New physics effects on the CP asymmetries in $B\to\phi K_S$ and $B\to\eta' K_S$ decays

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Within the standard model (SM), the time-dependent CP asymmetries in $B\to\psi K_S$, $B\to\eta' K_S$, and $B\to\phi K_S$ are expected to give the same result, i.e., $\sin 2\beta$. However, recent measurements of the mixing-induced CP asymmetries in $B\to\phi K_S$ and $B\to\eta' K_S$ modes give results whose central values differ from the SM expectations. We explore the effect of new physics in the two Higgs doublet model (THDM), which allows tree level flavor changing neutral currents (so-called model III), and the model with an extra vectorlike down quark (VLDQ). We find that the observed mixing-induced CP asymmetry for $B\to\phi K_S$ cannot be accommodated by the THDM but can be explained in the VLDQ model, and both models can explain the observed asymmetry for $B\to\eta' K_S$ mode.

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I. INTRODUCTION

A new era in B physics has just started with the advent of B factories. With the accumulation of huge data in the B system, the standard model (SM) will be subjected to a very stringent test. At the same time, the experiments at B factories are also potential sources for probing new physics. The BABAR [1] and Belle [2] measurements of the time-dependent asymmetries in the gold plated mode $B\to\psi K_S$ have provided the first evidence of CP violation in the B system. The observed world average of $\sin 2\beta$ [3],

$$\sin 2\beta_{\psi K_S}=0.734\pm 0.054, \tag{1}$$

agrees well with the SM prediction. This indicates that CP symmetry is significantly violated in nature and the Kobayashi-Maskawa (KM) mechanism [4] seems to be the dominant source of CP violation, in which the phase δ_{KM} is the only source of CP violation. However, this speculation does not exclude interesting CP violating new physics (NP) effects in other B decays. Since the decay $B\to\psi K_S$ ($b\to c\bar{c}s$) is a tree level process in the SM, the NP contributions to its amplitude are naturally suppressed. Moreover, at the loop level NP may give large contributions to the B^0-$\overline{B^0}$ mixing as well as to the loop-induced decay amplitudes. The former effects are universal to all B^0 decay modes and are constrained to be less than 20% compared to that of the SM contribution [3]. On the other hand, the effects of new physics in the decay amplitudes are nonuniversal, may vary from process to process, and can show up in the comparison of the CP asymmetries in different decay modes [5].

One of the most promising processes for NP searches widely considered in literature [6–14] is the decay $B\to\phi K_S$. Various NP scenarios have been presented to explain the data. Unfortunately, we do not know at present which is the correct one. Hopefully, careful study in future will rule out some of the scenarios, at least as far as the understanding of B physics and CP violation is concerned.

Unlike $B\to\psi K_S$, the process $B\to\phi K_S$ has no tree level amplitude, which makes inroads for NP to play an important role in this mode. In the SM the decay $b\to s\bar{s}s$, which contributes to $B\to\phi K_S$, is induced at the one-loop level. Thus, it is natural to expect that new physics contribution to this decay mode may be quite significant. According to the KM mechanism of CP violation, both CP asymmetries in $B\to\phi K_S$ and $B\to\psi K_S$ processes should measure the same quantity, namely $\sin 2\beta$, with negligible hadronic uncertainties [up to $O(\lambda^2)$, $\lambda\approx 0.2$] [5,12]. However, contrary to the SM expectations, the recent measurements of CP asymmetries in $B\to\phi K_S$ by BABAR [15] and Belle [2] Collaborations have registered significant deviation from the predictions, as

$$\sin(2\beta)_{\phi K_S}=-0.19^{+0.52}_{-0.50}\pm 0.09 \quad \text{BABAR},$$
$$\sin(2\beta)_{\phi K_S}=-0.73\pm 0.64\pm 0.18 \quad \text{Belle} \tag{2}$$

with an average

$$\sin((2\beta)_{\phi K_S})_{ave}=-0.39\pm 0.41. \tag{3}$$

The corresponding branching ratio is given (in units of 10^{-6}) as

$$BR(B\to\phi K^0)=8.7^{+1.7}_{-1.5}\pm 0.9 \quad \text{BABAR},$$
$$BR(B\to\phi K^0)=10.0^{+1.9+0.9}_{-1.7-1.3} \quad \text{Belle},$$
$$BR(B\to\phi K^0)=5.4^{+3.7}_{-2.7}\pm 0.7 \quad \text{CLEO [16]} \tag{4}$$

with an average

$$BR(B\to\phi K^0)_{ave}=8.67\pm 1.28. \tag{5}$$

One can see that there are large statistical errors associated with these measurements. Nevertheless, the data establish a 2.7σ deviation from the SM prediction $\sin(2\beta)_{\phi K_S}$

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$=\sin(2\beta)_{\psi K_S}$. Therefore, if the measurement of $\sin 2\beta$ in $B\to\psi K_S$ is considered as the first evidence of large CP violation in B system, then the difference between $\sin(2\beta)_{\phi K_S}$ and $\sin(2\beta)_{\psi K_S}$ is likely to be regarded as a potential hint for the presence of new physics. There are several attempts in the literature [6–14] with detail discussion on the possible implications of this result.

The second channel we are interested in is $B^0\to\eta' K_S$. This is another two-body decay mode which is similar to the two mentioned above. Since many alternative schemes have been presented in the literature recently to explain the $\sin(2\beta)_{\phi K_S}$ deviation, it is therefore very important to verify that each of the NP scenarios should successfully explain them all. At present it is difficult to say which is the correct description. In order to narrow down the same it is highly desirable that one should carefully study them. This will not only help us to narrow down the sources of NP but also provide important clues for hadronic B physics in general.

$B^0\to\eta' K_S$ also receives dominant contribution from the $b\to s\bar{s}s$ gluonic penguin, and therefore it is expected that the time-dependent mixing-induced CP asymmetry for this mode will also give the value $\sin 2\beta$ [3]. However, this decay mode also has a tiny CKM as well as color suppressed $b\to u\bar{u}s$ tree contributions along with $b\to s\bar{q}q(q=u,d)$ penguins, which induce deviation from the leading result. It has been shown in Ref. [17] that this deviation will be below a 2% level. Belle [18] and BABAR [19] Collaborations have recently measured the CP asymmetry for this mode which is given as

$$\sin(2\beta)_{\eta' K_S}=0.71\pm0.37^{+0.05}_{-0.06}\quad \text{Belle},$$
$$\sin(2\beta)_{\eta' K_S}=0.02\pm0.34\pm0.03\quad \text{BABAR} \tag{6}$$

with an average

$$\sin((2\beta)_{\eta' K_S})_{ave}=0.33\pm0.25, \tag{7}$$

whose central value also deviates significantly from SM expectations.

In this paper we would like to investigate the new physics effects on the CP asymmetry parameters of the decay $B^0\to\phi K_S$ and $B^0\to\eta' K_S$ modes, arising from some simple extensions of the SM. The model considered here is the two Higgs doublet model (THDM) which allows tree level flavor changing neutral currents, the so called model III and the model with an extra vectorlike down quark (VLDQ). We show that the observed data for $B^0\to\phi K_S$ can be easily accommodated in the VLDQ model whereas it cannot be explained in the THDM, and both the models can explain the data for $B^0\to\eta' K_S$ mode. It has already been discussed in Ref. [6] that THDM cannot explain the observed CP asymmetry in $B\to\phi K_S$ mode whereas VLDQ model can explain it. However, in this paper we have explicitly done the calculation for both the decay modes and confirm the result of Ref. [6] for the decay mode $B\to\phi K_S$.

The paper is organized as follows. In Sec. II, we present the basic formulas for CP violating parameters, in the presence of new physics. In Sec. III, we discuss CP violation effects in $B^0\to\phi K_S$ mode arising from the THDM and VLDQ model. The $B^0\to\eta' K_S$ process is discussed in Sec. IV. Section V contains our conclusion.

II. *CP* VIOLATION PARAMETERS

Here, we will present the basic formulas of CP asymmetry parameters, in the presence of new physics. Due to the contributions from new physics, these parameters deviate substantially from their standard model values. Let us consider the B^0 and $\overline{B^0}$ decay into a CP eigenstate f_{CP} (we consider $f_{CP}=\phi K_S$ or $\eta' K_S$). Here, we are presenting the formulas for $B^0\to\phi K_S$ mode, but the same results will also hold for $B^0\to\eta' K_S$ mode. The time-dependent CP asymmetry for $B\to\phi K_S$ can be described by [20]

$$\mathcal{A}_{\phi K_S}(t)=C_{\phi K_S}\cos(\Delta M_{B_d}t)+S_{\phi K_S}\sin(\Delta M_{B_d}t), \tag{8}$$

where we identify

$$C_{\phi K_S}=\frac{1-|\lambda|^2}{1+|\lambda|^2},\quad S_{\phi K_S}=-\frac{2\,\mathrm{Im}(\lambda)}{1+|\lambda|^2} \tag{9}$$

as the direct and the mixing-induced CP asymmetries. The parameter λ corresponds to

$$\lambda=\frac{q}{p}\frac{A(\overline{B}\to\phi K_S)}{A(B\to\phi K_S)}, \tag{10}$$

where q and p are the mixing parameters and represented by the CKM elements in the standard model as

$$\frac{q}{p}=\frac{V_{tb}^*V_{td}}{V_{tb}V_{td}^*}. \tag{11}$$

Using CKM unitarity, the amplitude for $\overline{B}\to\phi K_S$ is given as [12,21]

$$A(\overline{B}\to\phi K_S)=\lambda_c A^{cs}+\lambda_u A^{us}, \tag{12}$$

where $\lambda_q=V_{qb}V_{qs}^*$. The first term which is the dominant one is real. Thus if one neglects the subdominant amplitude, i.e., the doubly Cabibbo supressed second term which in general is expected to be very small, the mixing-induced CP asymmetry is given as $S_{\phi K_S}=\sin 2\beta$, same as the one for $B\to\psi K_S$ in the SM. It has beeen shown in Ref. [12] that the correction due to the second term is up to $O(\lambda^2)$, i.e.,

$$|S_{\phi K_S}-\sin 2\beta|\leqslant O(\lambda^2). \tag{13}$$

Adding a mild dynamical assumption to the SU(3) analysis, recently it has been shown in Ref. [21] that the upper bound of standard model pollution to the dominant amplitude of $B\to\phi K_S$ mode is of the order of 0.25 and for $B\to\eta' K_S$ is 0.35.

New physics could in principle contribute to both mixing and decay amplitudes. The new physics contribution to mixing is universal while it is nonuniversal and process dependent in the decay amplitudes. As the NP contributions to mixing phenomena are universal, it will still set $S_{\psi K_S} = S_{\phi K_S}$. Therefore, to explain the observed 2.7σ deviation in $S_{\psi K_S} - S_{\phi K_S}$, here we explore the NP effects only in the decay amplitudes. Thus including the NP contributions, we can write the decay amplitude for $B \to \phi K$ process as

$$A(B^0 \to \phi K) = A_{SM} + A_{NP} = A_{SM}[1 + r_{NP} \ \ e^{i\phi_{NP}}], \quad (14)$$

where $r_{NP} = |A_{NP}/A_{SM}|$ (A_{SM} and A_{NP} correspond to the SM and NP contributions to the $B \to \phi K_S$ decay amplitude) and $\phi_{NP} = \arg(A_{NP}/A_{SM})$, which contains both strong and weak phase components.

The branching ratio for $B \to \phi K$ decay process can be given as

$$BR(B \to \phi K) = BR^{SM}(1 + r_{NP}^2 + 2r_{NP}\cos\phi_{NP}), \quad (15)$$

where BR^{SM} represents the corresponding standard model value.

Now if we write $\phi_{NP} = \delta_{NP} + \theta_{NP}$, where δ_{NP} and θ_{NP} are the relative strong and weak phases between the new physics contributions to the decay amplitude and the SM part, one can then obtain the expressions for the CP asymmetries as

$$S_{\phi K} = \frac{\sin 2\beta + 2r_{NP}\cos\delta_{NP}\sin(2\beta + \theta_{NP}) + r_{NP}^2\sin(2\beta + 2\theta_{NP})}{1 + r_{NP}^2 + 2r_{NP}\cos\delta_{NP}\cos\theta_{NP}} \quad (16)$$

and

$$C_{\phi K} = \frac{-2r_{NP}\sin\delta_{NP}\sin\theta_{NP}}{1 + r_{NP}^2 + 2r_{NP}\cos\delta_{NP}\cos\theta_{NP}}. \quad (17)$$

In Eqs. (16) and (17) there are three unknowns, namely, r_{NP}, θ_{NP}, and δ_{NP}. So if somehow we could constrain the value of r_{NP} considering different new physics models, we could vary the θ_{NP} and δ_{NP} parameters to obtain the required value of $S_{\phi K}$.

III. CP VIOLATION IN $B \to \phi K_S$ PROCESS

To study the CP violation effects in $B^0 \to \phi K_S$ process, first we present the SM amplitude and then we consider the THDM and thereafter the model with an extra vectorlike down quark, in the following sections.

A. SM contributions

In the SM, the decay process $B \to \phi K_S$ proceeds through the quark level transition $b \to s\bar{s}s$, which is induced by the QCD, electroweak, and magnetic penguins. QCD penguins with the top quark in the loop contribute predominantly to such process. However, since we are looking for NP here we would like to retain all the contributions. The effective Hamiltonian describing the decay $b \to s\bar{s}s$ [22] is given as

$$H_{eff} = -\frac{G_F}{\sqrt{2}} V_{tb}V_{ts}^* \left(\sum_{j=3}^{10} C_j O_j + C_g O_g \right), \quad (18)$$

where $O_3, \ldots, O_6$ and $O_7, \ldots, O_{10}$ are the standard QCD and electrowork penguin operators, respectively, and O_g is the gluonic magnetic operator. Within the SM and at scale M_W, the Wilson coefficients $C_1(M_W), \ldots, C_{10}(M_W)$ at next to leading logarithmic order (NLO) and $C_g(M_W)$ at leading logarithmic order (LO) have been given in Ref. [23]. The corresponding QCD corrected values at the energy scale $\mu = m_b$ can be obtained using the renormalization group equation, as described in Ref. [24].

To calculate the B meson decay rate, we use the factorization approximation to evaluate the hadronic matrix element $\langle O_i \rangle \equiv \langle \bar{K}^0 \phi | O_i | \bar{B}^0 \rangle$. Since the hadronic matrix elements do not appear in the expressions for CP asymmetry parameters, they will not introduce any uncertainties in the results. In this approximation the matrix elements are given as $\langle O_3 \rangle = \langle O_4 \rangle = 4X/3$, $\langle O_5 \rangle = X$, $\langle O_6 \rangle = X/3$, $\langle O_7 \rangle = -X/2$, $\langle O_8 \rangle = -X/6$, and $\langle O_9 \rangle = \langle O_{10} \rangle = -2X/3$, where the factorizable hadronic matrix element X is given as $X = \langle \bar{K}^0(p_K) | \bar{s}\gamma_\mu(1-\gamma_5)b | \bar{B}^0(p_B) \rangle \langle \phi(q,\epsilon_\phi) | \bar{s}\gamma^\mu(1-\gamma_5)s | 0 \rangle = 2F_1^{BK}(M_\phi^2) f_\phi M_\phi \epsilon_\phi \cdot p_K$. For evaluating the matrix element of the most relevant operator, i.e., O_g, we use the procedure of [25], where it has been shown that the operator O_g is related to the matrix element of the QCD and electroweak penguin operators as

$$\langle O_g \rangle = -\frac{\alpha_S}{4\pi} \frac{m_b}{\sqrt{\langle q^2 \rangle}} \left[\langle O_4 \rangle + \langle O_6 \rangle - \frac{1}{N_C}(\langle O_3 \rangle + \langle O_5 \rangle) \right], \quad (19)$$

where q^μ is the momentum transferred by the gluon to the

$(\bar{s},s)$ pair. The parameter $\langle q^2\rangle$ introduces certain uncertainty into the calculation. In the literature its value is taken in the range $1/4 \lesssim \langle q^2\rangle/m_b^2 \lesssim 1/2$ [26], and we will use $\langle q^2\rangle/m_b^2 = 1/2$ [24], in our numerical calculations.

Thus, in the factorization approach the amplitude $A \equiv \langle \phi K^0 | H_{eff} | B^0\rangle$ of the decay $B^0 \to \phi K^0$ takes a form

$$A(\bar{B}^0 \to \phi \bar{K}^0) = -\frac{G_F}{\sqrt{2}} V_{tb} V_{ts}^* \left[a_3 + a_4 + a_5 - \frac{1}{2}(a_7 + a_9 + a_{10}) \right] X, \tag{20}$$

where X stands for the factorizable hadronic matrix element of which the exact form is irrelevant for us since it cancels out in the CP asymmetries. The coefficients a_i are given by

$$a_{2i-1} = C_{2i-1}^{eff} + \frac{1}{N_c} C_{2i}^{eff}, \quad a_{2i} = C_{2i}^{eff} + \frac{1}{N_c} C_{2i-1}^{eff}, \tag{21}$$

where N_C is the number of colors. The values of the QCD improved effective coefficients C_i^{eff} can be found in [24,27]. Now substituting the values of a_i for $N_C = 3$, from [27], the value of the form factor $F_1^{BK}(M_\phi^2) = 0.39$ and using the ϕ meson decay constant $f_\phi = 0.233$ GeV and $\tau_{B^0} = 1.542 \times 10^{-12}$ sec [28], we obtain the branching ratio in the SM as

$$BR^{SM}(B \to \phi K^0) = 10.5 \times 10^{-6} \tag{22}$$

which lies within the present experimental limits (5).

B. Two Higgs doublet model contributions

We now proceed to calculate the new physics effect in the THDM, which is one of the simplest extensions of the SM [29]. In such models, the tree level flavor changing neutral currents (FCNC's) are prevented by imposing one *ad hoc* discrete symmetry to constrain the THDM scalar potential and Yukawa Lagrangian and thus one obtains the so called model I and model II [30]. In model I both the up- and down-type quarks get mass from the Yukawa couplings to the same Higgs doublet ϕ_1 and in model II the up- and down-type quarks get their masses from Yukawa couplings to two different scalar doublets ϕ_1 and ϕ_2. Here we consider model III [31] of the THDM where no discrete symmetry is imposed and both up- and down-type quarks may have diagonal and/or off-diagonal flavor changing couplings with the two Higgs doublets ϕ_1 and ϕ_2.

The Yukawa Lagrangian of the quarks in model III is given in the form [27]

$$\mathcal{L}_Y^{\rm III} = \eta_{ij}^U \bar{Q}_{i,L} \tilde{\phi}_1 U_{j,R} + \eta_{ij}^D \bar{Q}_{i,L} \phi_1 D_{j,R} + \hat{\xi}_{ij}^U \bar{Q}_{i,L} \tilde{\phi}_2 U_{j,R} + \hat{\xi}_{ij}^D \bar{Q}_{i,L} \phi_2 D_{j,R} + \text{H.c.}, \tag{23}$$

where $\phi_i (i=1,2)$ are the two Higgs doublets of the THDM, $\tilde{\phi}_i = i\tau_2 \phi_i^*$, $Q_{i,L}$ with $i = 1,2,3$ are the left handed isodoublet quarks, and $U_{j,R}(D_{j,R})$ are the right handed isosinglet up- (down-) type quarks. $\eta_{i,j}^{U,D}$ correspond to the diagonal mass matrices of the up and down quarks, while the neutral and charged flavor changing couplings are

$$\xi_{ij}^{U,D} = \frac{\sqrt{m_i m_j}}{v} \lambda_{ij}, \quad \hat{\xi}_{neutral}^{U,D} = \xi^{U,D},$$

$$\hat{\xi}_{charged}^U = \xi^U V_{CKM}, \quad \hat{\xi}_{charged}^D = V_{CKM} \xi^D, \tag{24}$$

where V_{CKM} is the Cabibbo-Kobayashi-Maskawa mixing matrix [4]. The coupling constants λ_{ij} are the free parameters of the model to be determined from experimental data.

Recently Chao *et al.* [32] studied the $b \to s\gamma$ process and Xiao *et al.* [27] studied the charmless nonleptonic decays of B mesons using model III of the THDM where they have kept only the couplings $\lambda_{tt} = |\lambda_{tt}| e^{i\theta_t}$ and $\lambda_{bb} = |\lambda_{bb}| e^{i\theta_b}$ as nonzero. From the studies of [27,32], the following parameter space for model III is known:

$$\lambda_{ij} = 0 \quad \text{for} \quad ij \neq tt \quad \text{or} \quad bb,$$

$$|\lambda_{tt}| = 0.3, \quad |\lambda_{bb}| = 35, \quad \theta = (0° - 30°),$$

$$M_{H^+} = (200 \pm 100) \ \text{GeV}, \tag{25}$$

where $\theta = \theta_b - \theta_t$ is allowed by the available data. The advantage of keeping only these two couplings nonzero is that the neutral Higgs boson does not contribute at the tree level or one-loop level. The new contributions therefore come only from the charged Higgs penguin loop with heavy internal top quark.

The new physics will manifest itself by modifying the corresponding Inami-Lim [33] functions $C_0(x)$, $D_0(x)$, $E_0(x)$, and $E_0'(x)$ which determine the Wilson coefficients $C_3(M_W), \ldots, C_{10}(M_W)$ and $C_g(M_W)$ in SM. The new strong and electroweak penguin diagrams in the THDM can be obtained from the corresponding penguin diagrams in the SM by replacing the internal $W^\pm$ lines by the charged Higgs $H^\pm$ lines. Following the same procedure as in the SM, it is straightforward to calculate the new γ, Z^0, and gluonic penguin diagrams induced by the exchange of charged Higgs bosons in model III. These new Wilson coefficients $C_i^{H^\pm}(M_W)$, $i = 3, \ldots, 10$, at NLO level and C_g at the LO level can now be written as

$$C_3^{H^\pm}(M_W) = -\frac{\alpha_S(M_W)}{24\pi} E_0^{NP} + \frac{\alpha_{em}}{6\pi} \frac{1}{\sin^2\theta_W} C_0^{NP},$$

$$C_4^{H^\pm}(M_W) = \frac{\alpha_S(M_W)}{8\pi} E_0^{NP},$$

$$C_5^{H^\pm}(M_W) = -\frac{\alpha_S(M_W)}{24\pi} E_0^{NP},$$

$$C_6^{H^\pm}(M_W) = \frac{\alpha_S(M_W)}{8\pi} E_0^{NP},$$

$$C_7^{H^\pm}(M_W)=\frac{\alpha_{em}}{6\pi}[4C_0^{NP}+D_0^{NP}],$$

$$C_8^{H^\pm}(M_W)=C_{10}^{H^\pm}(M_W)=0,$$

$$C_9^{H^\pm}(M_W)=\frac{\alpha_{em}}{6\pi}\left[4C_0^{NP}+D_0^{NP}+\frac{1}{\sin^2\theta_W}4C_0^{NP}\right],$$

$$C_g^{H^\pm}(M_W)=-\tfrac{1}{2}E_0'^{NP}, \tag{26}$$

where the functions C_0^{NP}, D_0^{NP}, E_0^{NP}, and $E_0'^{NP}$ are the new physics contributions to the Wilson coefficients arising from the charged Higgs exchange penguin diagrams. These are given by

$$C_0^{NP}=-\frac{x_t}{16}\left[\frac{y_t}{1-y_t}+\frac{y_t}{(1-y_t)^2}\ln y_t\right]|\lambda_{tt}|^2,$$

$$D_0^{NP}=-\tfrac{1}{3}H(y_t)|\lambda_{tt}|^2,$$

$$E_0^{NP}=-\tfrac{1}{2}I(y_t)|\lambda_{tt}|^2,$$

$$E_0'^{NP}=\tfrac{1}{6}J(y_t)|\lambda_{tt}|^2-K(y_t)|\lambda_{tt}\lambda_{bb}|e^{i\theta}, \tag{27}$$

with

$$H(y)=\frac{38y-79y^2+47y^3}{72(1-y)^3}+\frac{4y-6y^2+3y^4}{12(1-y)^4}\ln y,$$

$$I(y)=\frac{16y-29y^2+7y^3}{36(1-y)^3}+\frac{2y-3y^2}{6(1-y)^4}\ln y,$$

$$J(y)=\frac{2y+5y^2-y^3}{4(1-y)^3}+\frac{3y^2}{2(1-y)^4}\ln y,$$

$$K(y)=\frac{-3y+y^2}{4(1-y)^2}-\frac{y}{2(1-y)^3}\ln y. \tag{28}$$

In the above use has been made of $x_t=m_t^2/M_W^2$ and $y_t=m_t^2/M_{H^+}^2$.

Since the charged Higgs bosons that appeared in model III have been integrated out at the scale M_W, the QCD running of Wilson coefficients $C_i^{H^\pm}(M_W)$ down to the scale $\mu=O(m_b)$ using the renormalization group equation can be done in the same way as in the SM. Including the new physics contributions the values of the effective Wilson coefficients at the scale $O(m_b)$ are explicitly given in Ref. [27]. Using the values for the Wilson coefficients from [27], we obtain the $B\to\phi K^0$ amplitude in the THDM as

$$A^{THDM}(B\to\phi K^0)=A^{SM}(1+0.28e^{i(\theta_{NP}+\delta_{NP})}). \tag{29}$$

Now taking $r_{NP}=0.28$ and varying the weak phase $\theta_{NP}=\{-\pi,\pi\}$ and strong phase $\delta_{NP}=\{0,2\pi\}$ according to Eq.

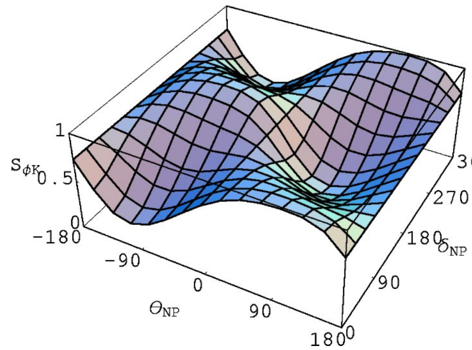

FIG. 1. 3D plot of $S_{\phi K}$ versus the weak phase θ_{NP} and strong phase δ_{NP} (in degrees) for $r_{NP}=0.28$.

(16), we find that the value of $S_{\phi K}$ cannot be negative as shown in Fig. 1. Thus the observed value of $S_{\phi K}$ cannot be accommodated in the THDM.

C. Contributions from the model with an extra vectorlike down quark

Now we consider the model with an additional vectorlike down quark [34]. It is a simple model beyond the SM with an enlarged matter sector with an additional vectorlike down quark D_4. The most interesting effects in this model concern CP asymmetries in neutral B decays into final CP eigenstates. At a more phenomenological level, models with isosinglet quarks provide the simplest self-consistent framework to study deviations of 3×3 unitarity of the CKM matrix as well as allow flavor changing neutral currents at the tree level. The presence of an additional down quark implies a 4×4 matrix $V_{i\alpha}$ $(i=u,c,t,4,\alpha=d,s,b,b')$, diagonalizing the down quark mass matrix. For our purpose, the relevant information for the low energy physics is encoded in the extended mixing matrix. The charged currents are unchanged except that the V_{CKM} is now the 3×4 upper submatix of V. However, the distinctive feature of this model is that the FCNC enters neutral current Lagrangian of the left handed down quarks:

$$\mathcal{L}_Z=\frac{g}{2\cos\theta_W}[\overline{u}_{Li}\gamma^\mu u_{Li}-\overline{d}_{L\alpha}U_{\alpha\beta}\gamma^\mu d_{L\beta}-2\sin^2\theta_W J^\mu_{em}]Z_\mu \tag{30}$$

with

$$U_{\alpha\beta}=\sum_{i=u,c,t}V^\dagger_{\alpha i}V_{i\beta}=\delta_{\alpha\beta}-V^*_{4\alpha}V_{4\beta}, \tag{31}$$

where U is the neutral current mixing matrix for the down sector which is given above. As V is not unitary, $U\neq\mathbf{1}$. In particular its nondiagonal elements do not vanish:

$$U_{\alpha\beta}=-V^*_{4\alpha}V_{4\beta}\neq0 \quad \text{for} \quad \alpha\neq\beta. \tag{32}$$

Since the various $U_{\alpha\beta}$ are nonvanishing, they would signal new physics and the presence of the FCNC at the tree level, and this can substantially modify the predictions for CP asymmetries. The new element U_{sb} which is relevant to our study is given as

$$U_{sb}=V_{us}^{*}V_{ub}+V_{cs}^{*}V_{cb}+V_{ts}^{*}V_{tb}\,. \tag{33}$$

The decay mode $B^0\to\phi K_S$ receives the new contributions both from color allowed and color suppressed Z-mediated FCNC transitions. The new additional operators are given as

$$O_1^{Z\text{-}FCNC}=[\overline{s}_\alpha\gamma^\mu(1-\gamma_5)b_\alpha][\overline{s}_\beta\gamma_\mu(C_V^s-C_A^s\gamma_5)s_\beta],$$

$$O_2^{Z\text{-}FCNC}=[\overline{s}_\beta\gamma^\mu(1-\gamma_5)b_\alpha][\overline{s}_\alpha\gamma_\mu(C_V^s-C_A^s\gamma_5)s_\beta], \tag{34}$$

where C_V^s and C_A^s are the vector and axial vector $Zs\overline{s}$ couplings. Using the Fierz transformation and the identity $(C_V^s-C_A^s\gamma_5)=[(C_V^s+C_A^s)(1-\gamma_5)+(C_V^s-C_A^s)(1+\gamma_5)]/2$, the matrix elements of the operators are given as

$$\langle\phi\overline{K}^0|O_1^{Z\text{-}FCNC}|\overline{B}^0\rangle=\left[\frac{4}{3}\,\frac{(C_V^s+C_A^s)}{2}+\frac{(C_V^s-C_A^s)}{2}\right]X,$$

$$\langle\phi\overline{K}^0|O_2^{Z\text{-}FCNC}|\overline{B}^0\rangle=\left[\frac{4}{3}\,\frac{(C_V^s+C_A^s)}{2}+\frac{1}{3}\,\frac{(C_V^s-C_A^s)}{2}\right]X. \tag{35}$$

The values for C_V^s and C_A^s are taken as

$$C_V^s=-\tfrac{1}{2}+\tfrac{2}{3}\sin^2\theta_W,\quad C_A^s=-\tfrac{1}{2}. \tag{36}$$

Thus the amplitude for $B\to\phi K$ arising from the Z-mediated FCNC tree diagram is given as

$$A^{VLDQ}(\overline{B}^0\to\phi\overline{K}_0)=\frac{G_F}{\sqrt{2}}\,\frac{4}{3}(-1+\sin^2\theta_W)U_{sb}X. \tag{37}$$

Using the experimental upper limit $Br(B\to X_S l^+l^-)<4.2\times10^{-5}$ [36], in Ref. [35] the bound on $|U_{bs}|$ is found to be $|U_{bs}|\leqslant2\times10^{-3}$. Recently Belle Collaboration [37] has measured the branching ratio for the process $B\to X_S l^+l^-$ as

$$Br(B\to X_S l^+l^-)=(6.1\pm1.4^{+1.4}_{-1.1})\times10^{-6}. \tag{38}$$

Using the above result one can obtain the value [35,38]

$$|Y_0(x_t)\ \lambda_t^{bs}+C_{U2Z}\ \ U_{bs}|=0.06\pm0.03, \tag{39}$$

where all the parameters in Eq. (39) are given in [35]. Thus one obtains the value of U_{bs} as

$$|U_{bs}|\simeq1\times10^{-3}. \tag{40}$$

Now using $\sin^2\theta_W$=0.23, we find

$$r_{NP}\simeq0.58. \tag{41}$$

The variation of $S_{\phi K}$ with respect to the strong phase δ_{NP} and weak phase θ_{NP} according to Eq. (16) in the VLDQ model is shown in Fig. 2 and the variation of branching ratio (15) is shown in Fig. 3. It can be seen from the figures that the observed asymmetry $S_{\phi K}$ and the branching ratio can be easily accommodated in this model.

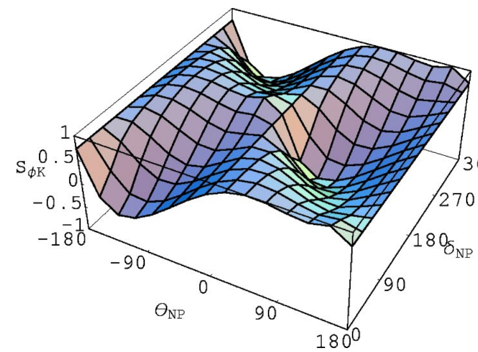

FIG. 2. 3D plot of $S_{\phi K}$ versus the weak phase θ_{NP} and strong phase δ_{NP} (in degrees) for r_{NP}=0.58.

IV. CP VIOLATION IN $B\to\eta' K_S$ PROCESS

At this stage we are in a position to test, as mentioned earlier, whether the above two models (model III of the THDM and VLDQ) can accommodate the result for another similar mode, which seems to be in agreement with the SM. In doing so, now we consider the $B^0\to\eta' K^0$ process.

A. Contributions from the SM and THDM

In the SM, in addition to $b\to s\overline{q}q$ $[q=(u,d,s)]$ penguins, the $B^0\to\eta' K_S$ process also receives a small contribution from color suppressed $b\to u\overline{u}s$ tree diagram. We first find out the standard model contribution. The matrix element in the SM is given as

$$A(\overline{B}^0\to\eta'\overline{K}^0)=\frac{G_F}{\sqrt{2}}\left[V_{ub}V_{us}^{*}\sum_{i=1}^{2}C_i^{eff}\langle\eta'\overline{K}^0|O_i|\overline{B}^0\rangle-V_{tb}V_{ts}^{*}\sum_{i=3}^{10}C_i^{eff}\langle\eta'\overline{K}^0|O_i|\overline{B}^0\rangle\right], \tag{42}$$

where $O_{1,2}$ are the tree and $O_{3\text{-}6(7\text{-}10)}$ are the QCD (electroweak) penguin operators. The matrix elements of these operators are given in the factorization approximation as [14]

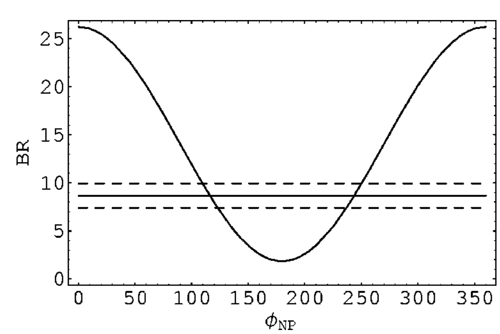

FIG. 3. Branching ratio of $B\to\phi K^0$ process (in units of 10^{-5}) versus the phase ϕ_{NP} (in degrees). The horizontal solid line is the central experimental value whereas the dashed horizontal lines denote the error limits.

$$\langle \eta' \overline{K^0}|O_1|\overline{B^0}\rangle=\tfrac{1}{3}X_2, \quad \langle \eta' \overline{K^0}|O_2|\overline{B^0}\rangle=X_2,$$

$$\langle \eta' \overline{K^0}|O_3|\overline{B^0}\rangle=\tfrac{1}{3}X_1+2X_2+\tfrac{4}{3}X_3,$$

$$\langle \eta' \overline{K^0}|O_4|\overline{B^0}\rangle=X_1+\tfrac{2}{3}X_2+\tfrac{4}{3}X_3,$$

$$\langle \eta' \overline{K^0}|O_5|\overline{B^0}\rangle=\frac{R_1}{3}X_1-2X_2-\left(1-\frac{R_2}{3}\right)X_3,$$

$$\langle \eta' \overline{K^0}|O_6|\overline{B^0}\rangle=R_1X_1-\frac{2}{3}X_2-\left(\frac{1}{3}-R_2\right)X_3,$$

$$\langle \eta' \overline{K^0}|O_7|\overline{B^0}\rangle=\frac{1}{2}\left[-\frac{R_1X_1}{3}-X_2+\left(1-\frac{R_2}{3}\right)X_3\right],$$

$$\langle \eta' \overline{K^0}|O_8|\overline{B^0}\rangle=\frac{1}{2}\left[-R_1X_1-\frac{X_2}{3}+\left(\frac{1}{3}-R_2\right)X_3\right],$$

$$\langle \eta' \overline{K^0}|O_9|\overline{B^0}\rangle=\frac{1}{2}\left[-\frac{X_1}{3}+X_2-\frac{4}{3}X_3\right],$$

$$\langle \eta' \overline{K^0}|O_{10}|\overline{B^0}\rangle=\frac{1}{2}\left[-X_1+\frac{X_2}{3}-\frac{4}{3}X_3\right], \tag{43}$$

where

$$X_1=i(m_B^2-m_{\eta'}^2)F_0^{B\to\pi}(m_{K^0}^2)\frac{X_{\eta'}}{\sqrt{2}}f_K,$$

$$X_2=i(m_B^2-m_{K^0}^2)F_0^{B\to K}(m_{\eta'}^2)\frac{X_{\eta'}}{\sqrt{2}}f_\pi,$$

$$X_3=i(m_B^2-m_{K^0}^2)F_0^{B\to K}(m_{\eta'}^2)Y_{\eta'}\sqrt{2f_K^2-f_\pi^2},$$

$$R_1=\frac{2m_{K^0}^2}{(m_b-m_d)(m_S+m_d)},$$

$$R_2=\frac{2(2m_{K^0}^2-m_\pi^2)}{(m_b-m_S)(m_S+m_S)}. \tag{44}$$

$X_{\eta'}=0.57$ and $Y_{\eta'}=0.82$ are the mixing parameters of the $u\bar{u}+d\bar{d}$ and $s\bar{s}$ components in the η' meson [39], which correspond to $\theta_P=-20°$. Thus the amplitude is given as

$$A(B^0\to\eta' K^0)=i\frac{G_F}{2}\Bigg(V_{ub}V_{us}^*a_2X_2-V_{tb}V_{ts}^*\Bigg\{\Bigg[\left(a_4-\frac{a_{10}}{2}\right.$$
$$\left.+\left(a_6-\frac{a_8}{2}\right)R_1\right]X_1+\Bigg(2(a_3-a_5)$$
$$-\frac{1}{2}(a_7-a_9)\Bigg)X_2+\Bigg[a_3+a_4-a_5$$
$$+\frac{1}{2}(a_7-a_9-a_{10})+\left(a_6-\frac{a_8}{2}\right)R_2\Bigg]X_3\Bigg\}\Bigg). \tag{45}$$

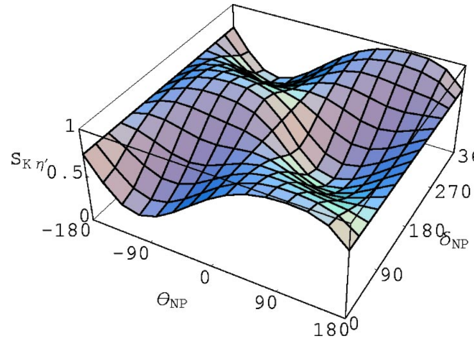

FIG. 4. 3D plot of $S_{\eta' K}$ versus the weak phase θ_{NP} and strong phase δ_{NP} (in degrees) for $r_{NP}=0.27$.

The decay width can be given by

$$\Gamma(B^0\to\eta' K^0)=\frac{|\vec{p}|}{8\pi m_B^2}|A(B^0\to\eta' K^0)|^2. \tag{46}$$

Using $F_0^{(B\to\pi)}(m_{K^0}^2)=0.335$, $f_{K(\pi)}=0.16(0.13)$ GeV, the quark masses as $(m_d,m_S,m_b)=(0.0076,0.122,4.88)$ GeV, and the values of the coefficients a_i's for $N_C=3$ from Ref. [27] we obtain the branching ratio in the standard model as

$$BR(B^0\to\eta' K^0)|_{SM}=3.24\times10^{-5}, \tag{47}$$

which is slightly less than the current experimental data [28],

$$BR(B^0\to\eta' K^0)=(5.8^{+1.4}_{-1.3})\times10^{-5}. \tag{48}$$

Now we consider the contributions arising from the THDM. As discussed earlier in this case due to the presence of new charged Higgs penguin diagrams, the values of the effective Wilson coefficients a_i's get modified. Again substituting their values from [27] in Eq. (45), we obtain the transition amplitude as

$$A^{THDM}(B^0\to\eta' K^0)=A^{SM}(1+0.27\ \ e^{i(\theta_{NP}+\delta_{NP})}). \tag{49}$$

Now taking $r_{NP}=0.27$ and varying the weak phase $\theta_{NP}=\{-\pi,\pi\}$ and strong phase $\delta_{NP}=\{0,2\pi\}$ we can see from Fig. 4 that the observed value of $S_{\eta' K}$ can be accommodated in the THDM. Furthermore, the observed branching ratio can also be explained in this model as seen from Fig. 5. If we take a crude assumption that the THDM and SM amplitudes interfere constructively, the maximum value of branching ratio is found to be

$$BR^{THDM}(B^0\to\eta' K^0)=5.22\times10^{-5}, \tag{50}$$

which lies within the present experimental limits [28].

B. Contributions from the VLDQ model

Now we consider the contributions arising from the extra vectorlike down quark model. In this case the $B^0\to\eta' K^0$ process proceeds through both color allowed and color suppressed tree level Z-mediated FCNC diagrams. The corresponding operators are given as

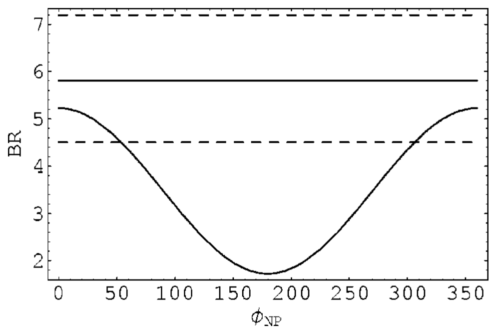

FIG. 5. Branching ratio of $B^0\to\eta' K^0$ (in units of 10^{-5}) process in the THDM versus the phase ϕ_{NP} (in degrees). The horizontal solid line is the central experimental value whereas the dashed horizontal lines denote the error limits.

$$O_1^{Z\text{-}FCNC}=[\overline{s}_\alpha\gamma^\mu(1-\gamma_5)b_\alpha][\overline{q}_\beta\gamma_\mu(C_V^q-C_A^q\gamma_5)q_\beta]$$
$$\equiv[\overline{s}_\alpha\gamma^\mu(1-\gamma_5)b_\alpha]\left[\overline{q}_\beta\gamma_\mu\left\{\frac{(C_V^q+C_A^q)}{2}(1-\gamma_5)+\frac{(C_V^q-C_A^q)}{2}(1+\gamma_5)\right\}q_\beta\right],$$

$$O_2^{Z\text{-}FCNC}=[\overline{s}_\beta\gamma^\mu(1-\gamma_5)b_\alpha][\overline{q}_\alpha\gamma_\mu(C_V^q-C_A^q\gamma_5)q_\beta]$$
$$\equiv[\overline{s}_\beta\gamma^\mu(1-\gamma_5)b_\alpha]\left[\overline{q}_\alpha\gamma_\mu\left\{\frac{(C_V^q+C_A^q)}{2}(1-\gamma_5)+\frac{(C_V^q-C_A^q)}{2}(1+\gamma_5)\right\}q_\beta\right]. \tag{51}$$

Using the Fierz transformation and equation of motion, the matrix elements of these operators are given as

$$\langle\eta'\overline{K}^0|O_1^{Z\text{-}FCNC}|\overline{B}^0\rangle=(C_A^u+C_A^d)X_2+C_A^sX_3+\tfrac{1}{6}\{[C_V^d(1+R_1)+C_A^d(1-R_1)]X_1+[C_V^s(1+R_2)+C_A^s(1-R_2)]X_3\}, \tag{52}$$

$$\langle\eta'\overline{K}^0|O_2^{Z\text{-}FCNC}|\overline{B}^0\rangle=\tfrac{1}{3}(C_A^u+C_A^d)X_2+\tfrac{1}{3}C_A^sX_3+\tfrac{1}{2}\{[C_V^d(1+R_1)+C_A^d(1-R_1)]X_1+[C_V^s(1+R_2)+C_A^s(1-R_2)]X_3\}. \tag{53}$$

So the amplitude for $B^0\to\eta' K^0$ in the VLDQ model is given as

$$A^{VLDQ}(\overline{B}^0\to\eta'\overline{K}^0)=\frac{G_F}{\sqrt{2}}U_{sb}\left(\tfrac{4}{3}\{(C_A^u+C_A^d)X_2+C_A^sX_3\}+\tfrac{2}{3}\{[C_V^d(1+R_1)+C_A^d(1-R_1)]X_1+[C_V^s(1+R_2)+C_A^s(1-R_2)]X_3\}\right). \tag{54}$$

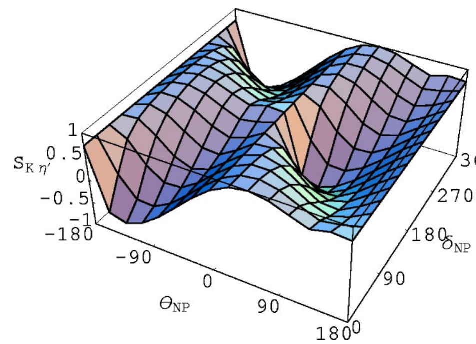

FIG. 6. 3D plot of $S_{\eta' K^0}$ versus the weak phase θ_{NP} and strong phase δ_{NP} (in degrees) for $r_{NP}=0.72$.

Substituting the values of $C^q_{V(A)}$ as

$$C_V^u=\tfrac{1}{2}-\tfrac{4}{3}\sin^2\theta_W,\quad C_A^u=\tfrac{1}{2},$$
$$C_V^{(s,d)}=-\tfrac{1}{2}+\tfrac{2}{3}\sin^2\theta_W,\quad C_A^{(s,d)}=-\tfrac{1}{2}, \tag{55}$$

we find

$$r_{NP}\simeq 0.72. \tag{56}$$

The variation of $S_{\eta' K^0}$ and the branching ratio according to Eqs. (16) and (15) in the vectorlike down quark model are shown in Figs. 6 and 7. It can be seen that the observed asymmetry and branching ratio for $B^0\to\eta' K^0$ mode can be easily accommodated in this model.

V. CONCLUSIONS

To summarize, the time-dependent CP asymmetry measurements in $B\to\phi K_S$ give $\sin 2\beta$, which is 2.7σ deviation from the corresponding value in $B\to\psi K_S$. According to the SM expectation they should measure the same. Unlike the $B\to\psi K_S$, which is a tree level process, $B\to\phi K_S$ occurs at the one-loop level, which allows room for new physics to play an important role. In this paper, we have explored two simple scenarios beyond the SM, the two Higgs doublet model (model III), and a model with an extra vectorlike down quark. We found that model III of the THDM is unable

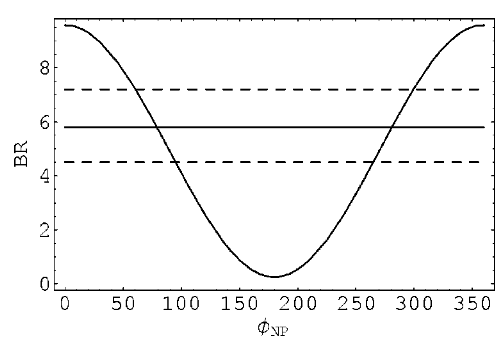

FIG. 7. Branching ratio of $B^0\to\eta' K^0$ process (in units of 10^{-5}) in the VLDQ model versus the phase ϕ_{NP} (in degree). The horizontal solid line is the central experimental value whereas the dashed horizontal lines denote the error limits.

to explain, whereas the vectorlike down quark model can easily explain the result.

It is important to note here that any new physics scenario that explains the ϕK_S discrepancy must also explain another similar two-body decay $B \to \eta' K_S$, which is also expected to give the same value as of ψK_S or ϕK_S, i.e., $\sin 2\beta$. In doing so it will be easy to rule out or narrow down the various NP scenarios. We found that both the models (model III of the THDM and VLDQ) can explain the $\eta' K_S$ result. This in turn gives us the clue that the VLDQ model may possibly be a strong contender for the NP effects responsible in $B \to \phi K_S$. It is worthwhile to emphasize that various supersymmetric models (as can be found in the literature) can explain the ϕK_S discrepancy. But apart from [13,14] none of the scenarios so far explained the simultaneous explanation of ϕK_S and $\eta' K_S$. On the other hand, our findings indicate that the simple nonsupersymmetric extension of the SM in terms of the matter content should not be ignored for possible NP candidature. Regardless of the sources of NP, if in future the ϕK_S result continues to be different from the SM expectation, then it will certainly establish the presence of NP.

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[1] BABAR Collaboration, B. Aubert *et al.*, Phys. Rev. Lett. **89**, 201802 (2002).

[2] Belle Collaboration, K. Abe *et al.*, hep-ex/0207098.

[3] Y. Nir, in *Proceedings of the ICHEP 2002*, Amsterdam, 2002 [Nucl. Phys. B (Proc. Suppl.) **117**, 111 (2003)].

[4] N. Cabibbo, Phys. Rev. Lett. **10**, 531 (1963); M. Kobayashi and T. Maskawa, Prog. Theor. Phys. **49**, 652 (1973).

[5] Y. Grossman and M.P. Worah, Phys. Lett. B **395**, 241 (1997).

[6] G. Hiller, Phys. Rev. D **66**, 071502(R) (2002).

[7] G. Barenboim, J. Bernabeu, and M. Raidal, Phys. Rev. Lett. **80**, 4625 (1998); M. Raidal, *ibid.* **89**, 231803 (2002).

[8] M. Ciuchini and L. Silverstini, Phys. Rev. Lett. **89**, 231802 (2002); C.-W. Chiang and J.L. Rosner, Phys. Rev. D **68**, 014007 (2003).

[9] E. Lunghi and D. Wyler, Phys. Lett. B **521**, 320 (2001); S. Khalil and E. Kou, Phys. Rev. D **67**, 055009 (2003); G.L. Kane, P. Ko, H. Yang, C. Kolda, J.-H. Park, and L.T. Yang, Phys. Rev. Lett. **90**, 141803 (2003); K. Agashe and C. D Carone, hep-ph/0304229.

[10] A. Datta, Phys. Rev. D **66**, 071702(R) (2002).

[11] S. Baek, Phys. Rev. D **67**, 096004 (2003).

[12] Y. Grossman, G. Isidori, and M. Worah, Phys. Rev. D **58**, 057504 (1998).

[13] B. Dutta, C.S. Kim, and S. Oh, Phys. Rev. Lett. **90**, 011801 (2003); A. Kundu and T. Mitra, Phys. Rev. D **67**, 116005 (2003).

[14] S. Khalil and E. Kou, hep-ph/0303214.

[15] BABAR Collaboration, B. Aubert *et al.*, hep-ex/0207070.

[16] CLEO Collaboration, R.A. Briere *et al.*, Phys. Rev. Lett. **86**, 3718 (2001); CLEO Collaboration, C.P. Jessop *et al.*, *ibid.* **85**, 2881 (2000).

[17] D. London and A. Soni, Phys. Lett. B **407**, 61 (1997).

[18] Belle Collaboration, K. Abe *et al.*, Phys. Rev. D **67**, 031102(R) (2003).

[19] BABAR Collaboration, B. Aubert *et al.*, hep-ex/0303046.

[20] I. I. Bigi and A. I. Sanda, *CP Violation*, Cambridge Monographs on Particle Physics, Nuclear Physics and Cosmology Vol. 9 (Cambridge University Press, Cambridge, England, 2000); G. C. Branco, L. Lavoura, and J. P. Silva, *CP Violation*, International Series of Monographs on Physics Vol. 103 (Oxford University Press, New York, 1999).

[21] Y. Grossman, Z. Ligeti, Y. Nir, and H. Quinn, Phys. Rev. D **68**, 015004 (2003).

[22] A. J. Buras and R. Fleischer, in *Heavy Flavours II*, edited by A. J. Buras and M. Lindner (World Scientific, Singapore, 1998), p. 65.

[23] G. Buchalla, A.J. Buras, and M.E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).

[24] A. Ali and C. Greub, Phys. Rev. D **57**, 2996 (1998); A. Ali, G. Kramer, and C.D. Lü, *ibid.* **58**, 094009 (1998).

[25] G. Barenboim, J. Bernabeu, J. Matias, and M. Raidal, Phys. Rev. D **60**, 016003 (1999).

[26] N.G. Deshpande and J. Trampetic, Phys. Rev. D **41**, 2926 (1990); H. Simma and D. Wyler, Phys. Lett. B **272**, 395 (1991); J.-M. Gerard and W.-S. Hou, Phys. Rev. D **43**, 2909 (1991); Phys. Lett. B **253**, 478 (1991).

[27] Z. Xiao, C.S. Li, and K.-T. Chao, Phys. Rev. D **63**, 074005 (2001); **65**, 114021 (2002).

[28] Particle Data Group, K. Hagiwara *et al.*, Phys. Rev. D **66**, 010001 (2002).

[29] J. F. Gunion, H. E. Haber, G. Kane, and S. Dawson, *The Higgs Hunter's Guide* (Addison-Wesley, Reading, MA, 1990).

[30] D. Atwood, L. Reina, and A. Soni, Phys. Rev. D **55**, 3156 (1997).

[31] T.P. Cheng and M. Sher, Phys. Rev. D **35**, 3484 (1987); M. Sher and Y. Yuan, *ibid.* **44**, 1461 (1991); W.S. Hou, Phys. Lett. B **296**, 179 (1992); D. Chang, W.S. Hou, and W.Y. Keung, Phys. Rev. D **48**, 217 (1993); Y.L. Wu and L. Wolfenstein, Phys. Rev. Lett. **73**, 1762 (1994); D. Atwood, L. Reina, and A. Soni, *ibid.* **75**, 3800 (1995).

[32] D. B.-Chao, K. Cheung, and W.-Y. Keung, Phys. Rev. D **59**, 115006 (1999).

[33] T. Inami and C.S. Lim, Prog. Theor. Phys. **65**, 297 (1981); **65**, 1772(E) (1981).

[34] Y. Grossman, Y. Nir, and R. Rattazzi, in *Heavy Flavours II*, edited by A. J. Buras and M. Lindner (World Scientific, Singapore, 1998), p. 755.

[35] G. Barenboim, F.J. Bottella, and O. Vives, Phys. Rev. D **64**, 015007 (2001).

[36] CLEO Collaboration, S. Glenn *et al.*, Phys. Rev. Lett. **80**, 2289 (1998).
[37] Belle Collaboration, J. Kaneko *et al.*, Phys. Rev. Lett. **90**, 021801 (2003).
[38] G. Buchalla, G. Hiller, and G. Isidori, Phys. Rev. D **63**, 014015 (2001).
[39] J.L. Rosner, Phys. Rev. D **27**, 1101 (1983); E. Kou, *ibid.* **63**, 054027 (2001).