3, 2.2) \times 10^{-3}$ [94], which we will use in $c_{2,3}$ and $b_{2,3}$. Our Z' contributions of interest, in eq. (3.4), enter eq. (3.7) only via $C_{7,9} = C_{7,9}^{\text{SM}} + C_{7,9}^{Z'}$ and $\tilde{C}_{7,9} = \tilde{C}_{7,9}^{Z'}$.

For numerical computation of $\mathcal{A}_{\bar{B}_s \to M_1 M_2}$, in view of table 1, the meson decay constants which we need are only $f_\pi = 131\,\text{MeV}$ and $f_\rho = 209\,\text{MeV}$, and the integral in eq. (3.6) can

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[3]The formula for b_2 given in [39] contains typos which we have corrected here in eq. (3.7).

Decay mode	QCDF	PQCD	SCET	
			Solution 1	Solution 2
$\bar{B}_s \to \eta\pi^0$	$0.05^{+0.03+0.02}_{-0.01-0.01}$	$0.05^{+0.02+0.01+0.00}_{-0.02-0.01-0.00}$	$0.032 \pm 0.013 \pm 0.008$	$0.025 \pm 0.010 \pm 0.003$
$\bar{B}_s \to \eta'\pi^0$	$0.04^{+0.01+0.01}_{-0.00-0.00}$	$0.11^{+0.05+0.02+0.00}_{-0.03-0.01-0.00}$	$0.001 \pm 0.000 \pm 0.005$	$0.052 \pm 0.021 \pm 0.015$
$\bar{B}_s \to \phi\pi^0$	$0.12^{+0.02+0.04}_{-0.01-0.02}$	$0.16^{+0.06+0.02+0}_{-0.05-0.02-0}$	$0.074 \pm 0.030 \pm 0.009$	$0.091 \pm 0.036 \pm 0.016$
$\bar{B}_s \to \eta\rho^0$	$0.10^{+0.02+0.02}_{-0.01-0.01}$	$0.06^{+0.03+0.01+0.00}_{-0.02-0.01-0.00}$	$0.078 \pm 0.031 \pm 0.022$	$0.059 \pm 0.023 \pm 0.006$
$\bar{B}_s \to \eta'\rho^0$	$0.16^{+0.06+0.03}_{-0.02-0.03}$	$0.13^{+0.06+0.02+0.00}_{-0.04-0.02-0.01}$	$0.003 \pm 0.001 \pm 0.013$	$0.141 \pm 0.056 \pm 0.042$
$\bar{B}_s \to \phi\rho^0$	$0.18^{+0.01+0.09}_{-0.01-0.04}$	$0.23^{+0.09+0.03+0.00}_{-0.07-0.01-0.01}$	$0.36 \pm 0.14 \pm 0.04$	

Table 2. Branching fractions, in units of 10^{-6}, of $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$ decays in the SM. For the first five modes, the last two columns correspond to the two solutions of SCET parameters in eq. (3.8). The errors of the SCET predictions are due to assumed 20% flavor-SU(3)-breaking effects and the errors in the ζs from fits to data, respectively. For comparison, the second and third columns contain results calculated in the frameworks of QCDF [31] and PQCD [32].

be treated with the aid of the relations $\int_0^1 d\nu\, \phi_M(\nu)/\nu = \int_0^1 d\nu\, \phi_M(\nu)/(1-\nu) \equiv \langle\chi^{-1}\rangle_M$ for $M = \pi, \rho$, in which cases $\langle\chi^{-1}\rangle_\pi = 3.3$ and $\langle\chi^{-1}\rangle_\rho = 3.45$ [47, 48]. Moreover, for the ζ's we adopt the two solutions derived from the fit to data done in ref. [48]:

$$\left(\zeta^P, \zeta_J^P, \zeta^V, \zeta_J^V, \zeta_g, \zeta_{Jg}\right)_1 = (0.137, 0.069, 0.117, 0.116, -0.049, -0.027),$$
$$\left(\zeta^P, \zeta_J^P, \zeta^V, \zeta_J^V, \zeta_g, \zeta_{Jg}\right)_2 = (0.141, 0.056, 0.227, 0.065, -0.100, 0.051). \tag{3.8}$$

From these, we can obtain $\zeta_{(J)}^{B\eta_q} = \zeta_{(J)}^{B\eta_s} = \zeta_{(J)}^P$, $\zeta_{(J)}^{B\phi} = \zeta_{(J)}^V$, and $\zeta_{(J)g}^{B\eta_q} = \zeta_{(J)g}^{B\eta_s} = \zeta_{(J)g}$ under the assumption of flavor-SU(3) symmetry [48]. In eq. (3.8), we have not displayed the errors of the ζs from the fit to data, which are available from ref. [48]. Other input parameters that we will employ are the meson masses $m_{\pi^0} = 134.977$, $m_\eta = 547.862$, $m_{\eta'} = 957.78$, $m_{\rho^0} = 769$, $m_\phi = 1019.46$, and $m_{B_s} = 5366.89$, all in units of MeV, and the B_s lifetime $\tau_{B_s} = 1.505 \times 10^{-12}$ s, which are their central values from ref. [10].

Before addressing the Z' influence on $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$, we provide the SM predictions for their branching fractions, which are collected in table 2. For the first five modes, the SCET numbers have been evaluated with the preceding formulas and parameter values, and the last two columns correspond to the two solutions of SCET parameters in eq. (3.8). For the sixth ($\phi\rho^0$) mode, the SCET entry has been computed with the CKM and SCET parameters supplied very recently in ref. [33]. The two errors in each of the SCET predictions are due to flavor-SU(3)-breaking effects which we have assumed to be 20% and due to the errors in the ζs from the fits to data, respectively, the latter errors being given in refs. [33, 48]. The SCET numbers for $\bar{B}_s \to (\eta, \eta')\rho^0, \phi\pi^0$ ($\bar{B}_s \to \phi\rho^0$) are close to the corresponding ones determined in ref. [48] ([33]).[4] For comparison, in the second and third columns we quote numbers calculated with QCD factorization (QCDF) [31] and perturbative QCD (PQCD) [32]. Evidently, these two methods produce results comparable to those of SCET, especially with its Solution 2 in the case of the first five modes, considering the errors in the predictions. The entries for $\bar{B}_s \to \phi\rho^0$ are

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[4]The SCET predictions in table 2 differ from those obtained in [39] because some of the input parameters used in our two papers are not the same.

also compatible with the new measurement $\mathcal{B}(\bar{B}_s \to \phi\rho^0)_{\exp} = (0.27 \pm 0.08) \times 10^{-6}$ [10] mentioned earlier. An important implication of what we see in this table is that NP would not be easily noticeable in the rates of these decays unless it could enhance them by more than a factor of 2. This possibility may be unlikely to be realized in the case of $\bar{B}_s \to \phi\rho^0$ which has been detected having a rate consistent with SM expectations. Nevertheless, as we demonstrate below, substantial enhancement can still occur in some of the other channels.

Now we include the Z' contributions from eq. (3.4) in order to examine their impact on these decays. As table 2 indicates that the predictions of the SCET Solution 1 for $\bar{B}_s \to \eta'(\pi^0, \rho^0)$ are comparatively quite suppressed, from this point on we employ only Solution 2 parameters in our treatment of $\bar{B}_s \to (\eta, \eta')(\pi^0, \rho^0), \phi\pi^0$. Thus, summing the SM and Z' terms for $m_{Z'} = 1\,\text{TeV}$, with the central values of the input parameters, we find the amplitudes (in units of GeV) for the π^0 channels to be

$$
\begin{aligned}
10^9 \mathcal{A}_{\bar{B}_s \to \eta\pi^0} &\simeq 1.67 + 0.47i + (3.96 - 0.08i)(\rho_L + \rho_R)(\delta_L - \delta_R)\,, \\
10^9 \mathcal{A}_{\bar{B}_s \to \eta'\pi^0} &\simeq 0.48 - 2.48i - (1.90 - 0.04i)(\rho_L + \rho_R)(\delta_L - \delta_R)\,, \\
10^9 \mathcal{A}_{\bar{B}_s \to \phi\pi^0} &\simeq -2.88 - 1.69i - (7.85 - 0.15i)(\rho_L - \rho_R)(\delta_L - \delta_R)
\end{aligned}
\tag{3.9}
$$

and for the ρ^0 channels

$$
\begin{aligned}
10^9 \mathcal{A}_{\bar{B}_s \to \eta\rho^0} &\simeq 2.56 + 0.77i + (6.32 - 0.12i)(\rho_L + \rho_R)(\delta_L + \delta_R)\,, \\
10^9 \mathcal{A}_{\bar{B}_s \to \eta'\rho^0} &\simeq 0.78 - 4.12i - (3.03 - 0.06i)(\rho_L + \rho_R)(\delta_L + \delta_R)\,, \\
10^9 \mathcal{A}_{\bar{B}_s \to \phi\rho^0} &\simeq -6.53 - 1.47i - (15.3 - 0.3i)(\rho_L - \rho_R)(\delta_L + \delta_R)\,,
\end{aligned}
\tag{3.10}
$$

where $\delta_{L,R}$ are defined in eq. (3.5).

We notice that the amplitudes in eqs. (3.9) and (3.10) do not all have the same dependence on $\rho_{L,R}$ and $\delta_{L,R}$. Therefore, although $\mathcal{B}(\bar{B}_s \to \phi\rho^0)_{\exp}$ implies a restraint on the values of $(\rho_L - \rho_R)(\delta_L + \delta_R)$ in $\mathcal{A}_{\bar{B}_s \to \phi\rho^0}$, the amplitudes for the other channels, which have different combinations of $\rho_{L,R}$ and $\delta_{L,R}$, may generally still be altered considerably with respect to their SM parts. However, in our particular Z' case $\rho_{L,R}$ must satisfy $\rho_R \sim 0.1\,\rho_L$, as can be inferred from figure 2. Hence, based on eq. (3.10), we may expect that the amplitudes for $\bar{B}_s \to (\eta, \eta')\rho^0$ do not deviate hugely from their SM values. To look into this more concretely, for definiteness we take $\rho_R = 0.1\,\rho_L$ and impose

$$
0.11 \leq 10^6 \, \mathcal{B}(\bar{B}_s \to \phi\rho^0) \leq 0.43\,,
\tag{3.11}
$$

which is the 2σ range of $\mathcal{B}(\bar{B}_s \to \phi\rho^0)_{\exp}$. From the allowed values of the product $\rho_L(\delta_L + \delta_R)$ we can assess how much the branching fractions of $\bar{B}_s \to (\eta, \eta')\rho^0$ are modified compared to the central values of their respective SM predictions in table 2 under SCET Solution 2. We show the results in figure 3, which also depicts eq. (3.11) relative to the SM prediction. We further find that the ranges $\rho_L(\delta_L + \delta_R) \in [-0.99, -0.71]$ and $[-0.23, 0.048]$ fulfill eq. (3.11). Within these ranges, represented by the horizontal portions of the unshaded areas in this figure, we learn that $\mathcal{B}(\bar{B}_s \to \eta\rho^0)$ (red solid curve) can reach
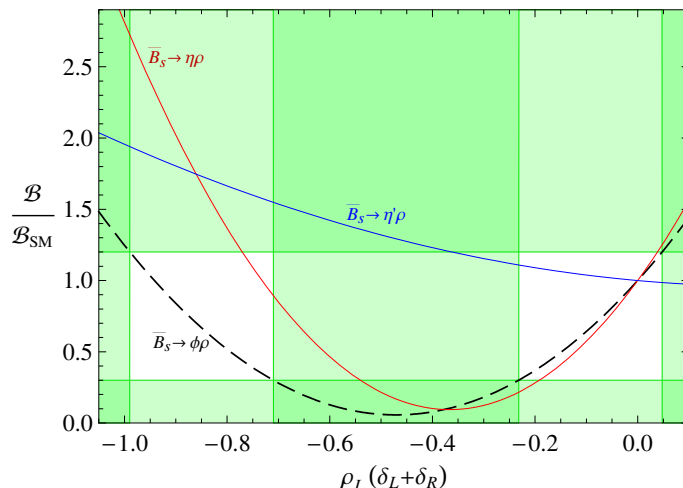
Figure 3. The calculated branching fractions of $\bar{B}_s \to \eta\rho^0$ (red solid curve), $\bar{B}_s \to \eta'\rho^0$ (blue solid curve), and $\bar{B}_s \to \phi\rho^0$ (black curve), normalized by their respective SM predictions listed in table 2, versus the product $\rho_L(\delta_L + \delta_R)$ in the case where $\rho_R = 0.1\,\rho_L$ and $m_{Z'} = 1\,\text{TeV}$. The vertical length of the unshaded areas and the $\rho_L(\delta_L + \delta_R)$ values within them satisfy the restriction in eq. (3.11).

up to ~ 2.7 times its SM value, whereas for $\bar{B}_s \to \eta'\rho^0$ (blue solid curve) the enhancement is at most about 1.9 times.[5]

For $\bar{B}_s \to (\eta, \eta', \phi)\pi^0$, the Z'-induced terms in eq. (3.9) are proportional to $\delta_L - \delta_R$. Therefore, these channels are not subject to the condition in eq. (3.11), and their amplitudes may be affected by the Z' contributions more than their ρ^0 counterparts. To examine this more quantitatively, we set $\rho_R = 0.1\,\rho_L$ as in the previous paragraph and subsequently compute the branching fractions of these π^0 modes for $-1 \leq \rho_L(\delta_L - \delta_R) \leq 1$. In figure 4 we present the results divided by the central values of their respective SM predictions in table 2 under SCET Solution 2. We observe that over most of the $\rho_L(\delta_L - \delta_R) > 0$ region covered in this plot the Z' effects can cause the branching fractions of $\bar{B}_s \to (\eta, \phi)\pi^0$ to exceed their SM counterparts by at least a factor of 2 and up to about an order of magnitude. Moreover, the $\bar{B}_s \to (\eta, \phi)\pi^0$ rates tend to be enhanced together with roughly similar enlargement factors. One also notices that the Z' impact could instead bring about substantial reduction of their rates. In table 3, we provide examples of the enhancement factors for some representative values of $\rho_L(\delta_L - \delta_R)$.

It is worth remarking that the Z'-generated coefficients $C_{3,5,7,9}^{Z'}$ and $\tilde{C}_{3,5,7,9}^{Z'}$ in eq. (3.4) enter the amplitudes for nonleptonic $b \to s$ decays which are not dominated by the contributions of the electroweak-penguin operators, such as $B_s \to K^{(*)}\bar{K}^{(*)}$ and $B \to \pi K^{(*)}$. Since these transitions have been observed, their data imply restrictions on the size of $\rho_L(\delta_L \pm \delta_R)$, as the requirement $\rho_R \sim 0.1\,\rho_L$ in our Z' scenario implies that the role of $\tilde{C}_{3,5,7,9}^{Z'}$ is minor. Our choices $|\rho_L(\delta_L \pm \delta_R)| \leq 1$ above correspond to $\left| C_7^{Z'} \pm C_9^{Z'} \right| \leq 0.0202 \simeq 2|C_9^{\text{SM}}|$, where $C_9^{\text{SM}} = -0.0103$ as quoted before. We have checked that the changes to the rates of those decays due to $|C_{7,9}^{Z'}| \lesssim |C_9^{\text{SM}}|$ are less than the uncertainties of the SCET es-

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[5]Before the LHCb detection of the $\bar{B}_s \to \phi\rho^0$ evidence [30], the possibilities of the $\bar{B}_s \to (\phi, \eta)\rho^0$ rates exceeding their SM predictions by an order of magnitude were entertained in [36, 39], respectively, for the $\delta_L = 0$ case.
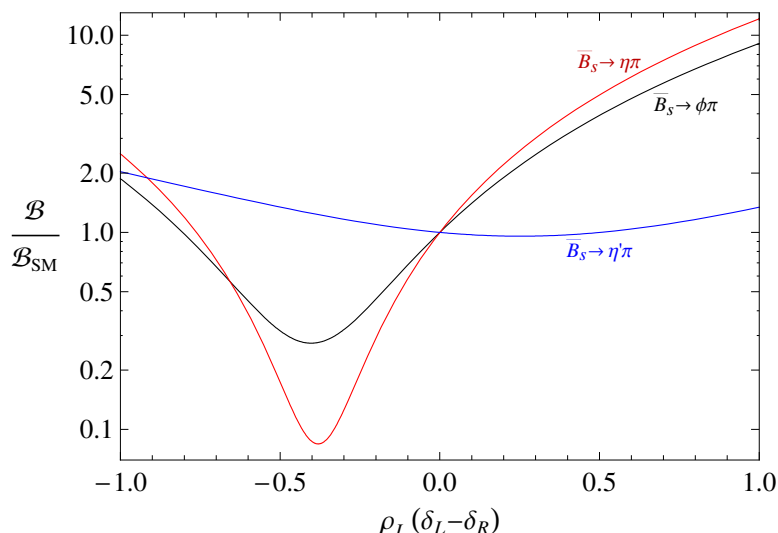
Figure 4. The calculated branching fractions of $\bar{B}_s \to \eta\pi^0$ (red curve), $\bar{B}_s \to \eta'\pi^0$ (blue curve), and $\bar{B}_s \to \phi\pi^0$ (black curve), normalized by their respective SM values listed in table 2 under SCET Solution 2, versus the product $\rho_L(\delta_L - \delta_R)$ in the case where $\rho_R = 0.1\,\rho_L$ and $m_{Z'} = 1\,\mathrm{TeV}$.

$\rho_L(\delta_L - \delta_R)$	$\bar{B}_s \to \eta\pi^0$	$\bar{B}_s \to \eta'\pi^0$	$\bar{B}_s \to \phi\pi^0$
-1	2.5	2.0	1.9
-0.5	0.17	1.3	0.32
0.5	5.0	1.0	3.9
1	12	1.3	9.1

Table 3. Enhancement factors of the branching fractions of $\bar{B}_s \to (\eta, \eta', \phi)\pi^0$ with respect to their SM predictions at a few representative values of $\rho_L(\delta_L - \delta_R)$ in the $\rho_R = 0.1\,\rho_L$ and $m_{Z'} = 1\,\mathrm{TeV}$ case.

timates in the SM, which are typically around 20% to 40% [33, 47, 48]. As for the influence of $C_{3,5}^{Z'}$, it can be minimized by adjusting the extra free parameters $\Delta_{L,R}^{dd}$ in eq. (3.4).[6] For comparison, earlier studies [36, 38, 40] concerning potential NP in $\bar{B}_s \to \phi(\pi^0, \rho^0)$ pointed out that rate enhancement factors of a few to an order of magnitude could still occur in the $\phi\pi^0$ mode and that $|C_j^{\mathrm{NP}}/C_9^{\mathrm{SM}}| \lesssim 2$ was not yet disfavored.

Finally, we would like to mention that the Z' coupling parameters of interest are separately consistent with constraints which may be pertinent from collider measurements. We illustrate this with the specific examples in table 4 for different sets of $\rho_L(\delta_L \pm \delta_R)$ and $\Delta_V^{\mu\mu}$ values in the aforesaid case where $\rho_R = 0.1\,\rho_L$ and $m_{Z'} = 1\,\mathrm{TeV}$. The choice $\left(\rho_L, \Delta_V^{\mu\mu}\right) = (0.8, 0.03)$ is evidently within the region covered in figure 2, while the points $\left(\rho_L, \Delta_V^{\mu\mu}\right) = \left((1.0, 1.2), 0.02\right)$ lie in the extension thereof. In this table, the displayed numbers for $\delta_{L,R}$ can comfortably comply with the condition $|\delta_{L,R}| \le 1.0\left[1 + (1.3\,\mathrm{TeV})^2/m_{Z'}^2\right]m_{Z'}/(3\,\mathrm{TeV})$

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[6]For instance, selecting $2\Delta_{L,R}^{dd} = -\Delta_{L,R}^{uu}$ would lead to $C_{3,5}^{Z'} = 0$, which was considered in [36, 39].

$\rho_L(\delta_L + \delta_R)$	$\rho_L(\delta_L - \delta_R)$	ρ_L	δ_L	δ_R	$\Delta_V^{\mu\mu}$
-0.85	0.5	0.8	-0.219	-0.844	0.03
-0.90	0.7	1.0	-0.1	-0.8	0.02
-0.95	0.9	1.2	-0.021	-0.771	0.02
-0.99	1.0	1.2	0.004	-0.829	0.02

Table 4. The quark-Z' coupling parameters ρ_L and $\delta_{L,R}$ corresponding to a few sample sets of $\rho_L(\delta_L \pm \delta_R)$ and $\Delta_V^{\mu\mu}$ in the $\rho_R = 0.1\,\rho_L$ and $m_{Z'} = 1\,\text{TeV}$ case.

inferred in ref. [97] from the study on LHC bounds in ref. [98]. For the lepton sector, the results of refs. [80, 100] imply that the selected $\Delta_V^{\mu\mu}$ values are compatible with LEP data on Z-boson decays into lepton pairs [10]. Furthermore, as the $\bar{e}eZ'$ interaction is supposed to be vanishing, restraints from LEP II measurements on $e^+e^- \to f\bar{f}$ can be evaded. Lastly, LHC searches for new high-mass phenomena in the dilepton final states have the potential for significantly probing $\delta_{L,R}$ and $\Delta_V^{\mu\mu}$ at the same time. Nevertheless, their samples values in table 4 can be checked to be consistent with the most recent $pp \to \ell^+\ell^- X$ results from the ATLAS experiment [99].[7]

4 Conclusions

We have explored the possibility that the recently observed anomalies in several $b \to s\mu^+\mu^-$ processes are attributable to the interactions of a Z' boson which also contribute to rare nonleptonic decays of the $\bar{B}_s$ meson, namely $\bar{B}_s \to (\eta, \eta', \phi)(\pi^0, \rho^0)$. Given that the amplitudes for these purely isospin-violating decays have CKM-suppressed tree components and tend to be controlled mainly by the electroweak-penguin operators, their decay rates are expected to be relatively small in the SM, making these modes potentially sensitive to signals beyond the SM. The Z' couplings are subject to various restrictions, particularly from the data on B_s-$\bar{B}_s$ mixing and the new experimental finding on $\bar{B}_s^0 \to \phi\rho^0$, besides the measurements of $b \to s\mu^+\mu^-$ transitions. We showed that, within the allowed parameter space, the Z' impact on $\bar{B}_s^0 \to (\eta, \phi)\pi^0$ can cause their rates to grow up to an order of magnitude greater than their expectations in the SM. On the other hand, the enhancement factors for $\bar{B}_s^0 \to \eta'\pi^0, (\eta, \eta')\rho^0$ are at most a few. The different enlargement factors of these different channels depend not only on the combinations of the Z' couplings occurring in their amplitudes, but also on how the SM and Z' terms in the amplitudes interfere with each other. Therefore, the observations of more of these decays in future experiments, together with improved upcoming data on $b \to s\mu^+\mu^-$, will test our Z' model more comprehensively.

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[7]We may test our Z' coupling choices with the latest ATLAS [99] constraint on a nonstandard quark-muon contact interaction of the form $\mathcal{L} = (4\pi/\Lambda^2)\eta_{\chi\chi'}\,\bar{q}_\chi\gamma^\beta q_\chi\,\bar{\mu}_{\chi'}\gamma_\beta\mu_{\chi'}$, where Λ is a heavy mass scale, $\eta_{\chi\chi'} = -1(1)$ if the new and SM contributions to $q\bar{q} \to \mu^+\mu^-$ interfere constructively (destructively), and $\chi, \chi' = L, R$. It turns out that the strongest restriction applies to the $\chi\chi' = RL$ or RR case and arises from the 95%-CL limit $\Lambda > 28\,\text{TeV}$ [99] corresponding to $4\pi/\Lambda^2 < 0.016\,\text{TeV}^{-2}$. This can be fulfilled by $\left|\Delta_R^{qq}\Delta_V^{\mu\mu}\right|$ for $q = u, d$ and the entries in the last two columns of table 4 with selections such as $\Delta_R^{uu} = -\Delta_R^{dd} = \delta_R/2$.

Acknowledgments

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