$B^0 \rightarrow \phi K_S$ in supergravity models with CP violations

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We examine the $B \rightarrow \phi K$ decays within the framework of SUGRA models making use of the improved QCD factorization method of Beneke *et al.* which allows calculations of nonfactorizable contributions. All other experimental constraints ($B \rightarrow X_s \gamma$, neutron and electron electric dipole moments, dark matter constraints, etc.) are imposed. We calculate the CP violating parameters $S_{\phi K_S}$, $C_{\phi K_S}$ and $\mathcal{A}_{\phi K^{\mp}}$ as well as the branching ratios (BRs) of B^0 and $B^{\pm}$, Br$[B \rightarrow \phi K]$. We find that, for the standard model (SM) and MSUGRA, it is not possible to account for the observed 2.7σ deviation between $S_{\phi K_S}$ and $S_{J/\psi K_S}$. In general the BRs are also in 3σ disagreement with experiment, except in the parameter region where the weak annihilation terms dominate the decay (and hence where the theory is least reliable). Thus if future data confirm the current numbers, this would represent the first significant breakdown of both SM and MSUGRA. We show then that, adding a SUGRA nonuniversal cubic soft breaking left-right term mixing the second and third generations in either the down or up quark sector, all data can be accommodated for a wide range of parameters. The full 6×6 quark mass matrices are used and the supersymmetric contributions calculated without approximation.

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I. INTRODUCTION

Rare decay modes of the B meson are important places to test the standard model (SM) and to look for new physics. However, large theoretical uncertainties in the calculations of exclusive nonleptonic B decays make it difficult to extract useful information from experimental data. Nevertheless, CP asymmetries of neutral B meson decays into final CP eigenstates, i.e., $B \rightarrow \phi K_S$ and $B \rightarrow J/\psi K_S$, are uniquely clean in their theoretical interpretations. Among these decay modes, $B^0 \rightarrow \phi K_S$ is induced only at the one loop level in the SM and hence is a very promising mode to see the effects of new physics. In the SM, it is predicted that the CP asymmetries of $B^0 \rightarrow \phi K_S$ and $B \rightarrow J/\psi K_S$ should measure the same $\sin 2\beta$ with a negligible $O(\lambda^2)$ difference [1]. On the other hand, the BaBar and Belle measurements [2–5] show a 2.7σ disagreement between $S_{\phi K_S}$ and $S_{J/\psi K_S}$ [5]:[1]

$$S_{J/\psi K_S} = 0.734 \pm 0.055,$$

$$S_{\phi K_S} = -0.38 \pm 0.41, \tag{1}$$

while $S_{J/\psi K_S} = \sin 2\beta_{J/\psi K_S}$ (which is a tree level process) is in excellent agreement with Buras' SM evaluation, $\sin 2\beta = 0.715^{+0.055}_{-0.045}$ from the Cabibbo-Kobayaski-Maskawa (CKM) matrix [8]. In addition, the branching ratios (BRs) and the direct CP asymmetries of both the charged and neutral modes of $B \rightarrow \phi K$ have also been measured [2–5]:[2]

$$\mathrm{BR}[B^0 \rightarrow \phi K_S] = (8.0 \pm 1.3)\times 10^{-6},$$

$$\mathrm{BR}[B^+ \rightarrow \phi K^+] = (10.9 \pm 1.0)\times 10^{-6}, \tag{2}$$

$$C_{\phi K_S} = -0.19 \pm 0.30,$$

$$\mathcal{A}_{CP}(B^+ \rightarrow \phi K^+) = (3.9 \pm 8.8 \pm 1.1)\%. \tag{3}$$

In general, any model should explain all these data. In particular, the relatively small uncertainties in the BRs of $B^+ \rightarrow \phi K^+$ and $B^0 \rightarrow \phi K_S$ need to be considered in the analysis since they are highly correlated and both are based on the $b \rightarrow s$ transition. In the SM, $\mathcal{A}_{CP}(B^+ \rightarrow \phi K^+)$ is small and agrees with Eq. (3). So this direct CP asymmetry result plays an important role in constraining the new physics contribution which might explain the discrepancy between $S_{\phi K_S}$ and $S_{J/\psi K_S}$.

The discrepancy between $S_{\phi K_S}$ and $S_{J/\psi K_S}$ has been discussed in some recent works [10–22] in the framework of

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[1]After submitting this work new data from Belle [6] gave a value of $S_{\phi K_S} = -0.96 \pm 0.5^{+0.09}_{-0.11}$ (a 3.5σ deviation from the standard model) and preliminary analysis of new data from BaBar [7] gave $S_{\phi K_S} = +0.45 \pm 0.43 \pm 0.07$. Belle and BaBar would then disagree by 2.1σ, and if one averages the new values one obtains [7] $S_{\phi K_S} = -0.15 \pm 0.33$ which is again 2.7σ from the standard model.

[2]Our average of BR$[B^+ \rightarrow \phi K^+]$ only includes BaBar and Belle since CLEO [9] is 2.3σ away. BR$[B^+ \rightarrow \phi K^+]$ would become $(9.4 \pm 0.9)\times 10^{-6}$ if CLEO is included.

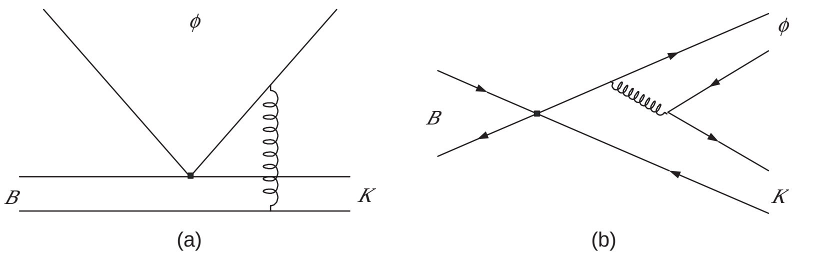

FIG. 1. Hard spectator scattering diagram (a) and weak annihilation diagram (b). In (a) the gluon can connect the spectator with either ϕ quark and in (b) the gluon can originate from any B quark or K quark.

supersymmetric (SUSY) models, especially in the minimal supersymmetric standard model (MSSM) with the mass insertion method [23]. Although these works can provide useful constraints on certain off-diagonal terms of squark mass matrices at low energy, we find that it is very interesting to investigate this problem in the context of grand unified theory (GUT) models. Since GUT models explain a number of phenomena at low energy by a few well motivated parameters at the GUT scale, various experimental measurements get correlated in this framework. Among supersymmetric GUT models, the R-parity conserved SUGRA model is one of the most favored models since it provides a natural explanation of the dark matter problem. The minimal SUGRA model (MSUGRA) with R-parity conservation has been investigated extensively because of its predictive power that comes from the fact that it depends on only a few new parameters. Unlike the MSSM, which is hard to be constrained due to its more than 100 new parameters (including 43 CP violating phases), the parameter space of the minimum SUGRA model has 4 parameters and 4 phases. This parameter space has several experimental constraints, i.e. $b\to s+\gamma$, neutron and electron electric dipole moments (EDM), Large Electron-Positron Collider (LEP) bounds and relic density measurements [24,25] etc. In this paper, we examine the observed discrepancy between $S_{\phi K_S}$ and $S_{J/\psi K_S}$ in the context of SUGRA models including MSUGRA and models with nonuniversalities. We consider all relevant experimental constraints in our calculation.

The calculation of $B\to\phi K$ decays involves the evaluation of the matrix elements of related operators in the effective Hamiltonian, which is the most difficult part in this calculation. However, the newly developed QCD improved factorization [Beneke-Bachalla-Neubert-Sachrajda (BBNS) approach] [26] provides a systematic way to calculate the matrix elements of a large class of B decays with significant improvements over the old factorization approach (naive factorization). It allows a QCD calculation of "nonfactorizable" contributions and model independent predictions for strong phases which are important in the theoretical evaluation of the direct CP asymmetries of B decays, e.g. for $B^-\to\phi K^-$, whose experimental result is given in Eq. (3). We adopt the QCD improved factorization in our $B\to\phi K$ calculations. Recently Du and co-workers [27–29] have published an improved calculation of $B\to PV$ decays. We followed here their calculational techniques [27] which are based on the original work [26] of Beneke *et al.* While the BBNS approach is an important advance in calculating B decays, it is not completely model independent. In the BBNS approach the hard gluon (H) and annihilation (A) diagrams (see Fig. 1) contain infrared divergences which are parametrized by an amplitude $\rho_{H,A}$ (with $\rho_{H,A}\leqslant 1$) and a phase $\phi_{H,A}$. (More details can be found in [26] and [27].) If the effects of these terms are small, the theoretical predictions are well defined. However, if these terms are large or dominant, the theory becomes suspect. We will see below that $S_{\phi K_s}$ is essentially independent of the infrared divergent terms, though the branching ratios can become sensitive to ρ_A and ϕ_A.

In this work, we first examine the MSUGRA model which is universal at the GUT scale and then consider nonuniversal terms. Although nonuniversalities may cause serious problems in some flavor changing processes, e.g. K^0-$\bar{K}^0$ mixing and $b\to s+\gamma$, we will show that the off-diagonal terms in the A parameter soft-breaking terms can satisfy all experimental data including the 2.7σ deviation between $S_{\phi K_S}$ and $S_{J/\psi K_S}$. We calculate the BR of $B^-\to\phi K^-$ which is highly correlated with $B^0\to\phi K^0$ and calculate the CP asymmetry of this mode in the allowed parameter space. We also calculate the CP asymmetry of the $b\to s\gamma$ decay mode.

This paper is organized as follows: In Sec. II, we give a brief description of the CP asymmetry of $B^0\to\phi K_S$ and the QCD factorization technique used in this paper. Then we discuss the SUGRA model and its contributions to $B^0\to\phi K_S$ in Sec. III. Some detailed discussions on experimental constraints implemented in our analysis are given in Sec. IV and the standard model predictions are discussed in Sec. V. Section VI is devoted to our results for the MSUGRA model, after which we proceed to the models with nonuniversalities in Sec. VII and we give conclusions in Sec. VIII.

II. *CP* ASYMMETRY OF $B\to\phi K$ DECAYS

The time dependent CP asymmetry of $B\to\phi K_S$ is described by

$$\mathcal{A}_{\Phi K_S}(t)\equiv\frac{\Gamma(\bar{B}^0_{\rm phys}(t)\to\phi K_S)-\Gamma(B^0_{\rm phys}(t)\to\phi K_S)}{\Gamma(\bar{B}^0_{\rm phys}(t)\to\phi K_S)+\Gamma(B^0_{\rm phys}(t)\to\phi K_S)}$$
$$=-C_{\phi K_S}\cos(\Delta m_B t)+S_{\phi K_S}\sin(\Delta m_B t) \tag{4}$$

where $S_{\phi K_S}$ and $C_{\phi K_S}$ are given by

$$S_{\phi K_S}=\frac{2\,{\rm Im}\,\lambda_{\phi K_S}}{1+|\lambda_{\phi K_S}|^2},\qquad C_{\phi K_S}=\frac{1-|\lambda_{\phi K_S}|^2}{1+|\lambda_{\phi K_S}|^2}, \tag{5}$$

and $\lambda_{\phi K_S}$ can be written in terms of decay amplitudes:

$$\lambda_{\phi K_S} = -e^{-2i\beta} \frac{\overline{\mathcal{A}}(\overline{B}^0 \to \phi K_S)}{\mathcal{A}(B^0 \to \phi K_S)}. \tag{6}$$

In our analysis, we find that the SUSY contributions to B_d-$\overline{B}_d$ mixing are small, and so from now on we will use the standard definition for β:

$$\beta \equiv \arg\left(\frac{V_{cd} V_{cb}^\star}{V_{td} V_{tb}^\star} \right). \tag{7}$$

Within the SM, $\sin 2\beta$ can be measured by $S_{J/\psi K_S}$. The current experimental result is given in Eq. (1). Since $B^0 \to J/\psi K_S$ decay is dominated by the SM tree level contribution, we expect that in our analysis the new physics will not affect the SM prediction for $\sin 2\beta$ from $B \to J/\psi K_S$. As a consequence, we further assume that the current SM fit for the CKM matrix will not be affected by models discussed in this paper.

The CP asymmetry of charged $B \to \phi K$ decay is defined as

$$\mathcal{A}_{\phi K^\mp} \equiv \frac{\Gamma(B^- \to \phi K^-) - \Gamma(B^+ \to \phi K^+)}{\Gamma(B^- \to \phi K^-) + \Gamma(B^+ \to \phi K^+)} = \frac{|\lambda_{\phi K^\mp}|^2 - 1}{|\lambda_{\phi K^\mp}|^2 + 1} \tag{8}$$

where

$$\lambda_{\phi K^\mp} = \frac{\overline{A}(B^- \to \phi K^-)}{A(B^+ \to \phi K^+)}. \tag{9}$$

From the above discussion, it is clear that our theoretical predictions for the experimental observables, e.g. $S_{\phi K_S}$, $C_{\phi K_S}$ and $\mathcal{A}_{\phi K^\mp}$, depend on the evaluation of decay amplitudes where the effective Hamiltonian plays an important role. The Effective Hamiltonian for $B \to \phi K$ in the SM is [26]

$$\mathcal{H}_{eff} = \frac{G_F}{\sqrt{2}} \sum_{q=u,c} V_{qb} V_{qs}^* \left[C_1 O_1^q + C_2 O_2^q + \sum_{k=3}^{10} C_k(\mu) O_k(\mu) + C_{7\gamma} O_{7\gamma} + C_{8g} O_{8g} \right] + \text{H.c.} \tag{10}$$

where the Wilson coefficients $C_i(\mu)$ can be obtained by running the renormalization group equation (RGE) from the weak scale down to scale μ. The definitions of the operator O_i's in the SM can be found in [26]. The SUSY contributions will bring in new operators $\tilde{O}_i$'s which can be obtained by changing $L \leftrightarrow R$ in the SM operators. We use $\tilde{C}_i$ to denote the Wilson coefficient of $\tilde{O}_i$. The decay amplitude of $B \to M_1 M_2$ can be expressed in terms of the matrix elements of O_i's, $\langle M_1 M_2 | O_i | B \rangle$. We evaluate these matrix elements in the QCD improved factorization technique. The necessary expressions can be found in [26,27].

Using the above Hamiltonian the amplitude of $B \to \phi K$ is

$$\mathcal{A}(B \to \phi K) = \mathcal{A}^f(B \to \phi K) + \mathcal{A}^a(B \to \phi K) \tag{11}$$

where $\mathcal{A}^f$ are factorized amplitudes which can be written as [27]

$$\mathcal{A}^f(B \to \phi K) = \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \sum_i V_{pb} V_{ps}^* a_i^p \langle \phi K | O_i | B \rangle_f, \tag{12}$$

and $\mathcal{A}^a$ is the weak annihilation decay amplitudes [27]:

$$\mathcal{A}^a(B \to \phi K) = \frac{G_F}{\sqrt{2}} f_B f_\phi f_K \sum V_{pb} V_{ps}^* b_i. \tag{13}$$

The matrix elements $\langle \phi K | O_i | B \rangle_f$ in Eq. (12) are the factorized hadronic matrix elements [30]. a_i's and b_i's contain the Wilson coefficients. Explicit expressions for them, as well as for $\mathcal{A}^a(B \to \phi K)$, can be found in [26] and [27].

In our discussion, the dominant SUSY contributions occur through $O_{7\gamma}$ and O_{8g} and we calculate these new SUSY contributions in the SUGRA framework.

III. SUGRA MODELS

The SUGRA model at the GUT scale can be described by its superpotential and soft-breaking terms:

$$W = Y^U Q H_2 U + Y^D Q H_1 D + Y^L L H_1 E + \mu H_1 H_2$$

$$\mathcal{L}_{\text{soft}} = -\sum_i m_i^2 |\phi_i|^2 - \left(\frac{1}{2} \sum_\alpha m_\alpha \overline{\lambda}_\alpha \lambda_\alpha + B\mu H_1 H_2 + (A^U Q H_2 U + A^D Q H_1 D + A^L L H_1 E) + \text{H.c.} \right). \tag{14}$$

Here Q, L are the left handed quark and lepton doublets, U, D and E are the right handed up, down and lepton singlets and $H_{1,2}$ are the Higgs boson doublets. In the minimal picture, the MSUGRA model contains a universal scalar mass m_0, a universal gaugino mass $m_{1/2}$ and the universal cubic scalar A terms:

$$m_i^2 = m_0^2, \quad m_\alpha = m_{1/2}, \quad A^{U,D,L} = A_0 Y^{U,D,L}. \tag{15}$$

This model contains four free parameters and a sign: m_0, $m_{1/2}$, A_0, $\tan\beta = \langle H_2 \rangle / \langle H_1 \rangle$ and the sign of μ.

However, the parameters $m_{1/2}$, μ and A can be complex and their phases can be $O(1)$. In order to accommodate the experimental bounds on the electron and neutron EDMs without fine tuning phases we extend MSUGRA by allowing the gaugino masses at M_G to have arbitrary phases. This model has been extensively studied in the literature [31–33]. Thus the SUSY parameters with phases at the GUT scale are $m_i = |m_{1/2}| e^{i\phi_i}$ i=1,2,3 [the gaugino masses are for the $U(1)$, $SU(2)$ and $SU(3)$ groups], $A_0 = |A_0| e^{i\alpha_A}$ and $\mu = |\mu| e^{i\phi_\mu}$. However, we can set one of the gaugino phases to zero and we choose $\phi_2 = 0$. Therefore, we are left with four phases. The EDMs of the electron and neutron can now al-low the existence of large phases in the theory [31–33]. In our calculation, we use $O(1)$ phases but calculate the EDMs to make sure that current bounds are satisfied.

We evolve the above parameters from the GUT scale down to the weak scale using full matrix RGEs. Since the $b\rightarrow s$ transition is a generation mixing process, it is necessary to use the full 6×6 matrix form of squark mass matrices in the calculation. We perform the calculation of SUSY contributions without any approximation.

We also include the one loop correction to bottom quark mass from SUSY [34], which is important in the calculation of SUSY contributions to the Wilson coefficients of the operators $O_{7\gamma}$ and O_{8g} and consequently affects the calculations of $B\rightarrow X_s\gamma$ and $B\rightarrow\phi K$ decays. We now discuss the experimental constraints in the next section.

IV. PARAMETER SPACE AND EXPERIMENTAL BOUNDS

In this section we review all the experimental constraints considered in our analysis, and briefly discuss their effects and importance.

A. $B\rightarrow X_s\gamma$

We use a relatively broad range for the branching ratio of $B\rightarrow X_s\gamma$ [35] to take into account the uncertainty in the theoretical calculation of $B\rightarrow X_s\gamma$ $(\pm 0.3\times 10^{-4})$:

$$2.2\times 10^{-4}<\mathrm{BR}(B\rightarrow X_s\gamma)<4.5\times 10^{-4}. \tag{16}$$

The SM prediction for the BR$[B\rightarrow X_s\gamma]$ is very close to the measured value [36], so the $b\rightarrow s$ transition in any new physics is strongly constrained. Since the $B\rightarrow\phi K$ decay also depends on the $b\rightarrow s$ transition, the BR$[B\rightarrow X_s\gamma]$ constraint needs to be implemented in any analysis of the $B\rightarrow\phi K$ decay. Besides the BR of $B\rightarrow X_s\gamma$, we also consider in this work the direct CP asymmetry of $B\rightarrow X_s\gamma$ for the existing of CP violating phases. The experimental measurement from CLEO gives [37]

$$\mathcal{A}_{b\rightarrow s+\gamma}=(-0.079\pm 0.108\pm 0.022)(1.0\pm 0.030) \tag{17}$$

or at 90% confidence level, $-0.27<\mathcal{A}_{b\rightarrow s+\gamma}<+0.10$.

B. The relic density

The recent Wilkinson Microwave Anisotropy Probe (WMAP) result gives [38]

$$\Omega_{CDM}h^2=0.1126^{+0.008}_{-0.009}. \tag{18}$$

We implement this bound at the 2σ level in our calculation: $0.094<\Omega h^2<0.129$. We also notice that when nonuniversal terms are present, new annihilation channels may arise and they are different from the usual MSUGRA $\tilde{\tau}-\tilde{\chi}^0$ coannihilation channel.

C. K^0-$\overline{K^0}$ mixing

It has been shown that K^0-$\overline{K^0}$ mixing can significantly constrain certain flavor changing sources in SUSY models [39]. The current experimental bound for ΔM_K is [40]

$$\Delta M_K=3.490\pm 0.006\times 10^{-12}\ \text{MeV}. \tag{19}$$

In nonuniversal SUGRA models, even if nonuniversal terms between the first two generations are not present at the GUT scale, ΔM_K can become large. This happens because the degeneracy between the first two generations can be broken by other nonuniversal terms via the RGEs. For example, the A_{32} terms in the trilinear coupling matrices A in Eq. (14) can give rise to new contribution to the m^2_{22} term via $\delta m^2_{22}\propto(1/16\pi^2)A_{23}A_{32}\log(M_{GUT}/M_{weak})$. Therefore, it is important to pay attention to ΔM_K even when there is no apparent direct source producing large contributions to ΔM_K.

D. Neutron and electron electric dipole moments

Neutron and electron EDMs can arise in any model with new CP violating phases. In SUSY models, an electron EDM arises from diagrams involving intermediate chargino-sneutrino states and intermediate neutralino-selectron states (for more details, see [31–33]). The current experimental bounds on neutron and electron EDMs are [40]

$$d_n<6.3\times 10^{-26}e\ \text{cm},\quad d_e<0.21\times 10^{-26}e\ \text{cm}. \tag{20}$$

There are other important phenomenological constraints considered, e.g. bounds on masses of SUSY particles and the lightest Higgs boson ($m_h\geqslant 114$ GeV).

V. $B\rightarrow\phi K$ DECAYS IN THE STANDARD MODEL

We first discuss $B\rightarrow\phi K$ decays in the SM. The largest theoretical uncertainties in this calculation come from weak annihilation diagrams which mostly depend on the divergent end-point integrals X_A parametrized in the form [26,27]

$$X_A=(1+\rho_A e^{i\phi_A})\ln\frac{m_B}{\Lambda_h},\quad \Lambda_h=\Lambda_{QCD},\quad \rho_A\leqslant 1. \tag{21}$$

Hard spectator processes contain similar integrals X_H which are parametrized in the same way. However, uncertainties from the hard spectator calculation are much smaller than those from the weak annihilation for this decay, so we will mainly concentrate on the latter. These weak annihilation contributions depend also on the strange quark mass, m_s, through the chirally enhanced factor κ_χ:

$$\kappa_\chi=\frac{2m_K^2}{m_b(m_s+m_q)} \tag{22}$$

where m_q is m_d or m_u.

In Fig. 2 we show the dependence of the branching ratio of $B^-\rightarrow\phi K^-$ mode on ϕ_A and m_s for $\rho_A=1$. Figure 3 shows the dependence of the BR on ρ_A. Similar graphs can be obtained for $B^0\rightarrow\phi K^0$. Since $\mathcal{A}_{\phi K^\mp}$ in the SM [Eq. (8)] is small, we can compare BR$[B^-\rightarrow\phi K^-]$ with the experimental measurement of BR$[B^+\rightarrow\phi K^+]$ given in Eq. (2). Before we discuss the graphs, we first list our parameters: $\rho_H=1$, $\phi_H=-68°$, $f_B=180$ MeV, $f_\phi=233$ MeV, $f_K=160$ MeV and $F^{BK}=0.34$. The CKM matrix elements can be obtained through the Wolfenstein parametrization [41]

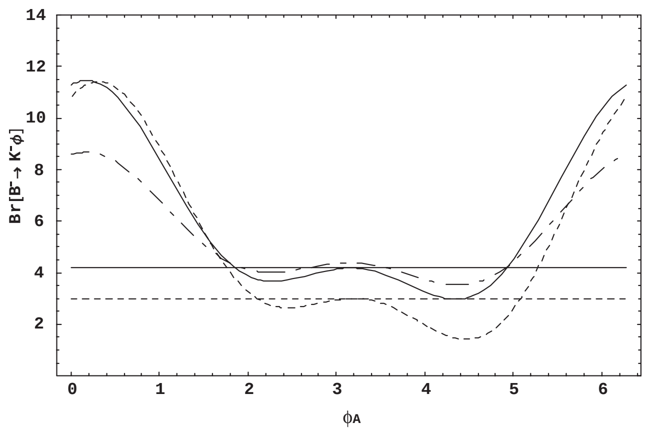

FIG. 2. Branching ratio of $B^- \to \phi K^-$ at $\rho_A = 1$. The solid curve corresponds to $\mu = m_b$, dashed curve for $\mu = 2.5$ GeV with m_s (2 GeV) = 96 MeV and the dot-dashed curve for $\mu = m_b$ with m_s (2 GeV) = 150 MeV. The two straight lines correspond to the cases without weak annihilation.

with $A = 0.819$, $\lambda = 0.2237$, $\rho = 0.224$ and $\eta = 0.324$. The integral $\int_0^1 [\Phi_B(\xi)/\xi] d\xi = m_B/\lambda_B$, where Φ_B is the B meson light-cone distribution amplitude, is parametrized by $\lambda_B = (0.35 \pm 0.15)$ GeV. For $\mu = 2.5$ GeV we use $\lambda_B = 0.2$ GeV, and for $\mu = m_b$ we use $\lambda_B = 0.47$ GeV. In addition, we always use asymptotic forms of the meson light-cone distribution amplitudes [26,27]. If not mentioned, we will use the above parameters in the following calculations.

In both figures, we give results for two different scales and two different m_s values, i.e., $\mu = m_b$ by solid lines [m_s (2 GeV) = 96 MeV] and the dot-dashed lines [m_s (2 GeV) = 150 MeV] and $\mu = 2.5$ GeV by dashed lines [m_s (2 GeV) = 96 MeV]. One can see that the scale dependence is not significant. The straight lines correspond to the branching ratios neglecting the weak annihilation contribution. Comparing Figs. 2 and 3, we see that a large branching ratio comparable to the experimental value is obtained in the region $\rho_A \simeq 1$ and $\phi_A \simeq 0$ (or 2π). However, in this region the weak annihilation diagrams dominate the branching ratio and thus the theory is most suspect. In the remaining part of the parameter space, where the weak annihilation effects are small and the theory is presumably reliable, the SM prediction of the branching ratio is small, about 3σ below the experimental value. We conclude therefore that where the theory is reliable the SM is in significant disagreement with the experimental value of the BR[$B^+ \to \phi K^+$], and in order to obtain a SM value in accord with the experiment one must use parameters where the theory is least reliable. A similar result holds for the BR[$B^0 \to \phi K_s$]. Here theory predicts a branching ratio about 10% smaller than for $B^+ \to \phi K^+$ [in accord with the experimental values of Eq. (2)] but the SM can achieve this only in the region where the weak annihilation processes dominate.

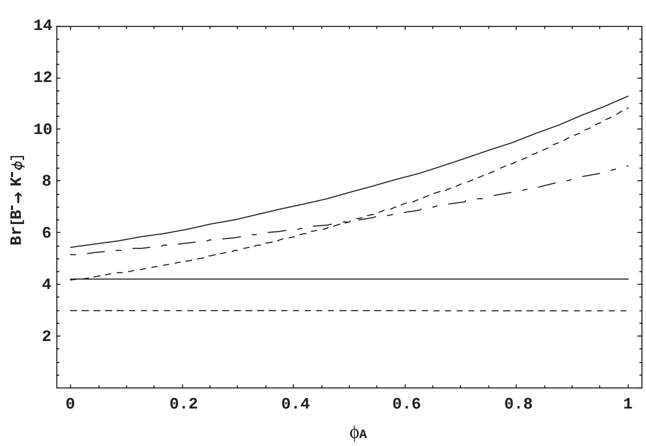

FIG. 3. Branching ratio of $B^- \to \phi K^-$ at $\phi_A = 0$. The solid curve corresponds to $\mu = m_b$, dashed curve for $\mu = 2.5$ GeV with m_s (2 GeV) = 96 MeV and the dot-dashed curve for $\mu = m_b$ with m_s (2 GeV) = 150 MeV. The two straight lines correspond to the cases without weak annihilation.

TABLE I. $S_{\phi K_S}$ with small weak annihilation at $\tan\beta = 10$ and 40 in MSUGRA. The values are the minimum that can be reached subject to all other experimental constraints.

$\tan\beta$	10				40			
$\lvert A_0 \rvert$	800	600	400	0	800	600	400	0
$m_{1/2} = 400$	0.74	0.74	0.73	0.73	0.71	0.70	0.70	0.69
$m_{1/2} = 500$	0.74	0.74	0.74	0.74	0.72	0.72	0.72	0.71

The dot-dashed line, in Fig. 2, corresponds to a larger value of m_s and we see that the BR is very sensitive to m_s only in the large annihilation region. The region with sufficiently large annihilation to accommodate the data decreases as m_s increases, since the annihilation amplitude then decreases, as can be seen from Eq. (22).

VI. MSUGRA

Before we proceed to present our results, let us first mention the values of the parameters used in our calculation of $B \to \phi K$ decay amplitudes in SUGRA models since the BRs depend sensitively on them. We use $\rho_{A,H} = 1$, $\phi_{A,H} = -68°$, $m_s(2\text{ GeV}) = 122.5$ MeV and a CKM fit giving $\sin 2\beta = 0.73$ and $\gamma = 59°$. (If we increase γ, the BR decreases, e.g. for $\gamma = 79°$, the SM BR decreases by $\sim 2\%$.) The SM BR is 4.72×10^{-6} and the weak annihilation contribution is small ($\sim 10\%$).

In Table I, we give the numerical results for two different values of the MSUGRA parameter $\tan\beta$ cases i.e. $\tan\beta = 40$ and 10 and for different $m_{1/2}$ and A_0 in the small weak annihilation case. For simplicity, we set $\alpha_A = \pi$ in the calculation of Table I. We use large phases for other parameters but still satisfy the EDM constrains. For example, for $m_{1/2} = 400$ GeV and $A_0 = 800$ GeV with $\tan\beta = 40$, we find that $\phi_1 = 70°$, $\phi_3 = 33°$, and $\phi_\mu = -13°$ (at the weak scale) satisfy the EDM constraints (for reasons discussed in detail in [33]). The phase α_A has a very small effect on $S_{\phi K_S}$ and this effect becomes smaller as the magnitude of A decreases. Thus a different value of α_A, in the above fit, can change $S_{\phi K_S}$ by $\pm 4\%$ for $A_0 = 800$ GeV. This change is even smaller for smaller A_0, e.g. for $A_0 = 200$ GeV, the change in

$S_{\phi K_S}$ is less than 2%. So far we have not specified m_0 for our results. The values of m_0, for different $m_{1/2}$ and A_0, are chosen such that the relic density constraint is satisfied. We also satisfy the BR[$b\rightarrow s+\gamma$] constraint and the Higgs mass constraint. The $S_{\phi K_S}$ values shown in the table are the minimum that can be reached for given $m_{1/2}$ and A_0.

It can be seen from Table I that the $S_{\phi K_S}$ values in MSUGRA differ only slightly from the SM prediction which is $\sin 2\beta$ evaluated using just the CKM phase. The branching ratios of $B\rightarrow\phi K$ decays also do not differ much from the SM prediction. Even if one went to the large weak annihilation region to accommodate the large branching ratios, $S_{\phi K_S}$ would still be similar to the numbers in Table I. Therefore, MSUGRA cannot explain the large BR and the 2.7σ difference between the $S_{\phi K_S}$ and the $S_{J/\psi K_S}$ experimental results. The reason is that, in MSUGRA, the only flavor violating source is in the CKM matrix, which cannot provide enough flavor violation needed for the $b\rightarrow s$ transition in $B\rightarrow\phi K$ decays. In the next section, we will search for the minimal extension of MSUGRA that can solve both the BR and CP problems of $B\rightarrow\phi K$ decays.

VII. SUGRA MODEL WITH NONUNIVERSAL A TERMS

In the preceding section, we showed that MSUGRA contributions to $B\rightarrow\phi K$ decays are negligible and thus MSUGRA needs to be extended if it is to explain the experimental results of $B\rightarrow\phi K$ decays. It is obvious that some nonuniversal soft breaking terms which can contribute to the $b\rightarrow s$ transition are necessary. There are two ways of enhancing the mixing between the second and the third generation: one can have nonuniversal terms in the squark mass matrices or in the $A^{U,D}$ matrices of Eq. (14). However, in a GUT model, at least the standard model gauge group must hold at M_G and hence the only squark m_{23}^2 that can occur is either left-left or right-right coupling. As discussed in [20], such nonuniversal terms produce only small effects on $B\rightarrow\phi K$ decays. Thus we are led to models with left-right mixing which can occur in the $A^{U,D}$ matrices as the simplest possible nonuniversal term relevant to $B\rightarrow\phi K$ decays. In this work then, we choose a model with nonzero (2,3) elements in the trilinear coupling A terms of Eq. (14) to enhance the left-right mixing of the second and the third generation. The A terms with nonzero (2,3) elements can be written as

$$A^{U,D}=A_0Y^{U,D}+\Delta A^{U,D} \tag{23}$$

where $\Delta A^{U,D}$ are 3×3 complex matrices and $\Delta A_{ij}^{U,D}=|\Delta A_{ij}^{U,D}|e^{\phi_{ij}^{U,D}}$. When $\Delta A^{U,D}=0$, MSUGRA is recovered. For simplicity, we will discuss first the case of non-zero ΔA_{23}^D and nonzero ΔA_{32}^D for $\tan\beta=10$ and 40. In both cases, all other entries in $\Delta A^{U,D}$ are set to zero. The other parameters are same as in the MSUGRA case. We will set the phases such that the EDM constraints are obeyed. At the GUT scale, we use a diagonal Yukawa texture for Y^U, while Y^D is constructed as VY_d^D where V is the CKM matrix and Y_d^D is the diagonalized matrix of the down-type Yukawa. The phenomenological requirements for the Yukawa matrices are that they produce the correct quark masses and the correct CKM matrix. Any other Yukawa texture which satisfies the same requirements can be obtained through unitary rotations. Therefore, our results can be recovered with other Yukawa textures if our A terms are rotated along with the Yukawas.

TABLE II. $S_{\phi K_S}$ at $\tan\beta=10$ with nonzero A_{23}^D and A_{32}^D.

$\lvert A_0\rvert$	800	600	400	0	$\lvert\Delta A_{23(32)}^D\rvert$
$m_{1/2}=300$	-0.50	-0.49	-0.47	-0.43	~ 50
$m_{1/2}=400$	-0.43	-0.40	-0.38	-0.36	~ 110
$m_{1/2}=500$	-0.46	-0.46	-0.44	-0.31	~ 200
$m_{1/2}=600$	-0.15	-0.13	-0.04	0.05	~ 280

In the calculations of decay amplitudes, we will use QCD parameters for the small weak annihilation region (see the last section) where the theory is reliable. In general it is possible that (see [27] for the calculational details of weak annihilation) the new physics can change the behavior of annihilation contributions when the relevant Wilson coefficients can be reduced or increased significantly. However, in our case with nonzero $\Delta A_{23,32}^{U,D}$ terms, the SUSY contribution mainly affects the Wilson coefficients $C_{8g(7\gamma)}$ (possibly also $\tilde{C}_{8g(7\gamma)}$) and these coefficients will not change the annihilation contributions compared to what we have in the SM calculation and thus our previous conclusion about the annihilation terms still holds.

A. Case I: $|\Delta A_{23}^D|=|\Delta A_{32}^D|$ and $\phi_{23}^D\neq\phi_{32}^D$

We show our results for $|\Delta A_{23}^D|=|\Delta A_{32}^D|$ but $\phi_{23}^D\neq\phi_{32}^D$ with $\tan\beta=10$ in Table II. We note that $|\Delta A_{23(32)}^D|$ is an increasing function of $m_{1/2}$. The phases ϕ_{23}^D and ϕ_{32}^D are approximately $-30°$ and $(75\sim115)°$, respectively. The other SUSY phases are $\phi_1\sim22°$, $\phi_3\sim31°$ and $\phi_\mu\sim-11°$. In addition, as mentioned above, the phase of A_0, i.e. α_A, is set to be π.

The BR[$B^-\rightarrow\phi K^-$] is $\sim10\times10^{-6}$ in the parameter space of Table II. We satisfy all other experiment constraints. We see that SUGRA models can explain the large BR and $S_{\phi K_S}$ of the $B\rightarrow\phi K$ decay modes even in the small annihilation region. Comparing with Eq. (1), one sees that the values of $S_{\phi K_S}$ in the table are within 1σ range of the experimental measurement. Reducing the BR[$B^-\rightarrow\phi K^-$] allows one to lower $S_{\phi K_s}$. For example, for $A_0=0$ and $m_{1/2}=600$ GeV, by adjusting ϕ_{32}^D, $S_{\phi K_s}$ can be reduced to -0.05 with BR[$B^-\rightarrow\phi K^-$]$\sim9\times10^{-6}$. In Table III we show the direct CP asymmetries of the $B^-\rightarrow\phi K^-$ decay, i.e. $\mathcal{A}_{\phi K^{\mp}}$, using the same parameters as in Table II. The CP asymmetry is around $-(2\sim3)$% and agrees with the experimental result shown in Eq. (3). This prediction depends on the choice of $\phi_{A,H}$ in Eq. (21). For example, if we use $\phi_{A,H}=28°$, we generate a large $\mathcal{A}_{\phi K^{\mp}}\sim27\%$. We find that there exists a reasonably large range of $\phi_{A,H}$ where we can satisfy the current bound on $\mathcal{A}_{\phi K^{\mp}}$. For example, at $m_{1/2}=300$ GeV and $A_0=800$ GeV where the SUSY contribution

TABLE III. $\mathcal{A}_{b\to s+\gamma}$ (left) and $\mathcal{A}_{\phi K^\mp}$ (right) at $\tan\beta=10$ with nonzero A_{23}^D and A_{32}^D.

$\lvert A_0 \rvert$	800		600		400		0	
$m_{1/2}=300$	1.2%	−3.7%	1.4%	−3.6%	1.7%	−3.6%	2.2%	−3.5%
$m_{1/2}=400$	1.9%	−3.5%	2.0%	−3.4%	2.2%	−3.3%	2.3%	−3.3%
$m_{1/2}=500$	2.6%	−3.5%	2.6%	−3.6%	2.5%	−3.5%	2.4%	−3.2%
$m_{1/2}=600$	2.0%	−2.8%	2.1%	−2.7%	2.1%	−2.5%	2.2%	−2.2%

is the largest, we find that $\mathcal{A}_{\phi K^\mp}$ varies from 9% to -4% when $\phi_{A,H}$ varies from $-100°$ to $-50°$ (for simplicity, we set $\phi_A=\phi_H$). In addition, since the annihilation contribution is small in that range, we find that the branching ratio is around $(9.5\sim 11)\times 10^{-6}$. The CP asymmetry of $b\to s\gamma$ is $\sim 1-3\%$. The present experimental errors for $C_{\phi K_S}$ are still large. For this model, $C_{\phi K_S}\sim -\mathcal{A}_{\phi K^\mp}$, which may be tested by future data.

In Tables IV and V, we give our results for $\tan\beta=40$ with nonzero $\Delta A_{23(32)}^D$. The phases ϕ_{23}^D and ϕ_{32}^D are $-(70\sim 0)°$ and $(80\sim 110)°$, respectively. $\phi_1\sim(25\sim 60)°$, $\phi_3\sim 25°$ and $\phi_\mu\sim -8°$. The off-diagonal elements $|\Delta A_{23(32)}^D|$ vary from 90 GeV to 250 GeV as $m_{1/2}$ increases. We compare Table IV with the results for $\tan\beta=10$ shown in Table II and we see that only low $m_{1/2}$ can satisfy experimental data for $\tan\beta=40$. The most important reason for this is that left-right mixing of the second and the third generation decreases significantly with increasing $\tan\beta$. This comes about as follows. The RGE running of $A_{23(32)}^D$ is not sensitive to $\tan\beta$. Therefore, for the same size of $A_{23(32)}^D$ input at the GUT scale, the weak scale values of $A_{23(32)}^D$ do not differ much for different $\tan\beta$. However, the A^D term enters into the down squark matrix after electroweak symmetry breaking when H_1 [see Eq. (14)] grows a vacuum expectation value proportional to $\cos\beta$. Hence the left-right mixing between the second and the third generation in the down squark matrix will be smaller for large $\tan\beta$. For low $m_{1/2}$ this reduction can be compensated by increasing the magnitude of $A_{23(32)}^D$. For example, for $m_{1/2}=300$ GeV, we use $|A_{23(32)}^D|\sim 90$ GeV in this case compared to 50 GeV in the case of $\tan\beta=10$ (see Table II). The chargino diagram contribution increases with $\tan\beta$ and can help to generate a large BR. But for large $m_{1/2}$, when the chargino contribution goes down, $|A_{23(32)}^D|$ must become much larger. However, as $A_{23(32)}$ increases, the pseudoscalar Higgs boson mass becomes small at the same time (but μ does not get smaller), which prevents $|A_{23(32)}^D|$ from having an unlimited increase. For example, for $m_{1/2}=600$ GeV and $A_0=800$ GeV, $|A_{23(32)}^D|=250$ GeV generates $m_A=580$ GeV which is still allowed for the dark matter constraint to be satisfied in the $\tilde{\tau}\leftrightarrow\tilde{\chi}^0$ co-annihilation channel. If $|A_{23(32)}^D|$ is increased more, the pseudoscalar mass gets smaller and the dark matter constraint can still be satisfied due to the available $\chi_1^0\chi_1^0\to A\to f\bar{f}$ channel. But with a further reduction of the pseudoscalar mass by increasing $|A_{23(32)}^D|$ further, this channel goes away when $m_A<2m_{\tilde{\chi}^0}$ and we must again satisfy the relic density using the stau-neutralino co-annihilation channel. However, the improvement of $S_{\phi K_s}$ in this scenario is small. For example, for the point mentioned above, $|A_{23(32)}^D|$ can be increased to around 480 GeV with relic density in the $\tilde{\tau}\leftrightarrow\tilde{\chi}^0$ channel but $S_{\phi K_S}$ can only be reduced from 0.37 (see Table IV) to 0.22 with the same branching ratio. Thus, the $S_{\phi K_S}$ and the branching ratio still cannot be satisfied.

If we use $\phi_{23}^D=\phi_{32}^D$ (equal phases) we have one less parameter, but that choice will not be able to satisfy experimental results. The reason is that the weak phase from the gluino contributions in the Wilson coefficients C_{8g} and the weak phase from $\tilde{C}_{8g}$ will cancel when $\phi_{23}^D=\phi_{32}^D$ because C_{8g} depends on A_{23}^D but $\tilde{C}_{8g}$ depends on $(A_{32}^D)^*$. For example, for $\tan\beta=10$ we find that $S_{\phi K_S}\sim 0.7$ since the gluino contribution dominates at lower $\tan\beta$. At $\tan\beta=40$, $S_{\phi K_S}$ can reach 0.45 since the chargino contribution is larger at higher $\tan\beta$, but this is not enough to satisfy the data.

B. Case II: $|\Delta A_{23}^U|=|\Delta A_{32}^U|$ and $\phi_{23}^U=\phi_{32}^U$

In this section we discuss the case $\Delta A_{23(32)}^D=0$ but $\Delta A_{23(32)}^U\neq 0$. The phases used are similar to those used in first two cases except $\phi_{23}^U=\phi_{32}^U$. This case is more complicated than the $A_{23(32)}^D\neq 0$ case. We find that it is easier to start by comparing them.

The first important change is that the ΔA_{32}^U contribution is much smaller than the ΔA_{32}^D contribution to the mixing between the second and the third generation in the down squark

TABLE IV. $S_{\phi K_S}$ (left) and BR$[B^-\to\phi K^-]\times 10^6$ (right) at $\tan\beta=40$ with nonzero ΔA_{23}^D and ΔA_{32}^D.

$\lvert A_0 \rvert$	800		600		400		0	
$m_{1/2}=300$	−0.40	10.0	−0.38	10.0	−0.33	10.1	−0.05	10.0
$m_{1/2}=400$	−0.11	8.0	−0.05	8.0	0.04	7.9	0.28	8.0
$m_{1/2}=500$	0.07	6.0	0.16	6.1	0.24	6.1	0.37	6.2
$m_{1/2}=600$	0.37	6.2	0.44	6.2	0.49	6.2	0.58	6.2

TABLE V. $\mathcal{A}_{b\to s+\gamma}$ (left) and $\mathcal{A}_{\phi K^{\mp}}$ (right) at $\tan\beta=40$ with nonzero A_{23}^D and A_{32}^D.

$\|A_0\|$	800		600		400		0	
$m_{1/2}=300$	-6.3%	-3.5%	-5.6%	-3.4%	-5.2%	-3.3%	-3.6%	-2.6%
$m_{1/2}=400$	-3.0%	-3.0%	-2.1%	-2.9%	-1.7%	-2.6%	-0.7%	-1.8%
$m_{1/2}=500$	-0.5%	-2.9%	-0.4%	-2.5%	-0.2%	-2.2%	0.2%	-1.7%
$m_{1/2}=600$	0.2%	-1.7%	0.3%	-1.4%	0.4%	-1.2%	0.6%	-0.8%

mass matrix due to the suppression by the second generation Yukawa coupling in the RGE of A_{32}^D. (Thus our choice of $\phi_{23}^U=\phi_{32}^U$ has no loss of generality.) Consequently, the size of the Wilson coefficient $\tilde{C}_{8g}$ is significantly reduced. Although ΔA_{23}^U still contributes, that contribution is also reduced (compared to ΔA_{23}^D) due to the RGE. Therefore, compared with the first case the total SUSY contributions are reduced especially for $\tan\beta=10$ and thus it becomes harder to fit the experimental results.

Another important change is the roles of some experimental constraints which are not important in the first case in the sense that they do not prevent the SUSY contributions from increasing, or at least their limits are not reached when we have solutions satisfying the *B*-decay data. Below are some comments concerning this.

(1) For $\tan\beta=40$ and low $m_{1/2}$, i.e. 300 GeV, the BR[$B\to X_s\gamma$] will constrain the size of $\Delta A_{23(32)}^U$. This is why the $S_{\phi K_S}$ and the branching ratio fit given in Table VI is not as good as the corresponding one shown in Table II for the $A_{23(32)}^D\neq 0$ case.

(2) When $m_{1/2}$ increases, the size of $\Delta A_{23(32)}^U$ also needs to be increased. But three other additional constraints are present, i.e., ΔM_K and ϵ_K from the K^0-$\overline{K}^0$ mixing and the mass of smallest up squarks (right-handed stop) $m_{\tilde{t}_R}$. For example, for $m_{1/2}=500$ and $A_0=600$ (and $m_0=431$ GeV by the relic density constraint) we get $m_{\tilde{g}}\sim m_{\tilde{q}}\sim 1000$ GeV (where $m_{\tilde{q}}$ is the average squark mass and $m_{\tilde{g}}$ is the mass of the gluino, see [39] for more details) and we find that $\sqrt{|\mathrm{Re}(\delta_{12}^d)_{LL}^2|}=7.1\times 10^{-2}$ and $\sqrt{|\mathrm{Im}(\delta_{12}^d)_{LL}^2|}=9.7\times 10^{-3}$ which are allowed by the experimental bounds on ΔM_K and ϵ_K [39] (the sizes of $(\delta_{12}^d)_{LR}$, $(\delta_{12}^d)_{RL}$ and $(\delta_{12}^d)_{RR}$ are around $10^{-8}\sim 10^{-7}$ and thus these bounds can be safely ignored in our case).

(3) The situation for the right-handed stop mass is similar to the pseudoscalar Higgs boson case we mentioned at the end of case I. We use the $\tilde{\tau}\leftrightarrow\tilde{\chi}^0$ channel to satisfy the dark matter constraints. Although it is possible to use a larger $A_{23(32)}^U$ which consequently reduces $m_{\tilde{t}_R}$ more and then opens the $\tilde{t}_R\leftrightarrow\tilde{\chi}^0$ channel, the room is small due to the smallness of $m_{\tilde{\chi}^0}$. In addition, the M_K and the ϵ_K bounds become harder to satisfy when $m_{\tilde{t}_R}$ is small. Therefore, as in the case where the pseudoscalar Higgs boson mass becomes small, possible improvements can not satisfy the experimental results of both $S_{\phi K_S}$ and BR[$B^-\to\phi K^-$].

A further difference is the $\tan\beta$ dependence. In case I, as was discussed above, gluino contributions depend inversely on $\tan\beta$ due to the way that $A_{23(32)}^D$ enters into the down squark mass matrix. But in this case, the gluino contributions are reduced significantly and the chargino plays a more important role, which will be enhanced by $\tan\beta$. Therefore, in this case, we see that larger $\tan\beta$ can satisfy the experimental results at small $m_{1/2}$, but small $\tan\beta$ cannot and that is why we have given results in Table VII only for two values of $m_{1/2}$ at $\tan\beta=10$ since higher $m_{1/2}$ cannot improve the situation.

We also comment concerning A_0 and its phase. So far, we have used the phase π for A_0. We find that using a different phase will not improve the results greatly. In general, the improvement is at a few percent level. (This holds also for case I.) For example, in case II, for $\tan\beta=40$, $m_{1/2}=400$ and $A_0=800$, we find that using $\alpha_A\sim -95°$ can improve $S_{\phi K_S}$ from -0.04 to -0.06.

Finally we note that the values of $\mathcal{A}_{b\to s+\gamma}$ and $\mathcal{A}_{\phi K^{\mp}}$ remain small, i.e. $\mathcal{A}_{b\to s+\gamma}$ and $\mathcal{A}_{\phi K^{\mp}}$ are $-(3\sim 0)\%$ and $-(3\sim 1)\%$ at $\tan\beta=10$, and $-(5\sim 0)\%$ and $-(3\sim 1)\%$ at $\tan\beta=40$.

VIII. CONCLUSION

In this paper, we have probed the $B\to\phi K$ decays in SUGRA models with *CP* violating phases to explain the discrepancy between the experimental measurements and the SM predictions of the *CP* asymmetry of $B^0\to\phi K_S$ and the branching ratios of the $B\to\phi K$ decays. We have calculated the *CP* asymmetries of $B^-\to\phi K^-$ and $B\to X_s\gamma$. In our analysis, we implemented all relevant experimental con-

TABLE VI. $S_{\phi K_S}$ (left) and BR[$B^-\to\phi K^-$]$\times 10^6$ (right) at $\tan\beta=40$ with nonzero ΔA_{23}^U and ΔA_{32}^U.

$\|A_0\|$	800		600		400		0		$\|\Delta A_{23(32)}^U\|$ (GeV)
$m_{1/2}=300$	0.03	8.4	0.04	9.0	0.01	8.0	0.17	8.0	~ 300
$m_{1/2}=400$	-0.07	8.5	-0.03	8.4	0	7.1	0.32	6.3	~ 600
$m_{1/2}=500$	0	6.5	0.07	6.4	0.18	6.0	0.44	6.1	~ 800
$m_{1/2}=600$	0.27	6.1	0.30	6.1	0.35	6.1	0.51	5.9	~ 1000

TABLE VII. $S_{\phi K_S}$ (left) and BR$[B^- \to \phi K^-] \times 10^6$ (right) at $\tan\beta = 10$ with nonzero ΔA_{23}^U and ΔA_{32}^U.

$\lvert A_0\rvert$	800		600		400		0		$\lvert\Delta A_{23(32)}^U\rvert$ (GeV)
$m_{1/2}=300$	0.17	6.5	0.16	6.3	0.32	6.1	0.60	5.2	~ 300
$m_{1/2}=400$	0.37	4.7	0.39	4.6	0.46	4.3	0.62	4.3	~ 550

straints, e.g. BR$[B \to X_s \gamma]$, relic density, $K^0 - \overline{K}^0$ mixing parameters and electron and neutron EDMs. We used the improved QCD factorization method [26,27] for the calculation of decay amplitudes.

The SM not only cannot explain the CP asymmetry of $B^0 \to \phi K_S$, it also fails to satisfy the BR$[B \to \phi K]$ data barring the region of large weak annihilation where the theory is most suspect. We then studied the MSUGRA model and found that it also has the same problem. Therefore, if the current experimental results continue to hold in the future, it will signal the first significant breakdown of the standard model and also of MSUGRA. This conclusion is important in the sense that one needs to construct a more complicated SUGRA model to satisfy experimental measurements which will provide important guidance to our future research on SUSY models and their signals at the accelerator experiments.

We have considered the extension of the MSUGRA model by adding nonuniversal A terms. For a GUT theory, the only natural choice is to have a left-right mixing between the second and the third generation in the up or down quark sectors i.e. $\Delta A_{23}^{U,D}$ and $\Delta A_{32}^{U,D}$ terms. We have examined thoroughly several different possibilities in this extension and their theoretical predictions and have found a large region of parameter space where all experimental results can be satisfied, including the CP asymmetries and branching ratios of the $B \to \phi K$ decays. This result is obtained without resorting to large weak annihilation amplitudes and so is based on reliable calculations of hadronic decays, and thus provides useful hints for the study of hadronic B decays. Further the size of $\Delta A_{23}^{U,D}$ needed is the same as the other soft breaking terms, and so is not anomalously small or large. Thus, this study also can provide important phenomenological information not only for accelerator physics but also for building models at the GUT scale and for exploring physics beyond it. In this connection, models based on Horava-Witten M theory can naturally give rise to nonzero values of ΔA_{23}. In previous work [42] it was shown that it was possible to construct a three generation model with $SU(5)$ symmetry using a nonstandard embedding based on a torus fibered Calabi-Yau threefold with a del Pezzo base dP_7. The model allowed Wilson line breaking to the standard model at M_G, and also had vanishing instanton charges on the physical orbifold plane. If in addition one assumed that the 5-branes in the bulk clustered around the distant plane, one could explain without undue fine tuning the general structure of the quark and lepton mass hierarchies and obtain the large mixing angle (LMA) solution for neutrino oscillations [42–44]. One can show that the model naturally gives rise to ΔA_{23} nonuniversal terms as required by the current B-factory data, while keeping the squark mass matrices essentially universal. This model will be discussed elsewhere [45].

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