Direct CP violation in semi-inclusive flavor-changing neutral current decays in the MSSM without $R$–parity

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Abstract

Semi-inclusive decays, $q_h \rightarrow q_ℓ X_{jj}$, are studied in the framework of the minimal supersymmetric standard model without $R$–parity, where $q_h$ ($q_ℓ$) are the second or third (first or second) generation quarks with the same charge and $X_{jj}$ is a vector meson formed by $q_jq_j$. The study is focused on the contributions of sfermions with $m_\tilde{f} < m_{\text{top}}$. In this mass region, CP asymmetries in top decays can be induced by taking into account the decay-widths of the exchanged-bosons, while in light-quark decays it can be generated due to the long-distance effects. The contributions of sfermions also alter the branching-ratios destructively or constructively depend on the phases of complex couplings of the $R$–parity violation interactions.

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The discovery of top quark \cite{1} have accomplished the particle content of quark sector predicted in the standard model (SM). The heavy top and also other light-quarks have been attracted by theorists and experimentalists to test the SM as well as open a window for the physics beyond the SM. Some valuable information are expected from some classes of its decays that should be observed in, namely, top and meson factories. Theoretically, the $BDK$ meson decays have been studied briskly from a few decades ago, while most of top decays are still under studied.

Comparing both of them, the study of $BDK$ decays are mostly confronted with theoretical difficulties like the non-perturbative effects. On the other hand, in top decays the difficulties are almost coming from the experimental side that is still far to carry out some precise measurements as will be achieved for $BDK$ decays in the meson factories, although top decays are clean of theoretical uncertainties because of its large mass scale. It is also well known that most of models beyond the SM contribute significantly to the rare $BDK$ decays, and its presence should be examined in the present or near future meson factories. On the contrary, although the rare top decays are also very sensitive to the new physics, e.g. \cite{2,3,4}, the rates are still at unreachable level even in the future top factories as the upgraded Tevatron or LHC. These facts encourage us to consider some modes with the order between the lowest charged-current and the rare decays. Then we consider the class of semi-inclusive decays $q_h \rightarrow q_{\ell} X_{jj}$. Here $q_h (q_{\ell})$ are the second or third (first or second) generation quarks and have the same charge ($Q_h = Q_{\ell}$), while $X_{jj}$ is any vector meson formed by $q_j \bar{q}_j$. Since diagrammatically both top and light-quark processes are same and the interactions work on them may be related each other, we are going to consider both top and light-quark decays simultaneously. Definitely, we will discuss the flavor-changing semi-inclusive decays : $t \rightarrow u(c) X_{jj}$, $c \rightarrow u X_{jj}$, $b \rightarrow d(s) X_{jj}$ and $s \rightarrow d X_{jj}$. The vector mesons are, for example, $X_{u_j \bar{u}_j} = \rho,\omega,J/\psi,\cdots$ and $X_{d_j \bar{d}_j} = \phi,\Upsilon,\cdots$.

Moreover, in the present paper, the interest is focused on the direct CP asymmetry defined as

$$ A_{CP} \equiv \frac{\Gamma - \bar{\Gamma}}{\Gamma + \bar{\Gamma}} \equiv \frac{\Delta}{\Sigma} \quad , $$

with $\bar{\Gamma}$ is the complex conjugate of decay-width $\Gamma$, while in general

$$ \Delta = - \sum_{x \neq y} \text{Im} (\alpha_x^* \alpha_y) \text{Im} (\mathcal{M}_x^* \mathcal{M}_y) \quad , $$

$$ \Sigma = \sum_x |\alpha_x|^2 |\mathcal{M}_x|^2 + \sum_{x \neq y} \text{Re} (\alpha_x^* \alpha_y) \text{Re} (\mathcal{M}_x^* \mathcal{M}_y) \quad , $$

if one describes the amplitude as $\mathcal{M} \equiv \sum_x \alpha_x \mathcal{M}_x$. Hence, the imaginary parts of $(\alpha_x^* \alpha_y)$ and $(\mathcal{M}_x^* \mathcal{M}_y)$ are required to be non-zero coincidently in order to have non-zero CP asymmetry.

Indeed, in the framework of SM, the CP violation in these decays have been studied in some papers. The decays $b \rightarrow d(s) X_{u_j \bar{u}_j}$ have been discussed in \cite{5} for the typical one $b \rightarrow d J/\psi$. It has been concluded that the CP asymmetry is tiny, i.e. $\sim O(10^{-3})$, generated due to strong or electromagnetic scattering in the final state. However, the size could be at a few percents level if one takes into account...
the long-distance effects of the intermediate states with same quark contents as the final state [10]. On the other hand, recently the decays $t \rightarrow u(c) X_{d_jd_j}$ have also been examined in [9]. Different with the bottom one, in the case of top decays, CP violation is induced only by the scattering in the final state. It gives the size to be less than $O(10^{-2})$. Therefore in the SM, CP asymmetries in the present class of decays are almost at unreachable level of experiments, but inversely it makes them to be good probes to detect new contributions beyond the SM.

Presently, one of the well-known candidates for models beyond the SM is the supersymmetric standard model (MSSM). The model is attractive because of solution of the naturalness problem and also a lot of interesting properties. Especially, one of them that is relevant with our interest is CP violation due to the broken $R$–parity ($R_p$). The $R_p$–conservation is imposed to prevent the terms which explicitly break the baryon ($B$) and lepton ($L$) numbers. In the SM, the gauge symmetry leads to the conservation of $B$ and $L$, while in the SUSY model it does not prevent the terms [9]. Without $R_p$, there will be additional terms in superpotential [9], that is

$$W_{R_p} = \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c + \lambda''_{ijk} U_i^c D_j^c D_k^c + \kappa_i L_i H_2.$$  (4)

Here, $L$ and $E^c$ ($Q$ and $U^c$, $D^c$) are the lepton doublet and anti-lepton singlet (quark doublet and anti-quark singlet) left-chiral superfields, while $H_{1,2}$ are the Higgs doublet chiral superfields. ($i, j, k$) are generation indices, while ($\lambda, \lambda', \lambda''$) are Yukawa coupling-strengths. In fact, up to now there is still no theoretical preference between conserved and violated $R_p$. Phenomenologically, some authors have paid attention to these terms, especially the stop contribution in above $\lambda'$ term, because of its possibility to explain the HERA anomaly [10]. If this is a true story, then one is required to confirm the $R_p$–violation in other modes. In this meaning, it is worthwhile to adopt the present model as a probe to test it.

It is obvious that the semi-inclusive decays, $q_h \rightarrow q_i X_{jj}$, can be induced by either $\lambda'$ or $\lambda''$ terms at tree-level. Then, the lagrangians that are relevant with the present processes are given by expanding the terms in Eq. (4),

$$\mathcal{L}_{\lambda'} = -\lambda_{ijk} \left[ \bar{\nu}_L^i \bar{d}_R^k \nu_L^j + \bar{d}_L^k \nu_L^j + \left( \bar{d}_R^k \right)^* \left( \bar{\nu}_L^i \right)^c \nu_L^j ight. \
\left. - \bar{e}_L^j \bar{d}_R^k u_L^i - \bar{u}_L^i \bar{d}_R^k e_L^i - \left( \bar{d}_R^k \right)^* \left( \bar{\nu}_L^i \right)^c u_L^j \right] + \text{h.c.},$$  (5)

$$\mathcal{L}_{\lambda''} = -\lambda''_{ijk} \left[ \bar{d}_R^k \left( \bar{u}_L^i \right)^c d_L^j + \bar{d}_R^k \left( \bar{d}_L^k \right)^c u_L^i + \bar{\nu}_L^i \left( \bar{d}_L^k \right)^c d_L^k \right] + \text{h.c.}.$$  (6)

Note that $\lambda''$ is anti-symmetric under the interchanges of $[j, k]$. These terms induce new contributions to the decays with same level as the standard $W$–boson mediated diagram. In point of view of the SM, there are two types of the decays that may occur in the present model,

1. The SM favored modes induced by $W$–boson, sleptons and down-type squarks exchange diagrams.
2. The SM forbidden modes induced by sneutrinos and up-type squarks exchange diagrams.
The second one above is known as the tree-level flavor-changing neutral current (FCNC) modes that are also allowed in some models with additional isosinglet charge (-1/3) quarks. Most important difference is, in the present model the unitarity of CKM matrix is not altered at all. In this meaning, the mode is very interesting, if experimentally the unitarity of CKM matrix is known to be conserved while, for example, the decay $b \to s \phi$ is observed at appropriate level. Again, there is no tree-level FCNC modes in up-type quark sector in the present model.

After performing Fierz transformation, the amplitude in the processes through multiple $B-$boson mediated diagrams, are governed by the following operator

$$\mathcal{M} = \sqrt{2} G_F f_{X_j} m_{X_j} m_W^2 \epsilon_{X_j}^{\mu*} \sum_B \left[ \left(C_{ijB}^* C_{hjB}\right) \mathcal{F}_2^{\ell X_j B} \right] [\bar{q}_\ell \gamma_\mu L q_h],$$

by taking the factor of SM-like contribution as normalization factor. Here $L = (1 - \gamma_5)/2$ and $\mathcal{F}_2^{\ell X_j B}$ is $B-$boson propagator that will be given later. The vector meson is factorized as

$$\langle 0 | \bar{q}_j \gamma^\mu q_j | X_{jj} \rangle = m_{X_j} f_{X_j} \epsilon_{X_j}^{\mu} ,$$

where $f_{X_j}$ is a constant with dimensions of mass and $\epsilon_{X_j}^{\mu}$ is the polarization vector.

The coupling constants $C_{ijB}^* C_{hjB}$ are given in Tab. 1, where $a \equiv \left(2 \sqrt{2} G_F m_W^2\right)^{-1}$ and $V$ denotes the CKM matrix respectively. In the table, non-zero conditions are derived from the anti-symmetric of interchanging the indices of $\lambda''$, while the allowed modes are determined from the kinematics.

First, let us consider the branching-ratio. In general, it is better to consider the charged-current decay normalized one, that is

$$\text{BR}(q_h \to q_\ell X_{jj}) = \frac{\Gamma(q_h \to q_\ell X_{jj})}{\Gamma(q_h \to q_\ell X^{\pm})} \times \text{BR}(q_h \to q_\ell X^{\pm}) ,$$

to eliminate some uncertainties in the overwhole factors. Here $q_\ell$ is any light-quark that has different charge with $q_h$, while $X^{\pm}$ is anything with charge $\pm 1$. Definitely, we normalize $t \to u(c) X_{jj}$ (other light-quark modes) with $t \to b W$ (its semi-leptonic decays) as usual. The decay-width in numerator can be written as

$$\Gamma(q_h \to q_\ell X_{jj}) = \frac{G_F^2 f_{X_j}^2 \hat{m}_h^4 m_h}{8 \pi} \sqrt{g_{X_j}^{\ell X_j}} \mathcal{F}_1^{\ell X_j} \sum_B \left| \left(C_{ijB}^* C_{hjB}\right) \mathcal{F}_2^{\ell X_j B} \right|^2 ,$$

in the heavy quark center-mass system under assumptions that $E_{X_{jj}} = 2 E_j$ and $m_{X_{jj}} = 2 m_j$, with $E$ denotes time-component of four-momentum. Here, a caret means normalization with $m_h$. Keeping the light-quark masses

$$g_{X_j}^{\ell X_j} \equiv 1 + \hat{m}_x^4 + \hat{m}_y^4 - 2 \left(\hat{m}_x^2 + \hat{m}_y^2 + \hat{m}_x^2 \hat{m}_y^2\right) ,$$

$$\mathcal{F}_1^{\ell X_j} \equiv \left(1 - \hat{m}_x^2\right)^2 + \left(1 + \hat{m}_x^2\right) \hat{m}_y^2 - 2 \hat{m}_y^4 ,$$

$\mathcal{F}_2^{\ell X_j B}$ is the $B-$boson propagator contribution,

$$\mathcal{F}_2^{\ell X_j B} = \left[ \left(1 + \frac{1}{4} \hat{m}_y^2 - \frac{1}{2} \sqrt{g_{X_j}^{\ell X_j} + \frac{3}{4} \hat{m}_y^2 - \hat{m}_x^2}\right) + i \hat{m}_z \hat{m}_x \right]^{-1} \text{ for } q_h = t ,$$

$$\approx -\hat{m}_z^2 \text{ for } q_h \neq t .$$
under the kinematical condition: $m_B < m_h$ for $q_h = t$. Beware of including the
decay-width in the propagator is essential for CP asymmetry. This point will be
discussed later. Further, the charged-current decays in denominator are given as
follows,

$$
\Gamma(q_h \to q\ell X^\pm) = \begin{cases}
\frac{G_F m_h}{8\sqrt{2}\pi}|V_{h\ell}|^2 \sqrt{g^{\ell W} F_1^{\ell W}} & \text{for } (q_h = t, X^\pm = W^\pm), \\
\frac{G_F m_h}{192\pi^3}|V_{h\ell}|^2 F_3^{\ell} & \text{for } (q_h \neq t, X^\pm = \ell\bar{\nu}),
\end{cases}
$$

(14)

where $F_3^{\ell}$ accounts the phase space function in semi-leptonic decays, i.e.

$$
F_3^x = 1 - 8\hat{m}_x^2 + 8\hat{m}_x^6 - \hat{m}_x^8 - 24\hat{m}_x^4 \ln \hat{m}_x,
$$

(15)

while the QCD corrections to the decays have been omitted. For $q_h = t$, the QCD
corrections are predicted to be tiny.
Now we are discussing the main interest in the paper. As shown in Eqs. (1) and (2), non-zero CP asymmetries arise from non-zero imaginary part of the interference terms between the amplitudes and its coupling-strengths as well. For top decays in the present model, these requirements are satisfied by taking into consideration the decay-width in the boson propagator and the complex couplings \((V, \lambda', \lambda'')\). This is the reason why one can not neglect the decay-width in Eq. (13) as pointed out before. Calculating \(\Delta\) defined in Eq. (2) gives

\[
\Delta = -\frac{G_F^2 f^2_{\lambda'\lambda''} m_W^4 m_h}{4 \pi} \sqrt{g^{\ell X_{jj}}} \mathcal{F}_{1}^{\ell X_{jj}} \]

\[
\times \sum_{B_x \neq B_y} \text{Im} \left[ \left( C^*_{\ell B_x} C_{h B_x} \right)^* \left( C^*_{\ell B_y} C_{h B_y} \right) \right] \text{Im} \left[ \left( \mathcal{F}_2^{\ell X_{jj} B_x} \right)^* \mathcal{F}_2^{\ell X_{jj} B_y} \right].
\]  

Using Eqs. (1) and (10), CP asymmetry in top decay is found to be

\[
\mathcal{A}_{CP} = -\sum_{B_x \neq B_y} \text{Im} \left[ \left( C^*_{\ell B_x} C_{h B_x} \right)^* \left( C^*_{\ell B_y} C_{h B_y} \right) \right] \text{Im} \left[ \left( \mathcal{F}_2^{\ell X_{jj} B_x} \right)^* \mathcal{F}_2^{\ell X_{jj} B_y} \right] \]

\[
\times \left\{ \sum_B \left| C^*_{\ell B} C_{h B} \right|^2 \left| \mathcal{F}_2^{\ell X_{jj} B} \right|^2 \right\}^{-1}
\]

\[
+ \sum_{B_x \neq B_y} \text{Re} \left[ \left( C^*_{\ell B_x} C_{h B_x} \right)^* \left( C^*_{\ell B_y} C_{h B_y} \right) \right] \text{Re} \left[ \left( \mathcal{F}_2^{\ell X_{jj} B_x} \right)^* \mathcal{F}_2^{\ell X_{jj} B_y} \right].
\]  

On the other hand, in light-quark decays CP asymmetries are largely affected by the hadrons in the initial, intermediate and final states (11). Especially, as mentioned first, it has been pointed out that CP asymmetries of the present class of B decays may be enhanced by the long-distance effects of the intermediate states with same quark contents as the final state, while other intermediate states with different quark contents are negligible (10). In this case, different amplitude is generated by the penguin operator that has different phase with the tree one. In our notation, for general hadronic decay \(X_{hm} \rightarrow X' \rightarrow X_{tn} X_{jj}\) with \(X'\) is an intermediate state that has the same quark content as \(X_{tn} X_{jj}\), the amplitude is expressed as

\[
\mathcal{M} = \sum_B \left\{ \left( C^*_{\ell B} C_{h B} \right) T_{X_{tn} X_{jj}} + \left( C^*_{\ell h' B} C_{h B} \right) P_{X_{tn} X_{jj}} \right. \]

\[
+ \frac{i}{2} \sum_{X'} \left[ \left( C^*_{\ell B} C_{h B} \right) T_{X'} + \left( C^*_{\ell h' B} C_{h B} \right) P_{X'} \right] T_{X'}, \]
\]  

under an assumption that the rescattering effects can be treated perturbatively. \(T\) (P) denotes the tree (penguin) operator, while \(T_{X'}\) denotes the scattering amplitude of \(X' \rightarrow X_{tn} X_{jj}\). Here, the penguin contribution is normalized by the heaviest inner-line particle \((q_{h'})\) contribution. Hence, \(\Delta\) reads

\[
\Delta = \text{Im} \left[ \left( C^*_{\ell B} C_{h B} \right)^* \left( C^*_{\ell h' B} C_{h B} \right) \right] \sum_{X'} \text{Im} \left[ \tilde{T}_{X_{tn} X_{jj}}^* \tilde{P}_{X_{tn} X_{jj}} + \tilde{T}_{X'}^* \tilde{P}_{X'} \right] \]

\[
+ \frac{i}{2} \left( \tilde{T}_{X_{tn} X_{jj}}^* \tilde{P}_{X'} - \tilde{T}_{X'}^* \tilde{P}_{X_{tn} X_{jj}} \right) T_{X'}. \]  

\[
(19)
\]
and CP asymmetry in light-quark mode becomes

$$\mathcal{A}_{CP} \approx \sum_{X'} \text{Im} \left[ \frac{(\tilde{T}_{X'mX'j} \tilde{T}_{X'}) T_{X'}}{|\tilde{T}_{X'mX'j}|^2 + |\tilde{T}_{X'mX'j}|} \right] \times \left( \frac{\bar{P}_{X'mX'j}}{T_{X'}} + \frac{P_{X'}}{T_{X'}} + \frac{i}{2} \left( \frac{\bar{P}_{X't} - \bar{P}_{X'mX'j}}{T_{X'}} \right) \right) \right]$$

for type 1,

$$\text{Im} \left[ \begin{bmatrix} V_{lh'}^* V_{hh'} \\ V_{lj}^* V_{hj} \end{bmatrix} \right]$$

for type 2, \hspace{1cm} (20)

if the total decay-width is approximately dominated by tree operator. Here \( \tilde{f}' \) denotes the lightest sfermion and \( \tilde{T} \) is a complex conjugate of \( T \). A tilde means

$$\tilde{T} \equiv \begin{cases} T^W \left( 1 + \sum_f \frac{C_{\ell f} C_{h f} F_{X_{fX_{fX}}}^2}{V_{hj}^* V_{hj}^W} \right) & \text{for type 1,} \\
T^P \left( 1 + \sum_{f \neq f'} \frac{C_{\ell f} C_{h f} F_{X_{f+f}X_{f+f}}^{2}}{C_{\ell f'} C_{h f'} F_{X_{f+f}X_{f+f}}^{2}} \right) & \text{for type 2,} \end{cases} \hspace{1cm} \text{(21)}$$

while

$$\tilde{P} \equiv P^W \left( 1 + \sum_f \frac{C_{\ell f} C_{h f} P_{f} F_{X_{fX_{fX}}}^2}{V_{hj}^* V_{hj}^W} \right),$$

for both types. \( T^B (P^B) \) denotes the \( B \)-boson mediated tree (penguin) operator. Since the matrix elements depend on the hadronic states, \( \langle \tilde{T}_{X'mX'j} \rangle \left( \langle \tilde{P}_{X'mX'j} \rangle \right) \) is in general different with \( \langle T_{X'} \rangle \left( \langle P_{X'} \rangle \right) \).

Remark that the results in Eqs. (11), (13) and (14), at least numerically, are not altered so much by the diagonalization of squarks \( \bar{q}_L, \bar{q}_R \), although in fact, squarks are essentially mixed each other due to large Yukawa coupling of their partner quarks in the MSSM. So for a rough order estimation and also reducing the model dependence on the diagonalization, it is better to use the weak eigenstate as it is. Imagine the process is through a squark mediated diagram, then we can appreciate this point in two extreme cases. First case is when the masses are almost decoupled, then the branching-ratio will be quadruple, while the CP asymmetry will be reduced by half. On the contrary, when the mass difference is extremely large, one can neglect the large ones because the contribution will be suppressed by inverse of its mass square.

Now we are ready to make numerical analysis for the branching-ratios and CP asymmetries. Many authors have extracted some direct and indirect bounds for the coupling strengths in Eqs. (1) and (2), i.e. \( \lambda' \) and \( \lambda'' \). The bounds can be seen in Tab. 1 of ref. (13). However, until now there is still no rigid constraints for \( \lambda'' \). Moreover, since one of the \( B- \) and \( L- \)parity is still possible solutions to maintain a
stable proton and allow for $R_p$-violation as well, we assume that only one of these symmetries has been violated. Next, we consider only the must-be lightest sfermion for each sector and neglect the other heavier sfermions because its contributions should be suppressed. Hence, the analysis is simplified and can be done in a general way, i.e. it is sufficient to consider $W$- and one $\tilde{f}$-mediated diagrams for type 1, or only single $\tilde{f}$-mediated diagram for type 2.

For the branching-ratios of charged-current decays, we put $\mathcal{BR}(b \rightarrow c \ell \bar{\nu}) = 0.103$ and $\mathcal{BR}(t \rightarrow b W) \sim 1$ by assuming the mode to be dominant in top quark decays. Also use the experimental results, $m_u = 6$(MeV), $m_c = 1.3$(GeV), $m_t = 180$(GeV), $m_d = 10$(MeV), $m_s = 200$(MeV), $m_b = 4.3$(GeV), $m_W = 80.33$(GeV), $\Gamma_W = 2.07$(GeV) and the Wolfenstein parameters of CKM matrix $(A, \lambda, \rho, \eta) = (0.86, 0.22, 0.3, 0.34)$ [11]. In figure captions, the couplings are redefined as $C_W \equiv V^*_{i\ell}V_{hj}$ and $C_f \equiv \lambda'_{ijk}^* \lambda'_{ij'k'}$ or $\lambda''_{ijk}^* \lambda''_{ij'k'}$ respectively. The size of $C_f$ in some figures is fixed to be $\sim 0.015$ that is reliable enough for most coupling-strengths listed in [13]. In all figures, for the must-be lightest sfermion mass we put $\frac{1}{2}m_t < m_{\tilde{f}} < m_t$ that satisfies the kinematical requirement mentioned below Eq. (13). This region is still above the lower bound from the LEP experiments. Further, we put $\Gamma_{\tilde{f}} \sim 2\Gamma_w$ for whole region of sfermion masses. For the narrow region of masses under consideration, this approximation is good although in general the decay-width must be dependent on the mass. The phase of complex coupling $C_f$ is defined as

$$C_f \equiv |C_f| e^{i\theta}.$$  

Since large branching-ratios of light-quark decays in the SM are favored, it is better to describe the ratio of SM and MSSM with $R_p$-violation cases. Then the unknown parameter $\hat{f}_{X_{ij}}$ will be eliminated as shown in Figs. 1 and 2 (the right one). However, in case of either type 1 modes, or type 2 modes with tiny $C_W$ ($C_f \gg C_W \sim 0$), one must plot the branching-ratios itself as depicted in the left figure in Fig. 4 with leaving $\hat{f}_{X_{ij}}$ as unknown. Note that in Fig. 2, it seems significant differences between neither $b \rightarrow d(s) \phi$ with $b \rightarrow d(s) \rho(\omega)$ nor $s \rightarrow d \rho(\omega)$ with $c \rightarrow u \phi$. On the other hand, for CP asymmetries in light-quark decays, a rough prediction can be performed simply by using Eqs. (13) and (21). In general CP asymmetry for type 1 will be changed by a factor of

$$\mathcal{A}_{CP} \approx \mathcal{A}_{CP}^{SM} \left(1 + \frac{a C_{\tilde{f}} m_W^2}{C_W m_{\tilde{f}}^2}\right)^{-1},$$  

since $\langle \tilde{P} \rangle \approx \langle P^W \rangle$ for $C_{\tilde{f}}$ around the present value $[2, 4]$. For example, let us consider CP asymmetry in the decay $b \rightarrow d J/\psi$. In the present model it will be changed to be $\mathcal{A}_{CP} \approx (-14 \sim 2) \times \mathcal{A}_{CP}^{SM}$ for $C_{\tilde{f}} = 0.015$ and various set of $(\theta, m_{\tilde{f}})$. On the other hand, in the framework of SM with including the long-distance effects, e.g. $B^- \rightarrow D^0 D^- \rightarrow \pi^- J/\psi$, the value has been predicted to be $\mathcal{A}_{CP}^{SM} \sim 1%$ [3]. The prediction for other light-quark modes can be accomplished by the same procedure respectively. Remark that the procedure here is not requiring any kinematical condition like before, i.e. $m_B < m_h$. 

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In top decays, the dependences on $m_\ell$ and $m_{Xjj}$ are drastically suppressed. It makes the discrepancies between different modes are almost coming from CKM matrix elements. So one can describe them generally as depicted in Figs. 3 and 4. From the figures, the sfermion contributions will be maximum near the resonances. An interesting behaviour appears in the figure (b) in Fig. 4, that is in the small coupling region light sfermions are favored to obtain large CP asymmetry and vice versa.

In conclusion, the class of decay modes $q_h \to q_\ell X_{jj}$ have been studied in the framework of the MSSM without $R_p$. The study has been done for top and light-quark decays simultaneously, and focused on the CP asymmetries for $\frac{1}{2}m_t < m_f < m_t$. It is shown that, CP asymmetries in top decays can be induced by taking into account the decay-widths of the exchanged-bosons, while in light-quark decays it can be generated due to the long-distance effects as usual. The sfermion contributions due to the new interactions in $R_p$-violation superpotential change the branching-ratios and CP asymmetries significantly. Both measurements are very sensitive to the coupling-strengths and sfermion masses as well, that makes the modes to be good probes to search for the $R_p$-violation in the MSSM. Finally, although the decays are suffered from the small Yukawa couplings compared with some supersymmetric productions, they have double kinematic reach that makes them to be better for achieving precise measurements. Therefore, combining the various production and decay modes will lead to a wide range of potential signals to search for the $R_p$-violation in the MSSM.

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Figure 1: Ratio of the branching-ratio of $b \to s J/\psi$ (left), and $b \to d J/\psi$ (right).

Figure 2: The branching-ratio of $b \to d(s) \phi$ or $b \to d(s) \rho(\omega)$ (left), and ratio of the branching-ratio of $s \to d \rho(\omega)$ or $c \to u \phi$ (right).
Figure 3: The branching-ratios of $t \to u\phi(\Upsilon)$ and $t \to c\phi(\Upsilon)$ for various (a) $(C_W, C_\tilde{f})$, (b) $C_\tilde{f}$ with $C_W = A\lambda^2$, and (c) $m_\tilde{f}$ with $C_W = A\lambda^2$.

Figure 4: The CP asymmetries in $t \to u\phi(\Upsilon)$ and $t \to c\phi(\Upsilon)$ for various (a) $(C_W, C_\tilde{f})$ with $\theta = \pi/6$, (b) $m_\tilde{f}$ with $\theta = \pi/6$, and (c) $\theta$ with $(|C_W|, |C_\tilde{f}|) = (A\lambda^2, 0.015)$. 