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

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The semi-inclusive decays $q_h \to q_l X_{jj}$ are studied in the framework of the minimal supersymmetric standard model without R parity, where q_h (q_l) are second or third (first or second) generation quarks with the same charge and X_{jj} is a vector meson formed by $q_j \overline{q_j}$. The study is focused on the contributions of sfermions with $m_{\tilde{f}} < m_{\rm top}$. In this mass region, CP asymmetries in top-quark 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 long-distance effects. The contributions of sfermions also alter the branching ratios destructively or constructively depending on the phase of the complex couplings of the R-parity violation interactions. [S0556-2821(98)02719-2]

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With the discovery of the top quark [1], the entire quark sector, as predicted by the standard model (SM), has been observed. The heavy top-quark and also other light quarks have been attractive to theorists and experimentalists to test the SM as well as open a window for physics beyond the SM. Some valuable information is expected from some classes of its decays that should be observed in, namely, top-quark and meson factories. Theoretically, the BDK meson decays were studied thoroughly a few decades ago, while most top-quark decays are still under study.

Comparing both of them, a study of BDK decays is mostly confronted with theoretical difficulties such as nonperturbative effects. On the other hand, in top-quark decays the difficulties stem from the experimental side that is still far from carrying out some precise measurements as will be achieved for BDK decays in meson factories, although top-quark decays are free of theoretical uncertainties because of its large mass scale. It is also well known that most models beyond the SM contribute significantly to rare BDK decays, and its presence should be examined in present or near-future meson factories. On the contrary, although rare top-quark decays are also very sensitive to new physics, e.g., [2–4], the rates are still at unreachable levels even in future top-quark factories such as the upgraded Fermilab Tevatron or CERN Large Hadron Collider (LHC). These facts encourage us to consider some modes with rates with order of magnitude between those of the lowest charged current and the rare decays. Then we consider the class of semi-inclusive decays $q_h \to q_l X_{jj}$. Here q_h (q_l) are second or third (first or second) generation quarks and have the same charge ($Q_h = Q_l$), while X_{jj} is any vector meson formed by $q_j \overline{q_j}$. Since diagrammatically both top- and light-quark processes are same and the interactions working on them may be related to each other, we are going to consider both top- and light-quark decays simultaneously. Specifically, we will discuss the flavor-changing semi-inclusive decays $t \to u(c) X_{jj}$, $c \to u X_{jj}$, $b \to d(s) X_{jj}$, and $s \to d X_{jj}$, where X_{jj} denotes the respective vector mesons, for example, $X_{u_j \overline{u_j}} = \rho, \omega, J/\psi, \ldots$ and $X_{d_j \overline{d_j}} = \phi, \Upsilon, \ldots$.

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

$$\mathcal{A}_{CP} \equiv \frac{\Gamma - \overline{\Gamma}}{\Gamma + \overline{\Gamma}} \equiv \frac{\Delta}{\Sigma}, \tag{1}$$

where $\overline{\Gamma}$ denotes the partial width for the CP conjugate to the partial decay width given by Γ. For the decays under consideration the right-hand side reads

$$\Delta = -\sum_{x \neq y} \mathrm{Im}(\alpha_x^* \alpha_y) \mathrm{Im}(\mathcal{M}_x^* \mathcal{M}_y), \tag{2}$$

$$\Sigma = \sum_x |\alpha_x|^2 |\mathcal{M}_x|^2 + \sum_{x \neq y} \mathrm{Re}(\alpha_x^* \alpha_y) \mathrm{Re}(\mathcal{M}_x^* \mathcal{M}_y), \tag{3}$$

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

Indeed, in the framework of the SM, the CP violation in these decays has been studied in some papers. The decays $b \to d(s) X_{u_j \overline{u_j}}$ have been discussed in [5] for the typical one $b \to d J/\psi$. 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 the few percent level if one takes into account the long-distance effects of the intermediate states with the same quark content as the final state [9]. On the other hand, re-

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cently the decays $t \to u(c) X_{d_j \bar{d}_j}$ have also been examined in [6]. Different from the bottom one, in the case of top-quark decays, CP violation is induced only by scattering in the final state. It gives a size of less than $O(10^{-2})$. Therefore in the SM, CP asymmetries in the present class of decays are almost at an unreachable level of experiment, but inversely it makes them 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 a solution of the naturalness problem and also a lot of interesting properties. Especially, one of them that is relevant to our interest is CP violation due to broken R parity ($R_{\rm p}$). $R_{\rm p}$ conservation is imposed to prevent 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 supersymmetric (SUSY) model it does not prevent the terms [7]. Without $R_{\rm p}$, there will be additional terms in the superpotential [8]; that is,

$$\mathcal{W}_{\not R_{\rm p}} = \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c + \lambda''_{ijk} U_i^c D_j^c D_k^c + \kappa_i L_i H_2 . \tag{4}$$

Here, L and E^c (Q and U^c, D^c) are the lepton doublet and antilepton singlet (quark doublet and antiquark 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 has still been no theoretical preference between conserved and violated $R_{\rm p}$.

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

$$\begin{aligned}\mathcal{L}_{\lambda'} = & -\lambda'_{ijk}[\tilde{\nu}_L^i \overline{d_R^k} d_L{}^j + \tilde{d}_L^k \overline{d_R^k} \nu_L{}^i + (\tilde{d}_R^j)^* (\overline{\nu_L^i})^c d_L{}^j \\ & - \tilde{e}_L^i \overline{d_R^k} u_L{}^j - \tilde{u}_L^j \overline{d_R^k} e_L{}^i - (\tilde{d}_R^k)^* (\overline{e_L^i})^c u_L{}^j] + \text{H.c.},\end{aligned} \tag{5}$$

$$\begin{aligned}\mathcal{L}_{\lambda''} = & -\lambda''_{ijk}[\tilde{d}_R^k (\overline{u_L^i})^c d_L{}^j + \tilde{d}_R^j (\overline{d_L^k})^c u_L{}^i + \tilde{u}_R^i (\overline{d_L^j})^c d_L{}^k] \\ & + \text{H.c.}\end{aligned} \tag{6}$$

Note that λ'' is antisymmetric under the interchanges of $[j,k]$. These terms induce new contributions to the decays with the same level as the standard W-boson-mediated diagram. From the point of view of the SM, there are two types of decays that may occur in the present model: (1) The SM-favored modes induced by W-boson, slepton, and down-type squark exchange diagrams; (2) The SM forbidden modes induced by sneutrinos and up-type squark 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. The most important difference is that in the present model the unitarity of the Cabibbo-Kobayashi-Maskawa (CKM) matrix is not altered at all. In this meaning, the mode is very interesting, if experimentally the unitarity of the CKM matrix is known to be conserved while, for example, the decay $b \to s\phi$ is observed at the appropriate level. Again, there are no tree-level FCNC modes in the up-type quark sector in the present model.

After performing a Fierz transformation, the amplitude in the processes through multiple $\mathcal{B}$-boson-mediated diagrams is governed by the following operator:

$$\begin{aligned}\mathcal{M} = & \sqrt{2} G_F f_{X_{jj}} m_{X_{jj}} m_W{}^2 \epsilon_{X_{jj}}{}^{\mu *} \\ & \times \sum_{\mathcal{B}} [(\mathcal{C}^*_{lj\mathcal{B}} \mathcal{C}_{hj\mathcal{B}}) \mathcal{F}_2{}^{lX_{jj}\mathcal{B}}][\overline{q_l} \gamma_\mu L q_h],\end{aligned} \tag{7}$$

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

$$\langle 0 | \overline{q_j} \gamma^\mu q_j | X_{jj} \rangle = m_{X_{jj}} f_{X_{jj}} \epsilon_{X_{jj}}{}^\mu , \tag{8}$$

where $f_{X_{jj}}$ is a constant with dimensions of mass and $\epsilon_{X_{jj}}{}^\mu$ is the polarization vector. The coupling constants $\mathcal{C}^*_{lj\mathcal{B}} \mathcal{C}_{hj\mathcal{B}}$ are given in Table I, where $a \equiv (2\sqrt{2} G_F m_W{}^2)^{-1}$ and V denotes the CKM matrix, respectively. In the table, nonzero conditions are derived from the antisymmetry of interchanging the indices of λ'', 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,

$$\mathcal{B}(q_h \to q_l X_{jj}) = \frac{\Gamma(q_h \to q_l X_{jj})}{\Gamma(q_h \to q_{l'} X^\pm)} \mathcal{B}(q_h \to q_{l'} X^\pm), \tag{9}$$

to eliminate some uncertainties in the overall factors. Here $q_{l'}$ is any light quark that has a different charge from q_h, while $X^\pm$ is anything with charge ± 1. Definitely, we normalize $t \to u(c) X_{jj}$ (other light-quark modes) with $t \to bW$ (its semileptonic decays) as usual. The decay width in the numerator can be written as

$$\begin{aligned}\Gamma(q_h \to q_l X_{jj}) = & \frac{\hat{G}_F^2 \hat{f}_{X_{jj}}^2 \hat{m}_W^4 m_h}{8\pi} \sqrt{g^{lX_{jj}}} \mathcal{F}_1{}^{lX_{jj}} \\ & \times \left| \sum_{\mathcal{B}} (\mathcal{C}^*_{lj\mathcal{B}} \mathcal{C}_{hj\mathcal{B}}) \mathcal{F}_2{}^{lX_{jj}\mathcal{B}} \right|^2 ,\end{aligned} \tag{10}$$

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

$$g^{xy} \equiv 1 + \hat{m}_x^4 + \hat{m}_y^4 - 2(\hat{m}_x^2 + \hat{m}_y^2 + \hat{m}_x^2 \hat{m}_y^2), \tag{11}$$

$$\mathcal{F}_1{}^{xy} \equiv (1 - \hat{m}_x^2)^2 + (1 + \hat{m}_x^2) \hat{m}_y^2 - 2\hat{m}_y^4 . \tag{12}$$

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

TABLE I. The coupling strength $\mathcal{C}^*_{lj\mathcal{B}}\mathcal{C}_{hj\mathcal{B}}$ for each mode.

Type	Decay mode	$\mathcal{B}$	$\mathcal{C}^*_{lj\mathcal{B}}\mathcal{C}_{hj\mathcal{B}}$	Nonzero condition	Allowed mode
1	$t\to uX_{d_j\bar{d}_j}$	W	$V_{uj}{}^*V_{tj}$	-	$j=1,2,3$
		$\tilde{d}_k$	$a(\lambda''_{1kj}{}^*\lambda''_{3kj})$	$k\neq j$	
		$\tilde{l}_k$	$a(\lambda'_{k1j}{}^*\lambda'_{k3j})$ -	-	
	$t\to cX_{d_j\bar{d}_j}$	W	$V_{cj}{}^*V_{tj}$	-	$j=1,2,3$
		$\tilde{d}_k$	$a(\lambda''_{2kj}{}^*\lambda''_{3kj})$	$k\neq j$	
		$\tilde{l}_k$	$a(\lambda'_{k2j}{}^*\lambda'_{k3j})$	-	
	$c\to uX_{d_j\bar{d}_j}$	W	$V_{uj}{}^*V_{tj}$	-	$j=1,2$
		$\tilde{d}_k$	$a(\lambda''_{1kj}{}^*\lambda''_{2kj})$	$k\neq j$	
		$\tilde{l}_k$	$a(\lambda'_{k1j}{}^*\lambda'_{k2j})$	-	
	$b\to dX_{u_j\bar{u}_j}$	W	$V_{jd}{}^*V_{jb}$	-	$j=1,2$
		$\tilde{d}_k$	$a(\lambda''_{j1k}{}^*\lambda''_{j3k})$	$k=2$	
		$\tilde{l}_k$	$a(\lambda'_{kj3}{}^*\lambda'_{kj1})$	-	
	$b\to sX_{u_j\bar{u}_j}$	W	$V_{js}{}^*V_{jb}$	-	$j=1,2$
		$\tilde{d}_k$	$a(\lambda''_{j2k}{}^*\lambda''_{j3k})$	$k=1$	
		$\tilde{l}_k$	$a(\lambda'_{kj3}{}^*\lambda'_{kj2})$	-	
	$s\to dX_{u_j\bar{u}_j}$	W	$V_{jd}{}^*V_{js}$	-	$j=1$
		$\tilde{d}_k$	$a(\lambda''_{j1k}{}^*\lambda''_{j2k})$	$k=3$	
		$\tilde{l}_k$	$a(\lambda'_{kj2}{}^*\lambda'_{kj1})$	-	
2	$b\to dX_{d_j\bar{d}_j}$	$\tilde{u}_k$	$a(\lambda''_{j1k}{}^*\lambda''_{j3k})$	$k=2$	$j=1,2$
		$\tilde{\nu}_k$	$a(\lambda'_{k1j}{}^*\lambda'_{k3j})$	-	
	$b\to sX_{d_j\bar{d}_j}$	$\tilde{u}_k$	$a(\lambda''_{j2k}{}^*\lambda''_{j3k})$	$k=1$	$j=1,2$
		$\tilde{\nu}_k$	$a(\lambda'_{k2j}{}^*\lambda'_{k3j})$	-	
	$s\to dX_{d_j\bar{d}_j}$	$\tilde{u}_k$	$a(\lambda''_{j1k}{}^*\lambda''_{j2k})$	$k=3$	$j=1$
		$\tilde{\nu}_k$	$a(\lambda'_{k1j}{}^*\lambda'_{k2j})$	-	

$$\mathcal{F}_2{}^{xyz}\begin{cases} =\left[\left(1+\frac{1}{4}\hat{m}_y^2-\frac{1}{2}\sqrt{g^{xy}+4\hat{m}_y^2}-\hat{m}_z^2\right)+i\hat{\Gamma}_z\hat{m}_z\right]^{-1} & \text{for } q_h=t, \\ \approx-\hat{m}_z^{-2} & \text{for } q_h\neq t, \end{cases} \tag{13}$$

under the kinematical condition $m_{\mathcal{B}}<m_h$ for $q_h=t$. Note that including the decay width in the propagator is essential for CP asymmetry. This point will be discussed later. Further, the charged-current decays in the denominator are given as follows:

$$\Gamma(q_h\to q_{l'}X^{\pm}) = \begin{cases} \dfrac{\hat{G}_F m_h}{8\sqrt{2}\pi}|V_{hl'}|^2\sqrt{g^{l'W}}\mathcal{F}_1{}^{l'W} & \text{for } (q_h=t, X^{\pm}=W^{\pm}), \\ \dfrac{\hat{G}_F^2 m_h}{192\pi^3}|V_{hl'}|^2\mathcal{F}_3{}^{l'} & \text{for } (q_h\neq t, X^{\pm}=l\bar{\nu}), \end{cases} \tag{14}$$

where $\mathcal{F}_3{}^{l'}$ accounts the phase space function in semileptonic decays, i.e.,

$$\mathcal{F}_3{}^{x}=1-8\hat{m}_x^2+8\hat{m}_x^6-\hat{m}_x^8-24\hat{m}_x^4\ln\hat{m}_x. \tag{15}$$

We remark that the QCD corrections to the light-quark decays have been omitted for simplicity. For a more precise estimation, the effects must be included; however, the size should be reduced because of the cancellation in the ratio defined in Eq. (1). For $q_h=t$, the QCD corrections are predicted to be tiny because the top quark should decay at once before hadronization.

Now we discuss the main interest of this paper. As shown in Eqs. (1) and (2), nonzero CP asymmetries arise from the nonzero imaginary part of the interference terms between the amplitudes and its coupling strengths as well. For top-quark 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,λ',λ''). This is the reason why one cannot neglect the decay width in Eq. (13) as pointed out before. Calculating Δ defined in Eq. (2) gives

$$\Delta = -\frac{\hat{G}_F^2 \hat{f}_{X_{jj}}^2 \hat{m}_W^4 m_h}{4\pi} \sqrt{g^{lX_{jj}}} \mathcal{F}_1{}^{lX_{jj}} \times \sum_{\mathcal{B}_x \neq \mathcal{B}_y} \mathrm{Im}[(\mathcal{C}_{lj\mathcal{B}_x}^* \mathcal{C}_{hj\mathcal{B}_x})^* (\mathcal{C}_{lj\mathcal{B}_y}^* \mathcal{C}_{hj\mathcal{B}_y})] \times \mathrm{Im}[(\mathcal{F}_2{}^{lX_{jj}\mathcal{B}_x})^* \mathcal{F}_2{}^{lX_{jj}\mathcal{B}_y}]. \tag{16}$$

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

$$\mathcal{A}_{CP} = -\sum_{\mathcal{B}_x \neq \mathcal{B}_y} \mathrm{Im}[(\mathcal{C}_{lj\mathcal{B}_x}^* \mathcal{C}_{hj\mathcal{B}_x})^* (\mathcal{C}_{lj\mathcal{B}_y}^* \mathcal{C}_{hj\mathcal{B}_y})] \times \mathrm{Im}[(\mathcal{F}_2{}^{lX_{jj}\mathcal{B}_x})^* \mathcal{F}_2{}^{lX_{jj}\mathcal{B}_y}] \left\{ \sum_{\mathcal{B}} |\mathcal{C}_{lj\mathcal{B}}^* \mathcal{C}_{hj\mathcal{B}}|^2 |\mathcal{F}_2{}^{lX_{jj}\mathcal{B}}|^2 + \sum_{\mathcal{B}_x \neq \mathcal{B}_y} \mathrm{Re}[(\mathcal{C}_{lj\mathcal{B}_x}^* \mathcal{C}_{hj\mathcal{B}_x})^* (\mathcal{C}_{lj\mathcal{B}_y}^* \mathcal{C}_{hj\mathcal{B}_y})] \times \mathrm{Re}[(\mathcal{F}_2{}^{lX_{jj}\mathcal{B}_x})^* \mathcal{F}_2{}^{lX_{jj}\mathcal{B}_y}] \right\}^{-1}. \tag{17}$$

On the other hand, in light-quark decays CP asymmetries are largely affected by hadrons in the initial, intermediate, and final states [9]. Especially, as mentioned first, it has been pointed out that CP asymmetries of the present class of B decays may be enhanced by long-distance effects of the intermediate states with same quark content as the final state, while other intermediate states with a different quark content are negligible [5]. In this case, a different amplitude is generated by the penguin operator that has a phase different from the tree one. In our notation, for general hadronic decay $X_{hm} \to X' \to X_{lm} X_{jj}$ with X' is an intermediate state that has the same quark content as $X_{lm} X_{jj}$, the amplitude is expressed as

$$\mathcal{M} = \sum_{\mathcal{B}} \left\{ (\mathcal{C}_{lj\mathcal{B}}^* \mathcal{C}_{hj\mathcal{B}}) T_{X_{lm}X_{jj}} + (\mathcal{C}_{lh'\mathcal{B}}^* \mathcal{C}_{hh'\mathcal{B}}) P_{X_{lm}X_{jj}} + \frac{i}{2} \sum_{X'} [(\mathcal{C}_{lj\mathcal{B}}^* \mathcal{C}_{hj\mathcal{B}}) T_{X'} + (\mathcal{C}_{lh'\mathcal{B}}^* \mathcal{C}_{hh'\mathcal{B}}) P_{X'}] \mathcal{T}_{X'} \right\}, \tag{18}$$

under the assumption that the rescattering effects can be treated perturbatively. T (P) denotes the tree (penguin) operator, while $\mathcal{T}_{X'}$ denotes the scattering amplitude of $X' \to X_{lm} X_{jj}$. Here, the penguin contribution is normalized by the heaviest virtual particle ($q_{h'}$) contribution. Hence, Δ reads

$$\Delta = \mathrm{Im}[(\mathcal{C}_{lj\mathcal{B}}^* \mathcal{C}_{hj\mathcal{B}})^* (\mathcal{C}_{lh'\mathcal{B}}^* \mathcal{C}_{hh'\mathcal{B}})] \sum_{X'} \mathrm{Im}\left[\tilde{T}_{X_{lm}X_{jj}}^* \tilde{P}_{X_{lm}X_{jj}} + \tilde{T}_{X'}^* \tilde{P}_{X'} + \frac{i}{2} (\tilde{T}_{X_{lm}X_{jj}}^* \tilde{P}_{X'} - \tilde{T}_{X'} \tilde{P}_{X_{lm}X_{jj}}^*) \right] \mathcal{T}_{X'}, \tag{19}$$

and CP asymmetry in light-quark mode becomes

$$\mathcal{A}_{CP} \approx \sum_{X'} \mathrm{Im} \left\{ \frac{(\tilde{T}_{X_{lm}X_{jj}}^* \tilde{T}_{X'}^*) \mathcal{T}_{X'}}{|\tilde{T}_{X_{lm}X_{jj}}|^2 + |\bar{\tilde{T}}_{X_{lm}X_{jj}}|^2} \times \left[\frac{\tilde{P}_{X_{lm}X_{jj}}}{\tilde{T}_{X'}^*} + \frac{\tilde{P}_{X'}}{\tilde{T}_{X_{lm}X_{jj}}^*} + \frac{i}{2} \left(\frac{\tilde{P}_{X'}}{\tilde{T}_{X'}^*} - \frac{\tilde{P}_{X_{lm}X_{jj}}}{\tilde{T}_{X_{lm}X_{jj}}^*} \right) \right] \right\} \times \begin{cases} \mathrm{Im}\left[\dfrac{V_{lh'}{}^* V_{hh'}}{V_{lj}{}^* V_{hj}} \right] & \text{for type 1,} \\ \mathrm{Im}\left[\dfrac{V_{lh'}{}^* V_{hh'}}{\mathcal{C}_{lj\tilde{f}'}^* \mathcal{C}_{hj\tilde{f}'}} \right] & \text{for type 2,} \end{cases} \tag{20}$$

if the total decay width is approximately dominated by the tree operator. Here $\tilde{f}'$ denotes the lightest sfermion and $\bar{\tilde{T}}$ is the CP conjugate of $\tilde{T}$. A tilde means

$$\tilde{T} \equiv \begin{cases} T^W \left(1 + \displaystyle\sum_{\tilde{f}} \frac{\mathcal{C}_{lj\tilde{f}}^* \mathcal{C}_{hj\tilde{f}}}{V_{lj}{}^* V_{hj}} \frac{\mathcal{F}_2{}^{lX_{jj}\tilde{f}}}{\mathcal{F}_2{}^{lX_{jj}W}} \right) & \text{for type 1,} \\ T^{\tilde{f}'} \left(1 + \displaystyle\sum_{\tilde{f} \neq \tilde{f}'} \frac{\mathcal{C}_{lj\tilde{f}}^* \mathcal{C}_{hj\tilde{f}}}{\mathcal{C}_{lj\tilde{f}'}^* \mathcal{C}_{hj\tilde{f}'}} \frac{\mathcal{F}_2{}^{lX_{jj}\tilde{f}}}{\mathcal{F}_2{}^{lX_{jj}\tilde{f}'}} \right) & \text{for type 2,} \end{cases} \tag{21}$$

while

$$\tilde{P} \equiv P^W \left(1 + \sum_{\tilde{f}} \frac{\mathcal{C}_{lj\tilde{f}}^* \mathcal{C}_{hj\tilde{f}}}{V_{lh'}{}^* V_{hh'}} \frac{P^{\tilde{f}}}{P^W} \right) \tag{22}$$

for both types. $T^{\mathcal{B}}$ ($P^{\mathcal{B}}$) denotes the $\mathcal{B}$-boson-mediated tree (penguin) operator. Since the matrix elements depend on the hadronic states, $\langle \tilde{T}_{X_{lm}X_{jj}} \rangle$ ($\langle \tilde{P}_{X_{lm}X_{jj}} \rangle$) is in general different from $\langle \tilde{T}_{X'} \rangle$ ($\langle \tilde{P}_{X'} \rangle$).

We remark that the results in Eqs. (10), (16), and (19), at least numerically, are not altered so much by the diagonalization of squarks ($\tilde{q}_L, \tilde{q}_R$), although, in fact, squarks are essentially mixed with each other due to a large Yukawa coupling of their partner quarks in the MSSM [11]. 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. The 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 by using Eqs. (10) and (16) or (19). On the contrary, when the mass difference is extremely large, one can neglect the large ones because the contribution will be suppressed by the inverse of its mass square.

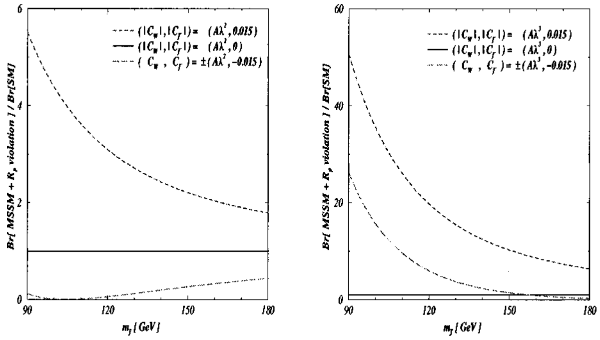

FIG. 1. Ratio of the branching ratio of $b \to sJ/\psi$ (left) and $b \to dJ/\psi$ (right).

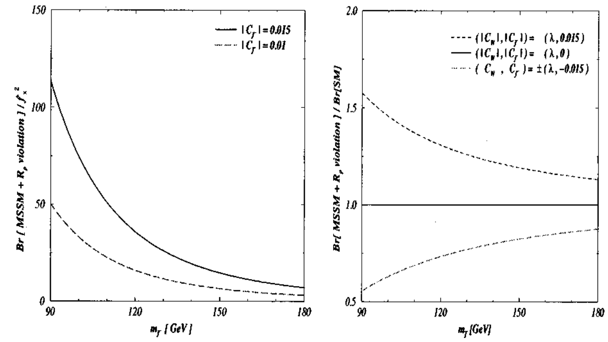

FIG. 2. The branching ratio of $b \to d(s)\phi$ or $b \to d(s)\rho(\omega)$ (left), and ratio of the branching ratio of $s \to d\rho(\omega)$ or $c \to u\phi$ (right).

Now we are ready to make a numerical analysis of the branching ratios and CP asymmetries. Many authors have extracted some direct and indirect bounds for the coupling strengths in Eqs. (5) and (6), i.e., λ' and λ''. The bounds can be seen in Table 1 of Ref. [12]. However, until now there has still been no rigid constraints for λ''. Moreover, since B or L parity is still a possible solution 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 take the other sfermions to be sufficiently heavy that their contributions may be neglected. Hence, the analysis is simplified; i.e., it is sufficient to take into account W- and one $\tilde{f}$-mediated diagrams for type 1 or only a single $\tilde{f}$-mediated diagram for type 2. This approximation is still good, at least, to the precision implied by not including QCD corrections.

For the branching ratios of charged-current decays, we put $\mathcal{B}(b \to cl\overline{\nu}) = 0.103$ and $\mathcal{B}(t \to bW) \sim 1$ by assuming the mode to be dominant in top-quark decays. Also we 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), and $\Gamma_W = 2.07$ (GeV) and the Wolfenstein parameters of the CKM matrix, $(A,\lambda,\rho,\eta) = (0.86,0.22,0.3,0.34)$ [10]. In the figure captions, the couplings are redefined as $C_W \equiv V_{lj}^* V_{hj}$ and $C_{\tilde{f}} \equiv \lambda'_{ijk}{}^* \lambda'_{i'j'k'}$ or $\lambda''_{ijk}{}^* \lambda''_{i'j'k'}$, respectively. The size of $C_{\tilde{f}}$ in some figures is $\leqslant 0.015$, which is the allowed limit for most coupling strengths listed in [12]. In all figures, for the must-be lightest sfermion mass we put $\frac{1}{2} m_t < m_{\tilde{f}} < m_t$, which 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 the 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 $C_{\tilde{f}}$ is defined as

$$C_{\tilde{f}} \equiv |C_{\tilde{f}}| e^{i\theta}. \tag{23}$$

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 $\hat{f}_{X_{jj}}$ 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 $(C_{\tilde{f}} \gg C_W \sim 0)$, one must plot the branching ratio itself as depicted in the left figure in Fig. 2, leaving $\hat{f}_{X_{jj}}$ as unknown. Note that in Fig. 2 it seems that there are no significant differences between either $b \to d(s)\phi$ with $b \to d(s)\rho(\omega)$ or $s \to d\rho(\omega)$ with $c \to u\phi$. 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

$$\mathcal{A}_{CP} \approx \mathcal{A}_{CP}^{\rm SM} \left(1 + \frac{aC_{\tilde{f}}}{C_W} \frac{m_W{}^2}{m_{\tilde{f}}{}^2} \right)^{-1}, \tag{24}$$

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 $b \to dJ/\psi$. In the present model it will be changed to $\mathcal{A}_{CP} \approx (-14 \sim 2) \times \mathcal{A}_{CP}^{\rm SM}$ for $C_{\tilde{f}} = 0.015$ and various sets of $(\theta, m_{\tilde{f}})$. On the other hand, in the framework of the SM including long-distance effects, e.g., $B^- \to D^0 D^- \to \pi^- J/\psi$, the value has been predicted to be $\mathcal{A}_{CP}^{\rm SM} \sim 1\%$ [5]. A prediction for other light-quark modes can be accomplished by the same procedure, respectively. We remark that the procedure here does not require any kinematical condition like before, i.e., $m_{\mathcal{B}} < m_h$.

In top quark decays, the dependences on m_l and $m_{X_{jj}}$ 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 behavior appears in Fig. 4(b); 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 $q_h \to q_l X_{jj}$ has 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 $\frac{1}{2} m_t < m_{\tilde{f}} < m_t$. It is shown that CP asymmetries in top quark 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 long-distance effects as

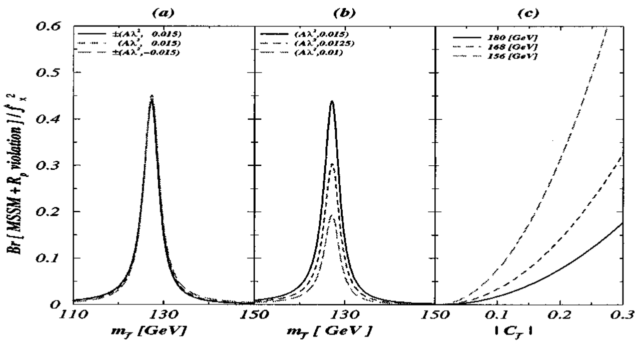

FIG. 3. The branching ratios of $t \to u\phi(\Upsilon)$ and $t \to c\phi(\Upsilon)$ for various (a) $(C_W, C_{\tilde{f}})$, (b) $C_{\tilde{f}}$ with $C_W = A\lambda^2$, and (c) $m_{\tilde{f}}$ with $C_W = A\lambda^2$.

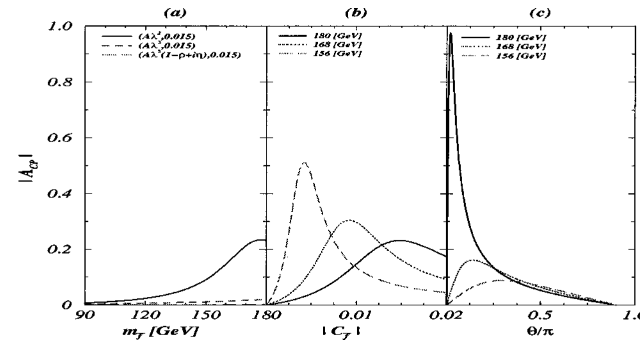

FIG. 4. The CP asymmetries in $t \to u\phi(\Upsilon)$ and $t \to c\phi(\Upsilon)$ for various (a) $(C_W, C_{\tilde{f}})$ with $\theta = \pi/6$, (b) $m_{\tilde{f}}$ with $\theta = \pi/6$, and (c) θ with $(|C_W|, |C_{\tilde{f}}|) = (A\lambda^2, 0.015)$.

usual. The sfermion contributions due to the new interactions in the 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, which makes the modes good probes to search for R_p violation in the MSSM. Finally, although the decays suffer from small Yukawa couplings compared to some supersymmetric productions, they have a double kinematic reach that makes them 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 R_p violation in the MSSM.

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