Dimuon CP Asymmetry in $B$ Decays and $Wjj$ Excess in Two Higgs Doublet Models

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Abstract

We analyze the puzzle of the dimuon CP asymmetry in $B_s$ decays in two Higgs doublet models. We show that the flavor changing neutral current (FCNC) induced by the Higgs coupling in a type III two Higgs doublet model provides a solution to the dimuon charge asymmetry puzzle by enhancing the absorptive part of the mixing amplitude $\Gamma_{12}$. We investigate different experimental constraints and show that it is possible to enhance $\Gamma_{12}$ in order to explain the dimuon asymmetry observed by D0. This enhancement requires large Higgs couplings to the first and second generations of quarks which may also explain the recent 3.2 $\sigma$ $Wjj$ excess observed by CDF.
I. INTRODUCTION

The D0 collaboration has measured the like-sign dimuon charge asymmetry in semileptonic $b$-hadron decays $A_{sl}^b$. The following result has been reported [1]:

$$A_{sl}^b = -0.00957 \pm 0.00251 \text{(stat)} \pm 0.00146 \text{(syst)}.$$  \hspace{1cm} (1)

The like-sign dimuon charge asymmetry $A_{sl}^b$ for semileptonic decay of $b$ hadrons is defined as

$$A_{sl}^b = \frac{N_{b}^{++} - N_{b}^{--}}{N_{b}^{++} + N_{b}^{--}}, \hspace{1cm} (2)$$

where $N_{b}^{++}$ and $N_{b}^{--}$ are the number of events containing two $b$-hadrons that decay semileptonically two positive or negative muons with the quark subprocesses: $b \to \mu^- \bar{\nu}X$ and $\bar{b} \to \mu^+ \nu X$.

This result indicates a 3.2 $\sigma$ deviation from the Standard Model (SM) prediction. A confirmation of this deviation would provide unambiguous evidence for new physics (NP) at low energy with a new source of CP violating phases. It is a common feature for any physics beyond the SM to possess additional sources of CP violation besides the SM phase in quark mixing matrix. These new phases can induce sizable contributions to direct and indirect CP asymmetries in $B_{d,s}$ decays and thereby resolve the apparent discrepancies between the observed results and the SM expectations.

The charge asymmetry $A_{sl}^b$ at the Tevatron can be expressed as $A_{sl}^b = (0.506 \pm 0.043)a_{sl}^d + (0.494 \pm 0.043)a_{sl}^s$ [1], where the charge asymmetry $a_{sl}^q (q = d, s)$ for “wrong-charge” semileptonic $B_q^0$-meson decay induced by the oscillation is defined by

$$a_{sl}^q = \frac{\Gamma(B_q^0 \to \mu^+ \nu X) - \Gamma(B_q^0 \to \mu^- \bar{\nu}X)}{\Gamma(B_q^0 \to \mu^+ \nu X) + \Gamma(B_q^0 \to \mu^- \bar{\nu}X)}. \hspace{1cm} (3)$$

This asymmetry can be written as [2]

$$a_{sl}^q = \frac{|\Gamma_{12}^q|}{|M_{12}^q|} \sin \phi_q = \frac{\Delta \Gamma_q}{\Delta M_q} \tan \phi_q, \hspace{1cm} (4)$$

where the mass and width differences between $B_q$ mass eigenstates are given by

$$\Delta M_{B_q} = M_{B_{u}} - M_{B_{d}} = 2|\Gamma_{12}^q|, \hspace{1cm} (5)$$

$$\Delta \Gamma_q = 2|\Gamma_{12}^q| \cos \phi_q, \hspace{1cm} (6)$$

and $M_{12}^q$ and $\Gamma_{12}^q$ are the dispersive and the absorptive parts of the mixing amplitudes, respectively. The CP violating phase $\phi_q$ is defined by $\phi_q = \text{arg} (-M_{12}^q/\Gamma_{12}^q)$. In the SM,
the arguments of \( M_{12}^q \) and \( \Gamma_{12}^q \) are aligned at the leading order due to the unitarity of the quark mixing matrix, and thus, the phase \( \phi_q \) is small irrespective of the individual phases, \( \beta_d \equiv \arg(-V_{cb}^*V_{cd}/V_{tb}^*V_{td}) \) and \( \beta_s \equiv \arg(-V_{tb}^*V_{ts}/V_{cb}^*V_{cs}) \). The ratio of \( \Gamma_{12}^q/M_{12}^q \) is roughly proportional to \( m_b^2/m_t^2 \) with a small phase, and as a result, the SM prediction of the charge asymmetry is small: \( a_{sl}^d(\text{SM}) \simeq -5 \times 10^{-4} \), \( a_{sl}^s(\text{SM}) \simeq 2 \times 10^{-5} \) and \( A_{sl}^b(\text{SM}) \simeq -0.00023 \), which is in clear contradiction with the D0 result in eq.(1).

In this respect, it is clear that a large new CP violation in \( B_s \) mixing is required to enhance \( a_{sl}^s \) if \( B_d \) mixing does not contain a new physics contribution. The phase \( \phi_s \) can be large in general in many new physics models because the phase alignment between \( M_{12}^s \) and \( \Gamma_{12}^s \) can be broken if a new particle, such as a supersymmetric particle, propagates in the loop diagram, which contributes to the mixing amplitude [3]. However, the absorptive part \( \Gamma_{12}^s \) is not necessarily modified by the propagation of new particles due to the on-shell condition of the intermediate states. The magnitude of \( M_{12}^s \) is determined by the mass difference \( \Delta M_{B_s} \), and thus eq.(4) tells us that \( a_{sl}^s \) has a maximal value if \( \Gamma_{12}^s \) is dominated by the SM tree level contribution [4]. One can easily show that if one simply extrapolates eq.(4), then one needs a new phase \( \phi_s \) with \( \sin \phi_s \gtrsim 1.6 \) (which is outside of the domain of \( \sin \phi_s \)) in order to account for the experimental result of \( A_{sl}^b \) [5] by the D0 collaboration. Also with a large \( \sin \phi_s \simeq 1 \), the decay width difference \( \Delta \Gamma_s = 2|\Gamma_{12}^s|\cos \phi_s \) is suppressed and becomes inconsistent with the experimental constraints. Therefore, one concludes that the sustainability of the D0 results of the like-sign dimuon charge asymmetry in semileptonic \( b \)-hadrons decay would be a clear hint of possible NP that modify the absorptive part of the mixing amplitude \( \Gamma_{12}^s \).

In this paper we show that the flavor changing neutral current (FCNC) induced by the Higgs coupling in two Higgs doublet models (THDM) can provide a solution to the dimuon charge asymmetry puzzle by enhancing \( \Gamma_{12}^s \). The THDM is classified by the selection of the Higgs couplings to fermions. We will consider a general type of coupling (so called type III THDM) to obtain an appropriate FCNC source to modify \( \Gamma_{12}^s \). We enumerate the experimental constraints and investigate if there is room to enhance \( \Gamma_{12}^s \) to achieve the D0 result of the dimuon asymmetry. We will obtain the operators generated by charged Higgs exchange which can enhance \( \Gamma_{12} \). To accomplish this, suitably large Higgs couplings to the first and second quark generations needed. Such couplings can, in addition, lead to an excess of Higgs decays into dijets. Indeed, the CDF collaboration has recently reported a 3.2\( \sigma \) excess in the 120-160 GeV range in the invariant mass distribution of the dijets in
association with a $W$ boson. The excess may be explained by the Higgs decays to dijets via the new Higgs couplings.

The presence of two Higgs doublets is required in supersymmetric models and in models with the left-right gauge symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. We briefly explore if the left-right symmetric model can explain the large dimuon asymmetry consistent with other experimental constraints.

II. DIMUON ASYMMETRY IN TWO HIGGS DOUBLET MODELS

The two Higgs Doublet Model is an extension of the SM that naturally introduces a new source of CP violation and FCNC. In this class of models, the most general renormalizable and gauge invariant Yukawa interactions are given by

$$- \mathcal{L}_Y = Y^u_{ij} q_i u_j H_u + Y^\nu_{ij} q_i u_j \tilde{H}_d + Y^d_{ij} q_i d_j H_d + Y^e_{ij} \ell_i e_j \tilde{H}_u,$$  

(7)

where the Higgs fields $H_{u,d}$ have hypercharges $Y = \pm 1/2$, and $\tilde{H}_{u,d}$ are defined as $\tilde{H}_{u,d} = -i \tau_2 H_{u,d}$. The Yukawa couplings $Y_f$ and $Y'_f$ ($f = u, d, e$) are $3 \times 3$ matrices with a generic flavor structure. In this case, the fermion masses are given as

$$M_u = Y^u_{ij} v_u + Y^{\nu}_{ij} v_d,$$  

$$M_d = Y^d_{ij} v_d + Y^{\nu}_{ij} v_u,$$  

$$M_e = Y^e_{ij} v_d + Y^{\nu}_{ij} v_u.$$  

(8)

(9)

(10)

In the basis where the mass matrix is diagonal, FCNC interactions through the neutral Higgs ($H$ and $A$) exchange are generated from the mismatch between the diagonalization of the mass matrices and the Yukawa interactions. For instance, couplings like $s_L b_R c d_d^0 H_d^0$ and $b_L s_R c d_d^0$ are induced by $(V_L^d V_d^d)^{ij} q_i d_j c d_d^0 H_d = (M^\text{diag}_d / v_d - V_L^d Y_d^d V_d^d \tan \beta)^{ij} q_i d_j c d_d^0 H_d$, where $\tan \beta = v_u / v_d$, and $V_L^d, V_R^d$ are the diagonalizing matrices of $M_d$. Hereafter, we will work in the basis where $M_d$ is diagonal, and we omit $V_L^d, V_R^d$ in the expressions. Then, by definition, $Y^d_{ij} = -Y^d_{ij} \tan \beta$ ($i \neq j$).

These couplings contribute to $B_s - \bar{B}_s$ mixing and modify the amplitude $M_{12}^s$ as follows:

$$M_{12}^s = (M_{12}^s)^{\text{SM}} + (M_{12}^s)^{2\text{HD}},$$  

(11)

where $(M_{12}^s)^{2\text{HD}}$ is given by

$$(M_{12}^s)^{2\text{HD}} \simeq \frac{Y^{23}_d Y^{32}_d}{m_H^2} (B_s^0 | \overline{b_L} s_R (\overline{b_R} s_L) | \overline{B}_s^0).$$  

(12)
We note that even in the minimal supersymmetric standard model (MSSM) this type of modification of $M^s_{12}$ via the neutral Higgs exchange can be obtained for large $\tan \beta$ through the finite correction of the Yukawa couplings due to soft SUSY breaking terms. In this case, if $\Gamma^s_{12} = (\Gamma^s_{12})_{\text{SM}}$, the charge asymmetry $a^s_{12}$ is given by

$$a^s_{12} = \frac{|(\Gamma^s_{12})_{\text{SM}}|}{(M^s_{12})_{\text{SM}}} \sin(\phi^s_{\text{SM}} + 2\theta_s),$$

where $r_s = |1 + M^2_{12}/M^S_{12}|$ and $2\theta_s = \arg(1 + M^2_{12}/M^S_{12})$. Thus, $2\theta_s$ can be large if $|Y^d_{23}Y^d_{32}|$ and $\arg(Y^d_{23}Y^d_{32})$ are large, satisfying the experimental constraint: $r_s \sim 1$. As mentioned above, a large value of $2\theta_s$ is necessary but not sufficient to account for the dimuon CP asymmetry $A^s_{12}$. A significant enhancement of $\Gamma^s_{12}$ is preferable.

In the 2HDM several $\Delta b = 1$ effective operators which may modify $\Gamma^s_{12}$, can be generated. For example, with non-vanishing $Y^u_{23} = (\bar{b}_{L,C,R})(C_{R,S,L})$ is generated through the charged Higgs exchange. Let us enumerate the possible $\Delta b = 1$ effective operators generated via the Higgs exchange, which are suitable for modifying $\Gamma^s_{12}$.

1. $$(\bar{b}_{L,C,R})(u^c_{R,S,L}) = -\frac{1}{2}(\bar{b}_{L,C,R})(u^c_{R,S,L})$$
2. $$(\bar{b}_{R,C,R})(u^c_{R,S,L}) = -\frac{1}{2}(\bar{b}_{R,C,R})(u^c_{R,S,L})$$
3. $$(\bar{b}_{R,C,R})(u^c_{R,S,L})$$
4. $$(\bar{b}_{R,C,R})(\tau_{L,R})$$

The operator $(\bar{b}_{L,C,R})(u^c_{R,S,L})$ is generated through the charged Higgs exchange with the coefficient:

$$\sum_{q,q'=u,c,t} \frac{V_{qb}V_{q's}}{m^2_{H^+}} \sin^2 \beta \left[ \frac{M_u \cos \beta}{v} - Y'_u \right]_{q'j} \left[ \frac{M_u \cos \beta}{v} - Y'_u \right]^{*}. \quad (14)$$

As we will see below, this is the preferred operator for modifying $\Gamma^s_{12}$.

The operator $(\bar{b}_{R,C,R})(u^c_{R,S,L})$ is generated through the charged Higgs exchange with the coefficient:

$$\sum_{q,q'=d,s,b} \frac{V_{iq}V_{q'j}}{m^2_{H^+}} \cos^2 \beta \left[ \frac{M_d \sin \beta}{v} - Y'_d \right]_{q3} \left[ \frac{M_d \sin \beta}{v} - Y'_d \right]^{*}. \quad (15)$$

This operator can modify $\Gamma^s_{12}$ through interference with $W$ boson exchange. The effect can be large if $Y^d_{23}$ and $(Y_d \sin \beta - Y'_d \cos \beta)_{22}$ are sizable. However, one needs fine-tuning to obtain a large contribution to $\Gamma^s_{12}$ since the strange mass is $m_s = (Y_d \cos \beta + Y'_d \sin \beta)_{22}v$. 

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Such fine tuning may also generate the operator $(\bar{b}_R s_L)(s_L s_R)$, which affects $B_d \rightarrow \phi K$ by an amount which causes disagreement with the experimental data.

The operators $(\bar{b}_R s_L)(u_R'u_L')$ and $(s_R b_L)(u_R'u_L')$ can be generated through the neutral Higgs exchange. We note that both operators are needed to modify $\Gamma_{12}^s$ since $(u_R'u_L')$ cannot be self-conjugate.

Finally, the operators $(\bar{b}_R s_L)(\tau_L \tau_R)$ and $(s_R b_L)(\tau_L \tau_R)$ can be generated from $Y_{23}^d$, $Y_{32}^d$ via the neutral Higgs exchange. It is remarkable that when $Y_{e'}_{22}^u$ generates the muon mass, $\text{Br}(B_s \rightarrow \tau\tau)$ is not constrained by experiment and $\Gamma_{12}$ can be modified. Note that in MSSM (or THDM with $Y_{e'} \rightarrow 0$), the following relation is obtained

$$\frac{\text{Br}(B_s \rightarrow \tau\tau)}{\text{Br}(B_s \rightarrow \mu\mu)} \simeq \left(\frac{m_\tau}{m_\mu}\right)^2. \quad (16)$$

Thus, due to the experimental bound on $\text{Br}(B_s \rightarrow \mu\mu)$, the quantity $\text{Br}(B_s \rightarrow \tau\tau)$ cannot be large enough to modify $\Gamma_{12}^s$.

In the case of neutral Higgs exchanges, $\Gamma_{12}$ is modified at one-loop level while $\Delta b = 2$ operator to modify $M_{12}^s$ is also generated through the neutral Higgs exchanges at tree level. Therefore, to obtain a sizable $\Gamma_{12}^s$ contribution, a large $M_{12}^s$ is also generated, which is unwanted.

Next we consider the following experimental constraints:

1. $B \rightarrow M_1 M_2$ and $B \rightarrow \ell\ell$ decays.

2. $b \rightarrow s\gamma$.

3. lifetime ratio $\tau_{B_s}/\tau_{B_d}$.

The constraints from two body decays into mesons and leptons are studied in [10], and the allowed operators for $\Delta b = 1$ are

$$\bar{s}b\bar{c}c, \quad \bar{s}b\tau\tau, \quad \bar{d}b\bar{c}u. \quad (17)$$

In the MSSM, it is difficult to enhance the $\Delta b = 1$ operators selectively because the interactions are related by known (or constrained) coupling constants.

The non-holomorphic Yukawa coupling $Y_u'$ can also contribute to the charged Higgs interaction and lead to an important effect on $b \rightarrow s\gamma$. In a non-SUSY type II 2HDM, the charged Higgs has to be heavier than $\sim 350$ GeV to agree with the experimental measurement of $b \rightarrow s\gamma$. However, in our case, this constraint can be relaxed by a small $Y_u'$ contribution. Note that $Y_{e'}_{22}^d$ ($Y_{d'}_{32}^d$) is constrained by $b_R \rightarrow s_L\gamma$ ($b_L \rightarrow s_R\gamma$) contribution.
The $b \to s\gamma$ process also constrains the large log contribution due to the renormalization group evolution below the $W$ boson mass [11]. The $b \to s\gamma$ constraint disfavors the neutral Higgs contributions. This is due to the fact that the operators $(\bar{s}_L b_R)(\bar{f}_L f_R)$ and $(\bar{s}_R b_L)(\bar{f}_R f_L)$ ($f = c, b, \tau$) give leading order large log corrections ($\delta C_{7,8} \sim m_f/m_b \ln M^2_W/m^2_b$) to the $b \to s\gamma$ operators, and hence they are dangerous operators. We note that the charged Higgs contribution can generate only next to leading order corrections.

Finally, the lifetime ratio $\tau_{B_s}/\tau_{B_d}$ provides a stringent constraint on the operator contributing to $\Gamma_{12}^s$. The modification of $\Gamma_{12}^s$ can induce a change ($\sim O(10)\%$) in the lifetime of $B_s$, which can be consistent with the large hadronic uncertainty. However, this uncertainty is cancelled in the lifetime ratio $\tau_{B_s}/\tau_{B_d}$. The current world average of the experimental result is [12]

$$\tau_{B_s}/\tau_{B_d} = 0.99 \pm 0.03, \quad (18)$$

Therefore, if $\Gamma_{12}^s$ is modified, the lifetime of $B_d$ should also modified, which provides a strong constraint. Actually, for $B_d$ system, it is known that the bound for $\text{Br}(B_d \to \tau\tau) < 4.1 \times 10^{-3}$. Consequently, the modification of $\Gamma_{12}^s$ via $B_s \to \tau\tau$ has a deficit due to this lifetime ratio constraint. The only allowed operator to modify the $B_d$ lifetime seems to be $\bar{d}bcu$, which can modify $\Gamma_{11}^d$. Fortunately, the operator $\bar{d}bcu$ alone modifies neither $M_{12}^d$ nor $\Gamma_{12}^d$ since $\bar{c}u$ is not self-conjugate.

The above constraints imply that the only possibility left is the charged Higgs exchange operator $(\bar{q}_L \gamma_{\mu} b_L) (u_R \gamma^{\mu} u_R) \ (q = d, s)$. The coefficients of these operators are obtained, as in eq. (14), and they are proportional to $X_{ib}X^*_{ij}$, where $X_{ij} = [(M^\text{diag}_u \cos \beta/v - Y_u')^TV]_{ij}$. In order to generate the operators $\bar{s}b\bar{c}c$ and $\bar{d}bcu$, we need $X_{cs}, X_{cb}$ and $X_{ud}$. The condition $X_{cs} \sim X_{ud}$ can make these effective operators comparable to keep the lifetime ratio $\tau_{B_s}/\tau_{B_d}$.

The condition can be satisfied when $Y_{u11}' \sim Y_{u22}'$. However, if $Y_{u11}'$ is sizable, one needs fine tuning to obtain the proper up quark mass. This can be relaxed if $\tan \beta$ is large.

The $\Gamma_{12}^s$ contribution is estimated as

$$\frac{\Gamma_{12}^s_{\text{THDM}}}{\Gamma_{12}^{s\text{SM}}} \sim \left( \frac{X_{cs}^* X_{cb}}{g^2 V_{ts}^* V_{tb}} \right)^2 \frac{M^4_W}{M^4_H} \gamma_{cc}, \quad (19)$$

where $\gamma_{cc} = \sqrt{1 - 4m^2_c/m^2_b(1 - 2/3(m^2_c/m^2_b))}$. The contribution to the dispersive part of the mixing can be written as

$$\frac{M^2_{12}^\text{THDM}}{M^2_{12}^{\text{SM}}} \sim \frac{X_{is}^* X_{ib} X_{js}^* X_{jb} M^2_W}{(g^2 V_{ts}^* V_{tb})^2} \frac{f(m^2_l/M^2_H, m^2_b/M^2_H, M^2_W/M^2_H)}{S(m^2_l/M^2_W)}, \quad (20)$$

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where $i, j = c, t$ and $f(x, y, z)$ and $S(x)$ are the Inami-Lim functions. If we choose $Y_u^{ij33}$ appropriately, the charged Higgs contribution to $M_{12}^s$ can vanish, keeping the contribution to $\Gamma_{12}^s$. (In fact, if $Y_u^{ij} \propto \delta_{ij}$ and $m_i \ll M_H$, a GIM-like mechanism works for the dispersive part of the meson mixing amplitudes). This is because the top-loop can contribute to $M_{12}^s$, but not to $\Gamma_{12}^s$. Therefore, one can avoid an excessive contribution to $M_{12}^s$ and still modify $\Gamma_{12}^s$ appropriately in this scenario.

In Fig. 1, we show $\Gamma_{12}^{\text{THDM}}/\Gamma_{12}^{\text{SM}}$ as a function of $y_u'$, where we assume $Y_u^{ij} = y_u' \delta_{ij}$ for simplicity. We choose the charged Higgs mass to be 160 GeV.

In order to modify $\Gamma_{12}^s$, we need $Y_u'$ whose magnitude depends on the charged Higgs mass $m_H$. As previously mentioned, we need to fine-tune in order to obtain the proper up quark mass in order to satisfy the lifetime ratio. The fine-tuning can be relaxed when the charged Higgs is light. However, the constraints from $B^- \rightarrow \tau \bar{\nu}$ and $B \rightarrow D\tau \bar{\nu}$ should be taken into account [13]. In a general 2HDM, the non-holomorphic coupling $\ell \tau^c H_u^*$ can contribute to the $\tau$ mass, and thus the $\ell \tau^c H_d$ coupling (which is important to $B \rightarrow \tau \bar{\nu}$) may have freedom to relax the $B \rightarrow \tau \bar{\nu}$ constraint.

The sizable Higgs coupling $Y_u^{111}$ can provide an interesting hadron collider signal. The coupling can cause a resonant production of the charged and neutral Higgs bosons (which mainly contain $H_d$). The Higgs bosons decay into two jets via the $Y_u^{n11}$ and $Y_u^{n22}$ couplings, which are needed to enlarge $\Gamma_{12}^s$. Thus, dijet excesses can be observed around the masses of the Higgs bosons. If the Higgs boson is light ($\lesssim 200$ GeV), the dijet events are buried under the QCD background. However, even if the Higgs boson is light, there is a chance to observe
dijet events produced by Higgs decays associated with $W/Z$ gauge bosons. Recently, the CDF collaboration has reported an excess of dijets associated with $W$ boson (which decay into $\ell \bar{\nu}$) [6]. The excess in the dijet mass distribution is in the 120-160 GeV range, and it can be explained if the Higgs boson mass is about 150-160 GeV (the peak of the dijet mass distribution shifts to lower mass due to cuts). We note that the bottom quark mass should be generated by $q_3 b^c H_u^*$ because $b$-quark excess is not observed.

The charged and neutral Higgs bosons associated with the $W$ boson are produced by the $t$-channel exchange of the left-handed quarks through the $Y_u'$ coupling. The resonant (or off-shell) production of the charged/neutral Higgs boson through $Y_u'$ can also contribute to the neutral/charged Higgs boson associated with the $W$ boson. Similar processes via a
Higgs doublet are also analyzed in [14, 15]. The CDF estimated production rate of $W jj$ excess is about 4 pb. We use MadGraph/MadEvent (version 5) to estimate the $WH$ cross-sections in this model. In Fig.2 we plot $\sigma(p\bar{p} \rightarrow WH)$ as a function of $y_u'$ for $m_{H^+} = 160$ GeV (neutral Higgs masses are also 160 GeV). In the figure we use a $K$ factor $\sim 1.35$. The $WH$ production cross-section is dominated by charged Higgs associated production (roughly 70%). The charged Higgs boson then decays 100% into jets. We find that the excess can be explained for Higgs couplings $Y_u'^{11} \sim Y_u'^{22} \sim 0.5$. We also have $ZH$ (includes both charged and neutral Higgs) productions in this model, but the cross-section is at least a factor 3 smaller compared to the $WH$ production. The $ZH$ production cross-section goes further down if the neutral Higgs masses are heavier than the charged Higgs masses.

In Fig.3 we plot $\Gamma_{12}^{THDM}/\Gamma_{12}^{SM}$ as a function of $y_u'^{23}$, where we choose $Y_u'^{11} = Y_u'^{22} = 0.5$ and $m_{H^+} = 160$ GeV. Since in this case we need $Y_u'^{23}$ to be $\sim 0.1$ to produce the desired $\Gamma_{12}^s$ modification. Due to this small $Y_u'^{23}$, we have a very small amount of single $b$ quark present in the $W jj$ signal. As mentioned before, $Y_u'^{23}$ is constrained by $b_R \rightarrow s_L \gamma$ operator. In order to turn on the $Y_u'^{23}$ term, one should make $Y_d'^{33}$ small and the bottom quark mass should come from the non-holomorphic term. The absorptive part of the mixing is insensitive to $Y_u'^{33}$, and thus $Y_u'^{33}$ can be used to adjust the dispersive part of the mixings.

III. DIMUON ASYMMETRY IN LEFT-RIGHT SYMMETRIC GAUGE THEORY

In SUSY models, the non-holomorphic terms ($Y_u', Y_d'$ couplings) can arise only from the finite corrections. A sizable contribution to $M_{12}^s$ is easily obtained, but in order to obtain sizable effects for $\Gamma_{12}^s$ consistence with other experimental constraints such as $b \rightarrow s \gamma$, $O(100)$ TeV scale physics should be considered.

In the left-right symmetric gauge theory (LR model), the gauge symmetry is extended to $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. In this model, a number of Higgs fields and couplings are usually needed in order to obtain the required quark masses and CKM mixings. The Higgs couplings can generate FCNC, which can affect the dimuon asymmetry, as described in the previous section. In addition, the $SU(2)_R$ gauge bosons can also contribute to $\Gamma_{12}^s$ and $M_{12}^s$. Let us estimate the contributions of these gauge bosons to the mixing amplitudes.

The RRRR operator $(\bar{q}_R \gamma_\mu b_R)(\bar{q}_R \gamma_\mu b_R)$ can contribute to $\Gamma_{12}^q (q = d, s)$:

$$\frac{\Gamma_{12}^{qLR}}{\Gamma_{12}^{SM}} \sim \left( \frac{V_{tq}^{R*} V_{tb}^R}{V_{tq}^{L*} V_{tb}^L} \right)^2 \left( \frac{g_{LR}}{g_L} \right)^4 \left( \frac{M_{W_L}}{M_{W_R}} \right)^4 (1 + O(m_c^2/m_b^2)), \quad (21)$$
where $V^L$ and $V^R$ are the left- and right-handed quark mixing matrices, $g_{L,R}$ are the gauge coupling constants, and $M_{W_{L,R}}$ are the masses of the $W_{L,R}$ gauge bosons (in general, these bosons can mix). We have used here the unitarity of the quark mixing matrices: $V_{uq}V_{ub} + V_{cq}V_{cb} + V_{tq}V_{tb} = 0$. Note that the $M_{12}$ contribution from the RRRR operator is tiny because of the unitarity of $V_R$ and $m_{u,c,t}/M_{W_R} \ll 1$.

The LLRR operator $(\bar{q}_L \gamma_\mu b_L)(\bar{q}_R \gamma_\mu b_R)$ can contribute to $M_{12}^q$:

$$
\frac{M_{12}^{q_{LR}}}{M_{12}^{q_{SM}}} \sim \left( \frac{V_{tq}^R V_{tb}^R}{V_{tq}^L V_{tb}^L} \right) \left( \frac{g_R}{g_L} \right)^2 \left( \frac{M_{W_L}}{M_{W_R}} \right)^2 \frac{2A_2(x_t^2, M_{W_L}^2/M_{W_R}^2)}{F(x_t^2)},
$$

(22)

where the loop functions $A_2$ and $F$ can be found in [17]. Note that the contribution to $\Gamma_{12}^q$ via the LLRR operator can be negligible because charm or up quark masses are inserted (or $W_L$-$W_R$ mixing is inserted twice).

As a result, we obtain

$$
\frac{M_{12}^{q_{LR}}}{M_{12}^{q_{SM}}} \sim \sqrt{\frac{\Gamma_{12}^{q_{LR}}}{\Gamma_{12}^{q_{SM}}}} \frac{2A_2(x_t^2, M_{W_L}^2/M_{W_R}^2)}{F(x_t^2)},
$$

(23)

The correction to $\Gamma_{12}^q$ should be less than about 30% when $|M_{12}^{q_{LR}}/M_{12}^{q_{SM}}| \lesssim 2$ and $M_{W_L}^2/M_{W_R}^2 \lesssim 10^{-3}$. Therefore, the contribution via the right-handed gauge boson is not a better choice to modify $\Gamma_{12}^q$ since modification to $M_{12}^q$ will be too large irrespective of the choice of the right-handed quark mixing matrix.

The merit of the right-handed gauge boson is to modify $\Gamma_{11}^d$ to adjust the lifetime ratio for $B_d$ and $B_s$. If the right-handed quark mixing matrix is

$$
V^R = \begin{pmatrix}
1 & 0 & 0 \\
0 & 0 & 1 \\
0 & 1 & 0
\end{pmatrix},
$$

(24)

the operator $(\bar{d}_R \gamma_\mu b_R)(\bar{c}_R \gamma_\mu u_R)$ is generated by $W_R$ exchange, and the lifetime of $B_d$ can be tuned depending on $g_{R}^2/M_{W_R}^2$. In the previous section, we found that the element $Y^\prime_{u11}$ can modify $\Gamma_{11}^d$ but it requires fine-tuning to obtain the proper up quark mass. The contribution from the right-handed gauge boson can relax the fine-tuning. We must adjust the contribution to make the lifetime ratio lie within the current experimental uncertainty in eq.(18). If we choose the elements (e.g. $V^R_{us}$) to be exactly zero, any unwanted contribution to meson mixings from the $W_R$ gauge boson can be avoided.
IV. CONCLUSION

In this paper we have investigated the dimuon CP asymmetry in $B_s$ decays in two Higgs doublet models. We find that the flavor changing neutral current (FCNC) induced by the Higgs couplings in type III two Higgs doublet model can enhance the decay width $\Gamma_{s12}$ and resolve the dimuon charge asymmetry puzzle, consistent with all the experimental constraints. The enhancement of $\Gamma_{s12}$ requires large Higgs couplings to the first and second generations of quarks, which may help explain the recent 3.2σ $Wjj$ excess observed at CDF.

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