CP violation in $B\to\phi K_S$ decay at large $\tan\beta$

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We consider the chargino contribution to CP violation in $B\to\phi K_S$ decay in the minimal supersymmetric standard model at large $\tan\beta$. It is shown that the Wilson coefficient C_{8g} of the chromomagnetic penguin operator can be significantly enhanced by the chargino-mediated diagrams while satisfying other direct or indirect experimental constraints. The enhanced C_{8g} allows a large deviation in the CP asymmetry from the standard model prediction, especially it can explain the apparent anomaly reported by BaBar and Belle.

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The measurement of CP asymmetries at B factories is a powerful probe of new physics (NP). In the standard model (SM), the origin of CP violation is the phase $\delta_{\rm CKM}$ in the 3×3 Cabibbo-Kobayashi-Maskawa (CKM) matrix of quark mixing. In general, there can be many new sources of CP violation in new physics models beyond the SM.

The time-dependent CP asymmetry in neutral B decays to CP eigenstates $B\to f_{CP}$ gives information on the two classes of CP violation C_f and S_f [1]:

$$A_{CP}(t)=\frac{\Gamma(\overline{B}(t)\to f_{CP})-\Gamma(B(t)\to f_{CP})}{\Gamma(\overline{B}(t)\to f_{CP})+\Gamma(B(t)\to f_{CP})}$$
$$=-C_f\cos(\Delta m_B t)+S_f\sin(\Delta m_B t), \qquad (1)$$

where Δm_B is the mass difference of the B system and $B(\overline{B})(t)$ is the state at time t which was $B(\overline{B})$ at $t=0$. The CP asymmetries C_f and S_f are determined by

$$\lambda_{CP}\equiv e^{-2i(\beta+\theta_d)}\frac{\overline{A}}{A}, \qquad (2)$$

where $\beta(\theta_d)$ is the contribution of SM (NP) to the phase in the B-$\overline{B}$ mixing and $\overline{A}(A)$ is the decay amplitude for $\overline{B}(B)\to f_{CP}$.

The time-dependent CP asymmetry in $B\to J/\psi K$ decay measured at BaBar [2] and Belle [3] shows the existence of CP violation in the B-meson system, and the current world average [1]

$$\sin2\beta_{J/\psi K_S}\equiv S_{J/\psi K_S}=0.734\pm0.054 \qquad (3)$$

is fully consistent with the CKM picture of CP violation in the SM. In the SM, the CP asymmetry in $B\to\phi K_S$ also measures the same angle β because CP violation in the decay amplitude is suppressed by the Cabibbo angle λ (≈0.22): $S_{\phi K_S}=S_{J/\psi K_S}+\mathcal{O}(\lambda^2)$.

Recently, BaBar [4] and Belle [5] have announced the first measurement of time-dependent CP asymmetries in $B\to\phi K_S$. The weighted average of the mixing-induced CP asymmetry,

$$\sin2\beta_{\phi K_S}\equiv S_{\phi K_S}=-0.39\pm0.41, \qquad (4)$$

differs from the SM prediction of Eq. (3) by 2.7σ. The Belle Collaboration has also reported the direct CP asymmetry

$$C_{\phi K}=0.56\pm0.43, \qquad (5)$$

which is consistent with zero as predicted in the SM.

Contrary to the $B\to J/\psi K_S$ decay where $\overline{A}/A=1$ to a good approximation even in the presence of NP, the loop-induced process $B\to\phi K_S$ generally allows large deviation from 1 in phase and modulus of $\overline{A}/A$. There are already many works which explain the apparent deviation (4) in various NP models [6–10].

The measured $\sin2\beta$'s from $B\to\eta' K_S$, $B\to(K^+K^-)K_S$ [5] and $\sin2\alpha$ from $B\to\pi^+\pi^-$ [11] are consistent with the SM although the errors are large. We note that, however, $B\to\phi K_S$ is unique in that it does not have the tree-level amplitude in the SM unlike most other B decays. Therefore, it is not unlikely that the NP manifests itself only in the $B\to\phi K_S$ decay [12].

The minimal supersymmetric (SUSY) standard model (MSSM) has many new CP violating phases besides the CKM phase of the SM. With CP violating phases of order one the electric dipole moments (EDMs) easily exceed the experimental upper bounds by several orders of magnitude. In addition, the general structure of sfermion mass matrices in the generation space leads unacceptable flavor changing neutral currents (FCNC) by gluino mediation. These SUSY CP and SUSY FCNC problems strongly constrain the MSSM parameters.

In the MSSM it has been shown that Eq. (4) can be accomodated if there exist new flavor structures in the up- or down-type squark mass matrices [6,10,13]. In this paper we will show that the chargino contribution can generate large deviation in $S_{\phi K}$ at large $\tan\beta$, even if CKM is the only source of flavor mixing.

Specifically, we adopt a decoupling scenario where the masses of the first two generation scalar fermions are very heavy [$\gtrsim\mathcal{O}(10\ \text{TeV})$], so that the SUSY FCNC and SUSY CP problems are solved without a naturalness problem [14]. It is known that if this pattern of soft terms are generated at high energy scale (such as the Planck scale), two-loop renormalization group (RG) evolution drives the scalar top quark mass squared negative, breaking the color and charge [15]. However, it was shown that the problem can be solved by introducing extra matters to make the sfermion masses RG invariant at the two-loop level [16]. We can also simply assume this scenario as a phenomenological solution to SUSY FCNC and SUSY CP problems, where the soft terms with little mixing between the first two generation squarks are generated at low energy scale by some unknown mechanism.

We also assume the flavor-changing off-diagonal elements of scalar fermions are vanishing to guarantee the absence of the gluino mediated FCNC. In this case the CKM matrix is the only source of flavor mixing, while there are new CP violating parameters μ, M_2 in the chargino matrix

$$M_C=\begin{pmatrix} M_2 & \sqrt{2}m_W \sin\beta \\ \sqrt{2}m_W\cos\beta & \mu \end{pmatrix} \qquad (6)$$

and A_t in the scalar top mass-squared matrix

$$M_{\tilde{t}}^2=\begin{pmatrix} m_{\tilde{Q}}^2+m_t^2+D_L & m_t(A_t^*-\mu\cot\beta) \\ m_t(A_t-\mu^*\cot\beta) & m_{\tilde{t}}^2+m_t^2+D_R \end{pmatrix}, \qquad (7)$$

where $D_L=(1/2-2/3\sin^2\theta_W)\cos 2\beta m_Z^2$ and $D_R=2/3\sin^2\theta_W\cos 2\beta m_Z^2$.

In this decoupling scenario, it has been shown that the light top squark and chargino contributions still allows large direct CP asymmetry in the radiative B decay up to $\pm 16\%$ [17]. It should be noted that the new contribution to B-$\overline{B}$ mixing is very small and this model is naturally consistent with Eq. (3) [17].

The decay $B\to\phi K_S$ is described by the $\Delta B=1$ effective Hamiltonian [18]. The chargino contribution to the Wilson coefficients of magnetic penguin operators is given by

$$\begin{aligned} C_{7\gamma}(\tilde{\chi}^{\pm})=&\sum_{I=1}^{2}\sum_{k=1}^{2}\frac{m_W^2}{m_{\tilde{t}_k}^2}\Bigg\{(\chi_I^{d_L})_{A2}^*(\chi_I^{d_L})_{A3} \\ &\times\left(I_1(x_{Ik})+\frac{2}{3}J_1(x_{Ik})\right)+(\chi_I^{d_L})_{A2}^*(\chi_I^{d_R})_{A3} \\ &\times\frac{m_{\tilde{\chi}_I^-}}{m_b}\left(I_2(x_{Ik})+\frac{2}{3}J_2(x_{Ik})\right)\Bigg\}, \\ C_{8g}(\tilde{\chi}^{\pm})=&\sum_{I=1}^{2}\sum_{k=1}^{2}\frac{m_W^2}{m_{\tilde{t}_k}^2}\Bigg\{(\chi_I^{d_L})_{k2}^*(\chi_I^{d_L})_{k3}J_1(x_{Ik}) \\ &+(\chi_I^{d_L})_{k2}^*(\chi_I^{d_R})_{k3}\frac{m_{\tilde{\chi}_I^-}}{m_b}J_2(x_{Ik})\Bigg\}, \end{aligned} \qquad (8)$$

where $x_{Ik}=m_{\tilde{\chi}_I^{\pm}}^2/m_{\tilde{t}_k}^2$ and $I_{1,2}(x), J_{1,2}(x)$ are the loop functions [19]. $(\chi_I^{d_{L(R)}})_{kq}$ is the top-squark(k)–chargino(I)–down-quark$(q_{L(R)})$ coupling:

$$(\chi_I^{d_L})_{kq}=-V_{I1}^*S_{\tilde{t}_k\tilde{t}_L}+V_{I2}^*S_{\tilde{t}_k\tilde{t}_R}\frac{m_t}{\sqrt{2}m_W\sin_\beta},$$

$$(\chi_I^{d_R})_{kq}=U_{I2}S_{\tilde{t}_k\tilde{t}_L}\frac{m_q}{\sqrt{2}m_W\cos\beta}, \qquad (9)$$

where $U,V(S)$ diagonalize(s) the chargino (top squark) mass matrix. In the analysis we have decoupled the charged Higgs contribution by assuming $m_{H^\pm}\gtrsim 1$ TeV to evade the very stringent two-loop EDM constraints at large $\tan\beta$ and relatively light Higgs pseudoscalar [24]. The effect of $H^\pm$ will be mentioned below.

We note that the $C_{7\gamma(8g)}^{\tilde{\chi}^\pm}$ has the enhancement factors $m_{\tilde{\chi}_I^\pm}/m_b$ by the chirality flip inside the loop and can dominate the SM contribution. On the other hand the Wilson coefficients of QCD penguin operators, $C_{3,\ldots,6}$, preserve the chirality and do not have such enhancement factors. In addition, due to the super-Glashow-Iliopoulos-Maiani (GIM) mechanism, the chargino contribution to $C_{3,\ldots,6}$ are much smaller than the SM values. The contribution to chirality flipped operators $C_{7\gamma(8g)}^{\prime\tilde{\chi}^\pm}$ are suppressed by m_s/m_b. We neglect the contribution of electroweak penguin operators which are also expected to be negligible in our scenario. Therefore the large deviation in $\overline{A}/A$ should be generated soley by C_{8g} in this scenario.

To calculate the hadronic matrix elements we use the QCD factorization method in Ref. [18]. In this approach it has been demonstrated that the strong phases of QCD penguin operators cancel out in the SM [7]. Since the NP contribution to $C_{3,\ldots,6}$ are negligible in our scenario, the strong phase is small as in the SM. The required $C_{8g}^{\tilde{\chi}^\pm}(m_b)$ [$C_{8g}^{\rm SM}(m_b)\approx -0.147$] for large deviation of CP asymmetries can be estimated from the approximate numerical expression for $\overline{A}$ [18,7]:

$$\begin{aligned}\overline{A}\propto&\sum_{p=u,c}V_{ps}^*V_{pb}(a_3+a_4^p+a_5)\\ &\approx -3.9\times 10^{-4}(3.7e^{0.21i}+4.5C_{8g}).\end{aligned} \qquad (10)$$

In the above expression the largest errors which are about 20% come from the decay constants and form factors. These uncertainties are largely canceled in the asymmetries in Eq. (1).

In Eq. (10) and the following analysis, we did not include weak annihilation contributions which are power suppressed but may be numerically important. These contributions are not derived in QCD factorization on first principles, and additional phenomenological parameter which can contain strong phase should be introduced to estimate the divergent integral at end point. The CP asymmetries $C_{\phi K}$, $S_{\phi K}$ can be increased depending on the size and phase of the parameter. Therefore our analysis is conservative and the full analysis, including the annihilation contributions, will be given in [12].

In Fig. 1 we show $S_{\phi K}$ as a function of arg(C_{8g}) for $|C_{8g}|=0.33,0.65$, and 1.0. From this figure we can see $|C_{8g}|\approx 0.33-0.65$ with a large positive phase can accomodate the deviation within 1σ.

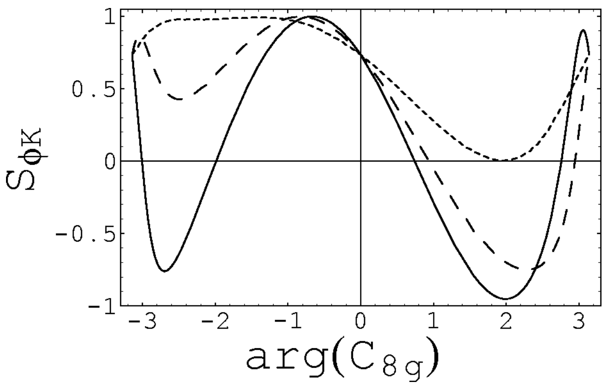

FIG. 1. $S_{\phi K}$ as a function of $\arg(C_{8g})$ for $|C_{8g}|=0.33$ (short dashed line), 0.65 (long dashed line), and 1.0 (solid line).

Since the chargino contributions to $C_{7\gamma}$ and C_{8g} have similar structures, one can think that large deviation in C_{8g} may result in the large deviation $C_{7\gamma}$. Too large deviation in $C_{7\gamma}$ will violate the measurement of $B(B\to X_s\gamma)$, for which we take $2\times10^{-4}<B(B\to X_s\gamma)<4.5\times10^{-4}$ [20], because it is already consistent with the SM prediction.

Due to the different loop functions, however, the two Wilson coefficients are not strongly correlated, and it is possible to have large deviation in C_{8g} while keeping $|C_{7\gamma}|$ constrained to satisfy $B(B\to X_s\gamma)$. Since χ^R in Eq. (9) is proportional to $\tan\beta$ for large $\tan\beta$, we need relatively large $\tan\beta$ to have sizable effects.

The neutral Higgs boson h^0 which is lighter than the Z boson at the tree level has large radiative corrections [21] which can be approximated at large $\tan\beta$ as

$$m_{h^0}^2 \approx m_Z^2 + \frac{3}{2\pi^2}\frac{m_t^4}{v^2}\log\left(\frac{m_{\tilde t_1}m_{\tilde t_2}}{m_t^2}\right), \tag{11}$$

where $v\approx 246$ GeV. The CERN e^+e^- collider LEP II sets the lower limit on the SM-like Higgs boson mass: $m_{h^0}\gtrsim 114.3$ (GeV) at 95% C.L. [22]. This limit gives very strong constraint on the top squark masses.

At large $\tan\beta$ ($\sim\mathcal{O}(50)$), the SUSY QCD and SUSY electroweak correction to the nonholomorphic couplings $H_u^* D^c Q$ become very important. This gives a large correction on the down-type quark masses and some CKM matrix elements [23]

$$m_b = \frac{\overline{m}_b}{1+\tilde\epsilon_3\tan\beta}, \qquad V_{JI} = V_{JI}^{\rm eff}\left[\frac{1+\tilde\epsilon_3\tan\beta}{1+\epsilon_0\tan\beta}\right] \tag{12}$$

where $\tilde\epsilon_3 \approx \epsilon_0 + \epsilon_Y y_t^2$ and $(JI)=(13)(23)(31)(32)$. $\overline{m}_b$, $V_{JI}^{\rm eff}$ are b-quark mass and CKM elements measured at experiments, respectively. At one-loop and $SU(2)_L\times U(1)_Y$ symmetric limit, the ϵ_0 and ϵ_Y are given by

$$\epsilon_0 = \frac{2\alpha_s}{3\pi}\mathrm{Re}\left(\frac{\mu^*}{m_{\tilde g}}\right) j(y_{\tilde b\tilde g}, y_{\tilde Q\tilde g}),$$

$$\epsilon_Y = \frac{1}{16\pi^2}\mathrm{Re}\left(\frac{A_t}{\mu}\right) j(y_{\tilde Q\mu}, y_{\tilde t\mu}), \tag{13}$$

where $y_{\tilde b\tilde g} = m_{\tilde b}^2/|m_{\tilde g}|^2$, $y_{\tilde Q\tilde g} = m_{\tilde Q}^2/|m_{\tilde g}|^2$, $y_{\tilde Q\mu} = m_{\tilde Q}^2/|\mu|^2$, $y_{\tilde t\mu} = m_{\tilde t}^2/|\mu|^2$. The loop function is given by $j(x,y)=[j(x)-j(y)]/(x-y)$ with $j(x)=x\log x/(x-1)$. Note that these SUSY threshold corrections are not easily decoupled even for very heavy superparticles. For degenerate SUSY parameters, $\epsilon_0\approx 0.013$ and $\epsilon_Y\approx 0.003$.

For the numerical analysis we fix the CKM matrix by using $|V_{us}^{\rm eff}|=0.2196$, $|V_{cb}^{\rm eff}|=4.12\times10^{-2}$, and $|V_{ub}^{\rm eff}|=3.6\times10^{-3}$ [22], leaving only $\delta_{\rm CKM}$ as a free parameter. For

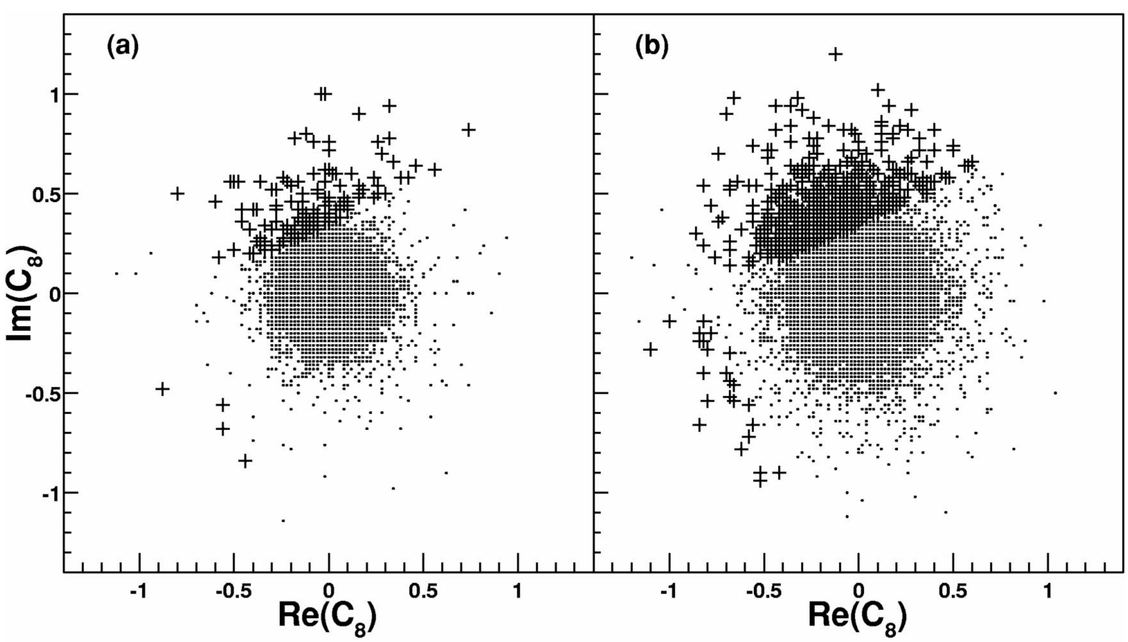

FIG. 2. Distribution of C_8 values for $\tan\beta=35$ (a) and $\tan\beta=60$ (b) in the complex plane. See the text for other parameters. The points with $S_{\phi K}<0$ are marked with a $+$ symbol.

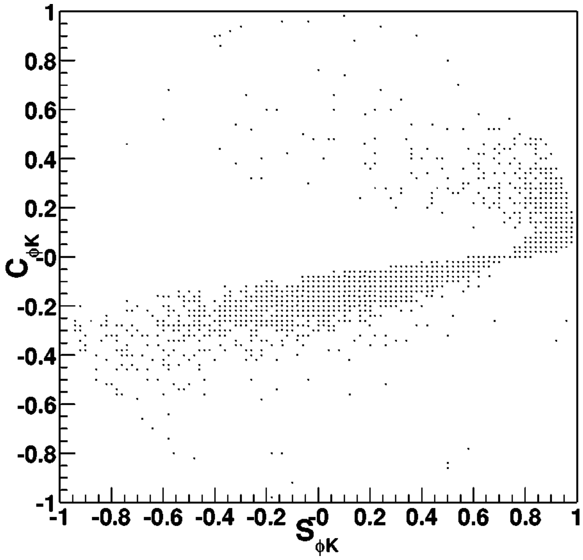

FIG. 3. Correlation between $C_{\phi K}$ and $S_{\phi K}$ for $\tan\beta = 60$.

example, for $\delta_{\rm CKM} = \pi/3$, we get $\Delta m_{B_d} = 0.491\ {\rm ps}^{-1}$ and $\sin 2\beta = 0.729$, which are close to the experimental central values [22]. The free SUSY parameters in our scenario are $\tan\beta$, M_2, μ, $m_{\tilde{Q}}$, $m_{\tilde{t}}$, and A_t ($m_{H^\pm} = 1$ TeV) (also $m_{\tilde{g}}$, $m_{\tilde{b}}$ are relevant for large $\tan\beta$). In Fig. 2 we show the distribution of C_{8g} in the complex plane for $\delta_{\rm CKM} = \pi/3$, $\tan\beta = 35(60)$, $m_{H^\pm} = 1$ TeV, $m_{\tilde{Q}} = 0.5$ TeV, $m_{\tilde{g}} = 1$ TeV, and $m_{\tilde{b}_R} = 0.5$ TeV. We scanned the other parameters as follows:

$$0 < m_{\tilde{t}} < 1\ \text{TeV}, \quad 0 < |\mu| < 1\ \text{TeV},$$
$$0 < |A_t| < 2\ \text{TeV}, \quad 0 < |M_2| < 1\ \text{TeV},$$
$$-\pi < \arg(\mu), \arg(A_t), \arg(M_2) < \pi. \tag{14}$$

We have fixed the phase on the gluino mass parameter such that $\arg(m_{\tilde{g}}) + \arg(\mu) = \pi$ to maximize the SUSY QCD correction. When scanning, we imposed the $B(B \to X_s\gamma)$ constraint and the direct search bounds on the (s)particle masses [22]: $m_{h^0} \geqslant 114.3$ GeV and $m_{\tilde{\chi}_1^\pm}, m_{\tilde{t}_1} \gtrsim 100$ GeV.

From Fig. 2 we can see that our scenario can easily accommodate the discrepancy (4). As mentioned above, we have chosen a large value for the charged Higgs boson mass $m_{H^\pm} = 1$ TeV to safely suppress the Barr-Zee type two-loop EDM constraints which are significant if $\tan\beta$ [24] is large and the pseudo-scalar Higgs boson (A^0) is relatively light.[1]

[1]The SUSY contribution to $(g-2)_\mu$ is also small in this case [25].

We have checked that for the smaller $m_{H^\pm}$ the larger deviation in $S_{\phi K}$ is possible due to the cancellation in the real part with the chargino contribution [12]. Therefore Fig. 2 is a rather conservative result for $S_{\phi K}$. We have checked that the allowed range for $S_{\phi K}$ is not sensitive to the change in $\delta_{\rm CKM}$ if we impose the constraint (3).

$C_{\phi K}$ is correlated with $S_{\phi K}$. In Fig. 3 the correlation is shown for $\tan\beta = 60$. We can see that it can also accommodate Eq. (5) although it favors small negative $C_{\phi K}$, which will be clarified by future experiments.

The direct CP asymmetry in $B \to X_s\gamma$ decay is also expected to be large when $S_{\phi K}$ has a large deviation from the SM expectations. We have checked that actually it can be very large but it is not correlated with $S_{\phi K}$. It is because $S_{\phi K}$ is determined by the complex phase on C_{8g} while $A_{CP}(B \to X_s\gamma)$ is mainly controlled by that on $C_{7\gamma}$ [26]. Because we have phases on three independent parameters μ, A_t, and M_2, the phases of $C_{7\gamma}$ and C_{8g} need not have strong correlations.

The $B(B \to \phi K)$ varies moderately over the parameter space we considered and seldom exceeds 15×10^{-6}, which is acceptable compared with the experimental measurements [4,5]. Also the mass difference Δm_s in the B_s-$\overline{B_s}$ system is close to the SM expectation $\Delta m_s \sim 14.5\ {\rm ps}^{-1}$ in most of the region of parameter space we considered, which may distinguish our scenario from other scenarios in Refs. [10,13].

For large $\tan\beta$ in the MSSM, through the Higgs-mediated FCNC the $B(B_s \to \mu^+\mu^-)$ can be enhanced by a few orders of magnitude over the SM prediction [27]. Observation of this leptonic decay mode, for example at Tevatron Run II, would be a clear signal of NP. However, this is possible only for relatively light A^0. Since the large deviation in the $S_{\phi K}$, although it needs new CP violating phase(s) of $\mathcal{O}(1)$, does not necessarily require light A^0 as we have shown, these two decay modes can be complementary to each other in searching for the MSSM at large $\tan\beta$.

In our scenario the sole source of large deviation in $S_{\phi K}$ is C_{8g}. It can have simultaneous effects on other decays such as $B \to \phi\phi$, $B \to \pi\pi$, $B \to K\pi$, etc. [12].

In conclusion we have considered the chargino contribution to the CP asymmetries $S_{\phi K}$ and $C_{\phi K}$ of $B \to \phi K_S$ decay in the CP violating MSSM scenario at large $\tan\beta$. We have shown that through the enhanced Wilson coefficient C_{8g} of a chromo-magnetic penguin operator by the large SUSY threshold corrections [23], there can be large deviations in the CP asymmetries.

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