Probing the high scale SUSY in CP violations of $K$, $B^0$ and $B_s$ mesons

Morimitsu Tanimoto* and Kei Yamamoto†

Department of Physics, Niigata University, Niigata 950-2181, Japan

Abstract

We probe the high scale SUSY at $10 - 50$ TeV in the CP violations of $K$, $B^0$ and $B_s$ mesons. In order to estimate the contribution of the squark flavor mixing to these CP violations, we discuss the squark mass spectrum, which is consistent with the recent Higgs discovery. Taking the universal soft parameters at the SUSY breaking scale, we obtain the squark mass spectrum at 10 TeV and 50 TeV, where the SM emerges. Then, the $6 \times 6$ mixing matrix between down-squarks and down-quarks is discussed by input of the experimental data of $K$, $B^0$ and $B_s$ mesons. It is found that $\epsilon_K$ is most sensitive to the high scale SUSY. The SUSY contributions for the time-dependent CP asymmetries $S_{J/\psi K_S}$ and $S_{J/\psi\phi}$ are $6 - 8\%$ at the SUSY scale of 10 TeV. We also discuss the SUSY contribution to the chromo-EDM of the strange quark.

*E-mail address: tanimoto@muse.sc.niigata-u.ac.jp
†E-mail address: yamamoto@muse.sc.niigata-u.ac.jp
1 Introduction

Although the supersymmetry (SUSY) is one of the most attractive candidates for the new physics, the SUSY signals have not been observed yet. Therefore, the recent searches for new particle at the LHC give us important constraints for SUSY. Since the lower bounds of the superparticle masses increase gradually, the squark and the gluino masses are supposed to be at the higher scale than 1 TeV \[1\]. Moreover, the SUSY model has been seriously constrained by the Higgs discovery, in which the Higgs mass is 126 GeV \[2\].

These facts suggest a class of SUSY models with heavy sfermions. If the SUSY is broken with the breaking scale \(10^{-100}\) TeV, the squark and slepton masses are expected to be also \(\mathcal{O}(10^{-100})\) TeV. Then, the lightest Higgs mass can be pushed up to 126 GeV, while all SUSY particle can be out of the reach of the LHC experiment. Therefore, the indirect search of the SUSY particles becomes important in the low energy flavor physics \[3, 4\].

The flavor physics is on the new stage in the light of LHCb data. The LHCb collaboration has reported new data of the CP violation of the \(B_s\) meson and the branching ratios of rare \(B_s\) decays \[5\]-\[16\]. For many years the CP violation in the \(K\) and \(B^0\) mesons has been successfully understood within the framework of the standard model (SM), so called Kobayashi-Maskawa (KM) model \[17\], where the source of the CP violation is the KM phase in the quark sector with three families. However, the new physics has been expected to be indirectly discovered in the precise data of \(B^0\) and \(B_s\) meson decays at the LHCb experiment and the further coming experiment, Belle II.

While, there are new sources of the CP violation if the SM is extended to the SUSY models. The soft squark mass matrices contain the CP violating phases, which contribute to the flavor changing neutral current (FCNC) with the CP violation \[18\]. We can expect the SUSY effect in the CP violating phenomena. However, the clear deviation from the SM prediction has not been observed yet in the LHCb experiment \[5\]-\[16\]. Therefore, we should carefully study the CP-violation phenomena.

The LHCb collaboration presented the time dependent CP asymmetry in the non-leptonic \(B_s \to J/\psi \phi\) decay \[8, 15, 16\], which gives a constraint of the SUSY contribution on the \(b \to s\) transition. In this work, we discuss the sensitivity of the high scale SUSY contribution to the CP violation of \(K^0, B_d\) and \(B_s\) mesons. For these decay modes, the most important process of the SUSY contribution is the gluino-squark mediated flavor changing process \[19\]-\[34\]. This FCNC effect is constrained by the CP violations in \(B^0 \to J/\psi K_s\) and \(B_s \to J/\psi \phi\) decays. The CP violation of \(K\) meson, \(\epsilon_K\), also provides a severe constraint to the gluino-squark mediated FCNC. In the SM, \(\epsilon_K\) is proportional to \(\sin(2\phi_1)\) which is derived from the time dependent CP asymmetry in \(B^0 \to J/\psi K_s\) decay \[35\]. The relation between \(\epsilon_K\) and \(\sin(2\phi_1)\) is examined by taking account of the gluino-squark mediated FCNC \[36\].

The time dependent CP asymmetry of \(B^0 \to \phi K_S\) and \(B^0 \to \eta' K^0\) decays are also considered as typical processes to search for the gluino-squark mediated FCNC because the penguin amplitude dominates this process. Furthermore, we discuss the semileptonic CP asymmetries of \(B^0\) and \(B_s\) mesons, which can probe the SUSY contribution.

In addition, it is remarked that the upper-bound of the chromo-EDM(cEDM) of the strange quark gives a severe constraint for the gluino-squark mediated \(b \to s\) transition \[37\]-\[40\]. The recent work shows us that the cEDM is sensitive to the high scale SUSY \[41\].
In order to estimate the gluino-squark mediated FCNC of the $K$, $B^0$ and $B_s$ meson for arbitrary squark mass spectra, we work in the basis of the squark mass eigenstate. There are three reasons why the SUSY contribution to the FCNC considerably depends on the squark mass spectrum. The first one is that the GIM mechanism works in the squark flavor mixing, and the second one is that the loop functions depend on the mass ratio of squark and gluino. The last one is that we need the mixing angle between the left-handed sbottom and right-handed sbottom, which dominates the $\Delta B = 1$ decay processes. Therefore, we discuss the squark mass spectrum, which is consistent with the recent Higgs discovery. Taking the universal soft parameters at SUSY breaking scale, we obtain the squark mass spectrum at the matching scale where the SM emerges, by using the Renormalization Group Equations (REG’s) of the soft masses. Then, the $6 \times 6$ mixing matrix between down-squarks and down-quarks is examined by input of the experimental data.

In section 2, we discuss the squark and gluino spectra. In section 3, we present the formulation of the CP violation in terms of the squark flavor mixing, and we present our numerical results in section 4. Section 5 is devoted to the summary. Relevant formulations are presented in Appendices A, B and C.

## 2 SUSY Spectrum

We consider the SUSY model with heavy sfermions. If the squark and slepton masses are expected to be also $\mathcal{O}(10)$ TeV, the lightest Higgs mass can be pushed up to 126 GeV.

Let us obtain the SUSY particle mass spectrum in the framework of the minimal supersymmetric standard model (MSSM), which is consistent with the observed Higgs mass. The numerical analyses have been given in refs. [42, 43].

At the SUSY breaking scale $\Lambda$, the quadratic terms in the MSSM potential is given as

$$V_2 = m_1^2 |H_1|^2 + m_2^2 |H_2|^2 + m_3^2 (H_1 \cdot H_2 + h.c.) ,$$

(1)

where we define $m_1^2 = m_{H_1}^2 + |\mu|^2$ and $m_2^2 = m_{H_2}^2 + |\mu|^2$ in terms of the soft breaking mass $m_{H_i}$ and the supersymmetric Higgsino mass $\mu$. The mass eigenvalues at the $H_1$ and $\tilde{H}_2 \equiv \epsilon H_2^*$ system are given

$$m_+^2 = \frac{m_1^2 + m_2^2}{2} \mp \sqrt{\left(\frac{m_1^2 - m_2^2}{2}\right)^2 + m_3^4} .$$

(2)

Suppose that the MSSM matches with the SM at the SUSY mass scale $Q_0 \equiv m_0$. Then, the smaller one $m_-^2$ is identified to be the mass squared of the SM Higgs $H$ with the tachyonic mass. On the other hand, the larger one $m_+^2$ is the mass squared of the orthogonal combination $\mathcal{H}$, which is decoupled from the SM at $Q_0$, that is, $m_{\mathcal{H}} \simeq Q_0$. Therefore, we have

$$m_-^2 = - m_0^2 (Q_0) , \quad m_+^2 = m_{\mathcal{H}}^2 (Q_0) = m_1^2 + m_2^2 + m^2 ,$$

(3)

with

$$m_3^4 = (m_1^2 + m_2^2)(m_2^2 + m^2) ,$$

(4)
which lead to the mixing angle between $H_1$ and $\tilde{H}_2$, $\beta$ as
\[ \tan^2 \beta = \frac{m_1^2 + m_2^2}{m_2^2 + m_2^4}, \] (5)
where
\[ H = \cos \beta H_1 + \sin \beta \tilde{H}_2, \]
\[ \mathcal{H} = -\sin \beta H_1 + \cos \beta \tilde{H}_2. \] (6)
Thus, the Higgs mass parameter $m^2$ is expressed in terms of $m_1^2$, $m_2^2$ and $\tan \beta$:
\[ m^2 = \frac{m_1^2 - m_2^2 \tan^2 \beta}{\tan^2 \beta - 1}. \] (7)
Below the energy scale $Q_0$, in which the SM emerges, the scalar potential is just the SM one as follows:
\[ V_{SM} = -m^2 |H|^2 + \frac{\lambda_H}{2} |H|^4. \] (8)
Here, the Higgs coupling $\lambda_H$ is given in terms of the SUSY parameters as
\[ \lambda_H(Q_0) = \frac{1}{4} (g^2 + g'^2) \cos^2 2\beta + \frac{3h_t^2}{8\pi^2} X_i^2 \left( 1 - \frac{X_i^2}{12} \right), \] (9)
where
\[ X_i = \frac{A_i(Q_0) - \mu(Q_0) \cot \beta}{Q_0}, \] (10)
and $h_t$ is the top Yukawa coupling of the SM. The parameters $m_2$ and $\lambda_H$ run with the SM Renormalization Group Equation (RGE) down to the electroweak scale $Q_{EW} = m_H$, and then give
\[ m_H^2 = 2m_H^2 = \lambda_H(m_H) v^2. \] (11)
It is easily seen that the VEV of Higgs, $\langle H \rangle$ is $v$, and $\langle \mathcal{H} \rangle = 0$, taking account of $\langle H_1 \rangle = v \cos \beta$ and $\langle \tilde{H}_2 \rangle = v \sin \beta$, where $v = 246\text{GeV}$.

Let us fix $m_H = 126\text{GeV}$, which gives $\lambda_H(Q_0)$ and $m^2(Q_0)$. This experimental input constrains the SUSY mass spectrum of MSSM. We consider the some universal soft breaking parameters at the SUSY breaking scale $\Lambda$ as follows:
\[ m_{Q_i}^2(\Lambda) = m_{U_i}^2(\Lambda) = m_{D_i}^2(\Lambda) = m_0^2 \quad (i = 1, 2, 3), \]
\[ M_1(\Lambda) = M_2(\Lambda) = M_3(\Lambda) = m_{1/2}, \]
\[ m_{H_1}^2(\Lambda) = m_{H_2}^2(\Lambda) = m_0^2, \]
\[ A_U(\Lambda) = A_0 y_U(\Lambda), \quad A_D(\Lambda) = A_0 y_D(\Lambda), \quad A_E(\Lambda) = A_0 y_E(\Lambda). \] (12)
Then, there is no flavor mixing at this scale if the universal soft masses are exactly satisfied. Different RGE effects for each flavor evolve the squark flavor mixing at the lower energy scale, which is controlled by the CKM mixing matrix. Since we take squark flavor mixing
as free parameters at the low energy, this universality condition has to be considered as an approximation and non-vanishing off diagonal squark mass matrix elements are introduced at the $\Lambda$ scale. We will show typical magnitudes of those off-diagonal elements in the numerical result to understand the level of our approximation in the numerical result.

Now, we have the SUSY five parameters, $\Lambda$, $\tan \beta$, $m_0$, $m_{1/2}$, $A_0$, where $Q_0 = m_0$. In addition to these parameters, we take $\mu = Q_0$. Inputing $m_H = 126$ GeV and taking $m_H \simeq Q_0$, we can obtain the SUSY spectrum for the fixed $Q_0$ and $\tan \beta$.

We consider the two case of $Q_0 = 10$ TeV and 50 TeV. The parameter set of the first case (a) is given as

$\Lambda = 10^{17}$ GeV, $Q_0 = m_0 = 10$ TeV, $m_{1/2} = 6.2$ TeV, $\tan \beta = 10$, $A_0 = 25.803$ TeV. (13)

Here $m_{1/2}$ and $A_0$ are tuned in order to obtain the proper $\lambda_H$ with the small $X_t(A_t)$, which gives $m_H = 126$ GeV at the electroweak $m_H$ scale. The parameter set of the second case (b) is given as

$\Lambda = 10^{16}$ GeV, $Q_0 = m_0 = 50$ TeV, $m_{1/2} = 63.5$ TeV, $\tan \beta = 4$, $A_0 = 109.993$ TeV. (14)

These parameter sets are easily found following from the numerical work in Ref. [12]. The obtained SUSY mass spectra at $Q_0$ are summarized in Table 1, where the top mass is sensitive to give the Higgs mass, and we use $m_t(m_t) = 163.5 \pm 2$ GeV [44, 45]. For the case (a), we show the running of SUSY masses in the MSSM from $\Lambda$ down to $Q_0$ in Figure 1 [46].

As seen in Table 1, the first and second family squarks are degenerate in their masses, on the other hand, the third ones split due to the large RGE’s effect. Therefore, the mixing angle between the first and second family squarks vanishes, but the mixing angles between the first-third and the second-third family squarks are produced at the $Q_0$ scale. The left-right mixing angle between $\tilde{b}_L$ and $\tilde{b}_R$ is given as

$$\theta \simeq \frac{m_b(A_b(Q_0) - \mu \tan \beta)}{m^2_{\tilde{b}_L} - m^2_{\tilde{b}_R}}.$$ (15)

It is noticed that the right-handed sbottom is heavier than the left-handed one. The lightest squark is the right-handed stop. Since we take the universal mass assumption for gauginos, $m_{1/2}$, the lightest gaugino is the Bino, $\tilde{B}$, whose mass is 2.9 TeV in the case of $Q_0 = 10$ TeV. That is the lightest supersymmetric particle (LSP) in our framework. Although these Wino and Bino mass values are consistent with the recent experimental result of searching for EW-gaugino [47], the Bino cannot be a candidate of the dark matter in this case [48, 49]. In order to get the Wino dark matter, we should relax the universal mass assumption for gauginos. However, this study does not affect our following numerical results of the CP violation, we do not discuss about the dark matter any more in this work.

3 Squark flavor mixing and CP violation

3.1 Squark flavor mixing

Let us consider the $6 \times 6$ squark mass matrix $M_{\tilde{q}}$ in the super-CKM basis. In order to move the mass eigenstate basis of squark masses, we should diagonalize the mass matrix by rotation
matrix $\Gamma^{(q)}_G$ as

$$m_q^2 = \Gamma^{(q)}_G M^2_q \Gamma^{(q)^*}_G,$$  \hspace{1cm} \text{(16)}$$

where $\Gamma^{(q)}_G$ is the $6 \times 6$ unitary matrix, and we decompose it into the $3 \times 6$ matrices as $\Gamma^{(q)}_G = (\Gamma^{(q)}_{GL}, \Gamma^{(q)}_{GR})^T$ in the following expressions:

$$\Gamma^{(d)}_{GL} = 
\begin{pmatrix}
  c_{13}^L & 0 & s_{13}^L e^{-i\phi_{13}^L} c_{\theta} & 0 & -s_{13}^L e^{-i\phi_{13}^L} s_{\theta} e^{i\phi} \\
  -s_{23}^L s_{13}^L e^{i(\phi_{13}^L - \phi_{23}^L)} & c_{23}^L & s_{23}^L c_{13}^L e^{-i\phi_{23}^L} & 0 & -s_{23}^L c_{13}^L s_{\theta} e^{i\phi} \\
  -s_{13}^L c_{23}^L e^{i\phi_{13}^L} & -s_{23}^L e^{i\phi_{23}^L} c_{\theta} & c_{13}^L c_{23}^L c_{\theta} & 0 & -c_{13}^L c_{23}^L s_{\theta} e^{i\phi}
\end{pmatrix},$$

$$\Gamma^{(d)}_{GR} = 
\begin{pmatrix}
  0 & 0 & s_{13}^R s_{\theta} e^{-i\phi_{13}^R} e^{i\phi} & 0 & s_{13}^R e^{-i\phi_{13}^R} c_{\theta} \\
  0 & 0 & s_{23}^R c_{13}^R s_{\theta} e^{-i\phi_{23}^R} e^{i\phi} & -s_{23}^R c_{13}^R e^{i(\phi_{23}^R - \phi_{23}^L)} & c_{23}^R & s_{23}^R c_{13}^R c_{\theta} e^{i\phi} \\
  0 & 0 & c_{13}^R c_{23}^R s_{\theta} e^{-i\phi} & -s_{23}^R c_{13}^R c_{\theta} e^{i\phi_{23}^R} & -s_{23}^R c_{13}^R e^{i\phi_{23}^R} & c_{13}^R c_{23}^R c_{\theta}
\end{pmatrix},$$  \hspace{1cm} \text{(17)}$$

where we use abbreviations $c_{ij}^{L,R} = \cos \theta_{ij}^{L,R}$, $s_{ij}^{L,R} = \sin \theta_{ij}^{L,R}$, $c_{\theta} = \cos \theta$ and $s_{\theta} = \sin \theta$ in Eq. (15). Here $\theta$ is the left-right mixing angle between $\bar{b}_L$ and $\bar{b}_R$. It is remarked that we take $s_{12}^{L,R} = 0$ due to the degenerate squark masses of the first and second families as discussed in the previous section.

The gluino-squark-quark interaction is given as

$$\mathcal{L}_{\text{int}}(\bar{g} q q) = -i \sqrt{2} g_s \sum_{\langle q \rangle} \bar{q}_i^a (T^a) \tilde{G}^a \left[ (\Gamma^{(q)}_{GL})_{ij} P_L + (\Gamma^{(q)}_{GR})_{ij} P_R \right] q_j + \text{h.c.},$$  \hspace{1cm} \text{(18)}$$

where $P_L = (1 - \gamma_5)/2$, $P_R = (1 + \gamma_5)/2$, and $\tilde{G}^a$ denotes the gluino field, $q^i$ are three left-handed ($i=1,2,3$) and three right-handed quarks ($i=4,5,6$). This interaction leads to the gluino-squark mediated flavor changing process with $\Delta F = 2$ and $\Delta F = 1$ through the box and penguin diagrams.
Table 1: Input parameters at Λ and obtained the SUSY spectra in the cases of (a) and (b).

| Case (a) | Case (b) |
|----------|----------|
| at Λ = 10^{17} \text{ GeV}, & at Λ = 10^{16} \text{ GeV}, |
| m_0 = 10 \text{ TeV}, & m_0 = 50 \text{ TeV}, |
| m_{1/2} = 6.2 \text{ TeV}, & m_{1/2} = 63.5 \text{ TeV}, |
| A_0 = 25.803 \text{ TeV}; & A_0 = 109.993 \text{ TeV}; |
| at Q_0 = 10 \text{ TeV}, & at Q_0 = 50 \text{ TeV}, |
| μ = 10 \text{ TeV}, & μ = 50 \text{ TeV}, |
| tan β = 10, & tan β = 4 |
| m_{\tilde{q}} = 12.8 \text{ TeV}, & m_{\tilde{q}} = 115.6 \text{ TeV}, |
| m_{\tilde{W}} = 5.2 \text{ TeV}, & m_{\tilde{W}} = 55.4 \text{ TeV}, |
| m_{\tilde{B}} = 2.9 \text{ TeV} |
| m_{\tilde{t}_L} = m_{\tilde{t}_L} = 12.2 \text{ TeV} |
| m_{\tilde{b}_R} = 14.1 \text{ TeV}, & m_{\tilde{b}_R} = 104.0 \text{ TeV}, |
| m_{\tilde{t}_R} = 8.4 \text{ TeV}, & m_{\tilde{t}_R} = 83.2 \text{ TeV} |
| m_{\tilde{\chi}_L \tilde{\chi}_L} = m_{\tilde{\chi}_L \tilde{\chi}_L} = 15.1 \text{ TeV} |
| m_{\tilde{\chi}_R \tilde{\chi}_R} ≃ m_{\tilde{\chi}_R \tilde{\chi}_R} = 14.6 \text{ TeV}, & m_{\tilde{\chi}_R \tilde{\chi}_R} = 110.7 \text{ TeV}, |
| m_{\tilde{\chi}_R \tilde{\chi}_R} = 110.7 \text{ TeV} |
| A_t = -1.2 \text{ TeV}, & A_t = -20.2 \text{ TeV}, |
| A_b = 5.1 \text{ TeV}, & A_b = 4.7 \text{ TeV}, |
| X_t = -0.22, & X_t = -0.65 |
| λ_H = 0.126, \ θ = 0.35^o |
| λ_H = 0.1007, \ θ = 0.05^o |

3.2 CP violation in ΔF = 2 and ΔF = 1 processes

Taking account of the gluino-squark interaction, the dispersive part of meson mixing \( M_{12}^D(P = K, B^0, B_s) \) are given as

\[
M_{12}^D = M_{12}^{SM} + M_{12}^{SU3},
\]

where \( M_{12}^{SU3} \) are written by SUSY parameters in Eq.(17) and its explicit formulation is given in Appendix A. The experimental data of ΔB = 2 process, the mass differences Δ\( M_{B^0} \) and Δ\( M_{B_S} \), and the CP-violating phases \( \phi_B \) and \( \phi_s \), give constraint to the SUSY parameters in Eq.(17). We also consider the constraint from the CP-violating parameter in the K meson, \( \epsilon_K \), and focus on the relation between \( \epsilon_K \) and \( \sin(2\beta) \), in which β is one angle of the unitarity triangle with respect to \( B^0 \).

The indirect CP asymmetry in the semileptonic decays \( B_q \rightarrow \mu^-X (q = d, s) \) leads to the nonzero asymmetry \( a_{sl}^q \) such as:

\[
a_{sl}^q = \frac{\Gamma(\bar{B}_q \rightarrow \mu^+X) - \Gamma(B_q \rightarrow \mu^-X)}{\Gamma(\bar{B}_q \rightarrow \mu^+X) + \Gamma(B_q \rightarrow \mu^-X)} = \text{Im} \left( \frac{\Gamma_{12}^q}{M_{12}^q} \right) \sim \frac{|\Gamma_{12}^q|}{|M_{12}^q|} \sin \phi_{sl}^q. \tag{20}
\]

The absorptive part of \( B_q - \bar{B}_q \) system \( \Gamma_{12}^q \) is dominated by the tree-level decay \( b \rightarrow c\bar{c}s \) etc in the SM. Therefore, we assume \( \Gamma_{12}^q = \Gamma_{12}^{SM} \) in our calculation. In the SM, the CP-violating phases are read \[{}^{[50]}\),

\[
\phi_{sl}^{SM} = (3.84 \pm 1.05) \times 10^{-3}, \quad \phi_{sl}^{SM} = -(7.50 \pm 2.44) \times 10^{-2},
\]

which correspond to

\[
a_{sl}^{SM} = (1.9 \pm 0.3) \times 10^{-5}, \quad \phi_{sl}^{SM} = -(4.1 \pm 0.6) \times 10^{-4}. \tag{22}
\]
The recent experimental data of these CP asymmetries are given as \[12, 45\]

\[a^s_{sl} = (-0.24 \pm 0.54 \pm 0.33) \times 10^{-2}, \quad a^d_{sl} = (-0.3 \pm 2.1) \times 10^{-3}. \quad (23)\]

The time dependent CP asymmetries in non-leptonic decays are also interesting to search for the SUSY effect. The \(\Delta B = 1\) transition amplitude is estimated by the effective Hamiltonian given as follows:

\[H_{\text{eff}} = \frac{4G_F}{\sqrt{2}} \left[ \sum_{q'=u,c} V_{q'b} V_{q's}^* \sum_{i=1,2} C_i O_i^{(q')} - V_{tb} V_{tq}^* \sum_{i=3-6,7\gamma;8G} (C_i O_i + \tilde{C}_i \tilde{O}_i) \right], \quad (24)\]

where \(q = s, d\). The local operators are given as:

\[O_1^{(q')} = (\bar{q}_i \gamma_\mu P_L q'_\beta)(\bar{q}'_\beta \gamma^\mu P_L b_\alpha), \quad O_2^{(q')} = (\bar{q}_i \gamma_\mu P_L q'_\beta)(\bar{q}'_\beta \gamma^\mu P_L b_\beta), \]

\[O_3 = (\bar{q}_i \gamma_\mu P_L b_\beta) \sum_{Q} (\bar{Q} Q) \gamma_\mu P_L b_\beta, \quad O_4 = (\bar{q}_i \gamma_\mu P_L b_\beta) \sum_{Q} (\bar{Q} Q) \gamma_\mu P_L Q_\alpha, \]

\[O_5 = (\bar{q}_i \gamma_\mu P_L b_\beta) \sum_{Q} (\bar{Q} Q) \gamma_\mu P_L R Q_\beta, \quad O_6 = (\bar{q}_i \gamma_\mu P_L b_\beta) \sum_{Q} (\bar{Q} Q) \gamma_\mu P_L R Q_\alpha, \]

\[O_{7\gamma} = \frac{e}{16\pi^2} m_b \bar{q}_i \sigma^{\mu \nu} P_R b_\beta F_{\mu \nu}, \quad O_{8G} = \frac{g_s}{16\pi^2} m_b \bar{q}_i \sigma^{\mu \nu} P_R T_{\alpha \beta} b_\beta G^a_{\mu \nu}, \quad (25)\]

where \(\alpha, \beta\) are color indices, and \(Q\) is taken to be \(u, d, s, c\) quarks. Here, the \(C_i\) is the Wilson coefficient and includes SM contribution and gluino-squark one, such as \(C_i = C_i^{\text{SM}} + C_i^{\tilde{g}}\). The \(C_i^{\text{SM}}\) is given in Ref. \[51\]. The terms \(\tilde{C}_i\) and \(\tilde{O}_i\) are obtained by replacing L(R) with R(L). The magnetic penguin contribution \(C_{7\gamma}\) and \(C_{8G}\) can be enhanced due to the left-right mixing. For the \(b \rightarrow s\) transition, the gluino contributions to these the Wilson coefficients, \(C_{7\gamma}\) and \(C_{8G}\), are given as follows:

\[C_{7\gamma}(m_{\tilde{g}}) = \frac{8 \sqrt{2} \alpha_s \pi}{3 2 G_F V_{tb} V_{ts}} \times \sum_{I=1}^{6} \frac{m_{\tilde{g}}}{m_{\tilde{d}_I}} \left\{ \left( \frac{\Gamma_{dI}^{(G)}}{\Gamma_{dI}^{(d)}} \right)^* \left( \frac{\Gamma_{dI}^{(G)}}{\Gamma_{dI}^{(d)}} \right) \right\}, \quad (26)\]

\[C_{8G}(m_{\tilde{g}}) = \frac{8 \sqrt{2} \alpha_s \pi}{3 2 G_F V_{tb} V_{ts}} \sum_{I=1}^{6} \frac{m_{\tilde{g}}}{m_{\tilde{d}_I}} \left\{ \left( \frac{\Gamma_{dI}^{(G)}}{\Gamma_{dI}^{(d)}} \right)^* \right\}, \quad (27)\]

where \(F_{I}(x_{\tilde{g}}^I)\) are the loop functions given in Appendix B with \(x_{\tilde{g}}^I = m_{\tilde{g}}^2/m_{\tilde{d}_I}^2\) \((I = 1 - 6)\). We estimate \(C_{7\gamma}\) and \(C_{8G}\) at the \(m_b\) scale including the effect of the leading order of QCD as follows \[51\]:

\[C_{7\gamma}(m_b) = \zeta C_{7\gamma}(m_{\tilde{g}}) + \frac{8}{3} (\eta - \zeta) C_{8G}(m_{\tilde{g}}), \quad (28)\]

\[C_{8G}(m_b) = \eta C_{8G}(m_{\tilde{g}}), \]

\[C_{7\gamma}(m_b) = \zeta C_{7\gamma}(m_{\tilde{g}}) + \frac{8}{3} (\eta - \zeta) C_{8G}(m_{\tilde{g}}), \quad (28)\]
where
\[
\zeta = \left( \frac{\alpha_s(m_{\tilde{g}})}{\alpha_s(m_{\tilde{b}})} \right)^{16} \left( \frac{\alpha_s(m_{\tilde{b}})}{\alpha_s(m_{\tilde{g}})} \right)^{15} \left( \frac{\alpha_s(m_{\tilde{g}})}{\alpha_s(m_{\tilde{t}})} \right)^{16} \left( \frac{\alpha_s(m_{\tilde{t}})}{\alpha_s(m_{\tilde{b}})} \right)^{21},
\]
\[
\eta = \left( \frac{\alpha_s(m_{\tilde{g}})}{\alpha_s(m_{\tilde{b}})} \right)^{14} \left( \frac{\alpha_s(m_{\tilde{b}})}{\alpha_s(m_{\tilde{g}})} \right)^{15} \left( \frac{\alpha_s(m_{\tilde{g}})}{\alpha_s(m_{\tilde{t}})} \right)^{14} \left( \frac{\alpha_s(m_{\tilde{t}})}{\alpha_s(m_{\tilde{b}})} \right)^{21},
\]
(29)

In the expression of Eq.(29), the QCD correction is taken into account for the case of the gluino mass being much smaller than the squark one [52].

Now that we discuss the time dependent CP asymmetries of \(B_0^0\) and \(B_s\) decaying into the final state \(f\), which are defined as [53]:
\[
S_f = \frac{2 \text{Im} \lambda_f}{1 + |\lambda_f|^2},
\]
(30)
where
\[
\lambda_f = \frac{q}{p} \frac{\bar{A}(B_q^0 \rightarrow f)}{A(B_q^0 \rightarrow f)}, \quad \frac{q}{p} \approx \sqrt{\frac{M_{q12}^*}{M_{q12}^2}},
\]
(31)
where \(A(B_q^0 \rightarrow f)\) is the decay amplitude in \(B_q^0 \rightarrow f\). The time-dependent CP asymmetries \(S_f\) are mixing induced CP asymmetry, where \(M_{q12}^*\) and \(A(B_q^0 \rightarrow f)\) include the SUSY contributions in addition to the SM one.

The time-dependent CP asymmetries in the \(B^0 \rightarrow J/\psi K_S\) and \(B_s \rightarrow J/\psi \phi\) decays are well known as the typical decay mode to determine the unitarity triangle. In this decays, we write \(\lambda_{J/\psi K_S}\) and \(\lambda_{J/\psi \phi}\) in terms of phase factors, respectively:
\[
\lambda_{J/\psi K_S} \equiv -e^{-i\phi_d}, \quad \lambda_{J/\psi \phi} \equiv e^{-i\phi_s}.
\]
(32)

In the SM, the phase \(\phi_d\) is given in terms of the angle of the unitarity triangle \(\phi_1\) as \(\phi_d = 2\phi_1\). On the other hand, \(\phi_s\) is given as \(\phi_s = -2\beta_s\), in which \(\beta_s\) is the one angle of the unitarity triangle in \(B_s\). Once \(\phi_d\) is input, \(\phi_s\) in the SM is predicted as [54]
\[
\phi_s = -0.0363 \pm 0.0017.
\]
(33)

If the SUSY contribution is non-negligible, \(\phi_d = 2\phi_1\) and \(\phi_s = -2\beta_s\) are not satisfied any more.

The recent experimental data of these phases are [8 55]
\[
\sin \phi_d = 0.679 \pm 0.020, \quad \phi_s = 0.07 \pm 0.09 \pm 0.01.
\]
(34)

These experimental values also constrain the mixing angles and phases in Eq.(17).

The \(b \rightarrow s\) transition is one-loop suppressed one in the SM, so the SUSY contribution to this process is expected to be sizable. In this point of view, we focus on the CP asymmetries in the \(b \rightarrow s\) transition, \(B^0 \rightarrow \phi K_S\) and \(B^0 \rightarrow \eta' K^0\). The CP asymmetries of \(B^0 \rightarrow \phi K_S\) and \(B^0 \rightarrow \eta' K^0\) have been studied for these twenty years [56 57 58]. In the SM, \(S_{\phi K_S}\) and \(S_{\eta' K^0}\) are same to \(S_{J/\psi K_S}\) within roughly 10% accuracy because the CP phase comes from mixing \(M_{12}^*\) in these mode. Once taking account of the new physics contribution, the \(S_{\phi K_S}\)
and $S_{\eta'K^0}$ are expected to be deviated from $S_{J/\psi K_S}$ because $B^0 \to J/\psi K_S$ is the tree-level decay whereas $B^0 \to \phi K_S$ and $B^0 \to \eta'K^0$ are one-loop suppressed one in the SM. Recent experimental fit results of these CP asymmetries are reported by HFAG as follows \cite{HFAG}:

$$S_{J/\psi K_S} = 0.679 \pm 0.020, \quad S_{\phi K_S} = 0.74^{+0.11}_{-0.13}, \quad S_{\eta'K^0} = 0.59 \pm 0.07. \quad (35)$$

These values may be regarded to be same within experimental error-bar and consistent with the SM prediction. In other words, these experimental results give severe constraints to the squark flavor mixing angle between the second-third families.

The CP asymmetries in $B^0 \to \phi K_S$ and $B^0 \to \eta'K^0$ containing the SUSY contribution are estimated in terms of $\lambda_f$ in Eq. (31):

$$\lambda_{\phi K_S, \eta'K^0} = -e^{-i\phi_4} \frac{\sum_{i=3,6,7,8G} \left( C_i \langle O_i \rangle + \tilde{C}_i \langle \tilde{O}_i \rangle \right)}{\sum_{i=3,6,7,8G} \left( C_i^* \langle O_i \rangle + \tilde{C}_i^* \langle \tilde{O}_i \rangle \right)}, \quad (36)$$

where $\langle O_i \rangle$ is the abbreviation for $\langle f | O_i | B^0 \rangle$. It is known that $\langle \phi K_S | O_i | B^0 \rangle = \langle \phi K_S | \tilde{O}_i | B^0 \rangle$ and $\langle \eta'K^0 | O_i | B^0 \rangle = -\langle \eta'K^0 | \tilde{O}_i | B^0 \rangle$, because these final states have different parities \cite{56,57,58}. Then, the decay amplitudes of $f = \phi K_S$ and $f = \eta'K^0$ are written in terms of the dominant gluon penguin ones $C_{SG}$ and $\tilde{C}_{SG}$ as follows:

$$\tilde{A}(B^0 \to \phi K_S) \propto C_{SG}(m_b) + \tilde{C}_{SG}(m_b),$$

$$\tilde{A}(B^0 \to \eta'K^0) \propto C_{SG}(m_b) - \tilde{C}_{SG}(m_b). \quad (37)$$

Since $\tilde{C}_{SG}(m_b)$ is suppressed compared to $C_{SG}(m_b)$ in the SM, the magnitudes of the time dependent CP asymmetries $S_f$ ($f = J/\psi K_S, \phi K_S, \eta'K^0$) are almost same in the SM prediction. If the squark flavor mixing gives the unsuppressed $\tilde{C}_{SG}(m_b)$, these CP asymmetries are expected to be deviated among them.

In order to obtain precise results, we also take account of the small contributions from other Wilson coefficients $C_i$ ($i = 3, 4, 5, 6$) and $\tilde{C}_i$ ($i = 3, 4, 5, 6$) in our calculations. We estimate each hadronic matrix element by using the factorization relations in Ref. \cite{59}:

$$\langle O_3 \rangle = \langle O_4 \rangle = \left( 1 + \frac{1}{N_c} \right) \langle O_5 \rangle, \quad \langle O_6 \rangle = \langle O_5 \rangle, \quad \langle O_{SG} \rangle = \frac{a_s(m_b)}{8\pi} \left( -\frac{2m_b}{\sqrt{\langle q^2 \rangle}} \right) \left( \langle O_4 \rangle + \langle O_6 \rangle - \frac{1}{N_c} (\langle O_4 \rangle + \langle O_5 \rangle) \right), \quad (38)$$

where $\langle q^2 \rangle = 6.3$ GeV$^2$ and $N_c = 3$ is the number of colors. One may worry about the reliability of these naive factorization relations. However this approximation has been justified numerically in the relevant $b \to s$ transition as seen in the calculation of PQCD \cite{60}.

We also consider the SUSY contribution for the $b \to s\gamma$ decay. The $b \to s\gamma$ is sensitive to the magnetic penguin contribution $C_{7\gamma}$. The branching ratio $BR(b \to s\gamma)$ is given as \cite{61}

$$\frac{BR(b \to s\gamma)}{BR(b \to ce\bar{\nu}_e)} = \frac{|V_{ts}V_{tb}|^2}{|V_{cb}|^2} \frac{6\alpha}{\pi f(z)} \left( |C_{7\gamma}(m_b)|^2 + |\tilde{C}_{7\gamma}(m_b)|^2 \right), \quad (39)$$

9
where
\[
f(z) = 1 - 8z + 8z^3 - z^4 - 12z^2 \ln z, \quad z = \frac{m_{c,\text{pole}}^2}{m_{b,\text{pole}}^2}.
\] (40)

The SM prediction including the next-to-next-to-leading order correction is given as [62]
\[
\text{BR}(b \to s\gamma)\text{(SM)} = (3.15 \pm 0.23) \times 10^{-4},
\] (41)
on the other hand, the experimental data are obtained as [45]
\[
\text{BR}(b \to s\gamma)\text{(exp)} = (3.53 \pm 0.24) \times 10^{-4}.
\] (42)

Therefore, we can examine the contribution of the gluino-squark mediated flavor-changing process to the \(b \to s\gamma\) process.

In our analysis we also discuss the relation between \(\epsilon_K\) and \(\sin 2\phi_1\), where \(\phi_1\) is the one angle of the unitarity triangle. The parameter \(\epsilon_K\) is given in the following theoretical formula
\[
\epsilon_K = e^{i\phi_1} \sin \phi_\epsilon \left( \frac{\text{Im}(M_{12}^K)}{\Delta M_K} + \xi \right), \quad \xi = \frac{\text{Im}A_0^K}{\text{Re}A_0^K}, \quad \phi_\epsilon = \tan^{-1} \left( \frac{2\Delta M_K}{\Delta \Gamma_K} \right),
\] (43)
with \(A_0^K\) being the isospin zero amplitude in \(K \to \pi\pi\) decays. Here, \(M_{12}^K\) is the dispersive part of the \(K^0 - \bar{K}^0\) mixing, and \(\Delta M_K\) is the mass difference in the neutral \(K\) meson. The effects of \(\xi \neq 0\) and \(\phi_\epsilon < \pi/4\) give suppression effect in \(\epsilon_K\), and it is parameterized as \(\kappa_\epsilon\) and estimated by Buras and Guadagnoli [35] as:
\[
\kappa_\epsilon = 0.92 \pm 0.02 .
\] (44)
The \(|\epsilon_{SM}^K|\) is given in terms of the Wolfenstein parameters \(\lambda, \bar{\rho}\) and \(\eta\) as follows:
\[
|\epsilon_{SM}^K| = \kappa_\epsilon C_\epsilon \hat{B}_K|V_{cb}|^2 \lambda^2 \bar{\eta} (|V_{cb}|^2 (1 - \bar{\rho})\eta_{tt} E(x_t) - \eta_{ct} E(x_c) + \eta_{ct} E(x_c, x_t))
\] (45)
with
\[
C_\epsilon = \frac{G_F^2 F_K^2 m_K M_W^2}{6\sqrt{2}\pi^2 \Delta M_K}.
\] (46)

It is easily found that \(|\epsilon_{SM}^K|\) is proportional to \(\sin(2\phi_1)\) because there is only one CP violating phase in the SM. Therefore, the observed value of \(S_{J/\psi K_S}\), which correspond to \(\sin(2\phi_1)\), should be correlated with \(|\epsilon_K|\) in the SM. According to the recent experimental results, it is found that the consistency between the SM prediction and the experimental data in \(\sin(2\phi_1)\) and \(|\epsilon_{SM}^K/\hat{B}_K|\) is marginal. This fact was pointed out by Buras and Guadagnoli [35] and called as the tension between \(|\epsilon_K|\) and \(\sin(2\phi_1)\). Note that \(|\epsilon_{SM}^K|\) also depends on the non-perturbative parameter \(\hat{B}_K\) in Eq.(43). Recently, the error of this parameter shrank dramatically in the lattice calculations [63]. In our calculation we use the updated value by the Flavor Lattice Averaging Group [64]:
\[
\hat{B}_K = 0.766 \pm 0.010 .
\] (47)
We can calculate $|\epsilon_K^{\text{SM}}|$ for the fixed $\sin(2\phi_1)$ by inputting this value.

Considering the effect of the squark flavor mixing in both $|\epsilon_K|$ and $S_{J/\psi K_S}$, this tension can be relaxed though the gluino-squark interaction. Then, $\epsilon_K$ is expressed as:

$$\epsilon_K = \epsilon_K^{\text{SM}} + \epsilon_K^{\text{SUSY}},$$

(48)

where $\epsilon_K^{\text{SUSY}}$ is induced by the imaginary part of the gluino-squark box diagram, which is presented in Appendix A. Since $s_{12}^{L(R)}$ vanishes in our scheme, $\epsilon_K$ is given in the second order of the squark mixing $s_{13}^{L(R)} \times s_{23}^{L(R)}$.

In addition to the above CP violating processes, the neutron EDM is also sensitive to the CP-violating phase of the squark mixing through cEDM of the strange quark. The experimental upper bound of the electric dipole moment of the neutron provides us the upper-bound of cEDM of the strange quark [37]-[40]. The cEDM of the strange quark $d_s^C$ comes from the gluino-squark interactions is given in Appendix C. The bound on the cEDM of the strange quark is estimated as [40] from the experimental upper bound of the neutron EDM as follows:

$$|d_s^C| < 0.5 \times 10^{-25} \text{ ecm.}$$

(49)

This bound also give severe constraints for phases of the mixing parameters of Eq.(17).

4 Numerical results

In this section we show our numerical results. At the first step, we constrain the squark flavor mixing parameters in Eq. (17) from the experimental data of the CP violation $\epsilon_K$, $\phi_q$ and $\phi_s$, and the mass difference $\Delta M_{B^0}$ and $\Delta M_{B_s}$ comprehensively. We have nine free parameters, in which there are four mixing angles $\theta_{13}^{L(R)}$ and $\theta_{23}^{L(R)}$, five phase $\phi_{13}^{L(R)}, \phi_{23}^{L(R)}, \phi$. In our analyses, we reduce the number of parameters by taking $\theta_{ij}^{\ell} = \theta_{ij}^{R}$ for simplicity, but we also discuss the case where this assumption is broken in the estimate of $\epsilon_K$ and the cEDM of the strange quark. Moreover, Wolfenstein parameters $\bar{\rho}, \bar{\eta}$ are free ones, which are determined by our numerical analyses. Other relevant input parameters such as quark masses $m_c, m_b$, the CKM matrix elements $V_{us}, V_{cb}$ and $f_B, f_K$, etc. are shown in our previous paper Ref. [33], which are referred from the PDG [45] and the UTfit Collaboration [44].

The uncertainties of these input parameters determine the predicted range of the SUSY contribution for the CP violations, $\Delta M_{B^0}$ and $\Delta M_{B_s}$. For example, the predicted range of the SUSY contribution for $\epsilon_K$ mainly comes from the uncertainties of $\hat{B}_K$, $|V_{cb}|$ and $m_t$ in addition to the observed error bar of $|\epsilon_K|$. If these uncertainties will be reduced in the future, the predicted range of the CP violation is improved.

At the second step, we predict the deviations of the time dependent CP asymmetries $S_f$ and the semileptonic CP asymmetries $a_q^q(q = d, s)$ from the SM taking account of the contribution of the gluino-squark interaction. The SUSY effect on the cEDM of the strange quark is also discussed.

In our analysis, we scan the mixing angles $s_{ij}^{L(R)}$ and phases in Eq. (17) in the region of $0 \sim 0.5$ and $0 \sim 2\pi$, respectively. At first, we show the analysis in the case of the SUSY scale $Q_0 = 10$ TeV in detail, and then, we also discuss the numerical results in the case $Q_0 = 50$ TeV.
Let us start with discussing the gluino-squark interaction effect on the $\Delta F = 2$ processes, $\epsilon_K$, $\Delta M_{B^0}$ and $\Delta M_{B_s}$, where the squark and gluino mass spectrum in Table 1 is input. We show the allowed region on the plane of $\sin(2\phi_1)$ and $|\epsilon_K^{SM}/\bar{B}_K|$ in Fig. 2. When we add the contribution of the gluino-squark interaction, $\epsilon_K^{SUSY}$, the allowed region of $\sin(2\phi_1)$ and $|\epsilon_K^{SM}/\bar{B}_K|$ converge within the experimental error-bar, where $\phi_d$ is not $2\phi_1$ any more as discussed below Eq.(32). The Figure 3 shows the $s_{13}^{L(R)} \times s_{23}^{L(R)}$ dependence of the SUSY contribution for $\epsilon_K$, that is $|\epsilon_K^{SUSY}/\epsilon_K|$. It is found that the SUSY contribution could be large up to 40%. It is remarked that $\epsilon_K$ is sensitive to the gluino-squark interaction even if the SUSY scale is 10 TeV.

We show the SUSY contribution to the CP violating phase $\phi_d$ versus $s_{13}^{L(R)}$ in Figure 4, where we define $\Delta \sin 2\phi_1 \equiv \sin \phi_d - \sin 2\phi_1$, which vanishes in the SM. The $\sin \phi_d$ could be
deviated from the SM in 6% as seen in this figure. We present the SUSY contribution to the mass difference $\Delta M_{B_s}$ versus $s_{23}^{L(R)}$. It is remarked that the SUSY contribution could be also 6% in the $\Delta M_{B_s}$.

We show the SUSY contribution to the CP violating phase $\phi_s$ versus $s_{23}^{L(R)}$ in Figure 6, where we define $\Delta \sin 2\beta_s \equiv \sin \phi_s - \sin 2\beta_s$, which vanishes in the SM. It is found that the deviation of $\sin \phi_s$ from $\sin 2\beta_s$ is at most 8%. As seen in Figure 7, the SUSY contribution for $\Delta M_{B_s}$ is very small, $\mathcal{O}(0.4)\%$.

Let us discuss the $b \to s$ transitions. Under the constraints of the experimental data $\epsilon_K$, $\phi_d$, and $\phi_s$, $\Delta M_{B_s}$ and $\Delta M_{B_s}$, we can predict the magnitude of the Wilson coefficients $C_8^i$ and $\tilde{C}_{8G}^i$, which give us the deviation from the SM predicted values. We show the ratio $|(C_{8G}^8 + \tilde{C}_{8G}^8)/C_{8G}|$ versus $s_{23}^{L(R)}$ in Figure 8. Thus $C_{8G}^8$ is at most 1% because of the small left-right mixing $\theta = 0.35^\circ$ as seen in Table 1. We also show the predicted $|(\tilde{C}_{7\gamma}^g + C_{7\gamma}^g)/C_{7\gamma}|$ in Figure 9. This magnitude is much smaller than the case of $C_{8G}^8$, about 0.15%. Thus $C_{7\gamma}^g$ do not affect the branching ratio of the $b \to s\gamma$ decay in Eq. (39).
Let us discuss the numerical results of $S_{\phi K_S}$ and $S_{\eta' K^0}$. Since $\tilde{C}_{\phi g}$ is small, the deviation from the SM prediction is also small. We show the ratio of $S_{\phi K_S}$ to $S_{\eta' K^0}$ versus $s_{23}^{L(R)}$ in Figure 10, where the SM predicts just one. The deviation from the SM is tiny, at most 0.2%. Thus, there is no chance to detect the SUSY contribution in these decay modes.

We discuss the magnitude of the SUSY contribution to the indirect CP violation $a_{sl}^d$ and $a_{sl}^s$. We show the predicted magnitudes in Figure 11. For the $B^0$ decay, the predicted region is $a_{sl}^d \simeq -0.001 \sim 0$, on the other hand, for the $B_s$ decay, $a_{sl}^s$ is predicted to be $a_{sl}^s \simeq 0 \sim 5 \times 10^{-5}$, where the SM gives $a_{sl}^{d\text{SM}} = -(4.1 \pm 0.6) \times 10^{-4}$ and $a_{sl}^{s\text{SM}} = (1.9 \pm 0.3) \times 10^{-5}$ as shown in Eq. (22).

At the last step, we discuss the cEDM of the strange quark, which depends on $s_{23}^{L(R)}$. Under the left-right symmetric assumption $s_{23}^L = s_{23}^R$, we show the predicted cEDM of the strange quark versus $s_{23}^{L(R)}$ in Figure 12. The predicted cEDM could be larger than the

---

**Figure 10:** The ratio of $S_{\phi K_S}$ to $S_{\eta' K^0}$ versus $s_{23}^{L(R)}$.

**Figure 11:** Predicted region of the semileptonic CP asymmetries $a_{sl}^d$ and $a_{sl}^s$. The SM prediction is shown by the pink region.

**Figure 12:** The predicted cEDM of the strange quark versus $s_{23}^L = s_{23}^R$. The horizontal line denote the experimental upper bound.

**Figure 13:** The predicted cEDM of the strange quark versus $|\epsilon_{K}^{\text{SUSY}}/\epsilon_{K}|$, where $R = s_{23}^R$. The horizontal line denote the experimental upper bound.
experimental bound of Eq. (49), $5 \times 10^{-26}$ cm, in the region of $s_{23}^{L(R)} \geq 0.17$.

In Figs. 2-11, we have not imposed the constraint of the cEDM of the strange quark. In order to see the effect of the cEDM constraint, we show the predicted magnitude of $d_s$ versus $|\epsilon_{\text{SUSY}}^{K}/\epsilon_{K}|$ in Fig. 13. Although some region in this plane is excluded by the experimental bound of the cEDM, the allowed region of $|\epsilon_{\text{SUSY}}^{K}/\epsilon_{K}|$ is not changed. This situation is understandable by considering the different phase dependence of $\phi_{23}^{L}$, $\phi_{23}^{R}$ and $\phi$ for $d_s$ and $\epsilon_{\text{SUSY}}^{K}$, respectively. Thus, the constraint of the cEDM of the strange quark does not change our predictions although some region of free phase parameters is excluded.

In addition, it is noticed that our result of $d_s$ depends on the assumption $s_{23}^{L} = s_{23}^{R}$ considerably. If we take the suppressed right-handed mixing $s_{23}^{R}/s_{23}^{L} = 0.1$, the predicted cEDM is just one order reduced, on the other hand, $\epsilon_{K}$ still have 40% contribution of the squark flavor mixing even in this case.

Let us discuss the typical mixing angles of $s_{13}^{L(R)}$ and $s_{23}^{L(R)}$ in our results. They are 0.1(0.2) for sizable SUSY contributions as seen in Fig. 3. These mixing angles are much larger than the CKM mixing elements $V_{cb}$ and $V_{ub}$. Therefore, non-vanishing off diagonal squark mass matrix elements are required at the $\Lambda$ scale as discussed below Eq. (12). For our squark mass spectrum, the mixing angle 0.1(0.2) corresponds to the off diagonal elements $(m_{\tilde{Q}}^{2})_{13}$ and $(m_{\tilde{Q}}^{2})_{23}$ to be $\sim 8(16)$ TeV$^2$ in the left-handed squark mass matrix. Due to the top-Yukawa coupling, the off diagonal element increases approximately 1.4 times at the $\Lambda$ scale compared with the one at the $Q_0$ scale by the RGE’s evolution, that is $\sim 10(20)$ TeV$^2$ while the diagonal component is 100 TeV$^2$. Thus, the universal soft masses should be considered in the approximation of 10(20)%.

Let us briefly discuss the case (b) $Q_0 = 50$ TeV. The CP violations sensitive to the SUSY contribution is only $\epsilon_{K}$. In the Figure 14, we show the $|\epsilon_{\text{SUSY}}^{K}/\epsilon_{K}|$ versus $s_{13}^{L(R)} s_{23}^{L(R)}$. The SUSY contribution could be also large up to 35%. Thus, $\epsilon_{K}$ is still sensitive to the gluino-squark interaction even if the SUSY scale is 50 TeV. This trend continue to the scale $Q_0 = 100$ TeV. On the other hand, cEDM is reduced to much smaller than the experimental upper bound, $5 \times 10^{-26}$ cm, as seen in Figure 15. The situation is different from the one in the case of $Q_0 = 10$ TeV. This result is understandable because the SUSY mass scale increases.
by five times and the left-right mixing angle $\theta$ is reduced from $0.35^\circ$ to $0.05^\circ$ compared with the case of $Q_0 = 10$ TeV as seen in Table 1.

We summarize our results in Table 2, where the sensitivity of the SUSY contribution is presented for the case of $Q_0 = 10$ TeV and 50 TeV. Most sensitive quantity of the SUSY contribution is $\epsilon_K$. However, more works are required to extract the SUSY contribution in $\epsilon_K$. The unitarity fit is needed to find any mismatch in the SM and single out the SUSY contribution. In order to obtain the more precise SM calculation for $\epsilon_K$, the uncertainties of $B_K$, $V_{cb}$ and $m_t$ must be reduced.

The SUSY contributions for $S_{J/\psi K_S}$, $S_{J/\psi \phi}$ and $\Delta M_{B^0}$ are at most 6−8%. Since the theoretical uncertainties in the SM is more than 10%, which mainly comes from $\bar{\rho}$ and $\bar{\eta}$, it is difficult to detect the deviations of 6−8% from the SM at present. We hope the precise determination of $\bar{\rho}$ and $\bar{\eta}$ in order to find the SUSY contribution of this level.

As seen in Table 2, the qualitative features at the 10 TeV and 50 TeV scale are almost same except for the cEDM of the strange quark. There is a big chance to observe the neutron EDM in the near future if the SUSY scale is at 10 TeV.

Before closing the presentation of the numerical results, we add a comment on the other gaugino contribution. Since left-handed squarks form SU(2) doublets, the mixing angle $\theta^L_{ij}$ also appear in the up-type squark mixing matrix. Consequently, there are additional contributions to the CP violations of $K$, $B^0$ and $B_s$ mesons induced by chargino exchanging diagrams. We have obtained the ratio of the chargino contribution to the gluino one for $\text{Im}M_{12}(K)$, $\text{Im}M_{12}^d(B^0)$ and $\text{Im}M_{12}^s(B_s)$ as 6%, 10% and 10%, respectively. Thus, the chargino contributions are the sub-leading ones.

![Table 2](image)

| $|\epsilon_K|$ | $S_{J/\psi K_S}$ | $S_{J/\psi \phi}$ | $\Delta M_{B^0}$ | $\Delta M_{B_s}$ | $|S_{\phi K_S}/S_{\eta' K^0}| - 1$ | $\text{BR}(b \rightarrow s \gamma)$ | $|a^d_{sl}|$ | $|a^s_{sl}|$ | $|\hat{d}_s^C|$ |
|---|---|---|---|---|---|---|---|---|---|
| | 40% | 6% | 8% | 6% | 0.4% | 0.2% | 0.3% | $\leq 1 \times 10^{-3}$ | $\leq 5 \times 10^{-5}$ | $\leq 4 \times 10^{-25}$ cm |
| (a) $Q_0 = 10$ TeV | 35% | 0.1% | 0.1% | 0.1% | 0.005% | 0.001% | | $\leq 8 \times 10^{-4}$ | $\leq 4 \times 10^{-5}$ | $\leq 1 \times 10^{-27}$ cm |
| (b) $Q_0 = 50$ TeV |

Table 2: The SUSY contribution in the cases (a) $Q_0 = 10$ TeV and (b) $Q_0 = 50$ TeV. The percents denote ratios of the SUSY contributions.
5 Summary

We have probed the high scale SUSY, which is at 10 TeV-50 TeV scale, in the CP violations of $K$, $B^0$ and $B_s$ mesons. In order to estimate the contribution of the squark flavor mixing to the CP violations, we discuss the squark mass spectrum, which is consistent with the recent Higgs discovery. Taking the universal soft parameters at the SUSY breaking scale, we obtain the squark mass spectrum at 10 TeV and 50 TeV, where the SM emerges, by using the RGE’s of MSSM. And then, the $6 \times 6$ mixing matrix between down-squarks and down-quarks is examined by input of the experimental data of $K$, $B^0$ and $B_s$ mesons.

It is found that $\epsilon_K$ is most sensitive to the SUSY even if the SUSY scale is at 50 TeV. Therefore, the estimate of $\epsilon_K$ should be improved by reducing uncertainties of the theoretical and experimental input in the SM. The SUSY contributions for $S_{J/\psi K^*_S}$, $S_{J/\psi \phi}$ and $\Delta M_{B^0}$ are $6-8\%$ at the SUSY scale of 10 TeV. The precise determination of $\tilde{\rho}$ and $\tilde{\eta}$ are required in order to find the SUSY contribution of this level.

We also discussed the high scale SUSY contribution in the semileptonic CP asymmetry of $B^0$ meson. We expect the Belle II experiment searching for the semileptonic CP asymmetry $a_{sl}^d$ to find the deviation from the one of the SM in future. Although the magnitude of cEDM of the strange quark depends on $s_{23}^U/s_{23}^L$ ratio and the left-right mixing angle of squarks considerably, there is a big chance to find the high scale SUSY by the observation of the neutron EDM.

In this work, we have discussed only the down quark-squark sector. We will study the up quark-squark and lepton-slepton sectors elsewhere.

Acknowledgment

This work is supported by JSPS Grand-in-Aid for Scientific Research, 21340055 and 24654062, 25-5222, respectively.

Appendix

A Squark contribution in $\Delta F = 2$ process

The $\Delta F = 2$ effective Lagrangian from the gluino-sbottom-quark interaction is given as

$$\mathcal{L}_{\text{eff}}^{\Delta F=2} = -\frac{1}{2} [C_{VLL}O_{VLL} + C_{VRR}O_{VRR}] - \frac{1}{2} \sum_{i=1}^{2} \left[ C_{SLL}^{(i)}O_{SLL}^{(i)} + C_{SRR}^{(i)}O_{SRR}^{(i)} + C_{SLR}^{(i)}O_{SLR}^{(i)} \right]$$

(50)

then, the $P^0-\bar{P}^0$ mixing, $M_{12}$, is written as

$$M_{12} = -\frac{1}{2m_P} \langle P^0 | \mathcal{L}_{\text{eff}}^{\Delta F=2} | \bar{P}^0 \rangle .$$

(51)
The hadronic matrix elements are given in terms of the non-perturbative parameters $B_i$ as:

\[
\begin{align*}
\langle P^0 | \mathcal{O}_{VLL} | P^0 \rangle &= \frac{2}{3} m_P^2 f_P^2 B_1, \\
\langle P^0 | \mathcal{O}_{VRR} | P^0 \rangle &= \langle P^0 | \mathcal{O}_{VLL} | P^0 \rangle, \\
\langle P^0 | \mathcal{O}_{SLL}^{(1)} | P^0 \rangle &= -\frac{5}{12} m_P^2 f_P^2 R_P B_2, \\
\langle P^0 | \mathcal{O}_{SLL}^{(2)} | P^0 \rangle &= \langle P^0 | \mathcal{O}_{SRR}^{(1)} | P^0 \rangle, \\
\langle P^0 | \mathcal{O}_{SLL}^{(2)} | P^0 \rangle &= \frac{1}{12} m_P^2 f_P^2 R_P B_3, \\
\langle P^0 | \mathcal{O}_{SLR}^{(1)} | P^0 \rangle &= \frac{1}{2} m_P^2 f_P^2 R_P B_4, \\
\langle P^0 | \mathcal{O}_{SLR}^{(2)} | P^0 \rangle &= \frac{1}{6} m_P^2 f_P^2 R_P B_5, \\
\end{align*}
\]

where

\[
R_P = \left( \frac{m_P}{m_Q + m_q} \right)^2,
\]

with $(P, Q, q) = (B_d, b, d)$, $(B_s, b, s)$, $(K, s, d)$.

The Wilson coefficients for the gluino contribution in Eq. (50) are written as

\[
C_{VLL}(m_\tilde{g}) = \frac{\alpha^2}{m_\tilde{g}^2} \sum_{i, j=1}^{6} (\lambda^{(d)}_{GLL})^{i j} (\lambda^{(d)}_{GLL})^{i j} \left[ \frac{11}{18} g_{2[1]}(x_i^g, x_j^g) + \frac{2}{9} g_{1[1]}(x_i^g, x_j^g) \right],
\]

\[
C_{VRR}(m_\tilde{g}) = C_{VLL}(m_\tilde{g})(L \leftrightarrow R),
\]

\[
C_{SRR}^{(1)}(m_\tilde{g}) = C_{SRR}^{(1)}(m_\tilde{g})(L \leftrightarrow R),
\]

\[
C_{SLL}^{(1)}(m_\tilde{g}) = C_{SRR}^{(1)}(m_\tilde{g})(L \leftrightarrow R),
\]

\[
C_{SLL}^{(2)}(m_\tilde{g}) = C_{SRR}^{(2)}(m_\tilde{g})(L \leftrightarrow R),
\]

\[
C_{SLR}^{(1)}(m_\tilde{g}) = \frac{\alpha^2}{m_\tilde{g}^2} \sum_{i, j=1}^{6} \left\{ (\lambda^{(d)}_{GLR})^{i j} (\lambda^{(d)}_{GR})^{i j} \left[ -\frac{11}{9} g_{2[1]}(x_i^g, x_j^g) \\
- \frac{14}{3} g_{1[1]}(x_i^g, x_j^g) - \frac{2}{3} g_{2[1]}(x_i^g, x_j^g) \right] \right\},
\]

\[
C_{SLR}^{(2)}(m_\tilde{g}) = \frac{\alpha^2}{m_\tilde{g}^2} \sum_{i, j=1}^{6} \left\{ (\lambda^{(d)}_{GLR})^{i j} (\lambda^{(d)}_{GR})^{i j} \left[ -\frac{14}{3} g_{1[1]}(x_i^g, x_j^g) + \frac{10}{9} g_{2[1]}(x_i^g, x_j^g) \right] \right\},
\]

where

\[
\begin{align*}
(\lambda^{(d)}_{GLL})^{i j}_K &= (\Gamma^{(d)}_{GL})^i_L (\Gamma^{(d)}_{GL})^j_R, \\
(\lambda^{(d)}_{GLL})^{i j}_K &= (\Gamma^{(d)}_{GL})^i_L (\Gamma^{(d)}_{GL})^j_R, \\
(\lambda^{(d)}_{GLR})^{i j}_K &= (\Gamma^{(d)}_{GL})^i_L (\Gamma^{(d)}_{GR})^j_K, \\
(\lambda^{(d)}_{GLR})^{i j}_K &= (\Gamma^{(d)}_{GL})^i_L (\Gamma^{(d)}_{GR})^j_K.
\end{align*}
\]
Here we take \((i,j) = (1,3), (2,3), (1,2)\) which correspond to \(B^0, B_s,\) and \(K^0\) mesons, respectively. The loop functions are given as follows:

- If \(x_i^g \neq x_j^g (x_{i,j}^g = m_{\tilde{g}_{i,j}}^2/m_{\tilde{g}}^2)\),

\[
g_{1[1]}(x_i^g, x_j^g) = \frac{1}{x_i^g - x_j^g} \left( \frac{x_i^g \log x_i^g}{(x_i^g - 1)^2} - \frac{1}{x_i^g - 1} - \frac{x_j^g \log x_j^g}{(x_j^g - 1)^2} + \frac{1}{x_j^g - 1} \right),
\]

\[
g_{2[1]}(x_i^g, x_j^g) = \frac{1}{x_i^g - x_j^g} \left( \frac{(x_j^g)^2 \log x_j^g}{(x_j^g - 1)^2} - \frac{1}{x_j^g - 1} - \frac{(x_j^g)^2 \log x_j^g}{(x_j^g - 1)^2} + \frac{1}{x_j^g - 1} \right). \tag{56}
\]

- If \(x_i^g = x_j^g\),

\[
g_{1[1]}(x_i^g, x_j^g) = -\frac{(x_i^g + 1) \log x_i^g}{(x_i^g - 1)^3} + \frac{2}{(x_i^g - 1)^2},
\]

\[
g_{2[1]}(x_i^g, x_j^g) = -\frac{2x_i^g \log x_i^g}{(x_i^g - 1)^3} + \frac{x_i^g + 1}{(x_i^g - 1)^2}. \tag{57}
\]

Taking account of the case that the gluino mass is much smaller than the squark mass scale \(Q_0\), the effective Wilson coefficients are given at the leading order of QCD as follows:

\[
C_{VLL}(m_b(\Lambda = 2 \text{ GeV})) = \eta_{VLL}^{B(K)} C_{VLL}(Q_0), \quad C_{VRR}(m_b(\Lambda = 2 \text{ GeV})) = \eta_{VRR}^{B(K)} C_{VLL}(Q_0),
\]

\[
\begin{pmatrix}
C_{SLL}^{(1)}(m_b(\Lambda = 2 \text{ GeV})) \\
C_{SLL}^{(2)}(m_b(\Lambda = 2 \text{ GeV}))
\end{pmatrix}
= \begin{pmatrix}
C_{SLL}^{(1)}(Q_0) \\
C_{SLL}^{(2)}(Q_0)
\end{pmatrix}
X_{LL}^{-1} \eta_{LL}^{B(K)} X_{LL},
\]

\[
\begin{pmatrix}
C_{SRR}^{(1)}(m_b(\Lambda = 2 \text{ GeV})) \\
C_{SRR}^{(2)}(m_b(\Lambda = 2 \text{ GeV}))
\end{pmatrix}
= \begin{pmatrix}
C_{SRR}^{(1)}(Q_0) \\
C_{SRR}^{(2)}(Q_0)
\end{pmatrix}
X_{RR}^{-1} \eta_{RR}^{B(K)} X_{RR},
\]

\[
\begin{pmatrix}
C_{SLR}^{(1)}(m_b(\Lambda = 2 \text{ GeV})) \\
C_{SLR}^{(2)}(m_b(\Lambda = 2 \text{ GeV}))
\end{pmatrix}
= \begin{pmatrix}
C_{SLR}^{(1)}(Q_0) \\
C_{SLR}^{(2)}(Q_0)
\end{pmatrix}
X_{LR}^{-1} \eta_{LR}^{B(K)} X_{LR}, \tag{58}
\]

where

\[
\eta_{VLL}^{B} = \eta_{VRR}^{B} = \left( \frac{\alpha_s(Q_0)}{\alpha_s(\tilde{g})} \right)^{\frac{6}{16}} \left( \frac{\alpha_s(m_{\tilde{g}})}{\alpha_s(m_t)} \right)^{\frac{6}{27}} \left( \frac{\alpha_s(m_t)}{\alpha_s(m_b)} \right)^{\frac{6}{27}},
\]

\[
\eta_{LL}^{B} = \eta_{RR}^{B} = S_{LL} \begin{pmatrix}
\frac{d_{LL}^{1}}{\eta_{b\tilde{g}}} & 0 \\
0 & \frac{d_{LL}^{2}}{\eta_{b\tilde{g}}}
\end{pmatrix} S_{LL}^{-1}, \quad \eta_{LR}^{B} = S_{LR} \begin{pmatrix}
\frac{d_{LR}^{1}}{\eta_{b\tilde{g}}} & 0 \\
0 & \frac{d_{LR}^{2}}{\eta_{b\tilde{g}}}
\end{pmatrix} S_{LR}^{-1},
\]

\[
\begin{pmatrix}
\eta_{b\tilde{g}} \\
\eta_{VLL}^{K} = \eta_{VRR}^{K} = \left( \frac{\alpha_s(Q_0)}{\alpha_s(m_{\tilde{g}})} \right)^{\frac{6}{16}} \left( \frac{\alpha_s(m_{\tilde{g}})}{\alpha_s(m_t)} \right)^{\frac{6}{27}} \left( \frac{\alpha_s(m_t)}{\alpha_s(m_b)} \right)^{\frac{6}{27}} \left( \frac{\alpha_s(m_b)}{\alpha_s(\Lambda = 2 \text{ GeV})} \right)^{\frac{6}{27}}.
\]

\eta^K_{LL} = \eta^K_{RR} = S_{LL} \left( \begin{array}{cc}
d_1^{LL} & 0 \\
0 & d_2^{LL}
\end{array} \right) S_{LL}^{-1},
\eta^K_{LR} = S_{LR} \left( \begin{array}{cc}
d_1^{LR} & 0 \\
0 & d_2^{LR}
\end{array} \right) S_{LR}^{-1},
\eta_{\lambda}\tilde{g} = \left( \frac{\alpha_s(Q_0)}{\alpha_s(m_{\tilde{g}})} \right)^{\frac{1}{10}} \left( \frac{\alpha_s(m_t)}{\alpha_s(m_t)} \right)^{\frac{3}{10}} \left( \frac{\alpha_s(m_b)}{\alpha_s(\Lambda = 2 \text{ GeV})} \right)^{\frac{3}{10}},
\eta_{\tilde{g}g} = (1 - \sqrt{241}),
d_1^{LL} = \frac{2}{3} (1 - \sqrt{241}),
d_2^{LL} = \frac{2}{3} (1 + \sqrt{241}),
d_1^{LR} = -16,
d_2^{LR} = 2,
S_{LL} = \left( \begin{array}{cc}
\frac{16 + \sqrt{241}}{60} & \frac{16 - \sqrt{241}}{60} \\
\frac{16 - \sqrt{241}}{60} & \frac{16 + \sqrt{241}}{60}
\end{array} \right),
S_{LR} = \left( \begin{array}{cc}
-2 & 1 \\
3 & 0
\end{array} \right),
X_{LL} = X_{RR} = \left( \begin{array}{cc}
1 & 0 \\
4 & 8
\end{array} \right),
X_{LR} = \left( \begin{array}{cc}
0 & -2 \\
1 & 0
\end{array} \right).

(59)

For the parameters \( B^{(d)}_1(i = 2 - 5) \) of B mesons, we use values in [66] as follows:

\[
B_2^{(B_d)}(m_b) = 0.79(2)(4), \quad B_3^{(B_d)}(m_b) = 0.92(2)(4),
B_4^{(B_d)}(m_b) = 1.15(3)\left( ^{+5}_{-7} \right), \quad B_5^{(B_d)}(m_b) = 1.72(4)\left( ^{+20}_{-6} \right),
B_2^{(B_s)}(m_b) = 0.80(1)(4), \quad B_3^{(B_s)}(m_b) = 0.93(3)(8),
B_4^{(B_s)}(m_b) = 1.16(2)\left( ^{+5}_{-7} \right), \quad B_5^{(B_s)}(m_b) = 1.75(3)\left( ^{+21}_{-6} \right).
\]

(60)

On the other hand, we use the most updated values for \( \hat{B}^{(d)}_1 \) and \( \hat{B}^{(s)}_1 \) as [43]

\[
\hat{B}^{(B_d)}_1 = 1.33 \pm 0.06, \quad \hat{B}^{(B_s)}_1 / \hat{B}^{(B_d)}_1 = 1.05 \pm 0.07.
\]

(61)

For the paremeters \( B^K_i(i = 2 - 5) \), we use following values [67],

\[
B_2^{(K)}(2\text{GeV}) = 0.66 \pm 0.04, \quad B_3^{(K)}(2\text{GeV}) = 1.05 \pm 0.12,
B_4^{(K)}(2\text{GeV}) = 1.03 \pm 0.06, \quad B_5^{(K)}(2\text{GeV}) = 0.73 \pm 0.10,
\]

(62)

and we take recent value of Eq. [47] for deriving \( B_4^{(K)}(2\text{GeV}) \).

\section{The loop functions \( F_i \)}

The loop functions \( F_i(x^I_g) \) are given in terms of \( x^I_g = m^2_g / m^2_{di} \) \((I = 3, 6)\) as

\[
F_1(x^I_g) = \frac{x^I_g \log x^I_g}{2(x^I_g - 1)^4} + \frac{(x^I_g)^2 - 5x^I_g - 2}{12(x^I_g - 1)^3}, \quad F_2(x^I_g) = -\frac{(x^I_g)^2 \log x^I_g}{2(x^I_g - 1)^4} + \frac{2(x^I_g)^2 + 5x^I_g - 1}{12(x^I_g - 1)^3},
F_3(x^I_g) = \frac{\log x^I_g}{(x^I_g - 1)^3} + \frac{x^I_g - 3}{2(x^I_g - 1)^2}, \quad F_4(x^I_g) = -\frac{x^I_g \log x^I_g}{(x^I_g - 1)^3} + \frac{x^I_g + 1}{2(x^I_g - 1)^2} = \frac{1}{2} \left( g_{21} \right)(x^I_g, x^I_g),
\]

(63)
C cEDM

The cEDM of the strange quark from gluino contribution is given by [65]

\[ d_s^C(Q_0) = -2\sqrt{4\pi\alpha_s(m_{\tilde{g}})}\text{Im}[A_{s}^{22}(Q_0)], \] (64)

where

\[ A_{s}^{22}(Q_0) = -\frac{\alpha_s(m_{\tilde{g}})}{4\pi} \frac{1}{3} \left\{ \frac{1}{2m_{d_3}^2} \right\} \left( m_s(\lambda_{GLL}^{(d)})_{22}^2 + m_s(\lambda_{GRR}^{(d)})_{22}^2 \right) \left( 9F_1(x_\tilde{g}) + F_2(x_\tilde{g}) \right) \]

\[ + m_g(\lambda_{GLR}^{(d)})_{22}^2 \left( 9F_3(x_\tilde{g}) + F_4(x_\tilde{g}) \right) \]

\[ + \frac{1}{2m_{d_6}^2} \left\{ \left( m_s(\lambda_{GLL}^{(d)})_{6}^2 + m_s(\lambda_{GRR}^{(d)})_{6}^2 \right) \left( 9F_1(x_\tilde{g}) + F_2(x_\tilde{g}) \right) + m_g(\lambda_{GLR}^{(d)})_{6}^2 \left( 9F_3(x_\tilde{g}) + F_4(x_\tilde{g}) \right) \right\}. \] (65)

Including the QCD correction, we get

\[ d_s^C(2\text{GeV}) = d_s^C(Q_0) \left( \frac{\alpha_s(Q_0)}{\alpha_s(m_{\tilde{g}})} \right)^{\frac{14}{21}} \left( \frac{\alpha_s(m_t)}{\alpha_s(m_{\tilde{g}})} \right)^{\frac{14}{23}} \left( \frac{\alpha_s(m_b)}{\alpha_s(m_{\tilde{g}})} \right)^{\frac{14}{25}} \left( \frac{\alpha_s(2\text{GeV})}{\alpha_s(m_{\tilde{g}})} \right)^{\frac{14}{26}}. \] (66)

References

[1] G. Aad et al. [ATLAS Collaboration], arXiv:1208.0949 [hep-ex];
G. Aad et al. [ATLAS Collaboration], Eur. Phys. J. C 73 (2013) 2362 arXiv:1212.6149;
S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 713 (2012) 408 arXiv:1205.0272 [hep-ex];
S. Chatrchyan et al. [CMS Collaboration], arXiv:1305.2390 [hep-ex];
A. Hoecker (CERN), Talk at Lepton-Photon Symposium, San Francisco, June 24, 2013.

[2] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716 (2012) 1 arXiv:1207.7214;
S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B 716 (2012) 30 arXiv:1207.7235.

[3] W. Altmannshofer, R. Harnik and J. Zupan, JHEP 1311 (2013) 202 arXiv:1308.3653 [hep-ph].

[4] T. Moroi and M. Nagai, Phys. Lett. B 723 (2013) 107 arXiv:1303.0668 [hep-ph].

[5] A. Bharucha et al. [LHCb Collaboration], arXiv:1208.3355 [hep-ex].

[6] RAaij et al. [LHCb Collaboration], arXiv:1303.7125 [hep-ex].

[7] RAaij et al. [LHCb Collaboration], arXiv:1304.6173 [hep-ex].
[8] R. Aaij et al. [LHCb Collaboration], arXiv:1304.2600 [hep-ex].
[9] R. Aaij et al. [LHCb Collaboration], arXiv:1305.2168 [hep-ex].
[10] R. Aaij et al. [LHCb Collaboration], arXiv:1304.6325 [hep-ex].
[11] R. Aaij et al. [LHCb Collaboration], New J. Phys. 15, 053021 (2013) arXiv:1304.4741.
[12] M. Vesterinen [on behalf of the LHCb Collaboration], arXiv:1306.0092 [hep-ex].
[13] R. Aaij et al. [LHCb Collaboration], Phys. Rev. Lett. 110 (2013) 021801 arXiv:1211.2674 [hep-ex].
[14] R. Aaij et al. [LHCb Collaboration], Phys. Rev. Lett. 110 (2013) 031801 arXiv:1210.4492 [hep-ex].
[15] R. Aaij et al. [LHCb Collaboration], Phys. Rev. Lett. 108 (2012) 101803 arXiv:1112.3183 [hep-ex].
[16] R. Aaij et al. [LHCb Collaboration], Phys. Lett. B 707 (2012) 497 arXiv:1112.3056.
[17] M. Kobayashi and T. Maskawa, Prog. Theor. Phys. 49 (1973) 652.
[18] F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, Nucl. Phys. B 477 (1996) 321 [hep-ph/9604387].
[19] S. F. King, JHEP 1009 (2010) 114 arXiv:1006.5895 [hep-ph].
[20] M. Endo, S. Shirai and T. T. Yanagida, Prog. Theor. Phys. 125 (2011) 921 arXiv:1009.3366 [hep-ph].
[21] M. Endo and N. Yokozaki, JHEP 1103 (2011) 130 arXiv:1012.5501 [hep-ph].
[22] J. Kubo and A. Lenz, Phys. Rev. D 82 (2010) 075001 arXiv:1007.0680 [hep-ph].
[23] Y. Kaburaki, K. Konya, J. Kubo and A. Lenz, Phys. Rev. D 84 (2011) 016007 arXiv:1012.2435 [hep-ph].
[24] J. Virto, JHEP 0911 (2009) 055 arXiv:0907.5376 [hep-ph].
[25] J. Virto, JHEP 1201 (2012) 120 arXiv:1111.0940 [hep-ph].
[26] P. Ko and J. -h. Park, Phys. Rev. D 80 (2009) 035019 arXiv:0809.0703 [hep-ph].
[27] P. Ko and J. -h. Park, Phys. Rev. D 82 (2010) 117701 arXiv:1006.5821 [hep-ph].
[28] R. -M. Wang, Y. -G. Xu, Q. Chang and Y. -D. Yang, Phys. Rev. D 83 (2011) 095010 arXiv:1102.2031 [hep-ph].
[29] J. K. Parry, Phys. Lett. B 694 (2011) 363 arXiv:1006.5331 [hep-ph].
[30] A. Hayakawa, Y. Shimizu, M. Tanimoto and K. Yamamoto, Phys. Lett. B 710 (2012) 446 [arXiv:1202.0486 [hep-ph]].

[31] Y. Shimizu, M. Tanimoto and K. Yamamoto, Prog. Theor. Phys. 128 (2012) 273 [arXiv:1205.1705 [hep-ph]].

[32] Y. Shimizu, M. Tanimoto and K. Yamamoto, Phys. Rev. D 87 (2013) 056004 [arXiv:1212.6486 [hep-ph]].

[33] Y. Shimizu, M. Tanimoto and K. Yamamoto, Phys. Lett. B 726 (2013) 747 [arXiv:1307.0374].

[34] A. Hayakawa, Y. Shimizu, M. Tanimoto and K. Yamamoto, PTEP 2014 (2014) 023B04, [arXiv:1311.5974 [hep-ph]].

[35] A. J. Buras and D. Guadagnoli, Phys. Rev. D 78 (2008) 033005 [arXiv:0805.3887].

[36] F. Mescia and J. Virto, Phys. Rev. D 86 (2012) 095004 [arXiv:1208.0534 [hep-ph]].

[37] J. Hisano and Y. Shimizu, Phys. Lett. B 581 (2004) 224 [hep-ph/0308255].

[38] J. Hisano and Y. Shimizu, Phys. Rev. D 70 (2004) 093001 [hep-ph/0406091].

[39] J. Hisano, M. Nagai and P. Paradisi, Phys. Rev. D 80 (2009) 095014 [arXiv:0812.4283].

[40] K. Fuyuto, J. Hisano and N. Nagata, Phys. Rev. D 87 (2013) 054018 [arXiv:1211.5228].

[41] K. Fuyuto, J. Hisano, N. Nagata and K. Tsumura, JHEP 1312 (2013) 010 [arXiv:1308.6493 [hep-ph]].

[42] A. Delgado, M. Garcia and M. Quiros, [arXiv:1312.3235 [hep-ph]].

[43] G. F. Giudice and R. Rattazzi, Nucl. Phys. B 757 (2006) 19 [hep-ph/0606105].

[44] UTfit Collaboration (http://www.utfit.org).

[45] J. Beringer et al. [Particle Data Group Collaboration], Phys. Rev. D 86 (2012) 010001.

[46] For example, see S. P. Martin, “A Supersymmetry primer,” in *Kane, G.L. (ed.): Perspectives on supersymmetry II* 1-153 [hep-ph/9709356].

[47] G. Aad et al. [ATLAS Collaboration], arXiv:1402.7029 [hep-ex].

[48] K. Harigaya, K. Kaneta and S. Matsumoto, arXiv:1403.0715 [hep-ph].

[49] K. Harigaya, M. Ibe, K. Ichikawa, K. Kaneta and S. Matsumoto, arXiv:1403.5880 [hep-ph].

[50] A. Lenz and U. Nierste, arXiv:1102.4274 [hep-ph]; G. Borissov, R. Fleischer and M. -Helen. Schune, arXiv:1303.5575 [hep-ph].
[51] G. Buchalla, A. J. Buras and M. E. Lautenbacher, Rev. Mod. Phys. 68 (1996) 1125 [hep-ph/9512380].
[52] J. Hisano, T. Kuwahara and N. Nagata, Phys. Lett. B 723 (2013) 324 [arXiv:1304.0343 [hep-ph]].
[53] T.Aushev, W. Bartel, A. Bondar, J. Brodzicka, T. E. Browder, P. Chang, Y. Chao and K. F. Chen et al., arXiv:1002.5012 [hep-ex].
[54] J. Charles et al. [CKMfitter Group], Eur. Phys. J. C 41 (2005) 1 [hep-ph/0406184].
[55] Y. Amhis et al. [Heavy Flavor Averaging Group], arXiv:1207.1158 [hep-ex].
[56] S. Khalil and E. Kou, Phys. Rev. Lett. 91 (2003) 241602 [hep-ph/0303214].
[57] M. Endo, S. Mishima and M. Yamaguchi, Phys. Lett. B 609 (2005) 95 [hep-ph/0409245].
[58] A. L. Kagan, lecture at the 30th SLAC Summer Institute on Particle Physics: Secrets of the B Meson, 5-16 August 2002, Stanford, USA. http://wwwconf.slac.stanford.edu/ssi/2002/ .
[59] R. Harnik, D. T. Larson, H. Murayama and A. Pierce, Phys. Rev. D 69 (2004) 094024 [hep-ph/0212180].
[60] S. Mishima and A. I. Sanda, Prog. Theor. Phys. 110 (2003) 549 [hep-ph/0305073].
[61] A. J. Buras, [hep-ph/9806471].
[62] M. Misiak, H. M. Asatrian, K. Bieri, M. Czakon, A. Czarnecki, T. Ewerth, A. Ferroglia and P. Gambino et al., Phys. Rev. Lett. 98 (2007) 022002 [hep-ph/0609232].
[63] T. Bae, Y. -C. Jang, H. Jeong, J. Kim, J. Kim, K. Kim, S. Kim and W. Lee et al., PoS LATTICE 2013, 476 (2013) [arXiv:1310.7319 [hep-lat]].
[64] S. Aoki, Y. Aoki, C. Bernard, T. Blum, G. Colangelo, M. Della Morte, S. Durr and A. X. E. Khadra et al., arXiv:1310.8555 [hep-lat].
[65] T. Goto, http://research.kek.jp/people/tgoto/ .
[66] D. Becirevic, V. Gimenez, G. Martinelli, M. Papinutto and J. Reyes, JHEP 0204 (2002) 025 [hep-lat/0110091].
[67] C. R. Allton, L. Conti, A. Donini, V. Gimenez, L. Giusti, G. Martinelli, M. Talevi and A. Vladikas, Phys. Lett. B 453 (1999) 30 [hep-lat/9806016].