Further Study of CP Violation and Branching Ratios for
\( \bar{B}^0 \rightarrow J/\psi K_s \) and \( \bar{B}^0 \rightarrow \phi K_s \) in the Standard Model and Beyond

Dong-Sheng Du and Mao-Zhi Yang
CCAST (WORLD LABORATORY) P.O. BOX 8730, BEIJING, 100080
and
Institute of High Energy Physics, Chinese Academy of Sciences,
P.O. Box 918(4), Beijing, 100039, People’s Republic of China

Abstract

In this work we study the CP violation for \( \bar{B}^0 \rightarrow J/\psi K_s \) and \( \bar{B}^0 \rightarrow \phi K_s \) up to leading and next-to-leading order QCD corrections in the standard model, two-Higgs-doublet model and the minimal supersymmetric extension of the standard model. We also study the effect of new physics on the branching ratios of these two decay modes. We find that within the parameter space constrained by the observation of the decay \( b \rightarrow s \gamma \), new physics does not affect the CP asymmetries greatly, and the prediction of new physics to the branching ratios of \( \bar{B}^0 \rightarrow J/\psi K_s \) and \( \bar{B}^0 \rightarrow \phi K_s \) is the same as that of the standard model up to a minor discrepancy as far as the Yukawa coupling

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†Mailing address
constants are perturbative.
I. INTRODUCTION

The decay of $\bar{B}^0 \rightarrow J/\psi K_s$ is expected to be one of the most promising channels to study the CP violation in the B decays at the B factory. In the standard model CP violation is produced via the complex phase of the Cabibbo-Kobayashi-Maskawa (CKM) matrix. Because the CKM matrix is unitary there is the unitarity equation

$$\sum_i V_{ij}V_{ik}^* = 0 \quad (j \neq k). \quad (1)$$

This equation can be represented as a closed triangle in the complex plane, which is called the unitarity triangle. Since the most poorly known entries in the CKM matrix are contained in the triangle equation

$$V_{ud}V_{ub}^* + V_{cd}V_{cb}^* + V_{td}V_{tb}^* = 0, \quad (2)$$

so this equation is most useful from the phenomenological point of view. This equation can be visualized as the unitarity triangle in Fig.1. The angles of $\alpha$, $\beta$, and $\gamma$ in Fig.1 can be determined through the detection of the CP violation of the decay modes $\bar{B} \rightarrow \pi \pi$, $\bar{B} \rightarrow \psi K_s$, $\bar{B} \rightarrow \phi K_s$, and $B_s \rightarrow \rho^0 K_s$ respectively. In the standard model CP violation is only produced through the CKM mechanism. It is believed that if the sum of the three angles $\alpha$, $\beta$, and $\gamma$ is measured to be $\pi$, the standard model gives the correct CP violating mechanism. If not, then there must be new physics beyond the standard model.

The decay $B^0 \rightarrow J/\psi K_s$ and $\bar{B}^0 \rightarrow \phi K_s$ are based on the transition $b \rightarrow q\bar{q}s$ in the quark level, where $q$ can be $c$ and $s$. Their CP conjugate processes are $B^0 \rightarrow J/\psi K_s$ and $\bar{B}^0 \rightarrow \phi K_s$ respectively. Thus both the state and CP conjugate state of the initial B meson can decay into the same final state. Its CP violating parameter, CP asymmetry, can be described by

$$A_{cp} = \frac{\int_0^\infty [\Gamma(B^0_{phys}(t) \rightarrow f) - \Gamma(\bar{B}^0_{phys}(t) \rightarrow f)] dt}{\int_0^\infty [\Gamma(B^0_{phys}(t) \rightarrow f) + \Gamma(\bar{B}^0_{phys}(t) \rightarrow f)] dt}$$
\[\begin{align*}
\xi &= \pm \frac{2q}{p_A}, \text{ here `+' for CP even } |f\rangle, \text{ `-' for CP odd } |f\rangle. \text{ For the present two decay modes that we study, } \\
\xi_f &= -\left(\frac{q}{p}\right)_B \left(\frac{q}{p}\right)_{K_s}^* \frac{\bar{A}}{A}. \text{ The parameters } p, q \text{ are defined in the physical state of } B \text{ and } K \text{ mesons. The } K_s \text{ meson is defined as } |K_s\rangle = p_K |K^0\rangle + q_K |\bar{K}^0\rangle, \text{ and the physical state of the neutral } B \text{ meson is defined as } \\
|B_L\rangle &= p_B |B^0\rangle + q_B |\bar{B}^0\rangle, \\
|B_H\rangle &= p_B |B^0\rangle - q_B |\bar{B}^0\rangle,
\end{align*}\]

where \( L, H \) denote light and heavy respectively. In the standard model the decay \( \bar{B}^0 \to J/\psi K_s \) is dominated by the tree level contribution, and \( \xi_{\psi K_s} \simeq -e^{-2i\beta} \) under the tree level approximation. The \( \bar{B}^0 \to \phi K_s \) is forbidden at the tree level, its relevant value of \( \xi \) is also equal to \( -e^{-2i\beta} \) approximately. However, in some non-standard models such as the two-Higgs-doublet model (THDM) and the minimal supersymmetric extension of the standard model, CP violation can either be produced through the CKM mechanism, or through the Higgs sector, or from both. So in order to investigate the very nature of CP violation, the \( B \) decays should be studied carefully in both standard and non-standard models.

In this work we study what impact the new mechanisms of non-standard model can give to the measurement of the CP violation of \( \bar{B}^0 \to J/\psi K_s \) and \( \bar{B}^0 \to \phi K_s \). To take into account the nonfactorization corrections, the inverse of the color number \( \frac{1}{N_c} \) should be replaced with \( \frac{1}{N_c} + \chi \), where \( \chi \) stands for the nonfactorization corrections, and this replacement amounts to changing the color number from \( N_c = 3 \) \[\text{[5]}. \] Here we consider three conditions with \( N_c = 2, 3 \text{ and } \infty \). We also investigate the new physics effects on the branching ratios of \( \bar{B}^0 \to J/\psi K_s \) and \( \bar{B}^0 \to \phi K_s \). The paper is organized as the following.

In section II we study the CP violation and the branching ratio of \( \bar{B}^0 \to J/\psi K_s \) in the
SM, THDM and the minimal supersymmetric extension of SM (MSSM). Section III is devoted to the study of $\bar{B}^0 \to \phi K_s$ in those three models. Section IV is for the discussion and conclusion.

II. The study of $\bar{B}^0 \to J/\psi K_s$

1) CP violation of $\bar{B}^0 \to J/\psi K_s$ in the standard model

The low energy effective Hamiltonian relevant to our study is\[6, 7\]

$$H_{\text{eff}} = \frac{G_F}{\sqrt{2}} \left[ \sum_{q=u,c} v_q \left\{ Q_1^u C_1(\mu) + Q_2^u C_2(\mu) + \sum_{k=3}^{10} Q_k C_k(\mu) \right\} \right] + \mathcal{H}.C., \quad (5)$$

where $C_k(\mu)$ (k=1, ... , 10) are Wilson Coefficients (WC) which are calculated in the renormalization group improved perturbation theory and include leading and next-to-leading order QCD corrections. $v_q$ is the product of Cabibbo-Kobayashi-Maskawa (CKM) matrix elements and defined as

$$v_q = \begin{cases} 
    v_{qd}^* v_{qb} & b \to d \text{ transitions}, \\
    v_{qs}^* v_{qb} & b \to s \text{ transitions}.
\end{cases}$$

The ten operators are taken in the following form\[6\]:

- $Q_1^u = (\bar{u}_\alpha b_\beta)_{V-A}(\bar{q}_\beta u_\alpha)_{V-A}$, $Q_2^u = (\bar{u}b)_{V-A}(\bar{q}u)_{V-A}$,
- $Q_3 = (\bar{q}b)_{V-A} \sum_{q'}(\bar{q}'q')_{V-A}$, $Q_4 = (\bar{q}_\beta b_\alpha)_{V-A} \sum_{q'}(\bar{q}'_\alpha q_\beta)_{V-A}$,
- $Q_5 = (\bar{q}b)_{V-A} \sum_{q'}(\bar{q}'q')_{V+A}$, $Q_6 = (\bar{q}_\beta b_\alpha)_{V-A} \sum_{q'}(\bar{q}'_\alpha q_\beta)_{V+A}$,
- $Q_7 = \frac{3}{2} (\bar{q}b)_{V-A} \sum_{q'} e_{q'}(\bar{q}'q')_{V+A}$, $Q_8 = \frac{3}{2} (\bar{q}_\beta b_\alpha)_{V-A} \sum_{q'} e_{q'}(\bar{q}'_\alpha q_\beta)_{V+A}$,
- $Q_9 = \frac{3}{2} (\bar{q}b)_{V-A} \sum_{q'} e_{q'}(\bar{q}'q')_{V-A}$, $Q_{10} = \frac{3}{2} (\bar{q}_\beta b_\alpha)_{V-A} \sum_{q'} e_{q'}(\bar{q}'_\alpha q_\beta)_{V-A}$,
where \( Q_1^q \) and \( Q_2^q \) are current-current operators, \( q=u, c \). For \( q = c \) case, \( Q_1^c \) and \( Q_2^c \) are obtained through making substitution \( u \rightarrow c \) in \( Q_1^u \) and \( Q_2^u \). \( Q_3 \sim Q_6 \) are QCD penguin operators, the sum \( \sum_{q'} \) is running over all flavors being active at \( \mu = m_b \) scale, \( q' = \{ u, d, s, c, b \} \). \( Q_7 \sim Q_{10} \) are electroweak penguin operators, \( e_{q'} \) are the electric charges of the relevant quarks in unit e which is the charge of the proton. The subscripts \( \alpha, \beta \) are \( SU(3)_c \) color indices. \((V \pm A)\) refers to \( \gamma_\mu (1 \pm \gamma_5) \).

For our calculation we use the renormalization scheme independent form of the Wilson Coefficients \( C'_i(\mu) \),

\[
\begin{align*}
C'_1 &= \bar{C}_1, & C'_2 &= \bar{C}_2, \\
C'_3 &= \bar{C}_3 - P_s/N_c, & C'_4 &= \bar{C}_4 + P_s, \\
C'_5 &= \bar{C}_5 - P_s/N_c, & C'_6 &= \bar{C}_6 + P_s, \\
C'_7 &= \bar{C}_7 + P_e, & C'_8 &= \bar{C}_8, \\
C'_9 &= \bar{C}_9 + P_e, & C'_{10} &= \bar{C}_{10},
\end{align*}
\]

(7)

where \( \bar{C}_i(\mu) \) is [3, 8, 9]

\[
\bar{C}(\mu) = (1 + \hat{r}_s^T \alpha_s(\mu)/4\pi + \hat{r}_e^T \alpha_{em}(\mu)/4\pi) \cdot C(\mu),
\]

(8)

and

\[
P_s = \frac{\alpha_s}{8\pi} \bar{C}_2(\mu) \left[ \frac{10}{9} - G(m_q, q, \mu) \right],
\]

\[
P_e = \frac{\alpha_{em}}{3\pi} \left( \bar{C}_1(\mu) + \frac{\bar{C}_2(\mu)}{N_c} \right) \left[ \frac{10}{9} - G(m_q, q, \mu) \right],
\]

\[
G(m_q, q, \mu) = -4 \int_0^1 x(1-x)dx \ln \frac{m_q^2}{\mu^2} - x(1-x)q^2,
\]

here \( q=u, c \). The renormalization scheme independent Wilson coefficients \( \bar{C}_i(\mu) \) has been given in Ref. [10],

\[
\bar{C}_1 = -0.313, \quad \bar{C}_2 = 1.150, \quad \bar{C}_3 = 0.017,
\]
\[ C_4 = -0.037, \quad C_5 = 0.010, \quad C_6 = -0.046, \]
\[ C_7 = -0.001 \cdot \alpha, \quad C_8 = 0.049 \cdot \alpha, \quad C_9 = -1.321 \cdot \alpha, \]
\[ C_{10} = 0.267 \cdot \alpha. \]  

(9)

With the effective Hamiltonian given in eq.(5) we obtain

\[
\tilde{A} \equiv \langle J/\psi K_s | H_{\text{eff}} | \bar{B}^0 \rangle = q_K^* \langle J/\psi \bar{K}^0 | H_{\text{eff}} | \bar{B}^0 \rangle
\]
\[
= q_K^* \frac{G_F}{\sqrt{2}} \left[ (C'_1 + \frac{C'_2}{N_c})v_c + \sum_{q=u,c} v_q \left[ C'_3 + \frac{C'_4}{N_c} + C'_5 + \frac{C'_6}{N_c} + \frac{3}{2} e_c (C'_7 + \frac{C'_8}{N_c} + C'_9 + \frac{C'_{10}}{N_c}) \right] \right]
\]
\[
\cdot \langle J/\psi | (\bar{c}c)_{V-A} | 0 \rangle \langle \bar{K}^0 | (\bar{s}b)_{V-A} | \bar{B}^0 \rangle
\]  

(10)

After we get eq.(10), we are now in the position to calculate \( \xi_{\psi K_s} \),

\[ \xi_{\psi K_s} = -\left( \frac{q}{p} \right)_B \left( \frac{q}{p} \right)^*_K \frac{\tilde{A}}{A}, \]  

(11)

and

\[
\left( \frac{q}{p} \right)_B \left( \frac{q}{p} \right)^*_K \approx \frac{v^*_b v_{td}}{v_b v^*_{td}} \cdot \frac{v_{cs} v^*_{cd}}{v_{cs} v^*_{cd}}
\]
\[
\approx \frac{1 - \rho - i\eta}{1 - \rho + i\eta}
\]
\[
\approx e^{-2i\beta},
\]

where \( \rho, \eta \) are the parameters in the Wolfenstein parameterization for the CKM matrix.

Substitute the numerical values of \( C'_i \) into eq.(10) and (11), we can get the numerical result of \( \xi_{\psi K_s} \). For the case without loop correction, only \( C'_1 \) and \( C'_2 \) contribute,

\[ \xi_{\psi K_s} \approx -e^{-2i\beta} \approx -\frac{1 - \rho - i\eta}{1 - \rho + i\eta}. \]  

(12)

For the case with loop corrections all \( C'_1 \cdots C'_{10} \) contribute,

\[ \xi_{\psi K_s} = -e^{-2i\beta} \frac{118.73 - \rho + i\eta}{118.73 - \rho - i\eta}
\]
\[ \approx \frac{1 - \rho - i\eta}{1 - \rho + i\eta} \frac{118.73 - \rho + i\eta}{118.73 - \rho - i\eta}. \]  

(13)
Experimental results have constrained ρ and η into a small area in the ρ-η plane\cite{1}, approximately, \(-0.2 < ρ < 0.3\), and \(0.2 < η < 0.4\). From eq.(12) and (13) we can see that loop diagram correction to ξ_{ψK_s} is only in the order of \(O(10^{-3})\). So the unitarity triangle angle \(β\) can be determined through the measurement of the CP asymmetry in \(B^0 \rightarrow J/ψK_s\) up to an approximation of the order \(O(10^{-3})\). The eq.(13) is obtained with \(N_c = 3\), when \(N_c = 2, \infty\) the loop corrections are also in the same order of \(O(10^{-3})\). This conclusion has been obtained through an estimate that the penguin contributions are proportional to \(V_{tb} V_{ts}^* \simeq λ^2\), \(V_{cb} V_{cs}^* \simeq λ^2\) and \(V_{ub} V_{us}^* \simeq λ^4 (λ ≈ 0.22)\), so up to very small corrections the penguin contributions have the same weak phase as the tree diagram contribution, then the penguin contribution affects the CP violation extremely small \cite{1, 11}. Our calculation using the QCD corrected Hamiltonian confirms this conclusion.

2) CP violation of \(B^0 \rightarrow J/ψK_s\) in the two-Higgs-doublet model

In the two-Higgs-doublet model\cite{12} the diagrams contributing to \(B^0 \rightarrow J/ψK_s\) are shown in Fig.2. The relevant low energy effective Hamiltonian has the same form as eq.(5), but the Wilson coefficients at the scale \(μ = M_W\) contain an extra contribution from the charged Higgs loop diagrams. We calculated the new contributions at the scale \(μ = M_W\) at first, and then using the two-loop-renormalization group equation to evolve them down to the scale \(μ = m_b\). The initial conditions in the scale \(μ = M_W\) are

\[
C_i(M_W) = C_i^{SM}(M_W) + C_i^{H^±}(M_W),
\]

where the SM contribution \(C_i^{SM}(M_W)\)'s have already been given in Ref. \cite{7}. In this paper, we calculate the contributions of charged Higgs loop diagrams, the results are

\[
C_1^{H^±}(M_W) = 0,
\]

\[
C_2^{H^±}(M_W) = 0,
\]
\[ C_{3}^{H^\pm}(M_W) = -\frac{\alpha_s(M_W)}{24\pi} F_{H^\pm}(x'_t), \]

\[ C_{4}^{H^\pm}(M_W) = \frac{\alpha_s(M_W)}{8\pi} F_{H^\pm}(x'_t), \]

\[ C_{5}^{H^\pm}(M_W) = -\frac{\alpha_s(M_W)}{24\pi} F_{H^\pm}(x'_t), \]

\[ C_{6}^{H^\pm}(M_W) = \frac{\alpha_s(M_W)}{8\pi} F_{H^\pm}(x'_t), \]

\[ C_{7}^{H^\pm}(M_W) = \frac{\alpha}{6\pi} \left[ E_{H^\pm}(x'_t) - D_{H^\pm}(x'_t) \right], \]

\[ C_{8}^{H^\pm}(M_W) = 0, \]

\[ C_{9}^{H^\pm}(M_W) = \frac{\alpha}{6\pi} \left[ E_{H^\pm}(x'_t) - D_{H^\pm}(x'_t) + \frac{1}{2\sin^2\theta_w} D_{H^\pm}(x'_t) \right], \]

\[ C_{10}^{H^\pm}(M_W) = 0, \]

where

\[ x'_t = \frac{m_t^2}{M_{H^\pm}^2}, \]

\[ E_{H^\pm}(x) = \frac{F_u^2}{18} \left[ \frac{38x - 79x^2 + 47x^3}{6(1-x)^3} + \frac{4x - 6x^2 + 3x^4}{(1-x)^4} \ln x \right], \]

\[ F_{H^\pm}(x) = \frac{F_u^2}{6} \left[ \frac{16x - 29x^2 + 7x^3}{6(1-x)^3} + \frac{2x - 3x^2}{(1-x)^4} \ln x \right], \]

\[ D_{H^\pm}(x) = -\frac{F_u^2}{2} x_t \left[ \frac{x}{1-x} + \frac{x}{(1-x)^2} \ln x \right], \]

where \( F_u \) is the coupling constant between the charged Higgs and the up- and down-type quarks. We consider two distinct two-Higgs-doublet modes, model I and model II, which naturally avoid tree-level flavor changing neutral current (FCNC). In model I, one doublet \( \phi_2 \) gives masses to all fermions and the other doublet \( \phi_1 \) essentially decouples from the fermions, and here \( F_u = \cot\beta \). In model II, \( \phi_2 \) gives mass to the up-type quarks, while
the other ($\phi_1$) gives the down-type quarks masses, and in this model $F_u = \cot\beta$, where $\tan\beta \equiv v_2/v_1$, which is the ratio of the vacuum expectation values (VEV) of the two Higgs doublet.

After presenting the initial values of the Wilson coefficients $C_i(M_W)$ in the two-Higgs-doublet model, now we shall evolve them down to the scale $\mu = m_b$, so that to take into account the QCD corrections. The renormalization-group equation for $C(\mu)$ is given by

$$\left[\mu \frac{\partial}{\partial \mu} + \beta(g) \frac{\partial}{\partial g}\right] C(M_W^2, g^2, \alpha) = \hat{\gamma}^T(g^2, \alpha) C(M_W^2, g^2, \alpha),$$

(17)

where $\hat{\gamma}$ is the anomalous dimensions. A. Buras et al. have calculated the $10 \times 10$ anomalous dimension matrix involving current-current operators, QCD penguin operators, and electroweak penguin operators up to next-to-leading order QCD corrections in the standard model. Now in the two-Higgs-doublet model and the supersymmetric extension of the standard model, the anomalous dimension matrix $\hat{\gamma}$ is just the same as that in the standard model. In this section and the next we will use the anomalous dimension matrix $\hat{\gamma}$ calculated by Buras et al. to evolve the Wilson coefficients from scale $\mu = M_W$ down to the $\mu = m_b$ scale.

The solution of the renormalization-group equation (19) is given by

$$C(\mu) = \hat{U}(\mu, M_W, \alpha)C(M_W),$$

(18)

where $\hat{U}(\mu, M_W, \alpha)$ is the evolution matrix from $M_W$ down to $\mu < M_W$. By resolving the eq. (17), we can get the numerical result for $\hat{U}(m_b, M_W, \alpha)$, which is presented in Appendix A.

The difference between the phases of $B^0 - \bar{B}^0$ mixing in the standard model and in the two-Higgs-doublet model is trivial. There are two types of box diagrams which contribute to the $B^0 - \bar{B}^0$ mixing: the diagrams with one $W$ and one charged Higgs propagating in it and the diagrams with two charged Higgses. The charged Higgs contributions
to the mixing involve the same CKM factors and almost have the same phase as the W-box contributions[14, 15]. So we will take \((q/p)_B = \frac{v_2^u v_{ud}}{v_{ub} v_{td}}\), which is the same as in the standard model. The case of the \(K^0 - \bar{K}^0\) mixing is similar to the \(B^0 - \bar{B}^0\) mixing. The contributions of charged Higgs bosons to the \(K^0 - \bar{K}^0\) mixing present a negligible phase difference from the standard model.

With the Wilson coefficients calculated in the two-Higgs-doublet model, and using eq.(3), (10) and (11), we can get the results of the CP asymmetry of \(\bar{B}^0 \rightarrow J/\psi K_s\). We now discuss our numerical results.

In the parameter space constrained by the observed rate of \(b \rightarrow s\gamma\), the possibility that charged-Higgs exchange induces large CP-violating effects through the complex phase of \(v_2/v_1\) has been ruled out[13, 16]. In the continuing work we will drop the complex phase of the VEV, i.e., only take into account the cases that \(v_2/v_1\) is a real number.

The observation of the decay \(b \rightarrow s\gamma\) gives a correlated bound to \(\cot\beta\) and the charged Higgs mass \(M_{H^\pm}\)[15]

\[
|\cot\beta| (100\text{GeV}/M_{H^\pm}) < 0.8. \tag{19}
\]

We make our calculation within this constraint of \(\cot\beta\) and \(M_{H^\pm}\), and also keep \(\cot\beta < 3\), which is the region of \(\cot\beta\) that keeps the Yukawa coupling perturbative[17]. We calculate three cases with the number of color \(N_c = 2, 3\) and \(\infty\). The result is that within the parameter space constrained by the experiment, charged Higgs effect reduces the CP asymmetry. But the effect is too small to distinguish the two-Higgs-doublet model from the SM.

3) CP violation of \(\bar{B}^0 \rightarrow J/\psi K_s\) in the MSSM

In the minimal supersymmetric extension of the standard model (MSSM), except for the usual particle fields in the standard model, there are the relevant supersymmetric partners, and an extra Higgs doublet. They add the new contributions to the process
of $b \to c\bar{c}s$. They are the contributions from: i) the charged Higgs loop diagrams, ii) the up-type squarks and chargino loop diagrams, iii) the down-type squarks, gluino and neutralino loop diagrams, which are shown in Fig.3. For the basic structure of the minimal supersymmetric extension of the standard model, we refer the readers to the Refs. [18, 19].

Because the assumption is quite strong that all the parameters in the soft supersymmetric breaking terms, such as the masses and couplings of the MSSM scalars and fermions, are unified at the grand unification (GUT) scale $\mu = M_{GUT}$, we relax all the constraints to the parameters of MSSM from this assumption and investigate the effects of the more general structure of MSSM. So the free parameters are the following: the masses of the left and right-handed up and down-type squarks, $m_{\tilde{f}_L}$, $m_{\tilde{f}_R}$; the mixing angles of the right and left-handed squarks, $\alpha_\tilde{f}$; the wino and charged Higgs masses $M_{\tilde{W}}$ and $M_{H^\pm}$; the superpotential $\mu$ and $\cot\beta$. For simplicity the following features should be assumed:

(i) Supersymmetric loop diagram contributions to $B \to X_s\gamma$ decay are mainly from the charged Higgs and chargino exchange.

(ii) The first two generations of up and down-type squarks are almost degenerate.

The above features are inferred from the observations noted in Ref. [19]. With these simplicities the remaining free parameters are left to be: a common mass $m_{\tilde{u}_L}$ for the first two generations of the left and right-handed squarks; the top squark masses $m_{\tilde{t}_L}$, $m_{\tilde{t}_R}$ and the top squark mixing angle $\alpha_{\tilde{t}}$; the masses $M_{\tilde{\omega}}$ and $M_{H^\pm}$; the $\mu$ and $\cot\beta$.

The minimal supersymmetric extension of the standard model gives the new contributions to the effective Hamiltonian of eq.(5) by adding some new terms into the Wilson coefficients. We express the new terms as $C^{SUSY}(\mu)$. Thus

$$C_i(\mu) = C_i^{SM}(\mu) + C_i^{SUSY}(\mu).$$

As in the SM, the Wilson coefficients are calculated in the scale $\mu = M_W$ at first, then use the renormalization group equation to evolve them down to the scale $\mu = m_b$. The values
of $C_i^{SM}(M_W)$ have been given in Ref. [7]. We calculate the supersymmetric contributions to the Wilson coefficients, $C_3^{SUSY}$, $C_4^{SUSY}$, $C_5^{SUSY}$, $C_6^{SUSY}$, $C_7^{SUSY}$, $C_8^{SUSY}$, $C_9^{SUSY}$, $C_{10}^{SUSY}$, at $\mu = M_W$ scale. The authors of Ref. [20] have calculated the contributions of gluino diagrams to the process $b \to q'q\bar{q}$. We refer to their results for the gluino diagram contribution to $C_3^{SUSY}$, $C_4^{SUSY}$, $C_5^{SUSY}$, $C_6^{SUSY}$ in this work. The initial conditions of the evolution are listed in the following,

$$C_1^{SUSY} = 0,$$

$$C_2^{SUSY} = 0,$$

$$C_3^{SUSY} = - \sum_{j=d,s,b} \frac{2\alpha_s^2 M_W^2}{g_2^2 \bar{m}_g} \left( s_L(x_j, y) - \frac{1}{6} o_L(x_j, y) \right) - \frac{\alpha_s}{24\pi} Z^{SUSY},$$

$$C_4^{SUSY} = - \sum_{j=d,s,b} \frac{\alpha_s^2 M_W^2}{g_2^2 \bar{m}_g} o_L(x_j, y) + \frac{\alpha_s}{8\pi} Z^{SUSY}, \tag{21}$$

$$C_5^{SUSY} = - \sum_{j=d,s,b} \frac{2\alpha_s^2 M_W^2}{g_2^2 \bar{m}_g} \left( s_R(x_j, y) - \frac{1}{6} o_R(x_j, y) \right) - \frac{\alpha_s}{24\pi} Z^{SUSY},$$

$$C_6^{SUSY} = - \sum_{j=d,s,b} \frac{\alpha_s^2 M_W^2}{g_2^2 \bar{m}_g} o_R(x_j, y) + \frac{\alpha_s}{8\pi} Z^{SUSY},$$

where [20]

$$s_L(x_j, y) = s_R(x_j, y) = \frac{2}{9} g(x_j, y) + \frac{4}{9} f(x_j, y),$$

$$o_L = \frac{7}{6} g(x_j, y) - \frac{2}{3} f(x_j, y) + \frac{1}{2} F(x_j) + \frac{4}{9} f(x_j),$$

$$o_R = \frac{1}{3} g(x_j, y) + \frac{7}{3} f(x_j, y) + \frac{1}{2} F(x_j) + \frac{4}{9} f(x_j), \tag{22}$$

and

$$f(x, y) = \frac{1}{y - x} \left\{ \frac{x}{(x - 1)^2} \ln x - \frac{y}{(y - 1)^2} \ln y - \frac{1}{x - 1} + \frac{1}{y - 1} \right\},$$

$$g(x, y) = \frac{1}{x - y} \left\{ \frac{x^2}{(x - 1)^2} \ln x - \frac{y^2}{(y - 1)^2} \ln y - \frac{1}{x - 1} + \frac{1}{y - 1} \right\},$$
\[ F(x) = \frac{1}{(1-x)^4} \left\{ -\frac{3}{2}x^3 + \frac{15}{2}x^2 - \frac{21}{2}x + \frac{9}{2} + (2x^3 - 6x^2 + 3x + 1)\ln x \right\}, \]

\[ f(x) = \frac{1}{(1-x)^4} \left\{ \frac{1}{3}x^3 - \frac{3}{2}x^2 + 3x - \frac{11}{6} - \ln x \right\}, \tag{23} \]

\[ x_j = \frac{\tilde{m}_j^2}{\tilde{m}_g^2}, \quad y = \frac{\tilde{m}_c^2}{\tilde{m}_g^2}. \]

The values of \( C_{i}^{SUSY} \), \( i = 7, \cdots, 10 \) are

\[ C_7^{SUSY} = \frac{\alpha}{6\pi} Y_1^{SUSY}, \]

\[ C_8^{SUSY} = 0, \]

\[ C_9^{SUSY} = \frac{\alpha}{6\pi} Y_2^{SUSY}, \tag{24} \]

\[ C_{10}^{SUSY} = 0. \]

And \( Z^{SUSY}, Y_1^{SUSY}, Y_2^{SUSY} \) in eq.(21) and (24) are presented in Appendix B. QCD corrections up to the next-to-leading order are taken into account by evolving the initial values of the Wilson coefficients down to the scale \( \mu = m_b \). Because the anomalous dimension matrices of the ten operators, \( C_i \) \( i = 1, \cdots, 10 \), are the same as in the SM, we can still use the evolution matrix \( \hat{U}(m_b, M_W, \alpha) \) listed in Appendix A. We perform the evolution using eq.(18) as in section 2).

After obtaining the Wilson coefficients \( C_i \) \( i = 1, \cdots, 10 \) with the QCD corrections up to the next-to-leading order in MSSM, we can still use eq.(3), (10), (11) to calculate the CP violating parameter with the Wilson coefficients of the MSSM. We perform those calculations within the parameter space allowed by experimental limit. For example, the experiments give the lower bounds of squark and chargino mass are 176 GeV and 45 GeV, respectively[21].

Because the complex phase of \( B^0 - \bar{B}^0 \) mixing is the same as the standard model up to the minor correction of the order of \((m_c/m_t)^2\) or less[22], we do not pay much attention
to it here.

We performed our calculation within the very large ranges of all the parameters \((m_\tilde{\nu}, \mu, m_\tilde{t}_L, m_\tilde{b}_L, m_\tilde{t}_R, m_\tilde{b}_R, \alpha_i)\) which are allowed by the experiment. We find that if we keep the Yukawa coupling is perturbative \((\text{cot} \beta < 3)\), the discrepancy from the SM is extremely small, it only happens in the third decimal number after the zero point.

4) The branching ratio of \(\bar{B}^0 \rightarrow J/\psi K_s\) in the SM, THDM and MSSM

Because there are great theoretical uncertainties in the calculation of the branching ratio of \(\bar{B}^0 \rightarrow J/\psi K\), in this section we do not want to give an precise prediction of the branching ratio. We only want to give a comparison of the predictions of these three models, and to probe the possibility to distinguish them by measuring the branching ratio of this decay mode. We calculate the branching ratio according to the BSW method [23].

In our calculation we take \(f_\psi = 0.386\) GeV, \(F^{BK}_1(0) = 0.379\) and \(F^{BK}_1(q^2) = \frac{F^{BK}_1(0)}{1-q^2/m_{1-}^2}\)
where \(m_{1-} = 5.43\) GeV. The result of our calculation with the color number being taken as \(N_c = 2, 3\) and \(\infty\) is that the difference between the predictions of the SM and MSSM is small, it is only up to a few percent level.

III The Study of \(\bar{B}^0 \rightarrow \phi K_s\)

In this section we report the results of our study of \(\bar{B}^0 \rightarrow \phi K_s\) using the effective Hamiltonian given in eq.(5), (14), (18) and (20). Because there is no tree level diagram contribution to \(\bar{B}^0 \rightarrow \phi K_s\) in the SM, new physics may have an observable effect. Using the effective Hamiltonian given in the last section, the amplitude of \(\bar{B}^0 \rightarrow \phi K_s\) is

\[
\langle K_s \phi | H_{\text{eff}} | \bar{B}^0 \rangle = \frac{G_f}{\sqrt{2}} g_s \sum_{q=u,c} v_q \left( (1 + \frac{1}{N_c}) C_3' + (1 + \frac{1}{N_c}) C_4' + C_5' + \frac{1}{N_c} C_6' + \frac{3}{2} f_s [ C_7' + \frac{1}{N_c} C_8' ] \right)
\]
\[ +(1 + \frac{1}{N_c}) C'_9 + (\frac{1}{N_c} + 1) C'_{10} \right) \langle \phi | (\bar{s}s)_{V-A} | 0 \rangle \langle K^0 | (\bar{s}b)_{V-A} | \bar{B}^0 \rangle \] (25)

where \( C'_i (i = 1, \cdots, 10) \) should be the renormalization scheme independent Wilson Coefficients calculated in the SM, THDM and MSSM, respectively. \( e_s = -\frac{2}{3} \) is the charge of the s-quark.

The Decay width of a \( \bar{B}^0 \) meson at rest decaying into \( \phi \) and \( K_s \) is

\[
\Gamma(\bar{B}^0 \to \phi K_s) = \frac{1}{8\pi} \left| \langle \phi K_s | H_{eff} | \bar{B}^0 \rangle \right|^2 \frac{p}{M_B},
\] (26)

where

\[
|p| = \frac{\{[M_B^2 - (M_\phi + M_{K_s})^2][M_B^2 - (M_\phi - M_{K_s})^2]\}^{1/2}}{2M_B}
\] (27)

is the momentum of \( \phi \) or \( K_s \). The corresponding branching ratio is given by

\[
B_{br}(\bar{B}^0 \to \phi K_s) = \frac{\Gamma(\bar{B}^0 \to \phi K_s)}{\Gamma_{\bar{B}^0_{tot}}}.
\] (28)

We take \( \Gamma_{\bar{B}^0_{tot}} = 4.22 \times 10^{-13} \text{GeV} \) in our calculation. The hadronic matrix element \( \langle \phi | (\bar{s}s)_{V-A} | 0 \rangle \langle K^0 | (\bar{s}b)_{V-A} | \bar{B}^0 \rangle \) is calculated in the Bauer, Stech, and Wirbel (BSW) method. The one-body vector matrix element of (V-A) current is

\[
\langle \phi | (\bar{s}s)_{V-A} | 0 \rangle = f_\phi M_\phi \epsilon_\phi,
\] (29)

where \( f_\phi = 0.233 \text{ GeV}, M_\phi = 1.02 \text{GeV} \), \( \epsilon_\phi \) is the polarization of \( \phi \) meson. The two-body pseudoscalar-pseudoscalar matrix element of the vector current is

\[
\langle \bar{K}^0 | (\bar{s}g_\mu(1-\gamma_5)b) | \bar{B}^0 \rangle = \left( P_B + P_K - \frac{M_B^2 - M_K^2}{q^2} q^2 \right) \mu F_{BK}(q^2, 1^-) + \frac{M_B^2 - M_K^2}{q^2} q^2 q_\mu F_{BK}(q^2, 0^+),
\] (30)

where \( q = P_B - P_K, F_{BK}(q^2, 1^-) = F_{BK}(0, 1^-)/(1 - \frac{q^2}{M^2_{1^-}}) \). In our calculation, we take \( F_{BK}(0, 1^-) = 0.379, M_{1^-} = 5.43 \text{GeV} \). Now with eq.(3), (25), and (28), we can
calculate the CP asymmetry and branching ratio of $\bar{B}^0 \to \phi K_s$. Our calculation shows that if we take $\cot\beta < 3$ [17], the difference between the predictions of new physics (THDM, MSSM) and SM is too small to be detectable in the future B factory.

V. Conclusion and discussion

We studied the CP asymmetry and the branching ratio of $\bar{B}^0 \to J/\psi K_s$ and $\bar{B}^0 \to \phi K_s$ up to the leading and next-to-leading order QCD corrections in the standard model, the two-Higgs-doublet model and the minimal supersymmetric extension of the standard model. In our calculation we can neglect the charged Higgs tree diagram contribution in the $b \to c\bar{c}s$ decay. Because this diagram gives the contribution to the operator $\bar{c}\gamma_\mu(1 + \gamma_5)c\bar{s}\gamma_\mu(1 - \gamma_5)b$, whose coefficient is $-2\cot^2\beta \frac{m_s^2}{M_H^2}$, and eq.(19) means $\cot^2\beta \frac{m_s^2}{M_H^2} < (\frac{0.8m_t}{100GeV})^2$. This implies the contribution of the charged Higgs tree diagram is too small, and can be safely neglected. So in the two-Higgs-doublet model and the MSSM, only some loop diagrams are added. Because of the coupling constant $\alpha$ and $\alpha_s$ suppression, the W-boson tree diagram dominates. Thus there will be no large difference between the CP asymmetries, branching ratios of $\bar{B}^0 \to J/\psi K_s$ predicted by these three models.

For $\bar{B}^0 \to \phi K_s$ decay, new physics only contributes an overall strong phase, and it does not affect the weak phase much. So in the THDM and MSSM, new physics has no large effect in the CP asymmetry of $\bar{B}^0 \to \phi K_s$.

After finishing our work, we noticed that the next-to-leading QCD corrections to $B \to X_s\gamma$ decay are calculated in the THDM in Ref. [24, 27].

Finally we come to our conclusion: (i) The tree diagram calculation for $A_{cp}$ in $\bar{B}^0 \to J/\psi K_s$ is reliable, which is independent of the nonfactorization corrections. (ii) It is difficult to distinguish the SM from the two Higgs-doublet model and MSSM by measuring CP asymmetries $A_{cp}$ and the branching fraction of $\bar{B}^0 \to J/\psi K_s$ and $\bar{B}^0 \to \phi K_s$ at B
factories.

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Appendix A

The evolution matrix from $M_W$ down to $\mu = m_b$ scale, where we take $m_b = 5.0\text{GeV}$.

\[ \hat{U}(m_b, M_W, \alpha) = \]

\[
\begin{pmatrix}
1.115, & -0.246, & 0, & 0, & 0, & 0, & 0, & 0, & 0, & 0 \\
-0.246, & 1.12, & 0, & 0, & 0, & 0, & 0, & 0, & 0, & 0 \\
-0.007, & 0.012, & 1.10, & -0.203, & 0.030, & 0.078, & 0.003, & 0.008, & -0.012, & 0.012 \\
0.006, & -0.033, & -0.285, & 0.977, & -0.017, & -0.169, & -0.002, & -0.017, & 0.034, & -0.018 \\
0.003, & 0.009, & 0.035, & 0.043, & 0.903, & 0.096, & -0.004, & 0.004, & -0.007, & 0.001 \\
0.005, & -0.038, & -0.059, & -0.172, & 0.317, & 1.717, & -0.003, & -0.029, & 0.039, & -0.024 \\
-0.003, & -0.001, & -0.001, & 0.003, & -0.006, & -0.001, & 0.913, & 0.053, & -0.0096, & -0.004 \\
-0.0008, & 0.0001, & -0.0003, & 0.001, & -0.0026, & -0.0104, & 0.326, & 1.949, & -0.0028, & -0.0005 \\
-0.004, & -0.0007, & 0.0038, & 0.0016, & -0.0019, & -0.001, & -0.010, & -0.004, & 1.101, & -0.2497 \\
0.0003, & -0.0001, & -0.001, & 0.005, & 0.0002, & -0.0001, & 0.001, & 0.0009, & -0.244, & 1.112
\end{pmatrix}
\]

(A1)
Appendix B

The functions $Z^{SUSY}, Y_1^{SUSY}$ and $Y_2^{SUSY}$ appearing in the Wilson coefficient $C_7^{SUSY}$ in eq. (21) and (24). Our result is consistent with Ref. [19] after replacing the lepton charge with the relevant quark charge.

1) contribution from charged Higgs loops with a z-boson coupling to the quark pair:

\[
Y_1 = \frac{1}{2} \cot^2 \beta x_t f_1 \left( \frac{m_t^2}{m_{H^\pm}^2} \right),
\]

\[
Y_2 = -\frac{1}{4} \left( \frac{1}{\sin^2 \theta_W} - 2 \right) \cot^2 \beta x_t f_1 \left( \frac{m_t^2}{m_{H^\pm}^2} \right);
\]

(B1)

2) contribution from charged Higgs loops with a photon coupling to the quark pair:

\[
Y_1 = Y_2 = \frac{1}{18} \cot^2 \beta f_2 \left( \frac{m_t^2}{m_{H^\pm}^2} \right).
\]

(B2)

3) contribution from charged Higgs loops with a gluon coupling to the quark pair:

\[
Z = \frac{1}{6} \cot^2 \beta f_3 \left( \frac{m_t^2}{m_{H^\pm}^2} \right).
\]

(B3)

4) contribution from chargino loops with a z-boson coupling to the quark pair:

\[
Y_1 = -\frac{2}{g_2^2 v_{ts} v_{tb}} \sum_{A,B=1}^{6} \sum_{I,J=1}^{2} (X_I^{UL})^+_{2A} (X_J^{UL})_{B3} \left\{ c_2(m_{\tilde{\chi}^\pm_i, A}, m_{\tilde{u}^+_A, B}^2, m_{\tilde{u}^-_B}^2, m_{\tilde{u}^-_B}^2) \right\} \delta_{AB} \delta_{IJ} \\
- c_2(m_{\tilde{u}^-_A}, m_{\tilde{\chi}^+_j, B}, m_{\tilde{\chi}^-_j, B}^2) \delta_{AB} V^*_{i1} V_{j1} + \frac{1}{2} m_{\tilde{\chi}^-_j, B} m_{\tilde{\chi}^-_j, B} c_0(m_{\tilde{u}^-_A, B}^2, m_{\tilde{\chi}^-_j, B}^2, m_{\tilde{\chi}^-_j, B}^2) \delta_{AB} U_{11}^* U_{J1}^* \right\},
\]

\[
Y_2 = \left( \frac{1}{\sin^2 \theta_W} - 2 \right) \frac{1}{g_2^2 v_{ts} v_{tb}} \sum_{A,B=1}^{6} \sum_{I,J=1}^{2} (X_I^{UL})^+_{2A} (X_J^{UL})_{B3} \left\{ c_2(m_{\tilde{\chi}^\pm_i, A}, m_{\tilde{u}^+_A, B}^2, m_{\tilde{u}^-_B}^2, m_{\tilde{u}^-_B}^2) \right\} \delta_{AB} \delta_{IJ} \\
- c_2(m_{\tilde{u}^-_A}, m_{\tilde{\chi}^+_j, B}, m_{\tilde{\chi}^-_j, B}^2) \delta_{AB} V^*_{i1} V_{j1} + \frac{1}{2} m_{\tilde{\chi}^-_j, B} m_{\tilde{\chi}^-_j, B} c_0(m_{\tilde{u}^-_A, B}^2, m_{\tilde{\chi}^-_j, B}^2, m_{\tilde{\chi}^-_j, B}^2) \delta_{AB} U_{11}^* U_{J1}^* \right\}.
\]

(B4)
5) contribution from chargino loops with a photon coupling to the quark pair:

\[ Y_1 = Y_2 = - \frac{1}{9g_2^2 v_t^2 v_t^* v_t^*} \sum_{A=1}^{6} \sum_{I=1}^{2} \frac{m_w^2}{m_{\tilde{\chi}_I^\pm A}^2} (X_I^{U_L})^+_{2A} (X_I^{U_L})_{A3} f_4 \left( \frac{m_{\tilde{\chi}_I^\pm A}^2}{m_{\tilde{\chi}_I^\pm A}^2} \right). \]  

(B5)

6) contribution from chargino loops with a gluon coupling to the quark pair:

\[ Z = - \frac{1}{3g_2^2 v_t^2 v_t^* v_t^*} \sum_{A=1}^{6} \sum_{I=1}^{2} \frac{m_w^2}{m_{\tilde{\chi}_I^\pm A}^2} (X_I^{U_L})^+_{2A} (X_I^{U_L})_{A3} f_5 \left( \frac{m_{\tilde{\chi}_I^\pm A}^2}{m_{\tilde{\chi}_I^\pm A}^2} \right). \]  

(B6)

7) contribution from neutralino loops with a z-boson coupling to the quark pair:

\[ Y_1 = \frac{2}{g_2^2 v_t^2 v_t^* v_t^*} \sum_{A,B=1}^{6} \sum_{I,J=1}^{4} \frac{1}{Z_I^{D_L}} (Z_{I,J}^{D_L}) B_3 \left\{ c_2(m_{\tilde{\chi}_I^O}^2, m_{\tilde{\chi}_J^O}^2, m_{\tilde{\chi}_J^O}^2) (\Gamma^{D_R \Gamma^{D_R^+}})_{AB} \delta_{IJ} \right\} - \frac{1}{2} m_{\tilde{\chi}_I^O} m_{\tilde{\chi}_J^O} \left( c_0(m_{\tilde{\chi}_I^O}^2, m_{\tilde{\chi}_J^O}^2) - \frac{1}{2} m_{\tilde{\chi}_I^O} m_{\tilde{\chi}_J^O} \right) \delta_{AB} (N_{I3}^{*} N_{I3}^{*} - N_{I4}^{*} N_{I4}^{*}) = 0 \}

(B7)

8) contribution from neutralino loops with a photon coupling to the quark pair:

\[ Y_1 = Y_2 = \frac{1}{54g_2^2 v_t^2 v_t^* v_t^*} \sum_{A=1}^{6} \sum_{I=1}^{4} \frac{m_w^2}{m_{\tilde{\chi}_I^O A}^2} (Z_I^{D_L})^+_{2A} (Z_I^{D_L})_{A3} f_6 \left( \frac{m_{\tilde{\chi}_I^O A}^2}{m_{\tilde{\chi}_I^O A}^2} \right). \]  

(B8)

9) contribution from gluino loops with a z-boson coupling to the quark pair:

\[ Y_1 = \frac{16g_3^2}{3g_2^2 v_t^2 v_t^* v_t^*} \sum_{A,B=1}^{6} \frac{1}{(\Gamma^{D_L})^+_{2A} (\Gamma^{D_L})_{B3} c_2(m_{\tilde{\chi}_I^O}^2, m_{\tilde{\chi}_J^O}^2, m_{\tilde{\chi}_J^O}^2; \Gamma^{D_R \Gamma^{D_R^+}})_{AB}, \]  

\[ Y_2 = \frac{8g_3^2}{3g_2^2 v_t^2 v_t^* v_t^*} \left( \frac{1}{\sin^2 \theta_w} - 2 \right) \sum_{A,B=1}^{6} (\Gamma^{D_L})^+_{2A} (\Gamma^{D_L})_{B3} c_2(m_{\tilde{\chi}_I^O}^2, m_{\tilde{\chi}_J^O}^2, m_{\tilde{\chi}_J^O}^2; \Gamma^{D_R \Gamma^{D_R^+}})_{AB}. \]
10) contribution from gluino loops with a photon coupling to the quark pair:

\[ Y_1 = Y_2 = \frac{4g_3^2}{81g_2^2v_t^2v_b^2} \sum_{A=1}^{6} \frac{m_w^2}{m_\tilde{d}_A^2} (\Gamma_{D_{IL}}^{L})^{+}_{A2} (\Gamma_{U_{IL}}^{U})_{A3} f_6 \left( \frac{m_\tilde{g}^2}{m_\tilde{d}_A^2} \right). \]  

(B10)

The functions appearing in the eq.(B1)\cdots(B10) are given by

\[ f_1(x) = \frac{x}{1-x} + \frac{x}{(1-x)^2} \ln x, \]

\[ f_2(x) = \frac{38x - 79x^2 + 47x^3}{6(1-x)^3} + \frac{4x - 6x^2 + 3x^4}{(1-x)^4} \ln x, \]

\[ f_3(x) = \frac{16x - 29x^2 + 7x^3}{6(1-x)^3} + \frac{2x - 3x^2}{(1-x)^4} \ln x, \]

\[ f_4(x) = \frac{52 - 101x + 43x^2}{6(1-x)^3} + \frac{6 - 9x + 2x^3}{(1-x)^4} \ln x, \]

\[ f_5(x) = \frac{2 - 7x + 11x^2}{6(1-x)^3} + \frac{x^3}{(1-x)^4} \ln x, \]

\[ f_6(x) = \frac{2 - 7x + 11x^2}{(1-x)^3} + \frac{6x^3}{(1-x)^4} \ln x, \]

\[ c_0(m_1^2, m_2^2, m_3^2) = - \left( \frac{m_1^2 \ln(m_1^2/\mu^2)}{(m_1^2 - m_2^2)(m_1^2 - m_3^2)} + (m_1 \leftrightarrow m_2) + (m_1 \leftrightarrow m_3) \right), \]

\[ c_2(m_1^2, m_2^2, m_3^2) = \frac{3}{8} - \frac{1}{4} \left( \frac{m_1^2 \ln(m_1^2/\mu^2)}{(m_1^2 - m_2^2)(m_1^2 - m_3^2)} + (m_1 \leftrightarrow m_2) + (m_1 \leftrightarrow m_3) \right). \]

The symbol conventions in eq.(B1)\cdots(B10) are the same as in Ref. [19].
Figure Captions

Fig.1: The unitarity triangle of CKM matrix in the complex plane.

Fig.2: Diagrams contributing to the process $b \rightarrow c\bar{c}s$ in the two-Higgs-doublet model, which mediate $\bar{B}^0 \rightarrow J/\psi K_s$ decay.

Fig.3: Diagrams contributing to the process $b \rightarrow c\bar{c}s$ in the MSSM.
Fig. 1

\[ V_{ud} V_{ub} \quad \alpha \quad V_{cd} V_{cb} \quad \gamma \quad V_{td} V_{tb} \quad \beta \]
Fig. 2
Fig. 3