

# CP violating dimuon charge asymmetry in general left-right models

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The recently measured charge asymmetry of like-sign dimuon events by the D0 collaboration at Tevatron shows the 3.9  $\sigma$  deviation from the standard model prediction. In order to solve this mismatch, we investigate the right-handed current contributions to  $B_s - \bar{B}_s$  and  $B_d - \bar{B}_d$  mixings which are the major source of the like-sign dimuon events in  $b\bar{b}$  production in general left-right models without imposing manifest or pseudo-manifest left-right symmetry. We find the allowed region of new physics parameters satisfying the current experimental data.

## I. INTRODUCTION

Recently the D0 collaboration has measured a deviation from the standard model (SM) prediction in the CP violating like-sign dimuon charge asymmetry in semileptonic  $b$  hadron decay with the  $9 \text{ fb}^{-1}$  integrated luminosity of  $p\bar{p}$  data at Tevatron [1]:

$$A_{sl}^b = -0.00787 \pm 0.00172 \text{ (stat.)} \pm 0.00093 \text{ (syst.)}. \quad (1)$$

The like-sign dimuon events comes from direct semileptonic decays of one of  $b$  hadrons and a semileptonic decay of the other  $b$  hadron following the  $B^0 - \bar{B}^0$  oscillation in  $b\bar{b}$  pair production at Tevatron, defined by

$$A_{sl}^b \equiv \frac{\Gamma(b\bar{b} \rightarrow \mu^+ \mu^+ X) - \Gamma(b\bar{b} \rightarrow \mu^- \mu^- X)}{\Gamma(b\bar{b} \rightarrow \mu^+ \mu^+ X) + \Gamma(b\bar{b} \rightarrow \mu^- \mu^- X)}. \quad (2)$$

At Tevatron experiment, both decays of  $B_d$  and  $B_s$  mesons contribute to the asymmetry. If we define the charge asymmetry of semileptonic decays of neutral  $B_q^0$  mesons as

$$a_{sl}^q \equiv \frac{\Gamma(\bar{B}_q^0(t) \rightarrow \mu^+ X) - \Gamma(B_q^0(t) \rightarrow \mu^- X)}{\Gamma(\bar{B}_q^0(t) \rightarrow \mu^+ X) + \Gamma(B_q^0(t) \rightarrow \mu^- X)}, \quad (3)$$

the like-sign dimuon charge asymmetry can be expressed in terms of  $a_{sl}^q$  as [2]

$$A_{sl}^b = \frac{1}{f_d Z_d + f_s Z_s} (f_d Z_d a_{sl}^d + f_s Z_s a_{sl}^s), \quad (4)$$

assuming that  $\Gamma(B_d^0 \rightarrow \mu^+ X) = \Gamma(B_s^0 \rightarrow \mu^+ X)$  to a very good approximation, where  $f_q$  are the production fractions of  $B_q$  mesons, and  $Z_q = 1/(1 - y_q^2) - 1/(1 + x_q^2)$  with  $y_q = \Delta\Gamma_q/(2\Gamma_q)$ ,  $x_q = \Delta M_q/\Gamma_q$ . These parameters are measured to be  $f_d = 0.402 \pm 0.013$ ,  $f_s = 0.112 \pm 0.013$ ,  $x_d = 0.771 \pm 0.007$ ,  $x_s = 26.3 \pm 0.4$ , and  $y_d = 0$ ,  $y_s = 0.052 \pm 0.016$  [3]. With these values, Eq. (4) is rewritten by

$$A_{sl}^b = (0.572 \pm 0.030)a_{sl}^d + (0.428 \pm 0.030)a_{sl}^s. \quad (5)$$

The non-zero dimuon asymmetry implies a difference between the  $B^0 \leftrightarrow \bar{B}^0$  transitions and the CP violation in the  $B$  system. In the SM, the source of the CP violation in the neutral  $B_q^0$  system is the phase of the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements involved in the box diagram. Using the SM values for the semileptonic charge asymmetries  $a_{sl}^d$  and  $a_{sl}^s$  of  $B_d^0$  and  $B_s^0$  mesons, respectively [4], the prediction of the dimuon asymmetry in the SM is given by

$$A_{sl}^b = (-2.7_{-0.6}^{+0.5}) \times 10^{-4}, \quad (6)$$

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which shows that the D0 measurement of Eq. (1) deviates about  $3.9 \sigma$  from the SM prediction. If the deviation is confirmed with other experiments, it indicates the existence of the new physics beyond the SM. Recently several works are devoted to explaining the D0 dimuon asymmetry in the SM and beyond [5].

As an alternative model solution to the mismatch between the measurement and the SM prediction of the dimuon charge asymmetry, we consider the left-right model (LRM) based on the  $SU(2)_L \times SU(2)_R \times U(1)$  gauge symmetry which is one of the attractive extensions of the SM [6]. The current measurement of the dimuon charge asymmetry can be explained in the LRM due to the sizable right-handed current contributions to  $B^0 - \bar{B}^0$  mixing [7]. The manifest left-right symmetry provides an natural answer to the origin of the parity violation. Involving the triplet Higgs field  $\Delta_{L,R}$  to break the additional  $SU(2)_R$  symmetry, the lepton number violating Yukawa terms are introduced and the see-saw mechanism for light neutrino masses is exploited in the LRM. This model arises as an intermediate theory in the SO(10) grand unified theory (GUT). In the LRM, the right-handed fermions transform as doublets under  $SU(2)_R$  and singlets under  $SU(2)_L$ , and the left-handed fermions behave reversely. Thus a bidoublet Higgs field is required for the Yukawa couplings and also responsible for the electroweak symmetry breaking (EWSB). The scale of the masses of the new gauge bosons in the LRM is constrained by direct searches and indirect analysis [8–11], and we will discuss the constraints on the model in further detail.

This paper is organized as follows. In section 2, we briefly review the charged sector in the general LRM. We explicitly show the right-handed current contributions in the neutral  $B$  meson system in section 3, and present the numerical analysis of  $B^0 - \bar{B}^0$  mixing and the dimuon charge asymmetry of  $B$  mesons in the general LRM in section 4. Finally we conclude in section 5.

## II. THE LEFT-RIGHT MODEL

We briefly review the main features of the LRM, which are necessary for our analysis. The gauge group of the left-right symmetric model is  $SU(2)_L \times SU(2)_R \times U(1)$ . There exist a bidoublet Higgs field  $\phi(2, \bar{2}, 0)$  and two triplet Higgs fields,  $\Delta_L(3, 1, 2)$  and  $\Delta_R(1, 3, 2)$  in the minimal LR model represented by

$$\phi = \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix}, \quad \Delta_{L,R} = \frac{1}{\sqrt{2}} \begin{pmatrix} \delta_{L,R}^+ & \sqrt{2}\delta_{L,R}^{++} \\ \sqrt{2}\delta_{L,R}^0 & -\delta_{L,R}^+ \end{pmatrix}, \quad (7)$$

of which kinetic terms are given by

$$\mathcal{L} = \text{Tr} [(D_\mu \Delta_{L,R})^\dagger (D^\mu \Delta_{L,R})] + \text{Tr} [(D_\mu \phi)^\dagger (D^\mu \phi)], \quad (8)$$

where the covariant derivatives are defined by

$$\begin{aligned} D_\mu \phi &= \partial_\mu \phi - i \frac{g_L}{2} W_{L\mu}^a \tau^a \phi + i \frac{g_R}{2} \phi W_{R\mu}^a \tau^a, \\ D_\mu \Delta_{L,R} &= \partial_\mu \Delta_{L,R} - i \frac{g_{L,R}}{2} [W_{L,R\mu}^a \tau^a, \Delta_{L,R}] - i g' B_\mu \Delta_{L,R}. \end{aligned} \quad (9)$$

The gauge symmetries are spontaneously broken by the vacuum expectation values (VEV)

$$\langle \phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} k_1 & 0 \\ 0 & k_2 \end{pmatrix}, \quad \langle \Delta_{L,R} \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ v_{L,R} & 0 \end{pmatrix}, \quad (10)$$

where  $k_{1,2}$  are complex in general and  $v_{L,R}$  are real, which lead to the charged gauge boson masses

$$\mathcal{M}_{W^\pm}^2 = \frac{1}{4} \begin{pmatrix} g_L^2(k_+^2 + 2v_L^2) & -2g_L g_R k_1^* k_2 \\ -2g_L g_R k_1 k_2^* & g_R^2(k_+^2 + 2v_R^2) \end{pmatrix} = \begin{pmatrix} M_{W_L}^2 & M_{W_{LR}}^2 e^{i\alpha} \\ M_{W_{LR}}^2 e^{-i\alpha} & M_{W_R}^2 \end{pmatrix}, \quad (11)$$

where  $k_+^2 = |k_1|^2 + |k_2|^2$  and  $\alpha$  is the phase of  $k_1^* k_2$ . Since the  $SU(2)_R$  breaking scale  $v_R$  should be higher than the electroweak scale,  $k_{1,2} \ll v_R$ ,  $W_R$  is heavier than  $W_L$ . Note that  $v_L$  is irrelevant for the symmetry breaking and just introduced in order to manifest the left-right symmetry. If the neutrino mass is purely determined by the see-saw relation  $m_\nu \sim v_L + k_+^2/v_R$ ,  $v_R$  is typically very large  $\sim 10^{11}$  GeV. It indicates that the heavy gauge bosons are too heavy to be produced at the accelerator experiments and the direct search of the  $SU(2)_R$  structure is hardly achieved. Therefore we assume that  $v_R$  is only moderately large,  $v_R \sim \mathcal{O}(\text{TeV})$ , for the heavy gauge bosons to be found at the LHC, and the Yukawa couplings are suppressed in order that the neutrino masses are at the eV scale. We let  $v_L$  be very small or close to 0 without loss of generality. This is achieved when the quartic couplings of  $(\phi\phi\Delta_L\Delta_R)$ -type

terms in the Higgs potential are set to be zero [12, 13] and warranted by the approximate horizontal U(1) symmetry [14] as well as the see-saw picture for light neutrino masses. We adopt this limit here and note that the Higgs boson masses are not affected by taking this limit [13].

The general Higgs potential in the LRM has been studied in Refs. [12, 13, 15]. After the mass matrix is diagonalized by a unitary transformation, the mass eigenstates are written as

$$\begin{pmatrix} W^\pm \\ W'^\pm \end{pmatrix} = \begin{pmatrix} \cos \xi & e^{-i\alpha} \sin \xi \\ -\sin \xi & e^{-i\alpha} \cos \xi \end{pmatrix} \begin{pmatrix} W_L^\pm \\ W_R^\pm \end{pmatrix}, \quad (12)$$

with the mixing angle

$$\tan 2\xi = -\frac{2M_{W_{LR}}^2}{M_{W_R}^2 - M_{W_L}^2}. \quad (13)$$

For  $v_R \gg |k_{1,2}|$ , the mass eigenvalues and the mixing angle reduce to

$$M_W^2 \approx \frac{1}{4}g_L^2(|k_1|^2 + |k_2|^2), \quad M_{W'}^2 \approx \frac{1}{2}g_R^2v_R^2, \quad \xi \approx \frac{g_L|k_1^*k_2|}{g_Rv_R^2}. \quad (14)$$

Here, the Schwarz inequality requires that  $\zeta_g \equiv (g_R/g_L)^2\zeta \geq \xi_g \equiv (g_R/g_L)\xi$  where  $\zeta \equiv M_{W'}^2/M_W^2$ . From the global analysis of muon decay measurements [16], the lower bound on  $\zeta_g$  can be obtained without imposing discrete symmetry as follows:

$$\zeta_g < 0.031 \quad \text{or} \quad M_{W'} > (g_R/g_L) \times 460 \text{ GeV}. \quad (15)$$

The new gauge boson mass  $M_{W'}$  is severely constrained from  $K_L - K_S$  mixing if the model has manifest ( $V^R = V^L$ ) left-right symmetry ( $g_R = g_L$ ):  $M_{W'} > 2.5$  TeV [17], where  $V^L(V^R)$  is the left(right)-handed quark mixing matrix. But, in general, the form of  $V^R$  is not necessarily restricted to manifest or pseudomanifest ( $V^R = V^{L*}K$ ) symmetric type, where  $K$  is a diagonal phase matrix [6]. Instead, if we take the following form of  $V^R$ , the limit on  $M_{W'}$  may be significantly relaxed to approximately 300 GeV, and the  $W'$  boson contributions to  $B_{d(s)}\bar{B}_{d(s)}$  mixings can be large [18]:

$$V_I^R = \begin{pmatrix} e^{i\omega} & \sim 0 & \sim 0 \\ \sim 0 & c_R e^{i\alpha_1} & s_R e^{i\alpha_2} \\ \sim 0 & -s_R e^{i\alpha_3} & c_R e^{i\alpha_4} \end{pmatrix}, \quad V_{II}^R = \begin{pmatrix} \sim 0 & e^{i\omega} & \sim 0 \\ c_R e^{i\alpha_1} & \sim 0 & s_R e^{i\alpha_2} \\ -s_R e^{i\alpha_3} & \sim 0 & c_R e^{i\alpha_4} \end{pmatrix}, \quad (16)$$

where  $c_R$  ( $s_R$ )  $\equiv \cos \theta_R$  ( $\sin \theta_R$ ) ( $0^\circ \leq \theta_R \leq 90^\circ$ ). Here the matrix elements indicated  $\sim 0$  may be  $\lesssim 10^{-2}$  and the unitarity requires  $\alpha_1 + \alpha_4 = \alpha_2 + \alpha_3$ . From the  $b \rightarrow c$  semileptonic decays of the  $B$  mesons, we can get an approximate bound  $\xi_g \sin \theta_R \lesssim 0.013$  by assuming  $|V_{cb}^L| \approx 0.04$  [19].

### III. $B^0 - \bar{B}^0$ MIXING

The neutral  $B_q$  meson system ( $q = d, s$ ) is described by the Schrödinger equation

$$i \frac{d}{dt} \begin{pmatrix} B_q(t) \\ \bar{B}_q(t) \end{pmatrix} = \left( M - \frac{i}{2} \Gamma \right) \begin{pmatrix} B_q(t) \\ \bar{B}_q(t) \end{pmatrix}, \quad (17)$$

where  $M$  is the mass matrix and  $\Gamma$  the decay matrix. The  $\Delta B = 2$  transition amplitudes

$$\langle B_q^0 | \mathcal{H}_{\text{eff}}^{\Delta B=2} | \bar{B}_q^0 \rangle = M_{12}^q, \quad (18)$$

yields the mass difference between the heavy and the light states of  $B$  meson,

$$\Delta M_q \equiv M_H^q - M_L^q = 2|M_{12}^q|, \quad (19)$$

where  $M_H^q$  and  $M_L^q$  are the mass eigenvalues for the heavy and the light eigenstates, respectively. The decay width difference is defined by

$$\Delta \Gamma_q \equiv \Gamma_L^q - \Gamma_H^q = 2|\Gamma_{12}^q| \cos \phi^q, \quad (20)$$

where the decay widths  $\Gamma_L$  and  $\Gamma_H$  are corresponding to the physical eigenstates  $B_L$  and  $B_H$ , respectively, and the CP phase is  $\phi^q \equiv \arg(-M_{12}^q/\Gamma_{12}^q)$ . The charge asymmetry in Eq. (3) is expressed as

$$a_{sl}^q = \frac{|\Gamma_{12}^q|}{|M_{12}^q|} \sin \phi^q = \frac{\Delta\Gamma^q}{\Delta M_q} \tan \phi^q, \quad (21)$$

of which the SM predictions are given by [4]

$$\begin{aligned} a_{sl}^d &= (-4.8_{-1.2}^{+1.0}) \times 10^{-4}, & a_{sl}^s &= (2.1 \pm 0.6) \times 10^{-5}, \\ \phi^d &= (-9.1_{-3.8}^{+2.6}) \times 10^{-2}, & \phi^s &= (4.2 \pm 1.4) \times 10^{-3}. \end{aligned} \quad (22)$$

In the SM,  $\Delta\Gamma_d/\Gamma_d$  is less than 1%, while  $\Delta\Gamma_s/\Gamma_s \sim 10\%$  is rather large. The decay matrix elements  $\Gamma_{12}^q$  is obtained from the tree level decays  $b \rightarrow c\bar{c}q$  where the dominant right-handed current contribution is suppressed by the heavy right-handed gauge boson mass  $M_{W_R}$  [20]. Therefore, we ignore the contributions of our model to  $\Gamma_{12}^q$  in this work.

We first consider the right-handed current contributions in the  $B_d^0 - \bar{B}_d^0$  system. The  $\Delta B = 2$  transition amplitudes in Eq. (18) is given by the following effective Hamiltonian in the LRM [7]:

$$H_{eff}^{B\bar{B}} = H_{eff}^{SM} + H_{eff}^{RR} + H_{eff}^{LR}, \quad (23)$$

where

$$H_{eff}^{SM} = \frac{G_F^2 M_W^2}{4\pi^2} (\lambda_t^{LL})^2 S(x_t^2) (\bar{d}_L \gamma_\mu b_L)^2, \quad (24)$$

$$\begin{aligned} H_{eff}^{LR} &= \frac{G_F^2 M_W^2}{2\pi^2} \left\{ [\lambda_c^{LR} \lambda_t^{RL} x_c x_t \zeta_g A_1(x_t^2, \zeta) + \lambda_t^{LR} \lambda_t^{RL} x_t^2 \zeta_g A_2(x_t^2, \zeta)] (\bar{d}_L b_R) (\bar{d}_R b_L) \right. \\ &\quad \left. + \lambda_t^{LL} \lambda_t^{RL} x_b \xi_g^- [x_t^3 A_3(x_t^2) (\bar{d}_L \gamma_\mu b_L) (\bar{d}_R \gamma_\mu b_R) + x_t A_4(x_t^2) (\bar{d}_L b_R) (\bar{d}_R b_L)] \right\}, \end{aligned} \quad (25)$$

and

$$\lambda_i^{AB} \equiv V_{id}^{A*} V_{ib}^B, \quad x_i \equiv \frac{m_i}{M_W} \quad (i = u, c, t), \quad \xi_g^\pm \equiv e^{\pm\alpha} \xi_g, \quad (26)$$

with

$$\begin{aligned} S(x) &= \frac{x(4 - 11x + x^2)}{4(1-x)^2} - \frac{3x^3 \ln x}{2(1-x)^3}, \\ A_1(x, \zeta) &= \frac{(4-x) \ln x}{(1-x)(1-x\zeta)} + \frac{(1-4\zeta) \ln \zeta}{(1-\zeta)(1-x\zeta)}, \\ A_2(x, \zeta) &= \frac{4-x}{(1-x)(1-x\zeta)} + \frac{(4-2x+x^2(1-3\zeta)) \ln x}{(1-x)^2(1-x\zeta)^2} + \frac{(1-4\zeta) \ln \zeta}{(1-\zeta)(1-x\zeta)^2}, \\ A_3(x) &= \frac{7-x}{4(1-x)^2} + \frac{(2+x) \ln x}{2(1-x)^3}, \\ A_4(x) &= \frac{2x}{1-x} + \frac{x(1+x) \ln x}{(1-x)^2}. \end{aligned} \quad (27)$$

Note that  $S(x)$  is the usual Inami-Lim function,  $A_1(x, \zeta)$  is obtained by taking the limit  $x_c^2 = 0$ , and  $H_{eff}^{RR}$  is suppressed because it is proportional to  $\zeta^2$ . Also in the case of  $V_I^R$ , one can see from Eq. (16) that there is no significant contribution of  $H_{eff}^{LR}$  to  $B_d^0 - \bar{B}_d^0$  mixing, so we only consider the  $V_{II}^R$  type mixing matrix for  $B_d^0 - \bar{B}_d^0$  mixing. The dispersive part of the  $B_d^0 - \bar{B}_d^0$  mixing matrix element can then be written as

$$M_{12}^d = M_{12}^{SM} + M_{12}^{LR} = M_{12}^{SM} (1 + r_{LR}^d), \quad (28)$$

where

$$r_{LR}^d \equiv \frac{M_{12}^{LR}}{M_{12}^{SM}} = \frac{\langle \bar{B}_d^0 | H_{eff}^{LR} | B_d^0 \rangle}{\langle \bar{B}_d^0 | H_{eff}^{SM} | B_d^0 \rangle}. \quad (29)$$

For specific phenomenological estimates one needs the hadronic matrix elements of the operators in Eqs. (24,25) in order to evaluate the mixing matrix element. We use the following parametrization:

$$\begin{aligned}\langle \bar{B}_d^0 | (\bar{d}_L \gamma_\mu b_L)^2 | B_d^0 \rangle &= \frac{1}{3} B_1 f_B^2 m_B, \\ \langle \bar{B}_d^0 | (\bar{d}_L \gamma_\mu b_L)(\bar{d}_R \gamma_\mu b_R) | B_d^0 \rangle &= -\frac{5}{12} B_2 f_B^2 m_B, \\ \langle \bar{B}_d^0 | (\bar{d}_L b_R)(\bar{d}_R b_L) | B_d^0 \rangle &= \frac{7}{24} B_3 f_B^2 m_B,\end{aligned}\quad (30)$$

where

$$\langle 0 | \bar{d}_\beta \gamma^\mu \gamma_5 b_\alpha | B_d^0 \rangle = -\langle \bar{B}_d^0 | \bar{d}_\beta \gamma^\mu \gamma_5 b_\alpha | 0 \rangle = -\frac{i f_B p_B^\mu}{\sqrt{2} m_B} \frac{\delta_{\alpha\beta}}{3}, \quad (31)$$

and where  $f_B$  is the  $B$  meson decay constant and  $B_i$  ( $i = 1, 2, 3$ ) are the bag factors. In the vacuum-insertion method [21],  $B_i = 1$  in the limit  $m_b \simeq m_B$ . We will use  $f_B B_i^{1/2} = (216 \pm 15)$  MeV for our numerical estimates [22]. Using the standard values of the quark masses and  $|V_{cd}^L| \approx 0.225$ , one can express  $r_{LR}^d$  in terms of the mixing angle and phases in the case of  $V_{II}^R$  in Eq. (16) as

$$\begin{aligned}r_{LR}^d &\approx 17.5 \left( \frac{1 - \zeta_g - (4.08 - 16.3\zeta_g) \ln(1/\zeta_g)}{1 - 5.58\zeta_g} \right) \zeta_g s_R^2 e^{-i(2\beta - \alpha_2 + \alpha_3)} \\ &- 756 \left( \frac{1 - 5.03\zeta_g - (0.490 - 1.96\zeta_g) \ln(1/\zeta_g)}{1 - 10.2\zeta_g + 30.1\zeta_g^2} \right) \zeta_g s_R c_R e^{-i(\beta + \alpha_3 - \alpha_4)} - 7.94 \xi_g s_R e^{-i(\beta + \alpha_3)},\end{aligned}\quad (32)$$

where the mixing phase  $\alpha$  was absorbed in  $\alpha_i$  by redefining  $\alpha_i + \alpha \rightarrow \alpha_i$ , and we used the approximation  $A_i(x, \zeta) \simeq A_i(x, \zeta_g)$  ( $i = 1, 2$ ) because  $\zeta$  dependence on  $A_i$  in Eq. (27) is rather weak for  $M_{W'} > 100$  GeV unless  $g_R/g_L$  is drastically different from unity.

On the other hand, the right-handed current contributions to  $B_s^0 - \bar{B}_s^0$  mixing is sizable only in the case of  $V_I^R$  as one can see from Eq. (16). Similarly to  $r_{LR}^d$ , we obtain  $r_{LR}^s$  in the case of  $V_I^R$  as

$$\begin{aligned}r_{LR}^s &\approx -3.47 \left( \frac{1 - \zeta_g - (4.08 - 16.3\zeta_g) \ln(1/\zeta_g)}{1 - 5.58\zeta_g} \right) \zeta_g s_R^2 e^{-i(-\alpha_2 + \alpha_3)} \\ &+ 162 \left( \frac{1 - 5.03\zeta_g - (0.490 - 1.96\zeta_g) \ln(1/\zeta_g)}{1 - 10.2\zeta_g + 30.1\zeta_g^2} \right) \zeta_g s_R c_R e^{-i(\alpha_3 - \alpha_4)} + 1.70 \xi_g s_R e^{-i\alpha_3}.\end{aligned}\quad (33)$$

The charge asymmetry  $a_{sl}^q$  in Eq. (21) can then be written in terms of  $r_{LR}^q$  in the LRM as

$$a_{LR}^q = a_{SM}^q \frac{\cos \phi_{LR}^q}{|1 + r_{LR}^q|} \left( 1 + \frac{\tan \phi_{LR}^q}{\tan \phi_{SM}^q} \right), \quad \phi_{LR}^q \equiv \arg(1 + r_{LR}^q), \quad (34)$$

where we omitted the subscript ‘ $sl$ ’ and the SM values of  $a_{sl}^q$  and  $\phi^q$  are given in Eq. (22). We use the above results for our numerical investigation of the right-handed current contributions to the like-sign dimuon charge asymmetry in semi-leptonic  $B$  decays in the next section.

#### IV. RESULTS

For our numerical analysis, we use the following  $2\sigma$  bounds obtained from the deviation of the present experimental data from the SM predictions on  $B$  meson mixing [23]:

$$0.62 < |1 + r_{LR}^d| < 1.15, \quad 0.79 < |1 + r_{LR}^s| < 1.23. \quad (35)$$

Note from Eqs. (32,33) that we have six independent new parameters ( $\zeta_g, \xi_g, \theta_R, \alpha_{2,3,4}$ ), and it is beyond the scope of this paper to perform a complete analysis by varying all six parameters. For simple illustration of the possible effect of the new interaction on  $B$  meson mixing, instead, we set  $\xi_g = \zeta_g/2$  and  $\alpha_{2,4} = 0$  because  $\xi_g$  contributions to  $B$  meson mixing is expected to be much smaller than  $\zeta_g$ 's and  $\alpha_3$  is important as the overall phase of  $r_{LR}^q$ .

In the case of  $V_I^R$ , as discussed earlier, the right-handed current contributions to  $B_s - \bar{B}_s$  mixing could be sizable while those to  $B_d - \bar{B}_d$  mixing is negligible. With the present experimental bounds of the dimuon charge asymmetry and  $B_s - \bar{B}_s$  mixing given in Eqs. (1,35), we first plot the allowed region of  $\alpha_3$  and  $\theta_R$  for  $M_{W'} = 800$  GeV at  $2\sigma$

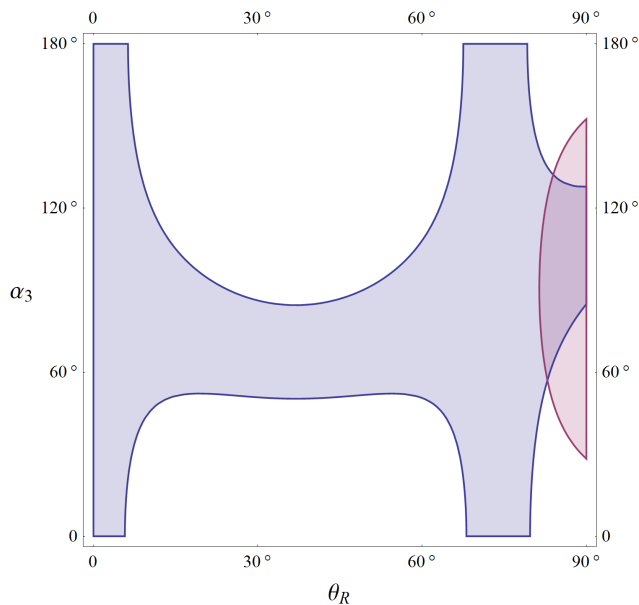


FIG. 1: Allowed regions for  $\alpha_3$  and  $\theta_R$  at  $2\sigma$  level for  $M_{W'} = 800$  GeV in the case of  $V_I^R$ . The red and blue regions are allowed by the current measurements of the like-sign dimuon charge asymmetry and  $B_s\bar{B}_s$  mixing, respectively.

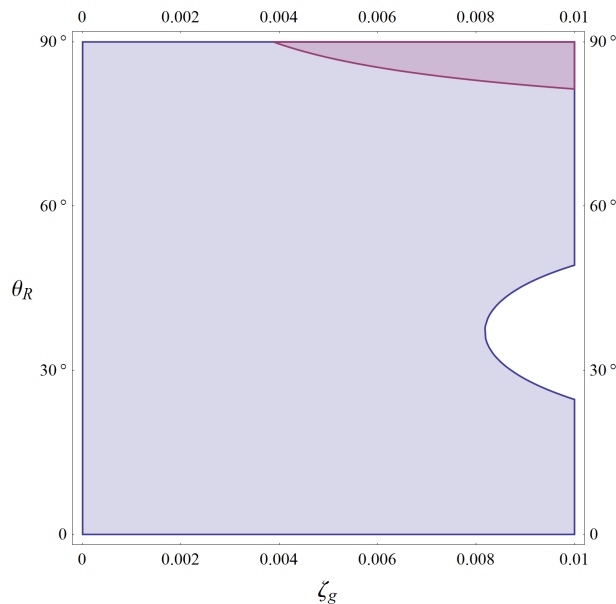


FIG. 2: Allowed regions for  $\theta_R$  and  $\zeta_g$  at  $2\sigma$  level for  $\alpha_3 = 90^\circ$  in the case of  $V_I^R$ . The red and blue regions are allowed by the current measurements of the like-sign dimuon charge asymmetry and  $B_s\bar{B}_s$  mixing, respectively.

level in Fig. 1. One can see from the overlapped allowed region in the figure that large values of  $\theta_R$  are preferred. This is the clear indication that manifest or pseudomanifest LRM is disfavored in this case. In Fig. 2, we plot the allowed region of  $\theta_R$  and  $\zeta_g$  for  $\alpha_3 = 90^\circ$  at  $2\sigma$  level. One can obtain the lower bound of  $\zeta_g \gtrsim 0.004$  from the figure which corresponds to the upper bound of  $W'$  mass  $M_{W'} \lesssim (g_R/g_L) \times 1.3$  TeV. For different values of  $\alpha_3$ , this mass bound can change, but not very much. In other words, if it happens that the mass of  $W'$  is much larger than the obtained upper bound, the right-handed contributions are not big enough to explain the present measurement of the dimuon charge asymmetry.

In the case of  $V_{II}^R$ , on the other hand, the right-handed current contributions to  $B_d - \bar{B}_d$  mixing could be sizable while those to  $B_s - \bar{B}_s$  mixing is negligible. Similarly to the  $V_I^R$  case, we plot the allowed region of  $\alpha_3$  and  $\theta_R$  for  $M_{W'} = 800$  GeV at  $2\sigma$  level in Fig. 3. The figure shows that small or large values of  $\theta_R$  are allowed unlike the  $V_I^R$

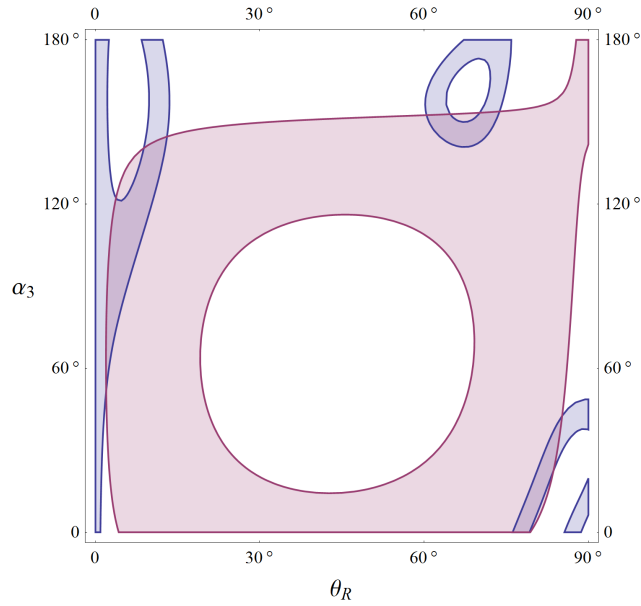


FIG. 3: Allowed regions for  $\alpha_3$  and  $\theta_R$  at  $2\sigma$  level for  $M_{W'} = 800$  GeV in the case of  $V_{II}^R$ . The red and blue regions are allowed by the current measurements of the like-sign dimuon charge asymmetry and  $B_d\bar{B}_d$  mixing, respectively.

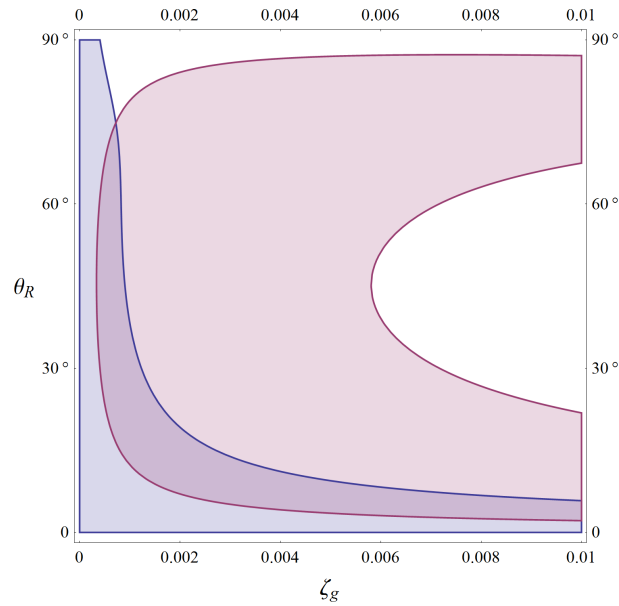


FIG. 4: Allowed regions for  $\theta_R$  and  $\zeta_g$  at  $2\sigma$  level for  $\alpha_3 = 90^\circ$  in the case of  $V_{II}^R$ . The red and blue regions are allowed by the current measurements of the like-sign dimuon charge asymmetry and  $B_d\bar{B}_d$  mixing, respectively.

case. In order for direct comparison with the  $V_L^R$  case, we plot again the allowed region of  $\theta_R$  and  $\zeta_g$  for  $\alpha_3 = 90^\circ$  at  $2\sigma$  level in Fig. 4. The figure shows that  $V_{II}^R$  scenario allows more wide range of allowed area of new parameter space and the lower bound of  $\zeta_g$  is approximately  $\zeta_g \gtrsim 0.0004$ . We obtain the corresponding upper bound of  $W'$  mass  $M_{W'} \lesssim (g_R/g_L) \times 4$  TeV. We found that this mass bound could be somewhat lower for different values of  $\alpha_3$ . It should also be noted that we have similar results for different  $\alpha_{2,4}$  in both scenarios.

## V. CONCLUDING REMARKS

In this paper, we studied the right-handed current contributions to the CP violating like-sign dimuon charge asymmetry in semi-leptonic  $B$  decays in general left-right models. Without imposing manifest or pseudomanifest left-right symmetry, we consider two types of mass mixing matrix  $V^R$  with which  $W'$  contributions are big enough to explain the current mismatch of the present measurements of the dimuon charge asymmetry and the SM prediction. We evaluated the sizes of  $W'$  contributions to  $B_d - \bar{B}_d$  and  $B_s - \bar{B}_s$  mixings which govern the dimuon charge asymmetry, and obtained the allowed regions of NP parameter spaces. With the given parameter sets, we have the following mass bounds of  $W'$ :  $M_{W'} \lesssim (g_R/g_L) \times 1.3$  TeV for Type I ( $V_I^R$ ) or  $M_{W'} \lesssim (g_R/g_L) \times 4$  TeV for Type II ( $V_{II}^R$ ), which represent the amount of NP effects enough to explain the present measurement of the dimuon charge asymmetry. If we consider the early LHC bound on  $W'$  [24], Type I model including manifest or pseudomanifest LRM is disfavored if  $g_R = g_L$ . This analysis can affect other  $B$  meson mixing related observables such as  $\sin 2\beta$  and mixing induced CP violation in  $B$  decays. A detailed discussion on such mixing induced CP asymmetries in general LRM can be found in Ref. [25], and a combined study including other decays with new experimental results will be discussed in the follow-up paper.

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