CP Violation in B Decays in a 2-Higgs Doublet Model for the Top Quark*

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In the absence of natural flavor conservation, multi-Higgs doublet models generally contain new sources of CP violation and anomalous charged Higgs Yukawa couplings. We present a charged-Higgs CP violation study of one such two-Higgs doublet model (2HDM) which treats the top quark differently from the other quarks. The phenomenological implications for the $K\bar{K}$ system and for B decays differ significantly from those of the standard model (SM) and of the 2HDM's with natural flavor conservation. In particular, the SM phase in this model could take a wide range of values, and the CP asymmetry in the "gold-plated" decay mode $B \to J/\psi K_S$ could be quite different from and even of opposite sign relative to the SM prediction. A new mechanism is also noted for generating a large neutron electric dipole moment which is close to the present experimental limit.

I. INTRODUCTION

Though CP violation has only been observed in the neutral kaon system, both the standard model (SM) and its many extensions predict large CP-violating effects in B decays. One of the main programs at the B factories is to test the SM Cabibbo-Kobayashi-Maskawa (CKM) paradigm of CP violation through, for example, measurement of the angles of unitarity triangle. In particular, the angle $\beta_{\rm CKM} \equiv \arg\left(-V_{cd}V_{cb}^*/V_{td}V_{tb}^*\right)$ can be cleanly related within the SM to the CP asymmetry in the "gold-plated" decay of $B \to J/\psi K_S$. The current SM fit already places quite a nontrivial constraint on this angle: $\sin(2\beta_{\rm CKM}) = .75 \pm .10$ [1]. Therefore direct experimental measurement of the angles of the unitarity triangle with values in gross disagreement with the SM expectation will be a clear signal for new sources of CP violation. It is thus important to explore the pattern of CP violation in B decays in different models of CP violation.

Two Higgs doublet models (2HDM) without natural flavor conservation [2,3] (NFC), also called Model III [4], contain in general both tree-level flavor-changing neutral Higgs (FCNH) couplings and new CP-violating phases beside the CKM phase. The phenomenology of this class of models thus differs significantly from that of the much studied 2HDM's with NFC (*i.e.,* Models I and II [5]) where neither tree level FCNH couplings nor new CP violating phases are present. It is worth noting that the abandonment of NFC is not against any fundamental principle and in fact it opens up many possibilities for flavor physics. Depending on the ansatz for FCNH couplings, type III 2HDM need not run into problems with the low energy data such as $K\bar{K}$ mixing. In this talk, we present a CP violation study of a specific type III 2HDM introduced in [6] where the top quark is treated differently than the other quarks. As will be shown, the charged Higgs sector of the model could lead to a distinctive pattern of CP violation in B decays even after imposing all the experimental constraints.

II. THE T2HDM

The top-quark 2HDM (T2HDM) [6] distinguishes itself from 2HDM's with NFC by the special status it gives to the top quark, as evidenced in the Yukawa couplings,

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$$\mathcal{L}_Y = -\overline{L}_L\phi_1 E\ell_R - \overline{Q}_L\phi_1 F d_R - \overline{Q}_L\widetilde{\phi}_1 G\mathbf{1}^{(1)}u_R - \overline{Q}_L\widetilde{\phi}_2 G\mathbf{1}^{(2)}u_R + \text{h.c.}, \tag{1}$$

where $\widetilde{\phi}_i = i\sigma^2\phi_i^*$ $(i = 1,2)$, and where E, F and G are 3×3 matrices in generation space; $\mathbf{1}^{(1)} \equiv \text{diag}(1,1,0)$; $\mathbf{1}^{(2)} \equiv \text{diag}(0,0,1)$; and Q_L and L_L are the usual left-handed quark and lepton doublets. Note that among the right-handed quarks, only t_R couples to ϕ_2, and that the mass hierarchy between the top quark and the other quarks is understood as a result of the hierarchy between the VEV's of ϕ_2 and ϕ_1. This form of the Yukawa interactions can be seen as a consequence of some discrete symmetry. Denoting the ratio of the two VEV's by $\tan\beta = v_2/v_1$, the model then requires a large $\tan\beta$ in accordance with the large mass ratio of the top and bottom quarks, m_t/m_b. In our analysis, we will always take $\tan\beta \geq 10$.

The absence of NFC in the model results in tree level FCNH couplings among the up-type quarks, whose contribution to $D\bar{D}$ mixing is dependent on the mixing among the right-handed up-type quarks [6]. For the charged Higgs sector, we obtain the following Yukawa Lagrangian [7],

$$\begin{aligned}\mathcal{L}_Y^C = (g/\sqrt{2}m_W)\left\{-\overline{u}_L V m_d d_R\left[G^+ - \tan\beta H^+\right] + \overline{u}_R m_u V d_L\left[G^+ - \tan\beta H^+\right]\right. \\ \left. + \overline{u}_R \Sigma^\dagger V d_L\left[\tan\beta + \cot\beta\right] H^+ + \text{h.c.}\right\},\end{aligned} \tag{2}$$

where $G^\pm$ and $H^\pm$ represent the would-be Goldstone bosons and the physical charged Higgs bosons, respectively. Here m_u and m_d are the diagonal up- and down-type quark mass matrices, V is the usual CKM matrix, and $\Sigma \equiv m_u U_R^\dagger \mathbf{1}^{(2)} U_R$ with U_R being the unitary rotation that brings the right-handed up-type quarks from gauge eigenstates to mass eigenstates. Unlike 2HDM's with NFC, the T2HDM depends on the U_R rotation. For our analysis, we take the simple parameterization for U_R [6,7],

$$U_R = \begin{pmatrix} \cos\phi & -\sin\phi & 0 \\ \sin\phi & \cos\phi & 0 \\ 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} 1 & 0 & 0 \\ 0 & \sqrt{1-|\epsilon_{ct}\xi|^2} & -\epsilon_{ct}\xi^* \\ 0 & \epsilon_{ct}\xi & \sqrt{1-|\epsilon_{ct}\xi|^2} \end{pmatrix}, \tag{3}$$

where $\epsilon_{ct} \equiv m_c/m_t$ and $\xi = |\xi|e^{-i\delta}$ is a complex number of order unity. The Σ matrix then depends only on the unknown parameter ξ,

$$\Sigma = \begin{pmatrix} 0 & 0 & 0 \\ 0 & m_c\epsilon_{ct}^2|\xi|^2 & m_c\epsilon_{ct}\xi^*\sqrt{1-|\epsilon_{ct}\xi|^2} \\ 0 & m_c\xi\sqrt{1-|\epsilon_{ct}\xi|^2} & m_t\left(1-|\epsilon_{ct}\xi|^2\right) \end{pmatrix}. \tag{4}$$

For numerical analysis, we will assume $|\xi| = 1$.

Several features of the model can be noted. First, it contains a new CP violating phase δ through its dependence on U_R. By comparison, 2HDM's with NFC (Models I and II) involve only the CKM phase. Secondly, charm quark Yukawa couplings of the type $H^+\bar{c}_R q_L$ $(q = d, s, b)$ are proportional to $\tan\beta$ and thus are enhanced in this model. In contrast, the corresponding couplings in Models I and II of the 2HDM are suppressed by $1/\tan\beta$. It is the presence of the new phase that leads to a novel CP-violating pattern in B decays, as will be shown later. On the other hand, the enhanced charm Yukawa's draw our immediate attention to the $K\bar{K}$ system, $b\to s\gamma$ and $B\to X_c\tau\nu$.

III. EXPERIMENTAL CONSTRAINTS

A. the $K\bar{K}$ system

Recall that there exists no tree level FCNH coupling among the down-type quarks in the present model. The most important contribution to the $K\bar{K}$ system from the charged Higgs sector is associated with the $HHcc$ box diagram which has a huge enhancement factor of $\tan^4\beta$. Therefore, both the $K\bar{K}$ mass difference Δm_K and the CP violation parameter ϵ_K are expected to place stringent constraints on the parameter space of the model.

1. Δm_K

In the SM the short-distance contribution to Δm_K is dominated by the charm quark and the long distance contribution can not be estimated reliably though is expected to be quite significant [8]. Including both the SM and the charged Higgs effects, we obtain in the T2HDM the total short distance contribution to the mass difference as,

$$(\Delta m_K)_{SD} = (G_F^2/6\pi^2) f_K^2 B_K m_K \lambda_c^2 \times \left[m_c^2 \eta_1 + (m_c^4 \tan^4 \beta / 4 m_H^2) \eta_1' \right] \quad (5)$$

where the first term is the SM charm contribution and the second term is due to the dominant $HHcc$ box diagram. Here B_K is the usual bag factor, $\lambda_c = V_{cs} V_{cd}^*$, and η_1 and η_1' are the QCD corrections to the two box diagrams. The SM top quark contribution is a few percent of the charm quark effect and is not included here. Similarly, in comparison to the $HHcc$ box diagram, the contributions from $WHcc$ and other box diagrams are small in the large $\tan\beta$ limit and are ignored.

To numerically deduce the allowed parameter space subject to the Δm_K constraint, we use the method described in [1] for error analysis. Assuming the magnitude of the long distance contribution to Δm_K to be no larger than 30% and simply taking $\eta_1' = \eta_1$, we find the 95% C.L. limit,

$$m_H / \tan^2 \beta > 0.48 \text{ GeV} \quad (6)$$

which is valid for $\tan\beta > 10$. As a result of the $\tan^4\beta$ dependence in Δm_K, this constraint puts a severe lower bound on the charged Higgs mass when $\tan\beta$ is large. By comparison, the charged Higgs effect on Δm_K in 2HDM's with NFC is suppressed by $\cot^4\beta$ and is thus negligible. Note also that the short distance effect is independent of the mixing parameter ξ of the T2HDM.

2. ϵ_K

Unlike Δm_K, the CP violation parameter ϵ_K is short distance dominated. In the SM, the dominant component of ϵ_K comes from the tt and tc box diagrams. The complete expression and its numerical evaluation, including the next-to-leading-order (NLO) QCD corrections, can be found in [8]. The charged Higgs contribution to ϵ_K is still dominated by the charm quark, with the imaginary part of the $HHcc$ box diagram now dependent on both the magnitude and the phase of ξ. More explicitly, this contribution from charged Higgs exchange is given by,

$$\epsilon_K^H = e^{i\frac{\pi}{4}} C_\epsilon B_K A \lambda^4 \eta_1' \sqrt{\rho^2 + \eta^2} \sin(\gamma + \delta) |\xi| (m_c \tan\beta)^4 / 4 m_W^2 m_H^2 \quad (7)$$

where A, λ, ρ, and η are the CKM parameters in the Wolfenstein parameterization [9], $\gamma \equiv \tan^{-1} \eta/\rho$ is the CKM phase, and $C_\epsilon = G_F^2 f_K^2 m_W^2 m_K / 6\sqrt{2}\pi^2 \Delta m_K = 3.78 \times 10^4$.

Due to its dependence on the new phase δ, ϵ_K no longer restricts the CKM angle γ. In fact, current data allows a wide range of values for γ. We can obtain bounds on the parameter $Y \equiv \sin(\gamma+\delta)|\xi|(\tan\beta/20)^4(200 \text{ GeV}/m_H)^2$ for any given value of γ by allowing $\sqrt{\rho^2+\eta^2}$ to vary within its 1σ uncertainties derived from $b \to ue\nu$. For a real CKM matrix ($\gamma = 0°$), we obtain the 95% C.L bound $0.08 < Y < 0.39$; the bound becomes $0.14 < Y < 0.65$ for $\gamma = -45°$. If γ takes its SM central value of 68° [1], the bound becomes nearly symmetric about zero as expected: $-0.085 < Y < 0.08$. Unlike Δm_K, the CP violation parameter ϵ_K imposes a constraint on the $(m_H, \tan\beta)$ plane that is dependent on $|\xi|$ and δ. As for Δm_K, the charged Higgs contribution to ϵ_K in 2HDM's with NFC is suppressed by $\cot^4\beta$ and is thus negligible when $\tan\beta$ is large.

B. the decay $b \to s\gamma$

The agreement between the experimental measurement and the SM prediction for the $b \to s\gamma$ decay rate places stringent constraints on possible flavor physics beyond the SM. In Model II for example, where the most important correction comes from the top quark and charged Higgs loop, a $\tan\beta$-independent lower bound of about 370 GeV can be imposed on the charged Higgs mass [10].

This picture changes dramatically in the T2HDM where the charm Yukawa's are greatly enhanced relative to those in Model II. Consequently, both top and charm loops involving charged Higgs exchange become important.

Furthermore, both amplitudes become complex due to their dependence on the new phase δ. As a result, the charged Higgs amplitude can interfere either constructively or destructively with the SM amplitude depending on the phase δ. The bound on the charged Higgs mass will be dependent on both $\tan\beta$ and ξ, and a relatively light Higgs can still be allowed.

For the numerical estimate, we will simply neglect the effect of the scalar operator $\bar{c}_R b_L \bar{s}_L c_R$ induced from Higgs exchange, and concentrate on the effect due to the charged Higgs correction to the matching condition for the leading-order (LO) Wilson coefficient $C_7^{(0)}$ at the scale m_W:

$$\delta C_7^{(0)}(m_W) = \sum_{u=c,t} \kappa^u \left[-\tan^2\beta + \left(\Sigma^T V^*\right)_{us}\left(\tan^2\beta+1\right)/m_u V^*_{us}\right] \\ \times \left\{B(y_u) + A(y_u)\left[-1+\left(\Sigma^\dagger V\right)_{ub}\left(\cot^2\beta+1\right)/m_u V_{ub}\right]/6\right\}. \tag{8}$$

In this expression $\kappa^u = \pm 1$ for $u = c, t$, $y_u = (m_u/m_H)^2$, and A and B are the standard expressions [11]. The effective Wilson coefficient at the scale $\mu \sim \mathcal{O}(m_b)$ is modified according to

$$C_7^{(0)\text{eff}}(\mu) \to C_7^{(0)\text{eff}}(\mu) + (\alpha_s(m_W)/\alpha_s(\mu))^{16/23}\delta C_7^{(0)}(m_W)\,. \tag{9}$$

The $b \to s\gamma$ constraint on the $(m_H, \tan\beta)$ plane, together with the constraints from Δm_K and ϵ_K, are presented in Fig. 1 for three representative choices of the CKM phase γ.

The experimental data on $B \to X_c \tau\nu$ and on other decays do not lead to new exclusion regions after imposing the $b \to s\gamma$ and $K\bar{K}$ constraints.

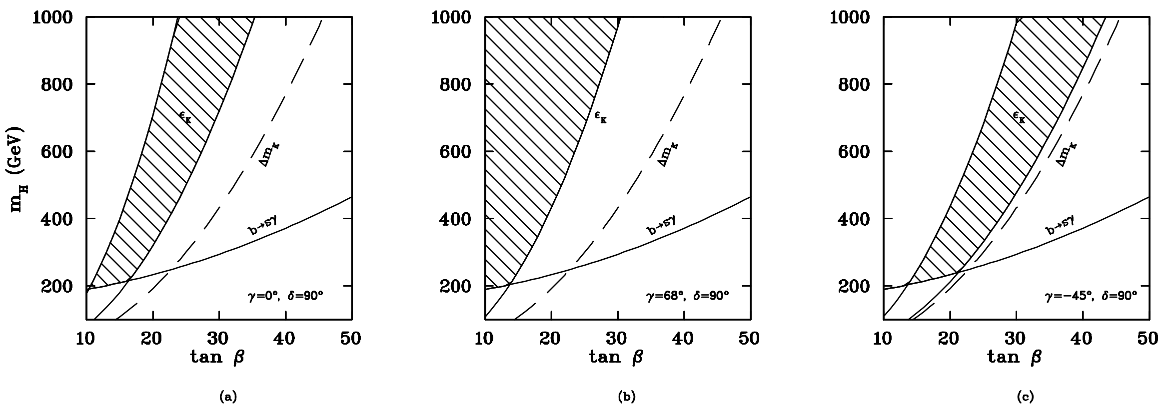

FIG. 1. Experimental constraints on the T2HDM from $b \to s\gamma$, ϵ_K and Δm_K for three representative choices of the CKM phase γ and for $\delta = 90°$ and $|\xi| = 1$: (a) $\gamma = 0°$, (b) $\gamma = 68°$ (SM central value), (c) $\gamma = -45°$. The allowed regions are shaded.

IV. CP ASYMMETRY IN $B \to J/\psi K_S$

In the SM, the time-dependent CP asymmetry $a(t) \equiv \left[\Gamma(B(t)) - \Gamma(\overline{B}(t))\right]/\left[\Gamma(B(t)) + \Gamma(\overline{B}(t))\right]$ for $B \to \psi K_S$ is free of hadronic uncertainties and is given by $a_{\rm SM}(t) = -\sin(2\beta_{\rm CKM})\sin(\Delta M t)$ [12], where $\Delta M = M_{B_H} - M_{B_L}$ is the mass difference between the neutral B mesons. Existing experimental data indirectly constrains $\sin(2\beta_{\rm CKM}) = .75 \pm .10$ [1] within the SM. Note that this CP asymmetry in the SM arises from the $B\bar{B}$ mixing [13] and not from the decay amplitudes, as can be seen easily from the Wolfenstein parameterization of the CKM matrix.

In the T2HDM, there is one additional diagram due to charged Higgs exchange that mediates the decay $B \to \psi K_S$. Integrating out the W and H^+ gives us the effective Lagrangian,

$$\mathcal{L}_{\rm eff} \simeq -2\sqrt{2}G_F V_{cb}V^*_{cs}\left[\bar{c}_L\gamma_\mu b_L \bar{s}_L\gamma^\mu c_L + 2\zeta e^{i\delta}\bar{c}_R b_L \bar{s}_L c_R\right] + \text{h.c.}, \tag{10}$$

where $\zeta e^{i\delta} \equiv (1/2)(V_{tb}/V_{cb})(m_c\tan\beta/m_H)^2\xi^*$ with ζ taken to be real and positive. Terms which are subdominant in the large $\tan\beta$ limit are not included in the above equation. Note that the charged-Higgs-mediated decay amplitude is complex and its interference with the SM amplitude could significantly modify the CP asymmetry in $B \to J/\psi K_S$ [14]. On the other hand, the charged Higgs has a vanishingly small effect on the $B\bar{B}$ mixing amplitude.

To proceed, we assume factorization and use Fierz rearrangement to evaluate the hadronic matrix elements of the two four-Fermi operators. The total amplitudes are then obtained,

$$\mathcal{A} \equiv \mathcal{A}(B \to \psi K_S) \simeq \mathcal{A}_{\rm SM}\left[1 - \zeta e^{-i\delta}\right] , \qquad \overline{\mathcal{A}} \equiv \overline{\mathcal{A}}(\bar{B} \to \psi K_S) \simeq \overline{\mathcal{A}}_{\rm SM}\left[1 - \zeta e^{i\delta}\right] \tag{11}$$

where the SM amplitudes satisfy $\mathcal{A}_{\rm SM} = \overline{\mathcal{A}}_{\rm SM}$. Therefore the ratio $\overline{\mathcal{A}}/\mathcal{A}$ gets a phase,

$$\overline{\mathcal{A}}/\mathcal{A} = \exp(-2i\vartheta) , \qquad \tan\vartheta = \zeta\sin\delta/(1 - \zeta\cos\delta) \tag{12}$$

Instead of measuring $\sin 2\beta_{\rm CKM}$, the CP asymmetry nows measures $\sin 2(\beta_{\rm CKM} + \vartheta)$

Although the magnitude of ϑ can be at most of order 10° after taking into account the experimental constraints, there could be large deviations of the CP asymmetry, $a_{\psi K_S} \equiv \sin(2\beta_{\rm CKM} + 2\vartheta)$, from the SM prediction. This can be understood as follows. The charged Higgs contribution to ϵ_K could be comparable in size to the SM effect, and this basically sets free the CKM angles $\beta_{\rm CKM}$ and γ. When $\beta_{\rm CKM}$ and γ take negative values, the CP asymmetry in $B \to J/\psi K_S$ can be of opposite sign relative to the SM expectation. The CP asymmetry in the present model is illustrated in Fig. 2 for three representative choices of γ. As can be seen from the figure, $a_{\psi K_S}$ can take almost any value in the T2HDM.

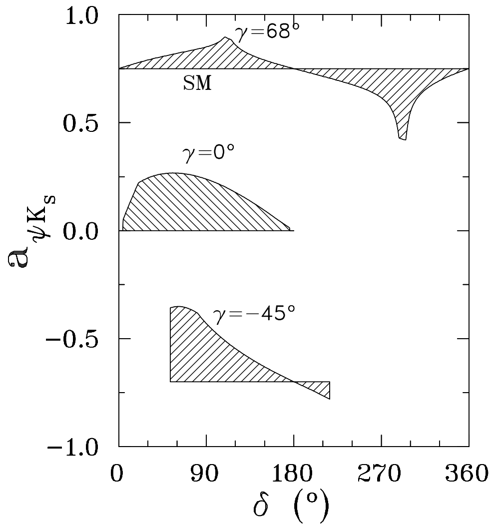

FIG. 2. Amplitude of the time-dependent CP asymmetry $a_{\psi K_S} \equiv \sin(2\beta_{\rm CKM} + 2\vartheta)$ versus the non-standard CP-odd phase δ for $B \to \psi K_S$. The top horizontal line is for the SM assuming the best fit value ($\sin 2\beta_{\rm CKM} = 0.75$) of Ref. [1]. The shaded regions correspond to the allowed ranges of the asymmetry in the T2HDM for three representative choices of γ: $\gamma = 68^\circ$ is the best fit of Ref. [1], $\gamma = 0^\circ$ corresponds to a real CKM matrix, and $\gamma = -45^\circ$ changes the sign of $a_{\psi K_S}$ relative to the SM expectation.

V. DISCUSSION

The analysis for $B \to J/\psi K_S$ can be straightforwardly extended to other B decay modes [15]. Of particular interest are those decays where the SM prediction for the CP asymmetry is zero or very small. These include $B_s \to J/\psi\phi$, $B_s \to J/\psi K_S$, and some others. In the T2HDM, it is not hard to see that the CP asymmetries for these two decays are $\sim \sin 2\vartheta$, and thus could be of order tens of percent. Indeed, a simple CP-violating pattern can be obtained for the various B decays that is quite distinctive from the SM prediction.

The complex, enhanced charm Yukawa for $H^+\bar{c}_R d_L$ provides an interesting scenario for generating a sizable neutron electric dipole moment (EDM) via a one loop diagram with virtual charm and H^+. Two vertices of the diagram involve charm quarks of opposite chirality in order for the diagram to have an imaginary piece. Using the non-relativistic quark model relation between the neutron EDM and quark EDM, one has,

$$d_n \simeq \frac{4}{3} d_d = \frac{\sqrt{2} G_F m_d}{9\pi^2} \frac{m_c^2 \tan^2\beta}{m_H^2} A|\xi| \sqrt{(1-\rho)^2+\eta^2}\lambda^4 \sin(\beta_{\rm CKM}-\delta) \left(\ln \frac{m_H^2}{m_c^2} - \frac{3}{4} \right) . \qquad (13)$$

Taking m_d as the current quark mass and imposing the experimental constraints on the model, we find that the neutron EDM can be as large as $10^{-27} - 10^{-26}$ $e \cdot {\rm cm}$, not too far below the present experimental limit. It is worth pointing out that this estimate is valid even if the new phase δ vanishes, as the $c_R - t_R$ mixing always gives the $H^+ \bar{c}_R d_L$ coupling a complex piece that is proportional to V_{td}.

The top-quark 2HDM provides one example for new CP-violating phases and non-standard charged Higgs Yukawa couplings which should be generally present in multi-Higgs doublet models when natural flavor conservation is not enforced. Although one need not take this particular model too seriously, the interesting phenomenological implications of flavor violation in general deserve further investigation.

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