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B_s-$\overline{B}_s$ Mixing in Z' Models with Flavor-Changing Neutral Currents

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Abstract

In models with an extra $U(1)'$ gauge boson, the family non-universal couplings to the weak eigenstates of the standard model fermions generally induce flavor-changing neutral currents. This phenomenon leads to interesting results in various B meson decays, for which recent data indicate hints of new physics involving significant contributions from $b \to s$ transitions. We analyze the B_s system, emphasizing the effects of Z' on the mass difference and CP asymmetries.

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I. INTRODUCTION

The study of B physics and the associated CP violating observables has been suggested as a good means to extract information of new physics at low energy scales [1–6]. Since B-$\overline{B}$ mixing is a loop-mediated process within the standard model (SM), it offers a good opportunity to see the footprints of physics beyond the SM. The currently observed $\Delta M_{B_d} = 0.489 \pm 0.008$ ps^{-1} [7] and its mixing phase $\sin 2\beta = 0.736 \pm 0.049$ extracted from the $J/\psi K_S$ mode [8] agree well with constraints obtained from other experiments [9]. However, no such information other than a lower bound $\Delta M_{B_s} > 14.4$ ps^{-1} [10] is available for the B_s meson yet.

Based upon the SM predictions, ΔM_{B_s} is expected to be slightly larger than 20 ps^{-1} and its mixing phase ϕ_s is only a couple of degrees. In contrast to the B_d system, its more than 25 times larger oscillation frequency and a factor of four less rate of hadronization from b quarks pose the primary challenges in the study of B_s oscillation and CP asymmetries. Since the $B_s \to J/\psi\phi$ decay is dominated by a CKM favored tree-level process, $b \to c\bar{c}s$, that does not involve any new weak phase in the SM, its asymmetry provides the most reliable information about the mixing phase ϕ_s. Although new physics contributions may not compete with the SM processes in most of the $b \to c\bar{c}s$ decays, they can play an important role in the B_s-$\overline{B}_s$ mixing because its loop nature in the SM. In particular, such a mixing can be significantly modified in models where a tree-level bottom-strange quark couping is allowed. It is thus seen that measuring the properties of B_s meson mixing is of high interest in future B physics studies and viewed as a means to reveal new physics [11, 12]. Since the current B factories do not run at the $\Upsilon(5S)$ resonance to produce B_s mesons, we therefore consider it one of the primary objectives of hadronic colliders to study B_s oscillation and decay in the coming years [13, 14].

In E_6 models, flavor changing neutral currents (FCNC) through an extra $U(1)'$ gauge boson can arise when the Z' couplings to physical fermion eigenstates are non-diagonal [15–17]. This is achieved through the introduction of the exotic fermions with different $U(1)'$ charges that mix with the SM fermions. However, to avoid inducing undesired FCNC mediated by the SM Z boson, one is restricted to schemes where only the right-handed fermions are mixed with the exotic fermions [18].

It is well-known that string models naturally give extra $U(1)'$ groups, at least one of which have family non-universal couplings with the SM fermions [19–22]. Generically, the physical and gauge eigenstates do not coincide. Therefore, unlike in the above-mentioned E_6 setup, the off-diagonal couplings of fermions with the Z' boson without mixing with additional states can be obtained. In such models, both left-handed and right-handed fermions can have family non-diagonal couplings with the Z' while their couplings with the Z remain family diagonal. Moreover, by a suitable construction of intersecting branes, it is possible to have diminishing $U(1)'$ charges for SM leptons [23]. In this paper, we will consider such leptophobic models with Z'-mediated FCNC in the quark sector.

Recently, we have studied the implications of a sizeable off-diagonal Z' coupling between the bottom and strange quark in the indirect CP asymmetry of $B \to \phi K_S$ decay [24], which is seen to have a significant deviation from the SM prediction [5, 6, 25, 26]. Here we want to extend our analysis to the B_s-$\overline{B}_s$ mixing where the Z' contribution also enters at the tree-level.

The paper is organized as follows. We review in Section II the basic formalism of B_s-$\overline{B}_s$ mixing. In Section III, we evaluate ΔM in the SM. In Section IV, we include the Z' contributions, both left-handed and right-handed couplings, in the mixing. Our main results are summarized in Section VI.

II. B_s-$\overline{B}_s$ MIXING

If we write the heavy and light eigenstates as

$$
\begin{aligned}
|B_s\rangle_L &= p|B_s^0\rangle + q|\overline{B}_s^0\rangle \ , \\
|B_s\rangle_H &= p|B_s^0\rangle - q|\overline{B}_s^0\rangle \ , \qquad (1)
\end{aligned}
$$

then the mixing factor

$$
\left(\frac{q}{p}\right)_{\rm SM} \simeq \sqrt{\frac{M_{12}^{\rm SM*}}{M_{12}^{\rm SM}}} \ , \qquad (2)
$$

has a phase

$$
\phi_s = 2\arg(V_{tb}V_{ts}^*) = -2\lambda^2\eta = \mathcal{O}(-2^\circ) \ , \qquad (3)
$$

where $\Gamma_{12}^{\rm SM} \ll M_{12}^{\rm SM}$ is used. Here the off-diagonal element of the decay matrix, $\Gamma_{12}^{\rm SM}$, is evaluated by considering decay channels that are common to both B_s and $\overline{B}_s$ mesons, and M_{12} is the off-diagonal element of the mass matrix. It is dominated by the charm-quark contributions over the up quarks in the box diagram due to the CKM enhancement. Unlike the kaon system, $\Gamma_{12}^{\rm SM}$ is much smaller than $M_{12}^{\rm SM}$ for B mesons. This is because the former is related to the B meson decays and thus set by the scale of its mass, whereas the latter is proportional to m_t^2. We can safely assume that Γ_{12} is not significantly modified by new physics, because Γ_{12} receives major contributions from CKM favored $b \to c\bar{c}s$ decays in SM. Therefore, the relation $\Gamma_{12} \ll M_{12}$ is unlikely to change.

The mass difference of the two physical states is

$$\Delta M \equiv M_H - M_L \simeq 2|M_{12}| \ . \tag{4}$$

The width difference is

$$\Delta\Gamma \equiv \Gamma_H - \Gamma_L = \frac{2{\rm Re}(M_{12}^*\Gamma_{12})}{|M_{12}|} = 2|\Gamma_{12}|\cos\theta \ , \tag{5}$$

where the relative phase $\theta = \arg(M_{12}/\Gamma_{12})$. Since Γ_{12} is dominated by the contributions from CKM favored $b \to c\bar{c}s$ decays, we have $\theta = \arg\left((V_{tb}V_{ts}^*)/(V_{cb}V_{cs}^*)\right) \simeq \pi$. Thus, in our convention, $\Delta\Gamma = -2|\Gamma_{12}|$ is negative in the SM. Although Γ_{12} is unlikely to be affected by new physics, the width difference always reduces as long as the weak phase of M_{12} gets modified [27].

The observability of B_s-$\overline{B}_s$ oscillation is often measured by the parameter

$$x_s \equiv \frac{\Delta M}{\Gamma_{B_s}} \ , \tag{6}$$

where $\Gamma_{B_s} = (4.51 \pm 0.18) \times 10^{-13}$ GeV, converted from the world average lifetime $\tau_s = 1.461 \pm 0.057$ ps [7]. Too large a value of x_s will be a challenge for experimental searches. Currently, the result from all ALEPH [28], CDF [29], DELPHI [30], OPAL [31], and SLD [32] studies of ΔM with a combined 95% CL sensitivity on ΔM_s of 17.8 ps^{-1} gives [10]

$$\Delta M > 14.4\,{\rm ps}^{-1} \ , \quad \text{and} \quad x_s > 20.6 \ . \tag{7}$$

It is also measured that $\Delta\Gamma_s/\Gamma_s = 0.16^{+0.15}_{-0.16} (< 0.54)$ with the upper bound in the parentheses quoted at 95% CL [10].

III. ΔM IN SM

First, let us define the $|\Delta B| = 2$ and $|\Delta S| = 2$ operators relevant for later discussions:

$$O^{LL} = [\bar{s}\gamma_\mu(1-\gamma_5)b][\bar{s}\gamma^\mu(1-\gamma_5)b] \ ,$$

$$
\begin{aligned}
O_1^{LR} &= [\bar{s}\gamma_\mu(1-\gamma_5)b][\bar{s}\gamma^\mu(1+\gamma_5)b] \ , \\
O_2^{LR} &= [\bar{s}(1-\gamma_5)b][\bar{s}(1+\gamma_5)b] \ , \\
O^{RR} &= [\bar{s}\gamma_\mu(1+\gamma_5)b][\bar{s}\gamma^\mu(1+\gamma_5)b] \ .
\end{aligned} \tag{8}
$$

Because of the $V-A$ structure in the SM weak interactions, only the operator O^{LL} contributes to B_s-$\overline{B}_s$ mixing. The other three operators appear in the Z' models because of the right-handed couplings and operator mixing through renormalization, as we will see in the next section.

In the SM,

$$
M_{12}^{\rm SM} \simeq \frac{1}{2m_{B_s}}\langle B_s^0|\mathcal{H}_{\rm eff}^{\rm SM}|\overline{B}_s^0\rangle \tag{9}
$$

is dominated by the top quark loop. The result accurate to the next-to-leading order (NLO) in QCD is given by [33]

$$
M_{12}^{\rm SM} = \frac{G_F^2}{12\pi^2}M_W^2 m_{B_s} f_{B_s}^2 (V_{tb}V_{ts}^*)^2 \eta_{2B} S_0(x_t)[\alpha_s(m_b)]^{-6/23}\left[1+\frac{\alpha_s(m_b)}{4\pi}J_5\right] B_{LL}(m_b) \ , \tag{10}
$$

where $x_t = (m_t(\mu_t)/M_W)^2$ and

$$
S_0(x) = \frac{4x-11x^2+x^3}{4(1-x)^2} - \frac{3x^3\ln x}{2(1-x)^3} \ . \tag{11}
$$

Using $m_t(\mu_t) = \mu_t = 170 \pm 5$GeV we find the numerical value of S_0 at x_t is $S_0(x_t) = 2.463$. The NLO short-distance QCD corrections are encoded in $\eta_{2B} \simeq 0.551$ and $J_5 \simeq 1.627$ [33]. The bag parameter $B_{LL}(\mu)$ is defined through the relation

$$
\langle \overline{B}_s|O_{LL}|B_s\rangle \equiv \frac{8}{3}m_{B_s}^2 f_{B_s}^2 B_{LL}(\mu) \ . \tag{12}
$$

Recent lattice analyses give the hadronic parameters $f_{B_s} = 230 \pm 30$ MeV and $B_{LL}(m_b) = 0.872 \pm 0.005$ [34–36].

For numerical estimates throughout this paper, we will use $G_F = 1.16639 \times 10^{-5}$ GeV^{-2}, $m_{B_s} = 5369.6 \pm 2.4$ MeV, and $M_W = 80.423 \pm 0.039$ GeV, and the Wolfenstein parameters [37] extracted various experiments are $\lambda = 0.2240 \pm 0.0036$, $A = 0.83 \pm 0.02$, $\rho = 0.216 \pm 0.079$, and $\eta = 0.341 \pm 0.028$ [38]. We find

$$
\begin{aligned}
\Delta M^{\rm SM} &= (1.32 \pm 0.35) \times 10^{-11}\ \text{GeV} \\
&= 20.0 \pm 5.4\ \text{ps}^{-1} \ , \\
x_s^{\rm SM} &= 29 \pm 8 \ .
\end{aligned} \tag{13}
$$

The main sources of errors are from m_t and f_{B_s} in ΔM. For As noted before, the central value is slightly larger than the current sensitivity based upon the world average.

Latest studies show that with one year of data, ΔM can be explored up to 30 ps^{-1} (ATLAS), 26 ps^{-1} (CMS), and 48 ps^{-1} (LHCb) (corresponding to x_s up to 46, 42, and 75) using exclusive hadronic modes at the LHC [14]. Assuming a luminosity of 2 fb^{-1} in a one-year run, the sensitivity of both BTeV and CDF on x_s can also reach up to 75 using the same modes [13].

IV. Z' CONTRIBUTIONS

For simplicity, we assume that the $U(1)'$ gauge group is orthogonal to the SM gauge group so that there is no mixing between the SM Z and the Z'. A purely left-handed off-diagonal Z' coupling to b and s quarks results in an effective $|\Delta B| = 2$, $|\Delta S| = 2$ Hamiltonian at M_W scale

$$\mathcal{H}_{\rm eff}^{Z'} = \frac{G_F}{\sqrt{2}} \left(\frac{g' M_Z}{g_1 M_{Z'}} B_{sb}^L \right)^2 O_{LL}(m_b) \equiv \frac{G_F}{\sqrt{2}} \rho_L^2 e^{2i\phi_L} O_{LL}(m_b) \ , \tag{14}$$

where g' is the $U(1)'$ gauge coupling, $g_1 = e/\sin\theta_W$, $M_{Z'}$ is the mass of the Z', and B_{sb}^L is the FCNC Z' coupling between the bottom and strange quarks. The parameters ρ_L and the weak phase ϕ_L in the Z' model are defined in the second part of the equation. Note that the Z' does not contribute to Γ_{12} at tree level because the intermediate Z' cannot be on shell. After evolving from the M_W scale to the m_b scale, the effective Hamiltonian becomes

$$\mathcal{H}_{\rm eff}^{Z'} = \frac{G_F}{\sqrt{2}} \left[1 + \frac{\alpha_s(m_b) - \alpha_s(m_w)}{4\pi} J_5 \right] R^{6/23} \rho_L^2 e^{2i\phi_L} O_{LL}(m_b) \ , \tag{15}$$

where $R = \alpha_s(M_W)/\alpha_s(m_b)$.

One immediately notices that although the above effective Hamiltonian is largely suppressed by the ratio $(g' M_Z)/(g_1 M_{Z'})$, it has only one power of G_F in comparison with the corresponding quadratic dependence in the SM. This is because the Z'-mediated process occurs at tree level.

The full description of the running of the Wilson coefficient from the M_W scale to m_b scale can be found in [33]. We only repeat the directly relevant steps here.

The renormalization group equation for the Wilson coefficients $\vec{C}$

$$\frac{d}{d\ \ln\mu} \vec{C} = \gamma^T(g) \vec{C}(\mu) \tag{16}$$

can be solved with the help of the U matrix

$$\vec{C}(\mu) = U(\mu, M_W)\vec{C}(M_W), \tag{17}$$

in which $\gamma^T(g)$ is the transpose of the anomalous dimension matrix $\gamma(g)$. With the help of $dg/d\ln\mu = \beta(g)$, U obeys the same equation as $\vec{C}(\mu)$. If we expand $\gamma(g)$ to the first two terms in the perturbative expansion,

$$\gamma(\alpha_s) = \gamma^{(0)}\frac{\alpha_s}{4\pi} + \gamma^{(1)}\left(\frac{\alpha_s}{4\pi}\right)^2 . \tag{18}$$

To this order the evolution matrix $U(\mu, m)$ is given by

$$U(\mu, m) = (1 + \frac{\alpha_s(\mu)}{4\pi}J)U^{(0)}(\mu, m)(1 - \frac{\alpha_s(m)}{4\pi}J) \tag{19}$$

$U^{(0)}$ is the evolution matrix in leading logarithmic approximation and the matrix J expresses the next-to-leading corrections to this evolution. We have

$$U^{(0)}(\mu, m) = V\left(\left[\frac{\alpha_s(m)}{\alpha_s(\mu)}\right]^{\frac{\vec{\gamma}^{(0)}}{2\beta_0}}\right)_D V^{-1} \tag{20}$$

where V diagonalizes $\gamma^{(0)T}$

$$\gamma_D^{(0)} = V^{-1}\gamma^{(0)T}V \tag{21}$$

and $\vec{\gamma}^{(0)}$ is the vector containing the diagonal elements of the diagonal matrix $\gamma_D^{(0)}$.

If we define

$$G = V^{-1}\gamma^{(1)T}V \tag{22}$$

and a matrix H whose elements are

$$H_{ij} = \delta_{ij}\gamma_i^{(0)}\frac{\beta_1}{2\beta_0^2} - \frac{G_{ij}}{2\beta_0 + \gamma_i^{(0)} - \gamma_j^{(0)}} \tag{23}$$

the matrix J is given by

$$J = VHV^{-1} \tag{24}$$

The operators O_{LL} and O_{RR} do not mix with others under renormalization. Their Wilson coefficients follow exactly the same RGE, where the above-mentioned matrices are all simple numbers. The factor

$$\left[1 + \frac{\alpha_s(m_b) - \alpha_s(M_W)}{4\pi}J_5\right]R^{6/23} \tag{25}$$

in Eqn. (15) reflects the RGE running. On the other hand, O_1^{LR} and O_2^{LR} form a sector mixed under RG running. Although the Z' boson only induces the operator O_1^{LR} at high energy scales, O_2^{LR} is generated after evolving down to low energy scales and, in particular, its Wilson coefficient C_2^{LR} is strongly enhanced by the RG effects [39].

With both contributions from the SM and the Z' boson with only left-handed FCNC couplings included, the mass difference

$$\begin{aligned}\Delta M &= \Delta M^{SM}\left(1+\frac{\Delta M^{Z'}}{\Delta M^{SM}}\right)\\ &= 20.0\,|1+3.566\times 10^5\rho_L^2 e^{2i\phi_L}|\ \mathrm{ps}^{-1}\ ,\end{aligned} \tag{26}$$

where ΔM^{SM} is the contribution in the SM and $\Delta M^{Z'}$ is the contribution from Z'. Similarly, the oscillation parameter

$$x_s(\rho_L,\phi_L) = 29.2\,|1+3.566\times 10^5\rho_L^2 e^{2i\phi_L}|\ . \tag{27}$$

It is noticed that with couplings of only one chirality are considered, the physical observables ΔM, x_s, and $\sin 2\phi_s$ to be considered below are periodic functions of the new weak phase with a period of 180°.

The effect of including a Z' with left-handed coupling is shown in Fig. 1 (a). It is noted that if ρ is too small, x_s is dominated by the SM contribution and has a value ~ 29. For ϕ_L around 90° and small enough ρ_L, the Z' contribution tends to cancel that of the SM and reduces x_s to be smaller than 29.2, the SM value. We show the contour plot of x_s in Fig. 1 (b) in the parameter space of the Z' model. We only show the range of $0^\circ \le \phi_L \le 180^\circ$ because x_s has a period of π in ϕ_L. The region inside of the $x_s = 20.6$ contour is ruled out by the currently measured lower bound on x_s. If the experimental value turns out to be significantly larger than that, one is led to a new physics explanation. From Fig. 1 we see that the Z' contribution dominates if $\rho \gtrsim 0.002$, independent of the actual value of ϕ_L. The planned resolution of Fermilab Run II and LHCb are both about 75 [13, 14]. From another point of view, if x_s is measured to fall within a range, one can read from the plot what is the allowed region for the Z'-model parameters.

In Fig. 2 (a), we show the sine of twice the mixing phase ϕ_s as a function of ρ_L and ϕ_L. The contours of fixed $\sin 2\phi_s$ with values of -0.5, 0.0 and 0.5 are shown as solid, dashed and dotted lines in Fig. 2 (b). Once x_s and $\sin 2\phi_s$ are extracted B_s decays, one can combine

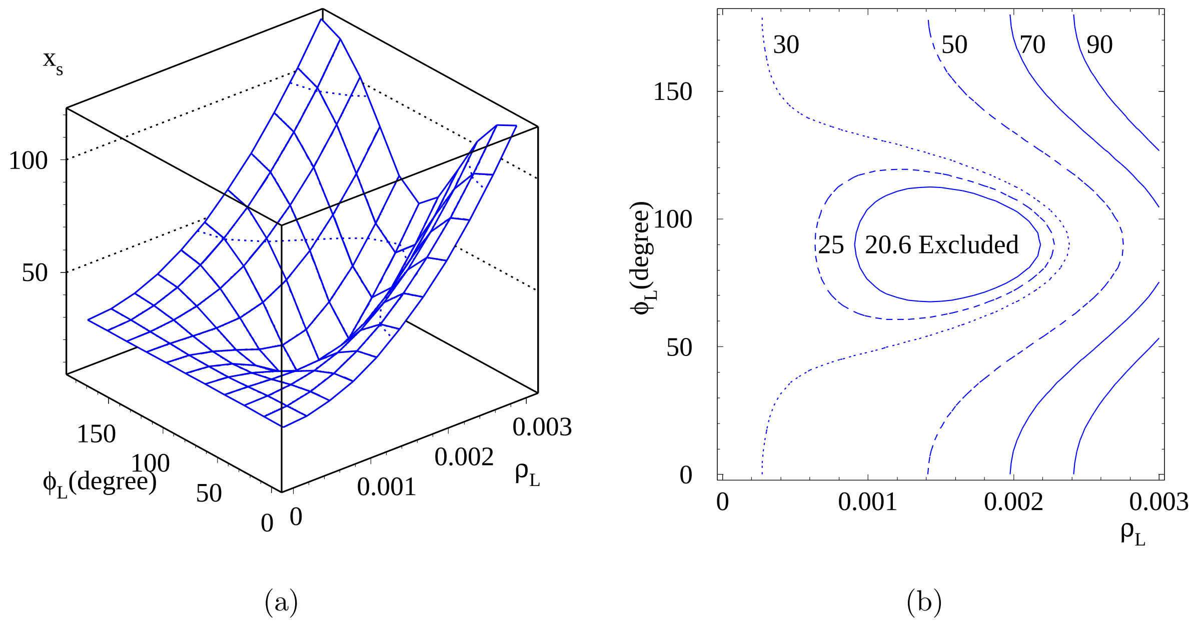


FIG. 1: Plot (a): Three-dimensional plot of x_s in the presence of Z'-mediated FCNC for left-handed b and s quarks as a function of ρ_L and ϕ_L, defined in Eq. (14). Plot (b): A contour plot of x_s in the presence of a Z'-mediated FCNC for left-handed b and s quarks.

Figs. 1 (b) and 2 (b) to determine ρ_L up to a two-fold ambiguity and ϕ_L up to a four-fold ambiguity in general, except for the special case when $\sin 2\phi_s \simeq 0$.

Once the right-handed Z' couplings are introduced, we immediately have the new $|\Delta B| = 2$ operators O_1^{LR}, O_2^{LR}, and O^{RR} defined in Eq. (8) in the effective Hamiltonian that contribute to B_s-$\overline{B}_s$ mixing. The matrix element of O^{RR} is the same as that of O^{LL}, while those of O_1^{LR} and O_2^{LR}, according to [35], are

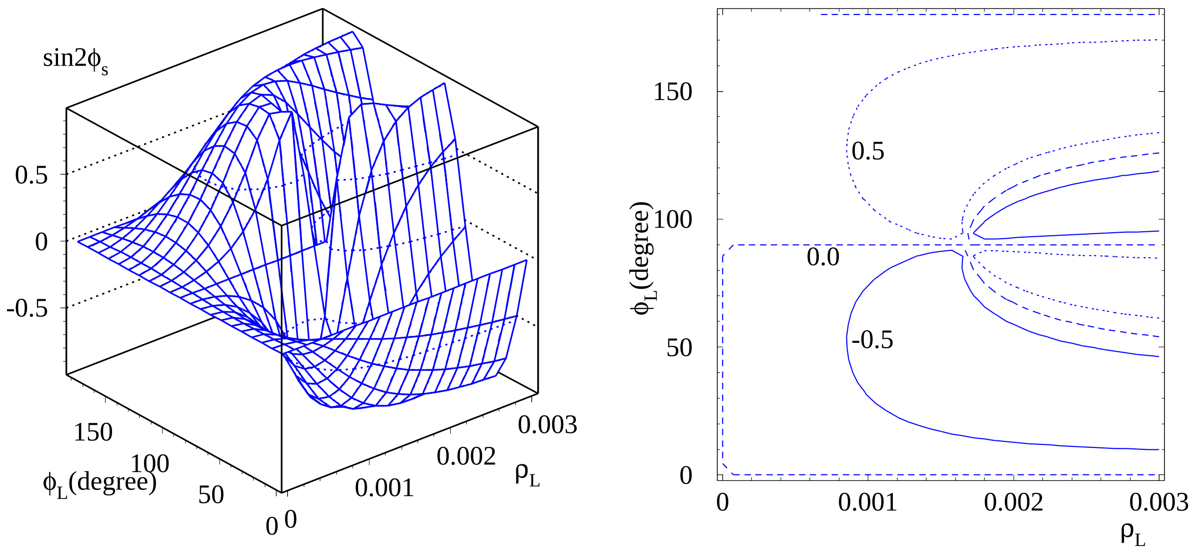
$$\langle \overline{B}_s | O_1^{LR} | B_s \rangle = -\frac{4}{3} \left(\frac{m_{B_s}}{m_b(m_b) + m_s(m_b)} \right)^2 m_{B_s}^2 f_{B_s}^2 B_1^{LR}(m_b) \tag{28}$$

$$\langle \overline{B}_s | O_2^{LR} | B_s \rangle = 2 \left(\frac{m_{B_s}}{m_b(m_b) + m_s(m_b)} \right)^2 m_{B_s}^2 f_{B_s}^2 B_2^{LR}(m_b) \tag{29}$$

For the Z' coupling to right handed currents, we define new parameters ρ_R and weak phase ϕ_R

$$\rho_R e^{i\phi_R} \equiv \frac{g' M_Z}{g_1 M_{Z'}} B_{sb}^R \ . \tag{30}$$

At the M_W scale, due to the right handed currents, we have addition contribution to the effective Hamiltonian that is similar to Eqn. (14). The effective Hamiltonian due to the



(a) (b)

FIG. 2: Plot (a): Three-dimensional plot of $\sin 2\phi_s$ in the presence of Z'-mediated FCNC for left-handed b and s quarks as a function of ρ_L and ϕ_L, defined in Eq. (14). Plot (b): A contour plot of $\sin 2\phi_s$ in the presence of a Z'-mediated FCNC for left-handed b and s quarks.

lift-right mixing is

$$\mathcal{H}_{\text{eff}}^{\mathcal{Z}'} = \frac{G_F}{\sqrt{2}}\rho_L\rho_R e^{-i(\phi_L-\phi_R)}(O_1^{LR}, O_2^{LR})\begin{pmatrix} 1 \\ 0 \end{pmatrix} . \tag{31}$$

In the RGE running, the Wilson coefficient for O_1^{LR} will mix with that of O_2^{LR} and the relevant anomalous dimension matrices are [39]

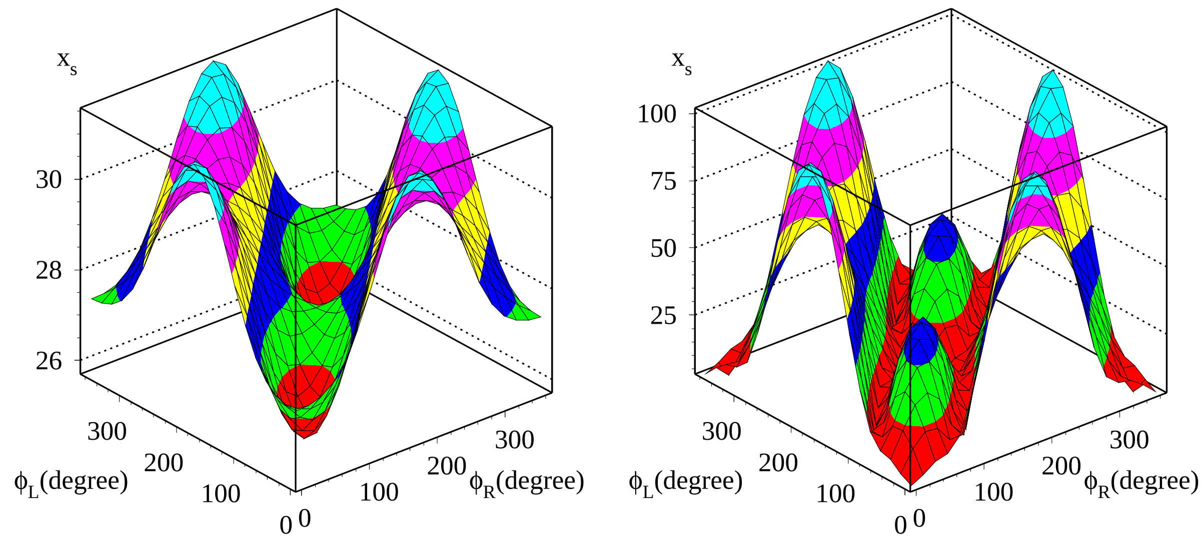
$$\gamma^{(0)} = \begin{pmatrix} \frac{6}{N_c} & 12 \\ 0 & -6N_c + \frac{6}{N_c} \end{pmatrix} , \quad \text{and} \tag{32}$$

$$\gamma^{(1)} = \begin{pmatrix} \frac{137}{6} + \frac{15}{2N_c^2} - \frac{22}{3N_c}f & \frac{200}{3}N_c - \frac{6}{N_c} - \frac{44}{3}f \\ \frac{71}{4} + \frac{9}{N_c} - 2f & -\frac{203}{6}N_c^2 + \frac{479}{6} + \frac{15}{2N_c^2} + \frac{10}{3}N_c f - \frac{22}{3N_c}f \end{pmatrix} , \tag{33}$$

where N_c is the number of colors and f is the number of active quarks. At the scale of the B meson, we have $f = 5$.

We will use $m_b(m_b) = 4.4$ GeV, $\Lambda_{\overline{MS}}^{(5)} = 225$ MeV and $m_b(m_b) + m_s(m_b) = 4.6$ GeV. Following Eqns. (20-24), we find the effective Hamiltonian at m_b for the operator $O_{1,2}^{LR}$ to be

$$\mathcal{H}_{\text{eff}}^{\mathcal{Z}'} = \frac{G_F}{\sqrt{2}}\rho_R\rho_L e^{-i(\phi_L-\phi_R)}(O_1^{LR}, O_2^{LR})\begin{pmatrix} 0.930 \\ -0.711 \end{pmatrix} . \tag{34}$$



(a) (b)

FIG. 3: Plot (a): x_s as a function of ϕ_L and ϕ_R for $\rho_L = \rho_R = 0.0002$. Plot (b): x_s as a function of ϕ_L and ϕ_R for $\rho_L = \rho_R = 0.001$.

Note that at M_W scale, the operator O_2^{LR} does not contribute. Through the operator mixing in RGE running, its effect becomes important at the m_b scale.

The bag parameters from Ref. [35] are $B_1^{LR}(m_b) = 1.753 \pm 0.021$ and $B_2^{LR}(m_b) = 1.162 \pm 0.007$ and the decay constant f_{B_s} is the same as before. The mass difference with all Z' contributions included is

$$\Delta M = 20.0\,1 + 3.566 \times 10^5 \rho_L^2 e^{2i\phi_L} + 3.566 \times 10^5 \rho_R^2 e^{2i\phi_R} - \\ 1.851 \times 10^6 \rho_L \rho_R \cos(\phi_L - \phi_R)|\ \mathrm{ps}^{-1}\ . \tag{35}$$

The overall contribution for the SM and Z' to x_s is,

$$x_s = 29.2\,|1 + 3.566 \times 10^5 \rho_L^2 e^{2i\phi_L} + 3.566 \times 10^5 \rho_R^2 e^{2i\phi_R} - \\ 1.851 \times 10^6 \rho_L \rho_R \cos(\phi_L - \phi_R)|\ . \tag{36}$$

To see the interference among different contributions, we set $\rho_L = \rho_R = 0.0002$ or 0.001 and plot x_s versus the weak phases ϕ_L and ϕ_R in Fig. 3.

First we note that after the RGE running, the operators O_1^{LR} and O_2^{LR} interfere constructively. After adding the corresponding contributions from the RL counterparts, they become the most dominant term. As one of the phase approaching 180° and the other phase 0°, the three terms from Z' all contribute positively. When the phases are close to be equal, the most dominant contribution approaches 0.

V. MISC

The inclusive branching ratio $b \to s\gamma$ is computed to the NLO in the SM to be $(3.79^{+0.39}_{-0.53}) \times 10^{-4}$ [41], where uncertainties of CKM parameters, charm mass and scale dependence are taken into account. Thus, the SM predicts the ratio

$$\left(\frac{BR(B_d \to X_s\gamma)}{\Delta M}\right)_{\rm SM} = (2.94 \pm 1.03) \times 10^7 \ {\rm GeV}^{-1} \ . \tag{37}$$

Possible things to discuss

1. **Include $Br(b \to s\gamma)^{SM}/\Delta M^{SM}_{B_s}$ and $Br(b \to s\mu^+\mu^-)^{SM}/\Delta M^{SM}_{B_s}$.**

2. **Discuss $\Delta\Gamma$.**

VI. CONCLUSION

Conclusions here.

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