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CP asymmetry in $B \to \phi K_S$ in a general two-Higgs-doublet model with fourth-generation quarks

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Abstract. We discuss the time-dependent CP asymmetry of the decay $B \to \phi K_S$ in an extension of the standard model with both a two Higgs doublet and additional fourth-generation quarks. We show that, although the standard model with a two Higgs doublet and the standard model with fourth-generation quarks *alone* are not likely to largely change the effective $\sin 2\beta$ from the decay $B \to \phi K_S$, the model with *both* an additional Higgs doublet and fourth-generation quarks can easily account for the possible large negative value of $\sin 2\beta$ without conflicting with other experimental constraints. In this model, additional large CP violating effects may arise from the flavor-changing Yukawa interactions between neutral Higgs bosons and the heavy fourth-generation down type quark, which can modify the QCD penguin contributions. With the constraints obtained from $b \to s\bar{s}s$ processes such as $B \to X_s\gamma$ and $\Delta m_{B_s^0}$, this model can lead to an effective $\sin 2\beta$ as large as -0.4 in the CP asymmetry of $B \to \phi K_S$.

1 Introduction

With the successful running of two B factories in KEK and SLAC, precise measurements of the time-dependent CP asymmetries as well as the direct CP asymmetries in rare B decays become available. Among those interesting decay modes, the most important one, the CP asymmetry of $B \to J/\psi K_S$, has been successfully measured, and a very good agreement with the standard model (SM) prediction on $\sin 2\beta$ was found.

However, the recent Belle results on $\sin 2\beta$ from $B \to \phi K_S$, although with significant errors, have indicated that the value of $\sin 2\beta$ from different decay modes could be significantly different. The most recent measurements give [1, 2]

$$\sin 2\beta = 0.47 \pm 0.34^{+0.08}_{-0.06} \text{ (Babar)},$$
$$\sin 2\beta = -0.96 \pm 0.5^{+0.09}_{-0.11} \text{ (Belle)}. \quad (1)$$

Of course, it is too early to draw any robust conclusion from the current preliminary data. Nevertheless, it opens a possibility that large new physics effects may show up in the $b \to sss$ processes, which has already triggered a large amount of theoretical efforts in examining the possible new physics contributions from various models. Besides the models related to supersymmetry which are the most promising ones, there are also a large class of models based on simple extensions of the matter contents of the SM, such as the standard models with a two Higgs doublet (S2HDM) [3–18] and the standard model with fourth-generation fermions (SM4) [19–23] etc. However, the most recent studies have pointed out that the contributions from the above mentioned two types of models to $B \to \phi K_S$ are in general not large enough to account for a large negative value of $\sin 2\beta$ in $B \to \phi K_S$ (for example $\sin 2\beta \approx -0.5$) [21, 24–26].

In this paper, we show that although due to the constraints from other experiments such as $b \to s\gamma$ and Δm_B etc., the general S2HDM and the SM4 alone are not likely to largely change the effective $\sin 2\beta$ in $B \to \phi K_S$, a model with *both* an additional Higgs doublet and 4th-generation quarks (denoted by S2HDM4) can significantly change the value of $\sin 2\beta$ without contradicting other experimental constraints. In this model, new large CP violating contributions may arise from the flavor-changing Yukawa interactions between the neutral Higgs boson and the fourth-generation down type quark b' (with $m_{b'} \gg m_b$), which changes the Wilson coefficients for QCD penguin operators and results in a large modification of effective $\sin 2\beta$. This mechanism is different from the case in the S2HDM in which the dominant contribution comes from changing the Wilson coefficients of the electro(chromo)-magnetic operators. The latter is subjected to a rather strong constraint from $b \to s\gamma$ and therefore cannot give enough contributions.

Let us begin with some model independent discussions. The definition of the effective $\sin 2\beta$ in $B \to \phi K_S$ is

$$\sin 2\beta_{\text{eff}} = \text{Im}\left[e^{2i\beta}\frac{\bar{\mathcal{A}}}{\mathcal{A}}\right]$$

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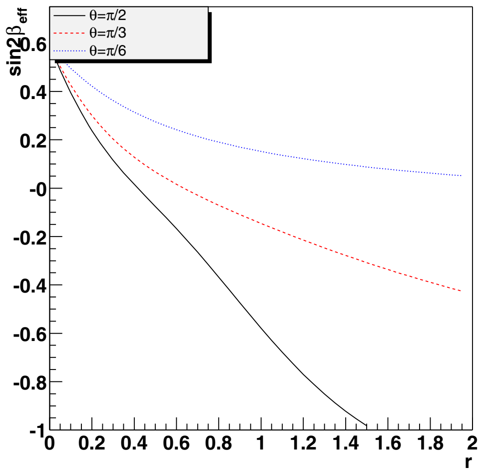

Fig. 1. The value of $\sin 2\beta_{\text{eff}}$ as a function of r. The solid, dashed and dotted curves correspond to $\theta = \pi/2, \pi/3$ and $\pi/6$ respectively

$$= \text{Im}\left[e^{2i\beta} \frac{\bar{\mathcal{A}}_{\text{SM}}(1 + re^{-i\theta})}{\mathcal{A}_{\text{SM}}(1 + re^{+i\theta})} \right], \tag{2}$$

where β is the SM value with $\sin 2\beta = 0.715^{+0.055}_{-0.045}$ [27]. $\bar{\mathcal{A}}_{\text{SM}}(\mathcal{A}_{\text{SM}})$ is the SM value of the decay amplitude of $\overline{B}^0(B^0) \to \phi K_{\text{S}}$. Here the two parameters r and θ parameterize the relative size and the additional CP violating phase of the new physics contributions. To get an idea of how $\sin 2\beta_{\text{eff}}$ is changed with the new physics contribution, we take some typical values of the phase θ, calculate the values of $\sin 2\beta_{\text{eff}}$, and show them in Fig. 1.

As shown in the figure, to explain the possibly large negative $\sin 2\beta_{\text{eff}}$, for instance, close to -0.5, in the case that θ is maximum ($\pi/2$), the value of r should be close to unity. For smaller θ such as $\pi/3$ and $\pi/6$, the value of r must be even larger. Therefore, to generate a large negative value of $\sin 2\beta_{\text{eff}}$ in the range of $-0.5 \sim -1.0$, the magnitude of the new physics contributions must be of the same order of magnitude as the one in the SM.

However, the new physics contributions must be constrained by other experiments, especially by the $b \to s$ transition related processes. The strictest constraint comes from the radiative decay of $B \to X_s\gamma$. The current data of $\text{Br}(B \to X_s\gamma; E_\gamma > 1.6\text{GeV}) = (3.28^{+0.41}_{-0.36}) \times 10^{-4}$ [28–31] are well reproduced in the framework of the next-to-leading order calculations in the SM (see e.g. [32,33]). Thus, if the new physics contribution carries no new phase, there is very little room for the new physics parameters. But in the case that new phases are present, the parameter space could be enlarged. This is because the data of $B \to X_s\gamma$ only constrains the absolute values of the Wilson coefficient $C_{7\gamma}$; if the new physics contribution does not change the absolute value of $C_{7\gamma}$, there will not be a serious problem. Thus the following relation must be satisfied for any new physics model:

$$|C_{7\gamma}| = \left|C_{7\gamma}^{\text{SM}} + C_{7\gamma}^{\text{NEW}}\right| \simeq \left|C_{7\gamma}^{\text{SM}}\right|, \tag{3}$$

with $C_{7\gamma}^{\text{SM}}$ and $C_{7\gamma}^{\text{NEW}}$ being the effective Wilson coefficient evaluated at the low energy scale ($\mu \approx m_b$) from SM and new physics models respectively. In this case, the absolute value of $C_{7\gamma}^{\text{NEW}}$ could vary largely from close to zero to about $-2C_{7\gamma}^{\text{SM}}$, which seems large enough for explaining the CP asymmetry in $B \to \phi K_{\text{S}}$. However, it follows from (3) that the data on $B \to X_s\gamma$ do strongly constrain the form of $C_{7\gamma}^{\text{NEW}}$; namely, the new physics must interfere in such a way that the total effect is roughly equivalent to adding a phase factor to $C_{7\gamma}^{\text{SM}}$, i.e $C_{7\gamma} \simeq |C_{7\gamma}^{\text{SM}}|e^{i\theta}$. Let us take an illustrative example in which the new physics contribution is purely electro(chromo)-magnetic and satisfy $C_{7\gamma} = |C_{7\gamma}^{\text{SM}}|e^{i\theta}$ and also $C_{8g} = |C_{8g}^{\text{SM}}|e^{i\theta}$ at the scale of m_W. Varying the value of θ from 0 to 2π and then running down to the low energy scale of $\mu \simeq m_b$ through the renormalization group equation, one finds that the value of $\sin 2\beta_{\text{eff}}$ in the decay $B \to \phi K_{\text{S}}$ only changes from 0.5 to 0.8. This naive discussion shows that if the dominant contribution from a new physics model is coming from $C_{7\gamma(8g)}^{\text{NEW}}$, the change to $\sin 2\beta_{\text{eff}}$ from the its SM value is limited. Unfortunately, the S2HDM belongs to this class of models. A recent analysis has confirmed that, within S2HDM, the value of $\sin 2\beta_{\text{eff}}$ can reach zero but is not likely to be largely negative [24–26].

For the model of SM4, there are additional up (t') and down (b') type quarks. The new phases may come from the extended Cabbibo–Kobayashi–Maskawa (CKM) matrix which is a four by four matrix in this model and contains undetermined matrix elements $V_{t'q}, V_{qb'}$ etc. To avoid the precise data of electro-weak (EW) processes, the mass of b' (t') has to be pushed to greater than ~200 GeV (~ 300 GeV). However, a phenomenological study showed that with the constraint of $B \to X_s\gamma$ and $B_s^0 - \overline{B^0}_s$ mixings being considered, its contribution to the CP violation of $B \to \phi K_{\text{S}}$ is not large enough either [21]. Thus, if the large negative value of $\sin 2\beta_{\text{eff}}$ in the decay $B \to \phi K_{\text{S}}$ is confirmed by future experiments, the above mentioned two models (i.e. S2HDM and SM4) will not be favored.

2 The model of S2HDM4

There are several directions in constructing models beyond the SM, such as enlarging the gauge groups to $SU(5)$, $SU(10)$ and E_6 etc., introducing new symmetries like various SUSY models, and expanding the matter contents, i.e., taking more fermions and Higgs bosons. The models of the last type can be regarded as simple extensions of the SM which keep the same gauge structure but still have rich sources of new contributions. The typical ones are the above mentioned S2HDM and SM4.

In this paper we would like to take a step further to consider a model with both a two Higgs doublet and fourth-generation quarks (S2HDM4). In this model, there are new Yukawa interactions between Higgs bosons and heavy fourth-generations quarks. Since in general the Yukawa interaction is expected to be proportional to the coupled quark mass, the new Yukawa couplings are much stronger than that in the S2HDM and SM4. Unlike in the case of

S2HDM, where the b quark contribution to the QCD penguin diagram through neutral Higgs boson loop is strongly suppressed by the small b quark mass, the same diagram with an intermediate b' quark may significantly contribute to the related processes [34]. This new feature only exists in this combined model and is of particular interest in studying the CP violation of $B\to\phi K_{\rm S}$ and other penguin dominant processes.

The Lagrangian for the S2HDM4 is given by

$$\mathcal{L}_Y=\bar\psi_{\rm L}Y_1^U\widetilde{\phi_1}u_{\rm R}+\bar\psi_{\rm L}Y_1^D\phi_1 d_{\rm R}+\bar\psi_{\rm L}Y_2^U\widetilde{\phi_2}u_{\rm R}$$
$$+\bar\psi_{\rm L}Y_2^D\phi_2 d_{\rm R}+{\rm H.c.} \qquad (4)$$

with the extended quark content of $u_{\rm L,R}=(u,c,t,t')_{\rm L,R}$ and $d_{\rm L,R}=(d,s,b,b')_{\rm L,R}$. The Yukawa coupling matrices $Y_i^{U(D)}$ are 4-dimensional matrices accordingly. The two Higgs fields ϕ_1,ϕ_2 have vacuum expectation values (VEVs) of $v_1{\rm e}^{{\rm i}\delta_1}$ and $v_2{\rm e}^{{\rm i}\delta_2}$ respectively, with $\sqrt{|v_1|^2+|v_1|^2}=v=246\,{\rm GeV}$. The relative phase $\delta=\delta_1-\delta_2$ between two VEVs is physical and provides a new source of CP violation [3–5]. In the mass eigenstates, the three physical Higgs bosons are denoted by H^0, A^0 and $H^\pm$ respectively. Due to the non-zero phase δ, all the Yukawa couplings become complex numbers in the physical mass basis, even if they are all real in the flavor basis. For simplicity, throughout this paper, we assume that the CKM matrix elements associated with t', i.e. $V_{t'q}$, are small enough to be ignored, and we will only focus on the neutral Higgs boson contributions.

In the mass basis, the Yukawa interactions between neutral Higgs bosons and quarks have the following general form:

$$\mathcal{L}_Y=\eta^q_{ij}\bar q_{i\rm L}q_{j\rm R}\phi+{\rm H.c.}, \qquad (5)$$

with $\phi=H^0$ or A^0. The Yukawa coupling $\eta^q_{\rm ij}$ is usually parameterized as

$$\eta^q_{ij}=\frac{\sqrt{m_{q_i}m_{q_j}}}{v}\xi_{q_iq_j}. \qquad (6)$$

In the Chen–Sher ansatz [35] motivated by a Fritzsch type of Yukawa coupling matrix the values of all $\xi_{q_iq_j}$s are of the same order of magnitude. However, from other contexts of the coupling matrix the relations among the $\xi_{q_iq_j}$ are different [36–38]. In the general case, they should be taken as free parameters to be determined or constrained by the experiments.

The effective Hamiltonian for $\Delta B=1$ charmless B decays reads

$$H_{\rm eff}=\frac{G_{\rm F}}{\sqrt2}\Bigg[V_{\rm ub}V_{\rm us}^*(C_1^uQ_1^u+C_2^uQ_2^u)+V_{\rm cb}V_{\rm cs}^*(C_1^cQ_1^c+C_2^cQ_2^c)$$
$$-V_{\rm tb}V_{\rm ts}^*\left(\sum_{i=3}^{10}C_iQ_i+C_{7\gamma}Q_{7\gamma}+C_{8g}Q_{8g}\right)\Bigg], \qquad (7)$$

where the operator basis Q_i can be found in [39]. In this model, the relevant Wilson coefficients at the scale of m_W from this model is given by

$$C_1(M_W)=\frac{11}{2}\frac{\alpha_{\rm s}(M_W)}{4\pi},$$
$$C_2(M_W)=1-\frac{11}{6}\frac{\alpha_{\rm s}(M_W)}{4\pi}-\frac{35}{18}\frac{\alpha_{\rm EM}}{4\pi},$$
$$C_3(M_W)=-\frac{\alpha_{\rm s}(M_W)}{24\pi}\times\left(\tilde E_0(x_t)+|\xi_{tt}|^2E_0^{\rm III}(y)+\frac{m_{b'}\sqrt{m_bm_s}}{2V_{tb}V_{ts}^*m_t^2}\xi^*_{bb'}\xi_{sb'}E_0^{\rm III}(y')\right),$$
$$+\frac{\alpha_{\rm EM}}{6\pi}\frac{1}{\sin^2\theta_{\rm W}}(2B_0(x_t)+C_0(x_t)),$$
$$C_4(M_W)=\frac{\alpha_{\rm s}(M_W)}{8\pi}\times\left(\tilde E_0(x_t)+|\xi_{tt}|^2E_0^{\rm III}(y)+\frac{m_{b'}\sqrt{m_bm_s}}{2V_{tb}V_{ts}^*m_t^2}\xi^*_{bb'}\xi_{sb'}E_0^{\rm III}(y')\right),$$
$$C_5(M_W)=-\frac{\alpha_{\rm s}(M_W)}{24\pi}\times\left(\tilde E_0(x_t)+|\xi_{tt}|^2E_0^{\rm III}(y)+\frac{m_{b'}\sqrt{m_bm_s}}{2V_{tb}V_{ts}^*m_t^2}\xi^*_{bb'}\xi_{sb'}E_0^{\rm III}(y')\right),$$
$$C_6(M_W)=\frac{\alpha_{\rm s}(M_W)}{8\pi}\times\left(\tilde E_0(x_t)+|\xi_{tt}|^2E_0^{\rm III}(y)+\frac{m_{b'}\sqrt{m_bm_s}}{2V_{tb}V_{ts}^*m_t^2}\xi^*_{bb'}\xi_{sb'}E_0^{\rm III}(y')\right),$$
$$C_{7\gamma}(M_W)=\frac{A(x_t)}{2}-\tfrac12\left(A(y)|\xi_t|^2+A(y')\frac{m_{b'}\sqrt{m_bm_s}}{2V_{tb}V_{ts}^*m_t^2}\xi^*_{bb'}\xi_{sb'}\right),$$
$$+B(y)|\xi_t\xi_b|{\rm e}^{{\rm i}\theta}-B(y')\frac{m_{b'}\sqrt{m_bm_s}}{2V_{tb}V_{ts}^*m_tm_b}\xi_{b'b}\xi_{sb'},$$
$$C_{8g}(M_W)=-\frac{D(x_t)}{2}-\frac12\left(D(y)|\xi_t|^2+D(y')\frac{m_{b'}\sqrt{m_bm_s}}{2V_{tb}V_{ts}^*m_t^2}\xi^*_{bb'}\xi_{sb'}\right),$$
$$+(y)|\xi_t\xi_b|{\rm e}^{{\rm i}\theta}-E(y')\frac{m_{b'}\sqrt{m_bm_s}}{2V_{tb}V_{ts}^*m_tm_b}\xi_{b'b}\xi^*_{sb'}, \qquad (8)$$

with $\alpha_{\rm s}(m_W)$ and $\alpha_{\rm EM}$ being the strong and electro-magnetic couplings at the scale m_W. The mass ratios x_t, y

and y' are defined by $x_t = m_t^2/m_W^2$, $y = m_t^2/m_{H^\pm}^2$ and $y' = m_{b'}^2/m_{H^0}^2$ respectively. The loop integration functions are standard and can be found in [6, 40–42]. Here we have ignored the coefficients for the electro-weak penguin diagrams since their effects are less significant in the decay of $B \to \phi K_\mathrm{S}$.

Note that the new contributions to QCD and electro-(chromo)-magnetic operators depend on different parameter sets. In the QCD penguin sector, the contribution depends on $\xi_{bb'}^* \xi_{sb'}$ where in the electro(chromo)-magnetic sector it depends on both $\xi_{b'b}\xi_{sb'}$ and $\xi_{bb'}^*\xi_{sb'}$. It is convenient to define two weak phases θ_1 and θ_2 with

$$\xi_{bb'}^* \xi_{sb'} = |\xi_{bb'}\xi_{sb'}| \mathrm{e}^{\mathrm{i}\theta_1} \text{ and } \xi_{b'b}\xi_{sb'} = |\xi_{b'b}\xi_{sb'}| \mathrm{e}^{\mathrm{i}\theta_2}. \quad (9)$$

Since in general $\xi_{b'b}$ and $\xi_{bb'}^*$ are complex numbers and $\xi_{b'b} \neq \xi_{bb'}^*$, the two phases are not necessarily equivalent. The presence of two independent phases rather than one is particular for this model, which gives different contributions to the QCD penguin and electro(chromo)-magnetic Wilson coefficients. The interference between them enlarges the allowed parameter space.

Note that the Wilson coefficient for QCD penguins may be complex numbers which provides additional sources of CP violation. To make a comparison, let us denote the Wilson coefficients in the SM by C_i^{SM}. Taking $\xi_{b'b} = \xi_{sb'} = 0.8$, $\theta_1 = 0.5$, $\theta_2 = -1.2$ and $m_{H^0} = m_{b'} = 200\,\mathrm{GeV}$ as an example, in the range of $40 < \xi_{bb'} < 60$, the ratio of $C_3/C_3^{\mathrm{SM}}(C_4/C_4^{\mathrm{SM}})$ has an imaginary part between -0.27 and -0.4 (-0.6 and -0.8). These large imaginary parts play an important role in CP violation.

3 Constraints from $B \to X_s\gamma$ and B_s^0-$\bar{B}_s^0$ mixing

Before making any predictions, one first needs to know how the new parameters in this model are constrained by other experiments. For the process we are concerned with, the strictest constraints come from $b \to s\bar{s}s$ processes such as $B \to X_s\gamma$ and B_s^0–$\bar{B}_s^0$ mixing, etc.

The expression for $B \to X_s\gamma$ normalized to $B \to X_c e\bar{\nu}_e$ reads

$$\frac{\mathrm{Br}(B \to X_s\gamma)}{\mathrm{Br}(B \to X_c e\bar{\nu}_e)} = \frac{6|V_{\mathrm{tb}}V_{\mathrm{ts}}^*|^2\alpha_{\mathrm{EM}}}{\pi|V_{\mathrm{cb}}|^2 f(m_c/m_b)}|C_{7\gamma}(\mu)|^2, \quad (10)$$

with $f(z) = 1 - 8z^2 - 24z^4 \ln z + 8z^6 - z^8$ and $\mathrm{Br}(B \to X_c e\bar{\nu}_e) = 10.45\%$. The low energy scale μ is set to m_b. Using the Wilson coefficients at the scale m_W and running down to the m_b scale through renormalization group equations, we obtain the predictions for $\mathrm{Br}(B \to X_s\gamma)$. For simplicity, we focus on the case in which the b' contribution dominates through the H^0 loop; namely, we push the masses of the charged Higgs bosons $H^\pm$ and the other pseudo-scalar boson A^0 to very high values, ($m_{H^\pm}, m_{A^0} > 500\,\mathrm{GeV}$), and ignore their contributions. We take the following typical values of the couplings:

$$|\xi_{bb'}| = 50, \quad |\xi_{b'b}| = 0.8, \quad |\xi_{sb'}| = 0.8,$$

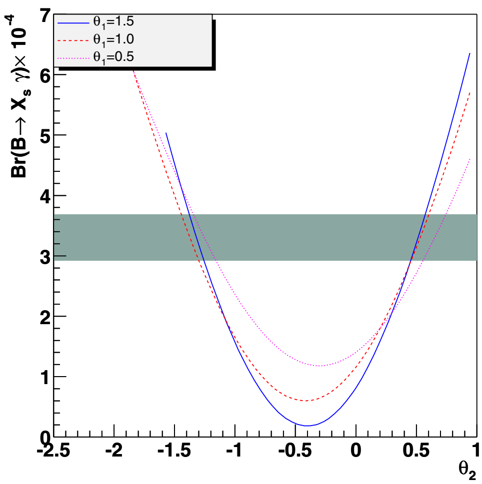

Fig. 2. The branching ratio of $B \to X_s\gamma$ as a functions of θ_2 in the model of S2HDM4 . The solid, dashed and dotted curves correspond to $\theta_1 = 1.5, 1.0$ and 0.5 respectively. The other parameters are taken from (11)

and

$$m_{H^0} = m_b' = 200\,\mathrm{GeV}, \quad (11)$$

and give in Fig. 2 the value of $\mathrm{Br}(B \to X_s\gamma)$ as a function of θ_1 with different values of θ_2.

From the figure, one finds that two separated ranges for the parameters θ_1 and θ_2 are allowed by the data:

$$-1.4 \lesssim \theta_2 \lesssim -1.2 \text{ and } 0.4 \lesssim \theta_2 \lesssim 0.7$$

for

$$0.5 \lesssim \theta_1 \lesssim 1.5. \quad (12)$$

Note that we do not make a scan for the full parameter space; nevertheless, the above obtained range is already enough for our purpose. Among the two allowed ranges, the one with $-1.4 \lesssim \theta_2 \lesssim -1.2$ is of particular interest. It will be seen below that in this range, the contribution to the CP asymmetry in $B \to \phi K_\mathrm{S}$ could be significant. In Fig. 3, we also give the allowed range of θ_1 with different values of θ_2. One finds that the allowed range for θ_1 is larger than for θ_2. In this figure, the interference between the two phases θ_1 and θ_2 is manifest. For θ_2 in the range of $(-1.0, -0.8)$, the allowed value for θ_1 is a narrow window around zero. But for θ_2 in the range of $(-1.4, -1.2)$, the allowed range for θ_1 could be between 0.5 and 2.0. Compared with the S2HDM in which only one phase appears, this interference effect for two phases enlarges the parameter space under the constraint of $B \to X_s\gamma$. Thus large contributions to the other processes are possible in this model.

The other $b \to s\bar{s}s$ process which could impose a strong constraint is the mass difference of the neutral B_s^0 meson. The measurements from LEP give a lower bound of $\Delta m_{B_s} > 14.9\,\mathrm{ps}^{-1}$. In this model, the b' contributes to

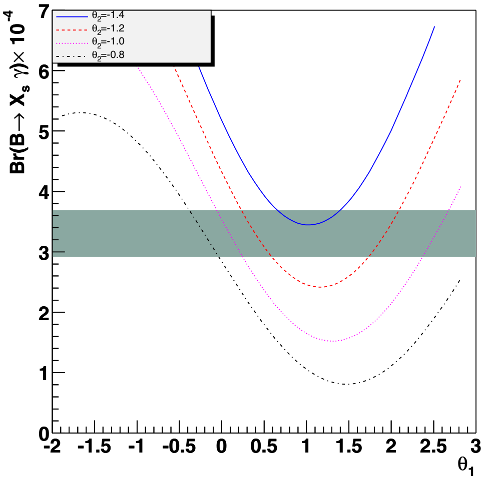

Fig. 3. The branching ratio of $B \to X_s\gamma$ as a function of θ_1 in the model of S2HDM4. The solid, dashed, dotted and dot-dashed curves correspond to $\theta_2 = 1.4, 1.2, 1.0$ and 0.8 respectively. The other parameters are taken from (11)

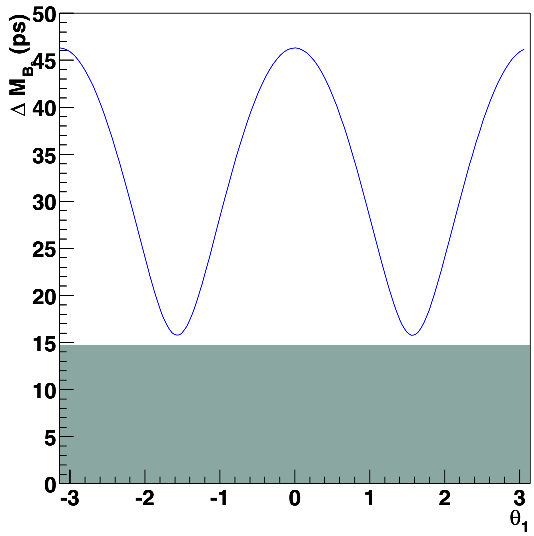

Fig. 4. The B_s^0 meson mass difference ΔB_s^0 as a function of θ_1 in the model of S2HDM4. Other parameters are taken from (11). The shadowed region is excluded by the data of ΔB_s^0

Δm_{B_s} only through box diagrams. The box diagram contribution to Δm_{B_s} is given by [40–43]

$$\begin{aligned}\Delta m_{B_s} &= \frac{G_{\rm F}^2}{6\pi^2}(f_{B_s}\sqrt{B_{B_s}})^2 m_{B_s} m_t^2 |V_{\rm ts}|^2 \\ &\times\Big\{\eta_{\rm tt}B^{\rm WW}(x_t) + \frac{1}{4}\eta_{\rm tt}^{\rm HH} y_t |\xi_{tt}|^4 B_V^{\rm HH}(y_t) \\ &+2\eta_{\rm tt}^{\rm HW} y_t |\xi_{\rm tt}|^2 B_V^{\rm HW}(y_t, y_w) \qquad (13)\\ &+\frac{1}{4}\eta_{\rm tt}^{\rm HH} y' \left(\frac{m_{b'}\sqrt{m_b m_s}}{2V_{tb}V_{ts}^* m_t^2}\xi_{bb'}^*\xi_{sb'}\right)^2 B_V^{\rm HH}(y')\Big\},\end{aligned}$$

where $G_{\rm F} = 1.16\times10^{-5}\,{\rm GeV}^{-2}$ is the Fermi constant. f_{B_s} and B_{B_s} are the decay constant and bag parameter for B_s^0. In the numerical calculations, we take the value of $f_{B_s}\sqrt{B_{B_s}} = 0.23\,{\rm GeV}$. The η_{ij} are the QCD correction factors. The loop integration functions of $B_{(V)}^{\rm HH,WW,HW}$ can be found in [41–43]. The mass ratios are defined as $y_t = m_t^2/m_{H^\pm}^2, y_w = m_t^2/m_W^2$ and $y' = m_{b'}^2/m_{H^0}^2$ respectively. Note that in the mass difference of the B_s^0 mesons, the contribution from S2HDM4 only depends on the parameter $\xi_{bb'}^*\xi_{sb'}$. So, only the phase θ_1 will be present in the expression.

Using the above obtained typical parameters in (11), the contribution to Δm_{B_s} is calculated and plotted as a function of θ_1 in Fig. 4. The figure shows that the current data of Δm_{B_s} do not impose a strong constraint on the value of θ_1.

The neutron electric dipole moment (EDM) is expected to give strong constraints on the new physics. In the SM, the neutron EDM is zero even at two-loop level. The current experimental upper limit gives EDM $< 1.1\times10^{-25}\,e$cm [44]. In general, the new physics contributes to the neutron EDM through one loop diagrams. In the presence of new scalars, additional significant contributions may arise, for example from the Weinberg gluonic operator [45] and also the two-loop Barr–Zee type diagrams [46, 47], etc.

However, we note that all the above three types of mechanisms are not related to $b \to s$ flavor-changing transitions and therefore will involve different parameters in this model. For the one-loop diagrams, the neutral EDM is mostly related to $\xi_{u(d)}$ and $\xi_{t(b')}$ through $u(d)$-quark EDM. For the Weinberg three gluonic operator, the dominant contribution is from an internal b' loop. Thus it is related to $\xi_{b'b'}$. Similarly, for the two-loop Barr–Zee diagram, the b'-quark loop will play the most important role and the couplings involve only $\xi_{u(d)}$, $\xi_{b'b'}$ etc.

Thus the neutron EDM will impose strong constraints on other parameters in this model and has less significance in the current study of the decay $B \to \phi K_{\rm S}$. This is significantly different from the S2HDM case in which the t-quark always dominates the loop contribution and the couplings ξ_{tt} and ξ_{bb} are subject to a strong constraint from neutron EDM.

Other constraints may come from K^0–$\overline{K^0}$ and B_d^0–$\overline{B_d^0}$ mixings. But those processes contain additional free parameters such as the Yukawa coupling of $\xi_{b'd}$ and $\xi_{sb'}$; the constraints from those processes are much weaker.

4 CP asymmetry in $B \to \phi K_{\rm S}$

Now we are in the position to discuss CP asymmetry in $B \to \phi K_{\rm S}$. The decay amplitude for $\bar{B} \to \phi\bar{K}^0$ reads

$$\begin{aligned}&\mathcal{A}(\bar{B}_d^0 \to \phi\bar{K}^0) \qquad (14)\\ &= -\frac{G_{\rm F}}{\sqrt{2}}V_{\rm ts}^*V_{\rm tb}\left(a_3 + a_4 + a_5 - \frac{1}{2}(a_7 + a_9 + a_{10})\right)X,\end{aligned}$$

with X being a factor related to the hadronic matrix elements. In the naive factorization approach $X = 2f_\phi m_\phi(\epsilon\cdot$

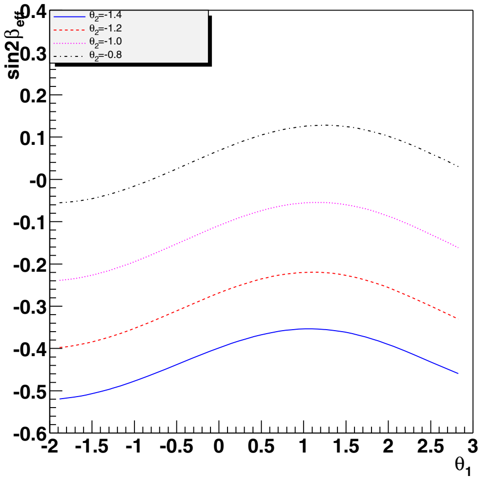

Fig. 5. The prediction for $\sin 2\beta_{\rm eff}$ as a function of θ_1 with different values of θ_2. The solid, dashed, dotted and dot-dashed curves corresponds to $\theta_2 = -1.4, -1.2, -1.0, -0.8$ respectively

$p_B)F_1(m_\phi)$, where ϵ, p_B, F_1 are the polarization vector of ϕ, the momentum of B meson and form factor, respectively. The coefficients a_i are defined through the effective Wilson coefficients $C_i^{\rm eff}s$ as follows:

$$a_{2i-1} = C_{2i-1}^{\rm eff} + \frac{1}{N_c} C_{2i}^{\rm eff}, \quad a_{2i} = C_{2i}^{\rm eff} + \frac{1}{N_c} C_{2i-1}^{\rm eff}, \quad (15)$$

Since the heavy particles such as $H^{\pm,0}$, A^0 and b' have been integrated out below the scale of m_W, the procedures to obtain the effective Wilson coefficients $C_i^{\rm eff}$ are exactly the same as in SM and can be found in [48] .

Using the above obtained parameters allowed by the current data, the predictions for the time-dependent CP asymmetry for $B \to \phi K_{\rm S}$ are shown in Fig. 5

In the figure, we give the value of $\sin 2\beta_{\rm eff}$ as a function of θ_1 with different values of $\theta_2 = 1.4, 1.2, 1.0$ and 0.8. Comparing with the constraints obtained from $B \to X_s\gamma$ and B_s^0–$\bar{B}_s^0$ mixings, one sees that in the allowed range of $-1.4 < \theta_2 < -1.2$ and $0.5 < \theta_1 < 1.5$, the predicted $\sin 2\beta_{\rm eff}$ can reach -0.4.

It is evident that the large negative value of $\sin 2\beta_{\rm eff}$ is a consequence of the interference effects between θ_1 and θ_2 and therefore is particular for this model. For a zero value of θ_1, there is no new phase in the QCD penguin sector. From Fig. 3, the allowed range for θ_2 is $-1.0 \le \theta_2 \le -0.8$. Then it follows from Fig. 5, that in this range the predicted $\sin 2\beta_{\rm eff}$ is at around zero. But for $\theta_1 \approx 0.5$, the allowed range for θ_2 is changed into $-1.4 \le \theta_2 \le -1.2$ and the predictions for $\sin 2\beta_{\rm eff}$ are much lower in the range of $(-0.4, -0.25)$.

5 Conclusions

In conclusion, we have discussed the CP asymmetry of decay $B \to \phi K_{\rm S}$, in the model of S2HDM4 which contains both an additional Higgs doublet and fourth-generation quarks. In this model, since the fourth-generation b' quark is much heavier than the b quark, the Yukawa interactions between a neutral Higgs boson and b' are greatly enhanced. This results in a significant modification of the QCD penguin diagrams. We have obtained the allowed range of the parameters from the process of $B \to X_s\gamma$ and Δm_{B_s}. Due to the more complicated phase effects, in this model the constraints from those process are weaker than that in S2HDM and SM4. The effective $\sin 2\beta_{\rm eff}$ in the decay $B \to \phi K_{\rm S}$ is predicted with the constrained parameters. We have found that this model can easily account for the possible large negative value of $\sin 2\beta$ without conflicting with other experimental constraints.

In this paper we focus on the case in which H^0 dominates. It is straightforward to find that the contribution from the other pseudo-scalar A^0 follows the same pattern. In the case of small mixing among the neutral scalars, the Yukawa couplings for H^0 and A^0 are directly related [12]. We find that for $m_{A^0} \approx 200\,{\rm GeV} \ll m_{H^0}$ its contribution to the decay amplitude of $B \to \phi K_{\rm S}$ is similar to the case of the H^0 dominance discussed above. For the case that m_{A^0} is close to m_{H^0}, the contribution from them are comparable, and the interference between the two could be important.

Since this model contributes new phases to QCD penguin diagrams, it remains to be seen if it has sizable effects on other penguin dominant processes, such as in the hadronic charmless B decays. Similarly, it is expected that in this model there are also significant contributions to the electro-weak penguin diagrams which deserve a further investigation (for recent discussions on EW penguin effects on $B \to \phi K$ see, e.g., [49–51].) It is well known that the EW penguin plays an important role in rare B decays. The current data on $B \to \pi\pi, \pi K$ have indicated some deviations from results based on the SM [52–57]. It is of interest to further investigate the new physics contributions to those decay modes within this model.

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References

1. K. Abe et al., Belle collaboration, Phys. Rev. Lett. **91**, 261602 (2003) [hep-ex/0308035]
2. B. Aubert et al., BABAR collaboration, hep-ex/0403026
3. Y.L. Wu, L. Wolfenstein, Phys. Rev. Lett. **73**, 1762 (1994) [hep-ph/9409421]
4. L. Wolfenstein, Y.L. Wu, Phys. Rev. Lett. **73**, 2809 (1994) [hep-ph/9410253]
5. Y.-L. Wu, hep-ph/9404241
6. Z.-J. Xiao, K.-T. Chao, C.S. Li, Phys. Rev. D **65**, 114021 (2002) [hep-ph/0204346]
7. S.L. Glashow S. Weinberg, Phys. Rev. D **15**, 1958 (1977)
8. M.J. Savage, Phys. Lett. B **266**, 135 (1991)
9. W.-S. Hou, Phys. Lett. B **296**, 179 (1992)

10. A. Antaramian, L.J. Hall, A. Rasin, Phys. Rev. Lett. **69**, 1871 (1992) [hep-ph/9206205]
11. L.J. Hall, S. Weinberg, Phys. Rev. D **48**, 979 (1993) [hep-ph/9303241]
12. D. Atwood, L. Reina, A. Soni, Phys. Rev. D **55**, 3156 (1997) [hep-ph/9609279]
13. Y.-B. Dai, C.-S. Huang, H.-W. Huang, Phys. Lett. B **390**, 257 (1997) [hep-ph/9607389]
14. D. Bowser-Chao, K.-m. Cheung, W.-Y. Keung, Phys. Rev. D **59**, 115006 (1999) [hep-ph/9811235]
15. R. Diaz, R. Martinez, J.A. Rodriguez, hep-ph/0010149
16. Y.-F. Zhou, Y.-L. Wu, Mod. Phys. Lett. A **15**, 185 (2000) [hep-ph/0001106]
17. Y.-L. Wu, Y.-F. Zhou, Phys. Rev. D **64**, 115018 (2001) [hep-ph/0104056]
18. Y.-F. Zhou, Y.-L. Wu, Eur. Phys. J. C **27**, 577 (2003) [hep-ph/0110302]
19. C.-S. Huang, W.-J. Huo, Y.-L. Wu, Phys. Rev. D **64**, 016009 (2001) [hep-ph/0005227]
20. C.-S. Huang, W.-J. Huo, Y.-L. Wu, Mod. Phys. Lett. A **14**, 2453 (1999) [hep-ph/9911203]
21. A. Arhrib, W.-S. Hou, Eur. Phys. J. C **27**, 555 (2003) [hep-ph/0211267]
22. K.-c. Chou, Y.-l. Wu, Y.-b. Xie, Chin. Phys. Lett. **1**, 47 (1984)
23. A. Datta, E.A. Paschos, CP violation, edited by C. Jarlskog, Adv. Ser. Direct. High Energy Phys. **3**, 292 (1989)
24. G. Hiller, Phys. Rev. D **66**, 071502 (2002) [hep-ph/0207356]
25. A.K. Giri, R. Mohanta, Phys. Rev. D **68**, 014020 (2003) [hep-ph/0306041]
26. C.-S. Huang, S.-h. Zhu, Phys. Rev. D **68**, 114020 (2003) [hep-ph/0307354]
27. A.J. Buras, hep-ph/0210291
28. R. Barate et al., ALEPH collaboration, Phys. Lett. B **429**, 169 (1998)
29. S. Chen et al., CLEO collaboration, Phys. Rev. Lett. **87**, 251807 (2001) [hep-ex/0108032]
30. K. Abe et al., Belle collaboration, Phys. Lett. B **511**, 151 (2001) [hep-ex/0103042]
31. B. Aubert et al., BaBar collaboration, hep-ex/0207076
32. P. Gambino, M. Misiak, Nucl. Phys. B **611**, 338 (2001) [hep-ph/0104034]
33. A.J. Buras, A. Czarnecki, M. Misiak, J. Urban, Nucl. Phys. B **631**, 219 (2002) [hep-ph/0203135]
34. Y.-L. Wu, Chin. Phys. Lett. **16**, 339 (1999) [hep-ph/9805439]
35. T.P. Cheng, M. Sher, Phys. Rev. D **35**, 3484 (1987)
36. Y.-F. Zhou, hep-ph/0307240, to appear in J. Phys. G
37. J.L. Diaz-Cruz, R. Noriega-Papaqui, A. Rosado, hep-ph/0401194
38. R.A. Diaz, R. Martinez, C.E. Sandoval, hep-ph/0311201
39. A.J. Buras, hep-ph/9806471
40. T. Inami, C.S. Lim, Prog. Theor. Phys. **65**, 297 (1981)
41. F.J. Gilman, M.B. Wise, Phys. Lett. B **93**, 129 (1980)
42. F.J. Gilman, M.B. Wise, Phys. Rev. D **27**, 1128 (1983)
43. Y.L. Wu, Y.F. Zhou, Phys. Rev. D **61**, 096001 (2000) [hep-ph/9906313]
44. K. Hagiwara et al., Particle Data Group collaboration, Phys. Rev. D **66**, 010001 (2002)
45. S. Weinberg, Phys. Rev. Lett. **63**, 2333 (1989)
46. J.D. Bjorken, S. Weinberg, Phys. Rev. Lett. **38**, 622 (1977)
47. S.M. Barr, A. Zee, Phys. Rev. Lett. **65**, 21 (1990)
48. G. Buchalla, A.J. Buras, Lauten, Rev. Mod. Phys. **68**, 1125 (1996) [hep-ph/9512380]
49. D. Atwood, G. Hiller, hep-ph/0307251
50. D.E. Morrissey C.E.M. Wagner, Phys. Rev. D **69**, 053001 (2004) [hep-ph/0308001]
51. N.G. Deshpande, D.K. Ghosh, hep-ph/0311332
52. B. Aubert et al., BABAR collaboration, Phys. Rev. Lett. **91**, 241801 (2003) [hep-ex/0308012]
53. K. Abe et al., Belle collaboration, Phys. Rev. Lett. **91**, 261801 (2003) [hep-ex/0308040]
54. K. Abe et al., Belle collaboration, hep-ex/0401029
55. Y.F. Zhou, Y.L. Wu, J.N. Ng, C.Q. Geng, Phys. Rev. D **63**, 054011 (2001) [hep-ph/0006225]
56. Y.-L. Wu, Y.-F. Zhou, Eur. Phys. J. Direct C **5**, 014 (2003) [hep-ph/0210367]
57. A.J. Buras, R. Fleischer, S. Recksiegel, F. Schwab, hep-ph/0312259