PHYSICAL REVIEW D **90,** 055003 (2014)

Distinguishing the right-handed up/charm quarks from the top quark via discrete symmetries in standard model extensions

Chao-Shang Huang,[1] Tianjun Li,[1,2] Xiao-Chuan Wang,[1] and Xiao-Hong Wu[3,1]

[1]*State Key Laboratory of Theoretical Physics, Institute of Theoretical Physics, Chinese Academy of Sciences, Beijing 100190, People's Republic of China*

[2]*School of Physical Electronics, University of Electronic Science and Technology of China, Chengdu 610054, People's Republic of China*

[3]*Institute of Modern Physics, East China University of Science and Technology, 130 Meilong Road, Shanghai 200237, People's Republic of China*

(Received 18 November 2013; revised manuscript received 22 July 2014; published 4 September 2014)

We propose a class of the two Higgs doublet standard models (SMs) with a SM singlet and a class of supersymmetric SMs with two pairs of Higgs doublets, where the right-handed up/charm quarks and the right-handed top quark have different quantum numbers under extra discrete symmetries. Thus, the right-handed up and charm quarks couple to one Higgs doublet field, while the right-handed top quark couples to another Higgs doublet. The quark Cabibbo-Kobayashi-Maskawa mixings can be generated from the down-type quark sector. As one of the phenomenological consequences of our models, we explore whether one can accommodate the observed direct CP asymmetry difference in singly Cabibbo-suppressed D decays. We how that it is possible to explain the measured values of CP violation under relevant experimental constraints.

DOI: 10.1103/PhysRevD.90.055003 PACS numbers: 11.25.Mj, 12.10.-g, 12.10.Kt

I. INTRODUCTION

Experimental data from the ATLAS [1,2], CMS [3,4], D0, and CDF [5] collaborations have confirmed the existence of the standard model (SM) Higgs boson. However, the quark Cabibbo-Kobayashi-Maskawa (CKM) mixing phase is not enough to explain the baryon asymmetry in the Universe and gives the contributions to electric dipole moments of electrons and neutrons much smaller than the experimental limits. Therefore, one needs new sources of CP violation, which has been one of the main motivations to search for new theoretical models beyond the SM for a long time.

The minimal extension of the SM is to enlarge the Higgs sector [6]. It has been shown that the two-Higgs-doublet models (2HDMs) naturally accommodate the electroweak precision tests, giving rise at the same time to many interesting phenomenological effects [7]. For a recent review on two-Higgs-doublet SMs, please see [8]. The generic scalar spectrum of the two-Higgs-doublet models consists of three neutral Higgs bosons and one charged Higgs boson pair. The direct searches for additional scalar particles at the LHC or indirect searches via precision flavor experiments will therefore continue being an important task in the upcoming years.

In this paper, we will propose a class of the two-Higgs-doublet SMs with a SM singlet and a class of the supersymmetric SMs with two pairs of Higgs doublets, where the right-handed up/charm quarks and the right-handed top quark have different quantum numbers under extra discrete symmetries. Therefore, the right-handed up and charm quarks couple to one Higgs doublet field, while the right-handed top quark couples to another Higgs doublet due to additional discrete symmetries. All the down-type quarks couple to the same Higgs doublet, and all the charged leptons couple to the same Higgs doublet. Also, the quark CKM mixings can be generated from the down-type quark sector. In particular, the first two-generation up-type quarks can have relatively large Yukawa couplings. As one of the phenomenological consequences of our models we explore if one can accommodate the experimental measurement of direct CP asymmetry difference in singly Cabibbo-suppressed D decays.

The CP asymmetry difference in $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$ decays has been measured by the LHCb Collaboration [9]. Combined with the results from the CDF [10], Belle [11], and previous *BaBar* [12] collaborations, the Heavy Flavor Averaging Group yields a world average of the difference of direct CP asymmetry in $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$ decays, $\Delta A_{CP} = (-0.656 \pm 0.154)\%$ in March 2012 [13]. However, the above results have not been confirmed by the latest experimental measurements. The updated LHCb result with pion-tagged analysis gives $\Delta A_{CP} = (-0.34 \pm 0.15 \pm 0.10)\%$ [14]. For the muon tagging, the measurements from LHCb using 1.0 fb^{-1} data at 7 TeV have $\Delta A_{CP} = (0.4 \pm 0.3 \pm 0.14)\%$ [15] and $\Delta A_{CP} = (+0.14 \pm 0.16 \pm 0.08)\%$ [16] with the latest 3 fb^{-1} data, which have an opposite sign compared to the pion-tagged results. In combination, the current world-averaged direct charm meson CP violation is $\Delta A_{CP} = (0.253 \pm 0.104)\%$ from the Heavy Flavor Averaging Group [13].

The CP asymmetry in charm meson decays has inspired a lot of theoretical discussions. The SM contributions to the

 © 2014 American Physical Society

direct CP asymmetry are discussed in Refs. [17–19]. Li et al. [18] showed that $\Delta A_{CP} = A_{CP}(K^+K^-) - A_CP(\pi^+\pi^-) = -1.00 \times 10^{-3}$, which is lower than the LHCb and CDF data. Based on the topological diagram approach for tree-level amplitudes and QCD factorization for a crude estimation of perturbative penguin amplitudes, Cheng and Chiang [19] showed that the CP asymmetry difference ΔA_{CP} is of order $-(0.14–0.15)\%$. Even with the maximal magnitude of QCD–penguin exchange amplitude $|PE| \sim T$ (T is the tree-level amplitude) and a maximal strong phase relative to T, one can only get $\Delta A_{CP} = -0.25\%$ which is still lower than the current world average. The $SU(3)$ effects have also been studied [20–24]. For recent discussions on the subjects, please see Ref. [25]. While the experiment is still not conclusive, there are some attempts to estimate the effects from new physics models, e.g., fourth generation [26], left-right model [27], diquark [28], supersymmetry [29,30], Randall-Sundrum model [31], compositeness [32,33], minimal flavor violation [34], other new physics models [35], and a χ^2 analysis of different measurements in the charm system [36].

We calculate the direct CP asymmetry difference in charm meson decays with experimental constraints satisfied in our models in the paper. The new feature of our work is that we consider the contributions from Higgs penguin induced operators and the mixing effect of Higgs penguin induced operator O_{13} into chromomagnetic operator O_{8g} at charm mass m_c scale. We find that it is possible to explain the measured values of CP violation under relevant experimental constraints.

This paper is organized as follows. We present a class of two-Higgs-doublet SMs and a class of the supersymmetric SMs in Secs. II and III. The effective Lagrangian of $c \to u$ transition, relevant Wilson coefficients, direct CP asymmetry in charm meson decays, and $\Delta c = 2$ and $\Delta c = 1$ constraints are given in Sec. IV. We conclude in Sec. V.

II. NONSUPERSYMMETRIC SMS

We consider the two-Higgs-doublet standard models [6]. First, let us explain the convention. We denote the left-handed quark doublets, the right-handed up-type quarks, the right-handed down-type quarks, the left-handed lepton doublets, and the right-handed leptons as q_i, u_i, d_i, l_i, and e_i, respectively, where $i = 1, 2, 3$. In addition, we introduce two pairs of the Higgs doublets as ϕ_1 and ϕ_2, and a SM singlet Higgs field S. Following the common convention, we assume that the $U(1)_Y$ charges for both ϕ_1 and ϕ_2 are $+1$.

Without loss of generality, we assume that ϕ_1 couples to the right-handed up and charm quarks, while ϕ_2 couples to the right-handed top quark. We classify the models as follows

(i) Model I: both the down-type quarks and the charged leptons couple to ϕ_2.

(ii) Model II: the down-type quarks couple to ϕ_1 while the charged letpons couple to ϕ_2.

(iii) Model III: the charged letpons couple to ϕ_1 while the down-type quarks couple to ϕ_2.

(iv) Model IV: both the down-type quarks and the charged leptons couple to ϕ_1.

To avoid the flavor changing neutral current (FCNC) constraints [37], we introduce a Z_3 symmetry. Under this Z_3 symmetry, the quark doublets, the up-type quarks, the Higgs fields, and the singlet transform as follows:

$$q_i \leftrightarrow q_i, \qquad u_k \leftrightarrow u_k, \qquad t \leftrightarrow \omega t, \qquad \phi_1 \leftrightarrow \phi_1, \\ \phi_2 \leftrightarrow \omega\phi_2, \qquad S \leftrightarrow \omega S, \tag{1}$$

where $\omega^3 = 1$, $i = 1, 2, 3$, and $k = 1, 2$. The transformation properties for down-type quarks, lepton doublets, and charged leptons will be given later for each model. By the way, to escape the FCNC constraints in the nonsupersymmetric SMs, we just need to consider Z_2 symmetry; i.e., we change each "ω^2" and "ω" into the "$-$" sign in our transformation equations. To match the supersymmetric SMs, we consider the Z_3 symmetry in this paper.

A. Model I

Under this Z_3 symmetry, the down-type quarks, the lepton doublets, and the charged leptons transform as follows:

$$d_i \leftrightarrow \omega^2 d_i, \qquad l_i \leftrightarrow l_i, \qquad e_i \leftrightarrow \omega^2 e_i. \tag{2}$$

Then, the SM fermion Yukawa Lagrangian is

$$-\mathcal{L} = y^u_{ki}\bar{u}_k q_i \phi_1 + y^t_i \bar{t} q_i \phi_2 + y^d_{ij}\bar{d}_i q_j \tilde{\phi}_2 + y^e_{ij}\bar{e}_i l_j \tilde{\phi}_2 + \text{H.c.}, \tag{3}$$

where y^u_{ij}, y^d_{ij}, and y^e_{ij} are Yukawa couplings and $\tilde{\phi}_i = i\sigma_2\phi_i^*$. Here, σ_2 is the second Pauli matrix. In particular, to avoid the FCNC constraints [37], we assume that the Yukawa couplings y^u_{13}, y^u_{23}, y^t_1, and y^t_2 are relatively small. It is clear that in the limit $y^u_{13} = y^u_{23} = y^t_1 = y^t_2 = 0$, there is no FCNC effect. Moreover, the quark CKM mixings are generated from the down-type quark sector. Let us define

$$\tan\beta \equiv \frac{\langle\phi_2\rangle}{\langle\phi_1\rangle}. \tag{4}$$

At large $\tan\beta$, the Higgs fields with dominant components from ϕ_1 will have large Yukawa couplings with the first two-generation up-type quarks.

The most general renormalizable Higgs potential at tree level, which is invariant under the $SU(2)_L \times U(1)_Y$ gauge symmetry and the Z_3 symmetry, is

$$
\begin{aligned}
V = {} & \frac{\lambda_1}{2}(\phi_1^\dagger\phi_1)^2 + \frac{\lambda_2}{2}(\phi_2^\dagger\phi_2)^2 + \frac{\lambda_S}{2}(S^\dagger S)^2 \\
& + \frac{\lambda_3}{2}(\phi_1^\dagger\phi_1)(\phi_2^\dagger\phi_2) + \frac{\lambda_4}{2}(\phi_1^\dagger\phi_2)(\phi_2^\dagger\phi_1) \\
& + \frac{\lambda_{S1}}{2}(S^\dagger S)(\phi_1^\dagger\phi_1) + \frac{\lambda_{S2}}{2}(S^\dagger S)(\phi_2^\dagger\phi_2) \\
& + [AS\phi_2^\dagger\phi_1 + \text{H.c.}] - \frac{1}{2}m_{11}^2\phi_1^\dagger\phi_1 \\
& - \frac{1}{2}m_{22}^2\phi_2^\dagger\phi_2 - \frac{1}{2}m_S^2 S^\dagger S, \qquad (5)
\end{aligned}
$$

where λ_i, λ_S, λ_{S1}, and λ_{S2} are dimensionless parameters, m_{11}^2, m_{22}^2, and m_S^2 are mass parameters, and A is a mass dimension-one parameter which is similar to the supersymmetry breaking trilinear soft term. λ_i for $i = 1, 2, 3, 4$, λ_S, λ_{S1}, λ_{S2}, m_{11}^2, m_{22}^2 and m_S^2 are real, while A is complex. In addition, the term $\lambda_5(\phi_1^\dagger\phi_2)^2$ and its Hermitian conjugate are forbidden by discrete Z_3 symmetry. Also, the terms $\lambda_6(\phi_1^\dagger\phi_1)(\phi_1^\dagger\phi_2)$ and $\lambda_6'(\phi_2^\dagger\phi_2)(\phi_1^\dagger\phi_2)$, as well as their Hermitian conjugates, which will induce the FCNC processes [37], are forbidden in our model, too. Interestingly, our model can be consistent with the constraints from the CP violation and FCNC processes even if A is not real [38–41].

For simplicity, we assume that the up-type quark Yukawa matrix is diagonal, and then there are no tree-level FCNC processes. Also, we assume that A is relatively small, and the vacuum expectation value (VEV) of S is much larger than the VEVs of ϕ_1 and ϕ_2, for example, $\langle S\rangle \simeq 3$ TeV. Thus, the mixings between S and ϕ_i are small and can be neglected. The Lagrangian of relevance for our discussion of direct CP violation in charm meson decays can be written as

$$
\begin{aligned}
-\mathcal{L} = {} & \frac{gm_{u_k}}{2m_W}\frac{c_\alpha}{c_\beta}H\bar{u}_k u_k - \frac{gm_{u_k}}{2m_W}\frac{s_\alpha}{c_\beta}h\bar{u}_k u_k + \frac{gm_t}{2m_W}\frac{s_\alpha}{s_\beta}H\bar{t}t \\
& + \frac{gm_t}{2m_W}\frac{c_\alpha}{s_\beta}h\bar{t}t - \frac{gm_{d_j}}{2m_W}\frac{s_\alpha}{s_\beta}H\bar{d}_j d_j - \frac{gm_{d_j}}{2m_W}\frac{c_\alpha}{s_\beta}h\bar{d}_j d_j \\
& + i\frac{gm_{u_k}}{2m_W}t_\beta A\bar{u}_k\gamma^5 u_k + i\frac{gm_t}{2m_W}ct_\beta A\bar{t}\gamma^5 t \\
& + i\frac{gm_{d_j}}{2m_W}ct_\beta A\bar{d}_j\gamma^5 d_j + \frac{gm_{u_k}}{2m_W}V_{kj}t_\beta H^+\bar{u}_k P_L d_j \\
& - \frac{gm_{d_j}}{2m_W}V_{kj}ct_\beta H^+\bar{u}_k P_R d_j - \frac{gm_t}{2m_W}V_{3j}ct_\beta H^+\bar{t}P_L d_j \\
& - \frac{gm_{d_j}}{2m_W}V_{3j}ct_\beta H^+\bar{t}P_R d_j + \cdots,
\end{aligned}
$$

where $s_\alpha = \sin\alpha$, $c_\alpha = \cos\alpha$, $s_\beta = \sin\beta$, $c_\beta = \cos\beta$, $t_\beta = \tan\beta$, and $ct_\beta = \cot\beta$, with α being the mixing angle between the real components of ϕ_1^0 and ϕ_2^0.

B. Model II

Under this Z_3 symmetry, the down-type quarks, the lepton doublets, and the charged leptons transform as follows:

$$
d_i \leftrightarrow d_i, \qquad l_i \leftrightarrow l_i, \qquad e_i \leftrightarrow \omega^2 e_i. \qquad (6)
$$

So the SM fermion Yukawa Lagrangian is

$$
-\mathcal{L} = y_{ki}^u\bar{u}_k q_i\phi_1 + y_i^t\bar{t}q_i\phi_2 + y_{ij}^d\bar{d}_i q_j\tilde{\phi}_1 + y_{ij}^e\bar{e}_i l_j\tilde{\phi}_2 + \text{H.c.} \qquad (7)
$$

Similar to model I, we assume that the Yukawa couplings y_{13}^u, y_{23}^u, y_1^t, and y_2^t are relatively small. The most general renormalizable Higgs potential at tree level, which is invariant under the $SU(2)_L \times U(1)_Y$ gauge symmetry and the Z_3 symmetry, is the same as that in Eq. (5) in model I. At large $\tan\beta$, the Higgs fields with dominant components from ϕ_1 will have large Yukawa couplings with the first two-generation up-type quarks, and all the down-type quarks.

With the same assumptions as in model I, the Lagrangian of relevance for our discussion can be written as

$$
\begin{aligned}
-\mathcal{L} = {} & \frac{gm_{u_k}}{2m_W}\frac{c_\alpha}{c_\beta}H\bar{u}_k u_k - \frac{gm_{u_k}}{2m_W}\frac{s_\alpha}{c_\beta}h\bar{u}_k u_k + \frac{gm_t}{2m_W}\frac{s_\alpha}{s_\beta}H\bar{t}t \\
& + \frac{gm_t}{2m_W}\frac{c_\alpha}{s_\beta}h\bar{t}t - \frac{gm_{d_j}}{2m_W}\frac{c_\alpha}{c_\beta}H\bar{d}_j d_j + \frac{gm_{d_j}}{2m_W}\frac{s_\alpha}{c_\beta}h\bar{d}_j d_j \\
& + i\frac{gm_{u_k}}{2m_W}t_\beta A\bar{u}_k\gamma^5 u_k + i\frac{gm_t}{2m_W}ct_\beta A\bar{t}\gamma^5 t \\
& - i\frac{gm_{d_j}}{2m_W}t_\beta A\bar{d}_j\gamma^5 d_j + \frac{gm_{u_k}}{2m_W}V_{kj}t_\beta H^+\bar{u}_k P_L d_j \\
& + \frac{gm_{d_j}}{2m_W}V_{kj}t_\beta H^+\bar{u}_k P_R d_j - \frac{gm_t}{2m_W}V_{3j}ct_\beta H^+\bar{t}P_L d_j \\
& + \frac{gm_{d_j}}{2m_W}V_{3j}t_\beta H^+\bar{t}P_R d_j + \cdots.
\end{aligned}
$$

C. Model III

Under this Z_3 symmetry, the down-type quarks, the lepton doublets, and the charged leptons transform as follows:

$$
d_i \leftrightarrow \omega^2 d_i, \qquad l_i \leftrightarrow l_i, \qquad e_i \leftrightarrow e_i. \qquad (8)
$$

So the SM fermion Yukawa Lagrangian is

$$
-\mathcal{L} = y_{ki}^u\bar{u}_k q_i\phi_1 + y_i^t\bar{t}q_i\phi_2 + y_{ij}^d\bar{d}_i q_j\tilde{\phi}_2 + y_{ij}^e\bar{e}_i l_j\tilde{\phi}_1 + \text{H.c.} \qquad (9)
$$

At large $\tan\beta$, the Higgs fields with dominant components from ϕ_1 will have large Yukawa couplings with the first two-generation up-type quarks, and all the charged leptons.

The rest of the discussion is similar to those in models I and II.

D. Model IV

Under this Z_3 symmetry, the down-type quarks, the lepton doublets, and the charged leptons transform as follows:

$$d_i \leftrightarrow d_i, \qquad l_i \leftrightarrow l_i, \qquad e_i \leftrightarrow e_i. \tag{10}$$

Then, the SM fermion Yukawa Lagrangian is

$$-\mathcal{L} = y^u_{ki}\bar{u}_k q_i \phi_1 + y^t_i \bar{t} q_i \phi_2 + y^d_{ij}\bar{d}_i q_j \tilde{\phi}_1 + y^e_{ij}\bar{e}_i l_j \tilde{\phi}_1 + \text{H.c.} \tag{11}$$

At large $\tan\beta$, the Higgs fields with dominant components from ϕ_1 will have large Yukawa couplings with the first two-generation up-type quarks, all the down-type quarks, and all the charged leptons. The rest of the discussion is similar to those in models I and II.

III. SUPERSYMMETRIC STANDARD MODELS

First, let us explain the convention. We denote the chiral superfields for the quark doublets, the right-handed up-type quarks, the right-handed down-type quarks, the lepton doublets, and the right-handed charged leptons as Q_i, U^c_i, D^c_i, L_i, and E^c_i, respectively, where $i = 1, 2, 3$. We also introduce two pairs of Higgs doublets, (H_u, H_d) and (H'_u, H'_d). In addition, we introduce three SM singlet Higgs fields: S, S', and T.

Without loss of generality, we assume that H_u couples to the right-handed up and charm quarks, H'_u couples to the right-handed top quark, and H_d couples to the right-handed down-type quarks. We classify the models as follows

(i) Model A: H'_d couples to the charged letpons.
(ii) Model B: H_d couples to the charged letpons.

To solve the μ problem, we consider a $Z_3 \times Z'_3$ discrete symmetry. Under the Z_3 symmetry, the SM quarks, the Higgs fields, and the singlet fields transform as follows:

$$\begin{aligned} &Q_i \leftrightarrow \omega Q_i, \quad U^c_k \leftrightarrow \omega U^c_k, \quad T^c \leftrightarrow \omega^2 T^c, \quad D^c_i \leftrightarrow \omega D^c_i, \\ &H_{u,d} \leftrightarrow \omega H_{u,d}, \quad H'_{u,d} \leftrightarrow H'_{u,d}, \quad S \leftrightarrow \omega S, \quad S' \leftrightarrow S', \\ &T \leftrightarrow \omega^2 T, \end{aligned} \tag{12}$$

where $\omega^3 = 1$. And under the Z'_3 symmetry, the SM quarks, the Higgs fields, and the singlet fields transform as below:

$$\begin{aligned} &Q_i \leftrightarrow Q_i, \quad U^c_i \leftrightarrow U^c_i, \quad T^c \leftrightarrow \omega'^2 T^c, \quad D^c_i \leftrightarrow D^c_i, \\ &H_{u,d} \leftrightarrow H_{u,d}, \quad H'_{u,d} \leftrightarrow \omega' H'_{u,d}, \quad S \leftrightarrow S, \quad S' \leftrightarrow \omega' S', \\ &T \leftrightarrow \omega'^2 T, \end{aligned} \tag{13}$$

where $\omega'^3 = 1$.

A. Model A

Under the $Z_3 \times Z'_3$ symmetry, the lepton doublets and the charged leptons, respectively, transform as follows:

$$\begin{aligned} &L_i \leftrightarrow L_i, \qquad E^c_i \leftrightarrow E^c_i, \\ &L_i \leftrightarrow \omega' L_i, \qquad E^c_i \leftrightarrow \omega' E^c_i. \end{aligned} \tag{14}$$

Then, the SM fermion Yukawa Lagrangian is

$$\begin{aligned} W_{\text{Yukawa}} = {} & y^u_{ik} Q_i H_u U^c_k + y^t_i Q_i T^c H'_u + y^d_{ij} Q_i H_d D^c_j \\ & + y^e_{ij} L_i H'_d E^c_j + \lambda_1 S H_d H_u + \lambda_2 S' H'_d H_u{}' \\ & + \lambda_3 T H_d H_u{}' + \lambda_4 T H_d{}' H_u + \lambda_5 S S' T + \frac{\kappa_1}{3} S^3 \\ & + \frac{\kappa_2}{3} S'^3 + \frac{\kappa_3}{3} T^3, \end{aligned} \tag{15}$$

where y^u_{ik}, y^t_i, y^d_{ij}, y^e_{ij}, λ_i, and κ_i are Yukawa couplings. To avoid the FCNC constraints, we assume that the Yukawa couplings y^u_{31}, y^u_{32}, y^t_1, and y^t_2 are relatively small, similar to the nonsupersymmetric models. In our model, we define

$$\tan\beta \equiv \frac{\langle H_d \rangle}{\langle H_u \rangle}, \tag{16}$$

which is different from the traditional minimal supersymmetric standard model. The VEV of H_u can be much smaller than that of H_d since H'_u couples to the top quark, i.e., the charm Yukawa coupling can be order 1. Note that the VEV of H_d can be about one order larger that of H'_d, and we obtain that the Yukawa couplings of down-type quarks can be about one order smaller than those of charged leptons compared to the SM.

B. Model B

Under the $Z_3 \times Z'_3$ symmetry, the lepton doublets and the charged leptons, respectively, transform as follows:

$$\begin{aligned} &L_i \leftrightarrow \omega L_i, \qquad E^c_i \leftrightarrow \omega E^c_i, \\ &L_i \leftrightarrow L_i, \qquad E^c_i \leftrightarrow E^c_i. \end{aligned} \tag{17}$$

Then, the SM fermion Yukawa Lagrangian is

$$\begin{aligned} W_{\text{Yukawa}} = {} & y^u_{ik} Q_i H_u U^c_k + y^t_i Q_i T^c H'_u + y^d_{ij} Q_i H_d D^c_j \\ & + y^e_{ij} L_i H_d E^c_j + \lambda_1 S H_d H_u + \lambda_2 S' H'_d H'_u \\ & + \lambda_3 T H_d H'_u + \lambda_4 T H'_d H_u + \lambda_5 S S' T + \frac{\kappa_1}{3} S^3 \\ & + \frac{\kappa_2}{3} S'^3 + \frac{\kappa_3}{3} T^3. \end{aligned} \tag{18}$$

To avoid the FCNC constraints, similar to model A, we assume that the Yukawa couplings y^u_{31}, y^u_{32}, y^t_1, and y^t_2 are relatively small.

IV. EFFECTIVE HAMILTONIAN AND DIRECT CP ASYMMETRIES IN D MESON DECAYS

The effective Hamiltonian for the $c \to u$ transition can be written as

$$\mathcal{H}^{\text{eff}}_{\Delta C=1} = \frac{G_F}{\sqrt{2}}\left\{\sum_{p=d,s}\lambda_p(C_1^p O_1^p + C_2^p O_2^p) + \lambda_b\left[\sum_{i=3}^{6} C_i O_i + C_{7\gamma}O_{7\gamma} + C_{8g}O_{8g} + \sum_{i=11}^{16}\sum_{q=u,d,s,c} C_i^q O_i^q\right]\right\}, \tag{19}$$

with $\lambda_p = V^*_{cp}V_{up}$ $(p = d, s)$ and $\lambda_b = V^*_{cb}V_{ub}$.

The complete list of operators is given as follows

$$\begin{aligned}
O_1^p &= (\bar{u}p)_{V-A}(\bar{p}c)_{V-A},\\
O_2^p &= (\bar{u}_\alpha p_\beta)_{V-A}(\bar{p}_\beta c_\alpha)_{V-A},\\
O_3 &= (\bar{u}c)_{V-A}\sum_q(\bar{q}q)_{V-A},\\
O_4 &= (\bar{u}_\alpha c_\beta)_{V-A}\sum_q(\bar{q}_\beta q_\alpha)_{V-A},\\
O_5 &= (\bar{u}c)_{V-A}\sum_q(\bar{q}q)_{V+A},\\
O_6 &= (\bar{u}_\alpha c_\beta)_{V-A}\sum_q(\bar{q}_\beta q_\alpha)_{V+A},\\
O_{7\gamma} &= \frac{e}{8\pi^2}m_c[\bar{u}\sigma_{\mu\nu}(1+\gamma^5)c]F^{\mu\nu},\\
O_{8g} &= \frac{g_s}{8\pi^2}m_c[\bar{u}\sigma_{\mu\nu}T^a(1+\gamma^5)c]G_a^{\mu\nu},\\
O_{11}^q &= (\bar{u}c)_{S+P}(\bar{q}q)_{S-P},\\
O_{12}^q &= (\bar{u}_\alpha c_\beta)_{S+P}(\bar{q}_\beta q_\alpha)_{S-P},\\
O_{13}^q &= (\bar{u}c)_{S+P}(\bar{q}q)_{S+P},\\
O_{14}^q &= (\bar{u}_\alpha c_\beta)_{S+P}(\bar{q}_\beta q_\alpha)_{S+P},\\
O_{15}^q &= [\bar{u}\sigma_{\mu\nu}(1+\gamma^5)c][\bar{q}\sigma^{\mu\nu}(1+\gamma^5)q],\\
O_{16}^q &= [\bar{u}_\alpha\sigma_{\mu\nu}(1+\gamma^5)c_\beta][\bar{q}_\beta\sigma^{\mu\nu}(1+\gamma^5)q_\alpha],
\end{aligned} \tag{20}$$

with $V \pm A = \gamma^\mu(1 \pm \gamma^5)$ and $S \pm P = (1 \pm \gamma^5)$.

The direct CP asymmetry of $D^0 \to K^+K^-$ can be written as

$$a_{K^+K^-} = 2\text{Im}\left(\frac{\lambda_b}{\lambda_s}R^s_{K,\text{SM}}\right) + 2\text{Im}\left(\frac{\lambda_b}{\lambda_s}R^s_{K,\text{NP}}\right), \tag{21}$$

where

$$\begin{aligned}
R^s_{K,\text{SM}} &= \frac{a_4^{\text{SM}} + r_\chi a_6^{\text{SM}}}{a_1},\\
R^s_{K,\text{NP}} &= \frac{1}{a_1}\left(a_4^{\text{NP}} - \frac{1}{12}a_{12}^s + r_\chi\left(a_6^{\text{NP}} + \frac{1}{4}a_{14}^s + 3a_{16}^s\right)\right),
\end{aligned} \tag{22}$$

where maximal strong phase is assumed, and only the weak phase is included in the above equation. The a_i coefficients are estimated in naive factorization:

$$\begin{aligned}
a_4^{\text{NP}} &= 3a_6^{\text{NP}} = -\frac{3C_F\alpha_s}{2\pi N_C}C_{8g}^{\text{NP}},\\
a_{12}^s &= C_{12}^s + C_{11}^s/N_C,\\
a_{14}^s &= C_{14}^s + C_{13}^s/N_C,\\
a_{16}^s &= C_{16}^s + C_{15}^s/N_C,
\end{aligned} \tag{23}$$

where the Wilson coefficients $C_{8g,11,12,13,14,15,16}$ are evaluated at charm quark mass m_c scale. For the direct CP asymmetry of $D^0 \to \pi^+\pi^-$, the upper index s should be replaced with d. In the flavor $SU(3)$ limit, we have $a_{\pi^+\pi^-} \simeq -a_{K^+K^-}$.

The Wilson coefficients can be evolved from W boson mass m_w scale to m_c scale through the intermediate bottom quark mass scale, m_b [42]. The main contribution in our case is $C_{8g}(m_c)$, which can be written as [43–46]

$$C_{8g}(m_c) \simeq 0.4983C_{8g}(m_w) - 0.1382C_2(m_w) + 0.4922C_{13}^c(m_w). \tag{24}$$

The direct CP asymmetry in the decays $D^0 \to K^+K^-$ and $D^0 \to \pi^+\pi^-$ can be estimated as

$$\Delta a_{CP} = a_{K^+K^-} - a_{\pi^+\pi^-} \simeq [-0.01676C_{8g}^{\text{NP}}(m_w) + 0.1142C_{13}(m_w)] \times 1\%. \tag{25}$$

For $\Delta a_{CP} \sim 0.1\%$, we should have $C_{8g}^{\text{NP}}(m_w) \sim 10$, or $C_{13}(m_w) \sim 1$.

We can further express $C_{11,13}^c$ as [46]

$$C_{11}^c = \frac{e^2}{16\pi^2}(C_{Q_1}^c - C_{Q_2}^c), \qquad C_{13}^c = \frac{e^2}{16\pi^2}(C_{Q_1}^c + C_{Q_2}^c). \tag{26}$$

To follow, we will calculate $C_{Q_{1,2}}^c$ and C_{8g} at m_w scale in models I, II, and A.

The contributions to C_{8g} from charged Higgs boson exchanges are

$$C_{8g} = -\cot^2_\beta\frac{1}{6}D(x_{H^\pm}) - E(x_{H^\pm}) \tag{27}$$

in model I and

$$C_{8g} = t_\beta^2\left[-\frac{1}{6}D(x_{H^\pm}) - E(x_{H^\pm})\right] \quad (28)$$

in model II, with $x_{H^\pm} = m_b^2/m_{H^\pm}^2$. The one-loop functions D and E are defined in Ref. [47].

In our calculations, we work in the limit of vanishing light quark masses, $m_u = m_d = m_s = 0$. The Wilson coefficients $C_{Q_{1,2}}$ at the leading order of $\mathcal{O}(\tan^2\beta)$ in model I are

$$\begin{aligned}
C_{Q_1}^c &= -\frac{m_b^2 m_c^2}{4m_w^2 s_w^2 c_\beta^2}\left(\frac{c_\alpha^2}{m_H^2}+\frac{s_\alpha^2}{m_h^2}\right)[f_{b0}(x_{H^\pm}) - f_{b0}(x_W)] \\
&\quad -\frac{3}{8}\frac{m_c^2 t_\beta}{s_w^2 c_\beta}\left(\frac{c_\alpha}{m_H^2}s_{\beta-\alpha}+\frac{s_\alpha}{m_h^2}c_{\beta-\alpha}\right)f_{c00}(x_W, x_{H^\pm}) \\
&\quad +\frac{m_c^2 t_\beta^2}{12 m_{H^\pm}^2 s_w^2}|V_{cb}|^2 f_{d00}(x_W, x_W, x_{H^\pm}), \\
C_{Q_2}^c &= \frac{m_b^2 m_c^2 t_\beta^2}{4m_w^2 s_w^2 m_A^2}[f_{b0}(x_{H^\pm}) - f_{b0}(x_W)] \\
&\quad +\frac{3}{8}\frac{m_c^2 t_\beta^2}{s_w^2 m_A^2} f_{c00}(x_W, x_{H^\pm}) \\
&\quad -\frac{m_c^2 t_\beta^2}{12 m_{H^\pm}^2 s_w^2}|V_{cb}|^2 f_{d00}(x_W, x_W, x_{H^\pm}).
\end{aligned} \quad (29)$$

where the one-loop functions $f_{b0,c00,d00}$ are defined in Ref. [48].

The Wilson coefficients $C_{Q_{1,2}}$ at the leading order of $\mathcal{O}(\tan^4\beta)$ in model II are

$$\begin{aligned}
C_{Q_1}^c &= -\frac{m_b^2 m_c^2}{4m_w^2 s_w^2}\frac{t_\beta^2}{c_\beta^2}\left(\frac{c_\alpha^2}{m_H^2}+\frac{s_\alpha^2}{m_h^2}\right) f_{b0}(x_{H^\pm}) \\
&\quad -\frac{m_b^2 m_c^2}{8m_w^2 s_w^2}\frac{t_\beta^2}{c_\beta^2}\left(\frac{c_\alpha^2}{m_H^2}+\frac{s_\alpha^2}{m_h^2}\right) \\
&\quad \times\left[3 f_{c00}(x_{H^\pm}) + \frac{m_b^2}{m_{H^\pm}^2} f_{c0}(x_{H^\pm})\right] \\
&\quad +\frac{m_b^4 m_c^2}{12 m_w^2 s_w^2 m_{H^\pm}^4} t_\beta^4 |V_{cb}|^2 f_{d0}(x_{H^\pm}), \\
C_{Q_2}^c &= \frac{m_b^2 m_c^2}{4m_w^2 s_w^2}\frac{t_\beta^4}{m_A^2} f_{b0}(x_{H^\pm}) \\
&\quad +\frac{m_b^2 m_c^2}{8m_w^2 s_w^2}\frac{t_\beta^4}{m_A^2}\left[3 f_{c00}(x_{H^\pm}) + \frac{m_b^2}{m_{H^\pm}^2} f_{c0}(x_{H^\pm})\right].
\end{aligned} \quad (30)$$

The leading contributions to the Wilson coefficients C_{8g} at the order of $\mathcal{O}(\tan^0\beta)$ and $C_{Q_{1,2}}$ at the order of $\mathcal{O}(\tan^2\beta)$ in model A from gluino exchanges are

$$\begin{aligned}
C_{8g} &= -\frac{1}{72\lambda_b}\frac{g_s^2}{g^2}\frac{m_W^2}{m_{\tilde g}^2}\bigg[F_{12}(x_{\tilde g})\delta_{12}^{LL} + F'_{12}(x_{\tilde g})\delta_{12}^{LR}\delta_{22}^{LR*} \\
&\quad - 8\frac{m_{\tilde g}}{m_c}F_{34}(x_{\tilde g})\delta_{12}^{LR} - 8\frac{m_{\tilde g}}{m_c}F'_{34}(x_{\tilde g})\delta_{12}^{LL}\delta_{22}^{LR}\bigg], \\
C_{Q_1}^c &= \frac{4}{3\lambda_b}\frac{g_s^2}{g^2 s_w^2}\frac{m_c m_{\tilde g}}{c_\beta^2}\left(\frac{c_\alpha^2}{m_H^2}+\frac{s_\alpha^2}{m_h^2}\right) f'_b(x_{\tilde g})\delta_{12}^{LL}\delta_{22}^{LR}, \\
C_{Q_2}^c &= -\frac{4}{3\lambda_b}\frac{g_s^2}{g^2 s_w^2}\frac{m_c m_{\tilde g}}{m_A^2} t_\beta^2 f'_b(x_{\tilde g})\delta_{12}^{LL}\delta_{22}^{LR},
\end{aligned} \quad (31)$$

where the one-loop functions are defined in Ref. [46].

The Higgs sector is subject to strong constraints from both the Higgs coupling measurements [49] and the direct heavier Higgs searches at LHC; in particular, $pp\to\Phi\to\tau^+\tau^-$ [50,51], $pp\to\Phi\to\mu^+\mu^-$ [52], and $pp\to b\Phi\to bbb$ [53] channels, with Φ as the neutral Higgs boson. The implications of the Higgs coupling measurements are studied in Refs. [54] and [55] with direct heavier Higgs searches within the 2HDMs. Besides the up and charm quark Yukawa couplings, the other Higgs couplings in model I are the same as in 2HDM 1, and in model II are the same as in 2HDM 4 [54]. We note that the constraints in the β and $\cos(\beta-\alpha)$ plane are much looser in model I than those in model II, while the latter are tightly around the alignment limit $\alpha = \beta - \pi/2$. In the numerical calculations, we consider the large $\tan\beta$ case. The direct heavier Higgs production channels through $\tau\tau$ and $\mu\mu$ are suppressed by $\sin^2\alpha$ from Yukawa couplings in both models I and II, while the bb channel is suppressed by $\sin^2\alpha$ in model I and enhanced by $\tan^2\beta$ in model II.

For numerical estimations, we choose the following parameters in the Higgs sector for model I: $t_\beta = 50$, $s_\alpha = -0.1$, $m_h = 126$ GeV, $m_H = 180$ GeV, $m_A = 220$ GeV, and $m_{H^\pm} = 250$ GeV. In model II, the measurement of $\mathrm{Br}(B\to X_s\gamma)$ puts a stringent bound on the lower limit of the mass of the charged Higgs, $m_{H^\pm}\geq 380$ GeV at 95% C.L. [56]. With a heavy charged Higgs pair, the Higgs sector quickly approaches the decoupling limits. For numerical studies, we choose the following parameters for model II: $t_\beta = 10$, $s_\alpha = -0.1$, $m_h = 126$ GeV, $m_H \simeq m_A \simeq m_{H^\pm} = 380$ GeV. In the supersymmetric version model A, the Yukawa couplings are similiar to those in model I. We also take the supersymmetric scale $m_{\tilde g} = m_{\tilde q} = 2$ TeV [49].

The charged Higgs contributions can be calculated as $C_{8g}^{H^\pm}\simeq -0.9\times 10^{-3}$ in model I, and $C_{8g}^{H^\pm}\simeq -0.047$ in model II. The contributions to $C_{13}(m_w)$ are suppressed in both models I and II, where we have $C_{13}^c(m_w)\sim -5.2\times 10^{-7}$ in model I, and $C_{13}^c(m_w)\sim -1.95\times 10^{-8}$ in model II. Therefore, due to the experimental constraints, the charged Higgs contributions cannot accommodate the direct *CP* measurement of charm decays.

In model A, for double insertion of $(\delta_{12}^{LL}\delta_{22}^{LR})$, we have $C_{8g}^{\tilde g}\sim 7.19\times\frac{(\delta_{12}^{LL}\delta_{22}^{LR})}{10^{-3}}$ and $C_{13}^c\sim -1.0\times\frac{(\delta_{12}^{LL}\delta_{22}^{LR})}{10^{-3}}$ from gluino exchange. For $(\delta_{12}^{LL}\delta_{22}^{LR})$ at the order of 10^{-3}, we can have both C_{8g} at the order of 10 and C_{13}^c of order 1, which are possible to accommodate the direct CP measurement of charm decays.

The constraint from the $D^0 - \bar{D}^0$ system can be found in Ref. [57]. The nonvanishing Wilson coefficients z_i $(i = 1, 2, \ldots 5)$ are

$$z_2 = \frac{g^4}{64\pi^2}\frac{\Lambda_{\rm NP}^2}{m_W^2}|\lambda_b|^2\frac{m_c^2}{m_W^2}x_W^2[I_2(x_W, x_W/x_{H^\pm}) - 2I_3(x_W, x_W/x_{H^\pm})] \quad (32)$$

at the leading order of $\mathcal{O}(t_\beta^0)$ in model I, and

$$z_2 = \frac{g^4}{64\pi^2}\frac{\Lambda_{\rm NP}^2}{m_W^2}|\lambda_b|^2 t_\beta^4 x_W^2 \times\left[\frac{1}{4}I_1(x_W, x_W/x_{H^\pm}) + \frac{m_c^2}{m_W^2}I_2(x_W, x_W/x_{H^\pm})\right] \quad (33)$$

at the leading order of $\mathcal{O}(t_\beta^4)$ in model II. The loop functions $I_{1,2,3}$ are defined in Ref. [58]. We can calculate z_2 for the above parameters, $z_2 \simeq -1.8\times 10^{-18}$ in model I, and $z_2 \simeq 7.7\times 10^{-13}(\frac{t_\beta}{10})^4$ in model II, which are below the experimental limits.

In model A, we obtain the gluino contributions

$$z_1 = -\frac{\alpha_s^2}{216}(\delta_{12}^{LL})^2[66\tilde{f}_6(m_{\tilde{q}}^2/m_{\tilde{g}}^2) + 24 f_6(m_{\tilde{q}}^2/m_{\tilde{g}}^2)],$$
$$\tilde{z}_2 = -\frac{\alpha_s^2}{216}(\delta_{12}^{LL}\delta_{22}^{LR})^2 204 f(x), \quad (34)$$

for $\Lambda_{\rm NP} = m_{\tilde{g}}$, where the functions f_6 and $\tilde{f}_6$ are given in Ref. [59], and f is defined as follows:

$$f(x) = \frac{60x^4(5+x)\ln(x) - 197x^5 - 25x^4 + 300x^3 - 100x^2 + 25x - 3}{60(x-1)^7}.$$

The leading order contributions from $(\delta_{12}^{LL})^2$ are included in z_1. In the numerical estimations, we take $\delta_{22}^{LR} = (m_c A_c - m_c\mu\tan\beta)/m_{\tilde{q}}^2 \simeq -m_c\mu\tan\beta/m_{\tilde{q}}^2 \simeq -0.015$ (with $\mu\sim 1.2 m_{\tilde{q}}$, $m_c \sim 0.5$ GeV when running to $m_{\tilde{q}}$ scale), and $\delta_{12}^{LL} \simeq 0.067$. With the parameter for model A, we have $z_1 \simeq 3.0\times 10^{-7}(\frac{\delta_{12}^{LL}}{0.067})^2$, and $\tilde{z}_2 \simeq -3.2\times 10^{-10}(\frac{\delta_{12}^{LL}\delta_{22}^{LR}}{10^{-3}})^2$, which are below the limits from the constraints of the $D^0 - \bar{D}^0$ system. However, due to the $SU(2)$ gauge invariance, the left-left up-type squark matrix is related to the down-type one. And we have $\delta_{12}^{LL} \simeq 0.067$ for down-type squarks, which does not satisfy the constraints from kaon system for the imaginary part ${\rm Im}(\delta_{12}^{LL}) \le 0.023$ with the supersymmetry scale at 2 TeV [60]. One way out is to consider the contributions of chirally opposite operators. We can get similar results if the above δ_{12}^{LL} is replaced with $\delta_{12}^{RR} \sim 0.067$, and δ_{22}^{LR} with $\delta_{22}^{LR*} \sim -0.015$. In this case, the up-type and down-type right-right squark matrixes are not related. Hence, the constraints from the kaon system are relaxed.

Recently, LHCb Collaboration has measured the leptonic and semileptonic decays of the charm meson; the upper limits are $B(D^0 \to \mu^+\mu^-) < 6.2(7.6)\times 10^{-9}$ at 90% (95%) C.L. [61] and $B(D^+ \to \pi^+\mu^+\mu^-) < 7.3(8.3)\times 10^{-8}$ at 90% (95%) C.L. [62]. The experimental bound on radiative charm decay is $B(D^0 \to \gamma\gamma) < 2.2\times 10^{-6}$ at 90% C.L. from the *BABAR* Collaboration [63], and $B(D^0 \to \gamma\gamma) < 4.7\times 10^{-6}$ at 90% C.L. from BESIII [64].

The corresponding Wilson coefficients are

$$C_{7\gamma} = G(x_{H^\pm}) + \frac{1}{6}\cot^2\beta A(x_{H^\pm}) \quad (35)$$

in model I,

$$C_{7\gamma} = t_\beta^2\left[G(x_{H^\pm}) + \frac{1}{6}A(x_{H^\pm})\right] \quad (36)$$

in model II, and

$$C_9 = -\frac{-1+4s_W^2}{s_W^2}\cot^2\beta\frac{x_W}{2}B(x_{H^\pm}) + \cot^2\beta x_{H^\pm}F(x_{H^\pm}),$$
$$C_{10} = -\frac{1}{s_W^2}\cot^2\beta\frac{x_W}{2}B(x_{H^\pm}) \quad (37)$$

in model I, while replacing $\cot^2\beta$ with t_β^2 in model II. The functions A, B, G, and F for the $c\to u$ transitions are defined as

$$A(x) = -\frac{x}{12}\left(\frac{5-10x-7x^2}{(1-x)^3} + \frac{6x(1-3x)\ln x}{(1-x)^4}\right),$$
$$B(x) = -\frac{x}{4}\left(\frac{1}{1-x} + \frac{\ln x}{(1-x)^2}\right),$$
$$F(x) = \frac{11-25x+40x^2}{54(1-x)^3} + \frac{2-3x+3x^3}{18(1-x)^4},$$
$$G(x) = -\frac{x}{6}\left(\frac{2}{(1-x)^2} - \frac{(1-3x)\ln x}{(1-x)^3}\right), \quad (38)$$

which differ from the ones in Ref. [65] for the $b\to s$ transitions.

The leading order contributions to the Wilson coefficients $C_{7\gamma,9,10}$ at the order of $\mathcal{O}(\tan^0\beta)$ in model A from gluino exchanges are

$$C_{7\gamma} = \frac{2}{72\lambda_b}\frac{g_s^2}{g^2}\frac{m_W^2}{m_{\tilde{g}}^2}\left[F_2(x_{\tilde{g}})\delta_{12}^{LL} + F_2'(x_{\tilde{g}})\delta_{12}^{LR}\delta_{22}^{LR*} - 4\frac{m_{\tilde{g}}}{m_c}F_4(x_{\tilde{g}})\delta_{12}^{LR} - 4\frac{m_{\tilde{g}}}{m_c}F_4'(x_{\tilde{g}})\delta_{12}^{LL}\delta_{22}^{LR}\right],$$

$$C_9 = \frac{4}{72\lambda_b}\frac{g_s^2}{g^2}\frac{m_W^2}{m_{\tilde{g}}^2}[f_6'(x_{\tilde{g}})\delta_{12}^{LL} + f_6''(x_{\tilde{g}})\delta_{12}^{LR}\delta_{22}^{LR*}], - \frac{1}{2\lambda_b s_W^2}\frac{g_s^2}{g^2}(-1+4s_W^2)[-f_{c00}^{(1)}(x_{\tilde{g}})\delta_{12}^{LR}\delta_{22}^{LR*} + f_{c00}^{(2)}(x_{\tilde{g}})\delta_{12}^{LL} + f_{c00}^{(3)}(x_{\tilde{g}})\delta_{12}^{LR}\delta_{22}^{LR*}],$$

$$C_{10} = -\frac{1}{2\lambda_b s_W^2}\frac{g_s^2}{g^2}[-f_{c00}^{(1)}(x_{\tilde{g}})\delta_{12}^{LR}\delta_{22}^{LR*} + f_{c00}^{(2)}(x_{\tilde{g}})\delta_{12}^{LL} + f_{c00}^{(3)}(x_{\tilde{g}})\delta_{12}^{LR}\delta_{22}^{LR*}], \tag{39}$$

where the one-loop functions are defined as follows: $f_6'(x) = x\frac{\partial f_6(x)}{\partial x}$, $f_6''(x) = \frac{x^2}{2}\frac{\partial^2 f_6(x)}{\partial x^2}$, $f_{c00}^{(1)}(x) = \frac{x^2}{2}\frac{\partial^2 f_{c00}(x,y)}{\partial x \partial y}|_{y->x}$, $f_{c00}^{(2)}(x) = x\frac{\partial f_{c00}(x,x)}{\partial x}$, $f_{c00}^{(3)}(x) = \frac{x^2}{2}\frac{\partial^2 f_{c00}(x,x)}{\partial x^2}$, and $F_{2(4)}$, $F'_{2(4)}$, and $f_{6(c00)}$ are defined in Ref. [46].

In model I, the short distance (SD) contribution from the charged Higgs exchange is negligible, $B(D^0 \to \gamma\gamma) \sim 10^{-14}$. In model II, the contribution can be estimated as $B(D^0 \to \gamma\gamma) = 2.8 \times 10^{-11}$. In model A with a double insertion of $(\delta_{12}^{LL}\delta_{22}^{LR})$, we have $C_{7\gamma}^{\tilde{g}} \sim -2.05 \times \frac{(\delta_{12}^{LL}\delta_{22}^{LR})}{10^{-3}}$ from gluino exchange. The SD contribution can be estimated as $B(D^0 \to \gamma\gamma) = 5.7 \times 10^{-7}$. In all three models, we have $B(D^0 \to \mu^+\mu^-)$ and $B(D^+ \to \pi^+\mu^+\mu^-)$ far below the current experimental bounds.

V. CONCLUSION

We proposed a class of the two-Higgs-doublet SMs with a SM singlet and a class of supersymmetric SMs with two pairs of Higgs doublets, where the right-handed up/charm quarks and the right-handed top quark have different quantum numbers under extra discrete symmetries. So the right-handed up and charm quarks couple to one Higgs doublet field, while the right-handed top quark couples to another Higgs doublet. We have studied the direct *CP* asymmetries in charm hadronic decays in models I, II, and A. We found that the large direct *CP* asymmetry difference cannot be accommodated within model I and II with the contributions of charged Higgs bosons. In model A, we can accommodate the experimental measurement of direct *CP* asymmetry with both O_{8g} and O_{13} operators, while the constraints from the $\Delta c = 2$ and $\Delta c = 1$ processes are satisfied.

We leave the detailed studies on phenomenological consequences of our models to the future.

ACKNOWLEDGMENTS

This research was supported in part by the Natural Science Foundation of China under Grants No. 10075069, No. 11375248, No. 10821504, No. 11075194, No. 11135003, No. 11135009, and No. 11275246.

[1] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[2] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **726**, 88 (2013); Report No. ATLAS-CONF-2013-079, 2013; Report No. ATLAS-CONF-2013-034, 2013; D. L. Mateos lecture at EPS 2013 for the ATLAS Collaboration.
[3] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[4] S. Chatrchyan *et al.* (CMS Collaboration), J. High Energy Phys. 06 (2013) 081; Report No. CMS-PAS-HIG-13-005, 2013.
[5] T. Aaltonen *et al.* (CDF and D0 Collaborations), Phys. Rev. Lett. **109**, 071804 (2012); Phys. Rev. D **88**, 052014 (2013).
[6] T. D. Lee, Phys. Rev. D **8**, 1226 (1973); Phys. Rep. **9C**, 143 (1974); P. Sikivie, Phys. Lett. **65B**, 141 (1976).
[7] J. F. Gunion, H. E. Haber, G. L. Kane, and S. Dawson, Front. Phys. **80**, 1 (2000); G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher, and J. P. Silva, Phys. Rep. **516**, 1 (2012).
[8] G. C. Branco, P. M. Ferreira, L. Lavoura, M. N. Rebelo, M. Sher, and J. P. Silva, Phys. Rep. **516**, 1 (2012).
[9] R. Aaij *et al.* (LHCb Collaboration), Phys. Rev. Lett. **108**, 111602 (2012).
[10] T. Aaltonen *et al.* (CDF Collaboration), Phys. Rev. Lett. **109**, 111801 (2012).
[11] B. R. Ko (Belle Collaboration), *Proc. Sci.*, ICHEP2012 (2013) 353 [arXiv:1212.1975].
[12] B. Aubert *et al.* (*BABAR* Collaboration), Phys. Rev. Lett. **100**, 061803 (2008).
[13] Y. Amhis *et al.* (Heavy Flavor Averaging Group Collaboration), arXiv:1207.1158; online update at http://www.slac.stanford.edu/xorg/hfag.
[14] R. Aaij *et al.* (LHCb Collaboration), Report No. LHCb-CONF-2013-003, 2013.
[15] R. Aaij *et al.* (LHCb Collaboration), Phys. Lett. B **723**, 33 (2013).
[16] R. Aaij *et al.* (LHCb Collaboration), J. High Energy Phys. 07 (2014) 041.
[17] G. Isidori, J. F. Kamenik, Z. Ligeti, and G. Perez, Phys. Lett. B **711**, 46 (2012).
[18] H.-n. Li, C.-D. Lu, and F.-S. Yu, Phys. Rev. D **86**, 036012 (2012).

[19] H.-Y. Cheng and C.-W. Chiang, Phys. Rev. D **85**, 034036 (2012); **85**, 079903(E) (2012).
[20] D. Pirtskhalava and P. Uttayarat, Phys. Lett. B **712**, 81 (2012).
[21] B. Bhattacharya, M. Gronau, and J. L. Rosner, Phys. Rev. D **85**, 054014 (2012).
[22] E. Franco, S. Mishima, and L. Silvestrini, J. High Energy Phys. 05 (2012) 140.
[23] H.-Y. Cheng and C.-W. Chiang, Phys. Rev. D **86**, 014014 (2012).
[24] G. Hiller, M. Jung, and S. Schacht, Phys. Rev. D **87**, 014024 (2013).
[25] J. Brod, arXiv:1302.3542.
[26] A. N. Rozanov and M. I. Vysotsky, JETP Lett. **95**, 397 (2012).
[27] C.-H. Chen, C.-Q. Geng, and W. Wang, Phys. Lett. B **718**, 946 (2013).
[28] C.-H. Chen, C.-Q. Geng, and W. Wang, Phys. Rev. D **85**, 077702 (2012).
[29] X. Chang, M.-K. Du, C. Liu, J.-S. Lu, and S. Yang, arXiv:1201.2565.
[30] G. F. Giudice, G. Isidori, and P. Paradisi, J. High Energy Phys. 04 (2012) 060.
[31] C. Delaunay, J. F. Kamenik, G. Perez, and L. Randall, J. High Energy Phys. 01 (2013) 027.
[32] B. Keren-Zur, P. Lodone, M. Nardecchia, D. Pappadopulo, R. Rattazzi, and L. Vecchi, Nucl. Phys. **B867**, 429 (2013).
[33] L. Da Rold, C. Delaunay, C. Grojean, and G. Perez, J. High Energy Phys. 02 (2013) 149.
[34] Y. Bai, J. Berger, J. L. Hewett, and Y. Li, J. High Energy Phys. 07 (2013) 166.
[35] W. Altmannshofer, R. Primulando, C.-T. Yu, and F. Yu, J. High Energy Phys. 04 (2012) 049.
[36] A. Dighe, D. Ghosh, and B. P. Kodrani, Phys. Rev. D **89**, 096008 (2014).
[37] S. L. Glashow and S. Weinberg, Phys. Rev. D **15**, 1958 (1977).
[38] A. G. Akeroyd, A. Arhrib, and E. M. Naimi, Phys. Lett. B **490**, 119 (2000).
[39] I. F. Ginzburg and I. P. Ivanov, arXiv:hep-ph/0312374.
[40] J. Horejsi and M. Kladiva, Eur. Phys. J. C **46**, 81 (2006).
[41] I. F. Ginzburg and I. P. Ivanov, Phys. Rev. D **72**, 115010 (2005).
[42] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
[43] F. Borzumati, C. Greub, T. Hurth, and D. Wyler, Phys. Rev. D **62**, 075005 (2000).
[44] G. Hiller and F. Kruger, Phys. Rev. D **69**, 074020 (2004).
[45] J.-F. Cheng, C.-S. Huang, and X.-h. Wu, Phys. Lett. B **585**, 287 (2004).
[46] J.-F. Cheng, C.-S. Huang, and X.-H. Wu, Nucl. Phys. **B701**, 54 (2004).
[47] B. Grinstein, R. P. Springer, and M. B. Wise, Nucl. Phys. **B339**, 269 (1990).
[48] C.-S. Huang and X.-H. Wu, Nucl. Phys. **B657**, 304 (2003).
[49] J. Beringer *et al.* (Particle Data Group Collaboration), Phys. Rev. D **86**, 010001 (2012) and 2013 partial update for the 2014 edition.
[50] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **713**, 68 (2012); CMS Collaboration, Report No. CMS-PAS-HIG-12-050, 2012.
[51] G. Aad *et al.* (ATLAS Collaboration), J. High Energy Phys. 02 (2013) 095.
[52] CMS Collaboration, Report No. CMS-PAS-HIG-12-011, 2012.
[53] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **722**, 207 (2013).
[54] N. Craig, J. Galloway, and S. Thomas, arXiv:1305.2424.
[55] B. Dumont, J. F. Gunion, Y. Jiang, and S. Kraml, arXiv:1405.3584.
[56] T. Hermann, M. Misiak, and M. Steinhauser, J. High Energy Phys. 11 (2012) 036.
[57] O. Gedalia, Y. Grossman, Y. Nir, and G. Perez, Phys. Rev. D **80**, 055024 (2009).
[58] V. D. Barger, J. L. Hewett, and R. J. N. Phillips, Phys. Rev. D **41**, 3421 (1990).
[59] Y. Grossman, A. L. Kagan, and Y. Nir, Phys. Rev. D **75**, 036008 (2007).
[60] M. Ciuchini, V. Lubicz, L. Conti, A. Vladikas, A. Donini, E. Franco, G. Martinelli, I. Scimemi *et al.*, J. High Energy Phys. 10 (1998) 008.
[61] R. Aaij *et al.* (LHCb Collaboration), Phys. Lett. B **725**, 15 (2013).
[62] R. Aaij *et al.* (LHCb Collaboration), Phys. Lett. B **724**, 203 (2013).
[63] J. P. Lees *et al.* (*BABAR* Collaboration), Phys. Rev. D **85**, 091107 (2012).
[64] H. Muramatsu, arXiv:1208.4744.
[65] B. Grinstein, M. J. Savage, and M. B. Wise, Nucl. Phys. **B319**, 271 (1989).