Impact of bosonic decays on the search for $\tilde{\tau}_2$ and $\tilde{\nu}_\tau$

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Abstract

We perform a detailed study of the decays of the heavier $\tau$ slepton ($\tilde{\tau}_2$) and $\tau$-sneutrino ($\tilde{\nu}_\tau$) in the Minimal Supersymmetric Standard Model (MSSM). We show that the decays into Higgs or gauge bosons, i.e. $\tilde{\tau}_2^- \to \tilde{\tau}_1^- + (h^0, H^0, A^0$ or $Z^0)$, $\tilde{\tau}_2^- \to \tilde{\nu}_\tau + (H^- or W^-)$, and $\tilde{\nu}_\tau \to \tilde{\tau}_1^- + (H^+ or W^+)$, can be very important due to the sizable $\tau$ Yukawa coupling and large mixing parameters of $\tilde{\tau}$. Compared to the decays into fermions, such as $\tilde{\tau}_2^- \to \tau^- + \tilde{\chi}_i^0$ and $\tilde{\tau}_2^- \to \nu_\tau + \tilde{\chi}_j^0$, these bosonic decay modes can have significantly different decay distributions. This could have an important influence on the search for $\tilde{\tau}_2$ and $\tilde{\nu}_\tau$ and the determination of the MSSM parameters at future colliders.
In the Minimal Supersymmetric Standard Model (MSSM) [1] supersymmetric (SUSY) partners of all Standard Model (SM) particles with masses less than O(1 TeV) are introduced. This solves the problems of hierarchy, fine-tuning and naturalness of the SM. Hence discovery of all SUSY partners and study of their properties are essential for testing the MSSM. Future colliders, such as the Large Hadron Collider (LHC), the upgraded Tevatron, $e^+e^-$ linear colliders, and $\mu^+\mu^-$ colliders will extend the discovery potential for SUSY particles to the TeV mass range and allow for a precise determination of the SUSY parameters.

In this article we focus on the sleptons of the third generation, i.e. staus ($\tilde{\tau}_{1,2}; m_{\tilde{\tau}_1} < m_{\tilde{\tau}_2}$) and tau-sneutrino ($\tilde{\nu}_\tau$). These particles may have properties different from the sleptons of the other two generations due to the sizable $\tau$ Yukawa coupling. Production and decays of $\tilde{\tau}_i$ and $\tilde{\nu}_\tau$ were studied in [2, 3, 4]. Like other sleptons, they can decay into fermions, i.e. a lepton plus a neutralino ($\tilde{\chi}_k^0$) or chargino ($\tilde{\chi}_j^\pm$):

\[
\begin{align*}
\tilde{\tau}_i^- & \rightarrow \tau^- \tilde{\chi}_k^0, & \tilde{\nu}_\tau & \rightarrow \nu_\tau \tilde{\chi}_k^0, \\
\tilde{\tau}_i^- & \rightarrow \nu_\tau \tilde{\chi}_j^+, & \tilde{\nu}_\tau & \rightarrow \tau^- \tilde{\chi}_j^+, \\
\end{align*}
\]

(1)

with $i, j = 1, 2$ and $k = 1, \ldots, 4$. In addition, the heavier stau $\tilde{\tau}_2$ and the tau-sneutrino can also decay into bosons [3, 4], i.e. a lighter slepton plus a gauge boson

\[
\begin{align*}
\tilde{\tau}_2^- & \rightarrow \tilde{\tau}_1^- Z^0, & \tilde{\nu}_\tau & \rightarrow \tilde{\tau}_1^- Z^0, \\
\tilde{\tau}_2^- & \rightarrow \tilde{\nu}_\tau W^-, & \tilde{\nu}_\tau & \rightarrow \tilde{\tau}_1^- W^+, \\
\end{align*}
\]

(2)

or a Higgs boson

\[
\begin{align*}
\tilde{\tau}_2^- & \rightarrow \tilde{\tau}_1^- (h^0, H^0, A^0), & \tilde{\nu}_\tau & \rightarrow \tilde{\tau}_1^- H^+, \\
\end{align*}
\]

(3)

The decays in Eqs. (2) and (3) are possible in case the mass splitting between the sleptons is sufficiently large.

In the present article we perform a more general analysis than [3, 4]. We point out that the $\tilde{\tau}_2^-$ and $\tilde{\nu}_\tau$ decays into gauge or Higgs bosons of Eqs. (2) and (3) can be very important in a large region of the MSSM parameter space due to the sizable $\tau$ Yukawa coupling and large $\tau$ mixing parameters. This importance of the bosonic modes relative to the fermionic modes of Eq. (1) could have a significant influence on searches for $\tilde{\tau}_2^-$ and $\tilde{\nu}_\tau$ at future colliders. An analogous study for the heavier stop ($\tilde{t}_2$) and sbottom ($\tilde{b}_2$) was performed in [5, 6].

First we summarize the MSSM parameters in our analysis. In the MSSM the stau sector is specified by the mass matrix in the basis ($\tilde{\tau}_L, \tilde{\tau}_R$) [7, 8]

\[
\mathcal{M}_\tilde{\tau}^2 = \begin{pmatrix} m_{\tilde{\tau}_L}^2 & a_\tau m_\tau \\ a_\tau m_\tau & m_{\tilde{\tau}_R}^2 \end{pmatrix}
\]

(4)

with

\[
m_{\tilde{\tau}_L}^2 = M_L^2 + m_Z^2 \cos 2\beta (\sin^2 \theta_W - \frac{1}{2}) + m_\tau^2.
\]

(5)
\[ m_{\tilde{\tau}_R}^2 = M^2_E - m_Z^2 \cos 2\beta \sin^2 \theta_W + m_\tau^2, \]
\[ a_\tau m_\tau = (A_\tau - \mu \tan \beta) m_\tau. \]

\( M^L, E \) and \( A_\tau \) are soft SUSY-breaking parameters, \( \mu \) is the higgsino mass parameter, and \( \tan \beta = v_2/v_1 \) with \( v_1, v_2 \) being the vacuum expectation value of the Higgs field \( H_1^0 \) \( (H_2^0) \).

Diagonalizing the matrix (4) one gets the mass eigenstates \( \tilde{\tau}_1 = \bar{\tau}_L \cos \theta_\tau + \bar{\tau}_R \sin \theta_\tau \), and \( \tilde{\tau}_2 = -\bar{\tau}_L \sin \theta_\tau + \bar{\tau}_R \cos \theta_\tau \) with the masses \( m_{\tilde{\tau}_1}, m_{\tilde{\tau}_2} (m_{\tilde{\tau}_1} < m_{\tilde{\tau}_2}) \) and the mixing angle \( \theta_\tau \). The stau mixing is large if \( |m_{\tilde{\tau}_1}^2 - m_{\tilde{\tau}_2}^2| \ll |a_\tau m_\tau| \), which may be the case for large \( \tan \beta \) and \( \mu \). The mass of \( \tilde{\nu}_\tau \) is given by

\[ m_{\tilde{\nu}_\tau}^2 = M^2_L + \frac{1}{2} m_Z^2 \cos 2\beta. \]

The properties of the charginos \( \tilde{\chi}_i^\pm (i = 1, 2; m_{\tilde{\chi}_1^\pm} < m_{\tilde{\chi}_2^\pm}) \) and neutralinos \( \tilde{\chi}_k^0 (k = 1, \ldots; 4; m_{\tilde{\chi}_1^0} < \ldots < m_{\tilde{\chi}_4^0}) \) are determined by the parameters \( M, M', \mu \) and \( \tan \beta \), where \( M \) and \( M' \) are the SU(2) and U(1) gaugino masses, respectively. Assuming gaugino mass unification we take \( M' = (5/3) \tan^2 \theta_W M \). The masses and couplings of the Higgs bosons \( h^0, H^0, A^0, H^\pm \), including leading radiative corrections, are fixed by \( m_A, \tan \beta, \mu, m_t, m_b, M_Q, M_U, M_D, A_t, \) and \( A_b \), where \( M_{Q,U,D} \) and \( A_{t,b} \) are soft SUSY-breaking parameters in the \((\tilde{t}, \tilde{b})\) sector. \( H^0 \) \( (h^0) \) and \( A^0 \) are the heavier (lighter) CP-even and CP-odd neutral Higgs bosons, respectively. For the radiative corrections to the \( h^0 \) and \( H^0 \) masses and their mixing angle \( \alpha \), we use the formulae of Ref. \[10\]; for those to \( m_{H^\pm} \) we follow Ref. \[11\]. We treat \( M_{L,E,Q,U,D} \) and \( A_{\tau,t,b} \) as free parameters.

The widths of the \( \tilde{\ell}_i \rightarrow \tilde{\tau}_2 \) or \( \tilde{\nu}_\tau \) decays into Higgs and gauge bosons are given by \((k = 1, \ldots, 4)\) \[9\]:

\[ \Gamma(\tilde{\ell}_i \rightarrow \tilde{\ell}_j^{(l)} H_k) = \frac{\kappa_{ijk}}{16\pi m_{\tilde{\ell}_i}^3} (G_{ij})^2, \quad \Gamma(\tilde{\ell}_i \rightarrow \tilde{\ell}_j^{(l)} V) = \frac{\kappa_{ijk}}{16\pi m_{\tilde{\ell}_i}^3} (c_{ijV})^2. \]

Here \( \tilde{\ell}_j^{(l)} = \tilde{\tau}_2 \) or \( \tilde{\nu}_\tau \) (with the indices \( i \) and \( j \) to be omitted for \( \tilde{\nu}_\tau \)), \( H_k = \{ h^0, H^0, A^0, H^\pm \} \) and \( V = \{ Z^0, W^\pm \} \). \( \kappa_{ijk} \equiv \kappa (m_{\tilde{\ell}_1}^2, m_{\tilde{\ell}_0}^2, m_{\tilde{\ell}_2}^2) \) is the usual kinematic factor, \( \kappa(x, y, z) = (x^2 + y^2 + z^2 - 2xy - 2xz - 2yz)^{1/2} \). Notice an extra factor \( \kappa^2/m_{\tilde{\ell}_i}^2 \) for the gauge boson modes. \( G_{ijk} \) denote the slepton couplings to Higgs bosons and \( c_{ijV} \) those to gauge bosons. Their complete expressions, as well as the widths of the fermionic decays, are given in \[9\].

Since \( m_\tau \) is rather small, we need a large difference between \( m_{\tilde{\tau}_L} \sim m_{\tilde{\nu}_\tau} \) and \( m_{\tilde{\tau}_R} \) in order to realize a mass splitting between \((\tilde{\tau}_1, \tilde{\tau}_2, \tilde{\nu}_\tau)\) large enough to allow the bosonic decays in Eqs. (2, 3). In this case the \( \tilde{\tau}_L - \tilde{\tau}_R \) mixing is rather small. Thus, in this article we consider two patterns of the mass spectrum of the sleptons: \( m_{\tilde{\tau}_1} < m_{\tilde{\tau}_2} \sim m_{\tilde{\nu}_\tau} \) with \((\tilde{\tau}_1, \tilde{\tau}_2) \sim (\tilde{\tau}_R, \tilde{\tau}_L)\) for \( m_{\tilde{\tau}_L} > m_{\tilde{\tau}_R} \), and \( m_{\tilde{\tau}_2} > m_{\tilde{\tau}_1} \sim m_{\tilde{\nu}_\tau} \) with \((\tilde{\tau}_1, \tilde{\tau}_2) \sim (\tilde{\tau}_L, \tilde{\tau}_R)\) for

\footnote{Notice that \[9\] \[10\] have a sign convention for the parameter \( \mu \) opposite to the one used here.}
m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R}$. The bosonic decays are therefore basically the decays of \((\tilde{\tau}_L, \tilde{\nu}_\tau)\) into \(\tilde{\tau}_R\) or vice versa. This is in strong contrast to the case of the \(\tilde{t}_2\) and \(\tilde{b}_2\) decays. In the latter case, large mass splittings can be obtained even for \(M_{\tilde{Q}} \sim M_{\tilde{L}} \sim M_{\tilde{D}}\) due to the large top and/or bottom Yukawa couplings, which can also cause large left-right mixings and a complex decay spectrum.

The leading terms of \(G_{ijk}\) and \(c_{ijV}\) which are relevant for the bosonic decays are given in Table 1. Here the Yukawa coupling \(h_\tau\) is proportional to the Yukawa coupling \(h_{ijV}\) which are relevant for the bosonic decays are suppressed by the small mixing. In contrast, the gauge couplings which are relevant for the bosonic decays are suppressed by the small mixing for \(\tilde{\tau}_i \sim \tilde{\tau}_R\). Hence the widths of the decays into Higgs bosons can be large for large tan \(\beta\), \(A_\tau\), and \(\mu\). Notice that the \(\tilde{\tau}_L - \tilde{\tau}_R\) mixing enhances the mass splitting between \(\tilde{\tau}_1\) and \(\tilde{\tau}_2\), which results in a larger phase space for the bosonic decays of \(\tilde{\tau}_2\). In contrast, the gauge couplings which are relevant for the bosonic decays are suppressed by the small mixing of sleptons, since the gauge interactions preserve the chirality of sleptons. However, this suppression is largely compensated by the extra factor \(\kappa^2/m_{\tilde{\nu}}^2\) in Eq. (10). In fact, since the gauge bosons in the decays are longitudinally polarized, the widths of the decays into gauge bosons for \(m_{\tilde{\tau}_i} - m_{\tilde{\nu}_{\tilde{j}}} \gg m_\nu\) are approximated by those into the corresponding Nambu-Goldstone (NG) bosons. As a result, the decays into gauge bosons are enhanced when the couplings of the sleptons to the NG bosons \((\propto a_\tau m_\tau/m_\nu)\) are large.

Table 1: Leading terms of the slepton couplings to Higgs and gauge bosons.

| Coupling                 | Expression                                      |
|-------------------------|-------------------------------------------------|
| \(\tilde{\tau}_1 \tilde{\tau}_2 h^0\) | \(h_\tau (\mu \cos \alpha + A_\tau \sin \alpha) \cos 2\theta_\tau\) |
| \(\tilde{\tau}_1 \tilde{\tau}_2 H^0\)   | \(h_\tau (\mu \sin \alpha - A_\tau \cos \alpha) \cos 2\theta_\tau\) |
| \(\tilde{\tau}_1 \tilde{\tau}_2 A^0\)   | \(h_\tau (\mu \cos \beta + A_\tau \sin \beta)\) |
| \(\tilde{\tau}_1 \tilde{\tau}_2 Z^0\)   | \(g \sin 2\theta_\tau\) |
| \(\tilde{\nu}_\tau \tilde{\tau}_1 H^\pm\) | \(h_\tau (\mu \cos \beta + A_\tau \sin \beta) \sin \theta_\tau\) |
| \(\tilde{\nu}_\tau \tilde{\tau}_2 H^\pm\) | \(h_\tau (\mu \cos \beta + A_\tau \sin \beta) \cos \theta_\tau\) |
| \(\tilde{\nu}_\tau \tilde{\tau}_1 W^\pm\) | \(-g \cos \theta_\tau\) |
| \(\tilde{\nu}_\tau \tilde{\tau}_2 W^\pm\) | \(-g \sin \theta_\tau\) |

Notice the factor \(\cos 2\theta_\tau\) in the \(\tilde{\tau}_1 \tilde{\tau}_2 h^0\) and \(\tilde{\tau}_1 \tilde{\tau}_2 H^0\) couplings, but not in the \(\tilde{\tau}_1 \tilde{\tau}_2 A^0\) coupling. In our case \(|\cos 2\theta_\tau| \sim 1\) due to the small \(\tilde{\tau}_L - \tilde{\tau}_R\) mixing, unlike the case of \(\tilde{t}_2\) and \(\tilde{b}_2\) decays. Similarly, the \(\tilde{\nu}_\tau \tilde{\tau}_1 H^\pm\) coupling is not suppressed by \(\tau\) mixing for \(\tilde{\tau}_i \sim \tilde{\tau}_R\).
We now turn to the numerical analysis of the $\tilde{\tau}_2$ and $\tilde{\nu}_\tau$ decay branching ratios. We calculate the widths of all possibly important two-body decay modes of Eqs. (1, 2, 3). Three-body decays are negligible in this study. We take $m_\tau = 1.78$ GeV, $m_t = 175$ GeV, $m_b = 5$ GeV, $m_Z = 91.2$ GeV, $\sin^2 \theta_W = 0.23$, $m_W = m_Z \cos \theta_W$ and $\alpha(m_Z) = 1/129$.

In order not to vary too many parameters we fix $M = 300$ GeV, $M_A = 150$ GeV, and $M_Q = M_U = M_D = A_t = A_b = 500$ GeV for simplicity. In our numerical study we take $m_{\tilde{\tau}_1}$, $m_{\tilde{\tau}_2}$, $A_\tau$, $\mu$, and $\tan \beta$ as input parameters. Note that for a given set of the input parameters we have two solutions for $(M_L, M_E)$ corresponding to the two cases $m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R}$ and $m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R}$. In the plots we impose the following conditions in order to respect experimental and theoretical constraints:

(i) $m_{\tilde{\chi}^\pm_1} > 100$ GeV, $m_{\tilde{\nu}_\tau} > 45$ GeV, $m_{h^0} > 90$ GeV, $m_{\tilde{\tau}_1, \tilde{\tau}_1, \tilde{b}_1} > m_{\tilde{\chi}^0_1}$, $m_{\tilde{\chi}^0_1} > 80$ GeV,

(ii) $A_\tau^2 < (M_L^2 + M_E^2 + m_{H_1}^2)$, $A_\tau^2 < 3 (M_Q^2 + M_U^2 + m_{H_2}^2)$, and $A_\tau^2 < 3 (M_Q^2 + M_D^2 + m_{H_1}^2)$, where $m_{H_1}^2 = (m_A^2 + m_Z^2) \sin^2 \beta - \frac{1}{2} m_Z^2$ and $m_{H_2}^2 = (m_A^2 + m_Z^2) \cos^2 \beta - \frac{1}{2} m_Z^2$,

(iii) $\Delta \rho \left(\tilde{t} - \tilde{b}\right) < 0.0012$ using the formula of [12].

Condition (i) is imposed to satisfy the experimental bounds on $\tilde{\chi}^\pm_1$, $\tilde{\chi}_1^0$, $\tilde{\tau}$, $\tilde{\nu}_\tau$, $\tilde{b}$, $\tilde{g}$, and $h^0$ from LEP [13] and Tevatron [14]. Note that $m_{\tilde{\chi}_1^0} > 80$ GeV is imposed in order to evade the experimental bounds on $m_{\tilde{t}_1}$ and $m_{\tilde{b}_1}$. Condition (ii) is approximately necessary to avoid color and charge breaking global minima [15] and to exclude unrealistically large $A_\tau$. Condition (iii) constrains $\mu$ and $\tan \beta$ in the squark sector. We note that the experimental data for the $b \to s \gamma$ decay give rather strong constraints [16] on the SUSY and Higgs parameters within the minimal supergravity model, especially for large $\tan \beta$. However, we do not impose this constraint since it strongly depends on the detailed properties of the squarks, including the generation-mixing.

In Fig. 1 we plot in the $A_\tau - \mu$ plane the contours of the branching ratios of the Higgs boson modes $\text{BR}(\tilde{\ell} \to \tilde{\ell} \ell H) \equiv \sum \text{BR}[\tilde{\ell} \to \tilde{\ell} + (h^0, A^0, H^\pm)]$, the gauge boson modes $\text{BR}(\tilde{\ell} \to \tilde{\ell} V) \equiv \sum \text{BR}[\tilde{\ell} \to \tilde{\ell} + (Z^0, W^\pm)]$, and the total bosonic modes $\text{BR}(\tilde{\ell} \to \tilde{\ell} B) \equiv \text{BR}(\tilde{\ell} \to \tilde{\ell} H) + \text{BR}(\tilde{\ell} \to \tilde{\ell} V)$, with $\ell = (\tilde{\tau}_2^-, \tilde{\nu}_\tau)$ and $\tilde{\ell} = (\tilde{\tau}_1^-, \tilde{\nu}_\tau)$. We take $m_{\tilde{\tau}_1} = 250$ GeV, $m_{\tilde{\tau}_2} = 500$ GeV, $\tan \beta = 30$ and show two cases $m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R}$ and $m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R}$, respectively. Note that $\tilde{\nu}_\tau$ decays into bosons and $\tilde{\tau}_2^-$ decays into $\tilde{\nu}_\tau$ are kinematically forbidden if $m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R}$ and $m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R}$, respectively.

We observe that $\text{BR}(\tilde{\ell} \to \tilde{\ell} H)$ increases with $|A_\tau|$ while $\text{BR}(\tilde{\ell} \to \tilde{\ell} V)$ increases with $|\mu|$. This dependence on $(A_\tau, \mu)$ is explained as follows. Note first that for large $\tan \beta$ the mixing between $H_1$ and $H_2$ is rather small. Hence, for $m_A > m_Z$, $(H^0, A^0, H^\pm)$ are mainly $H_1$ while $h^0$ and the NG bosons are mainly $H_2$. The couplings of $(\tilde{\tau}, \tilde{\nu}_\tau)$ to $H_1$ and $H_2$ are $\sim h_\tau A_\tau$ and $h_\tau \mu$, respectively. Therefore the decays to $(H^0, A^0, H^\pm)$ are enhanced for large $|A_\tau|$, whereas those to $h^0$ and to gauge bosons are enhanced for large $|\mu|$. This property can also be derived directly from Table 1 by noting that $|\sin \alpha| \ll 1$ and $\sin \beta \sim 1$. As a result, the total bosonic branching ratio $\text{BR}(\tilde{\ell} \to \tilde{\ell} B)$ becomes large,
and even dominant for \( m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R} \), in a wide region of the \( A_{\tau} - \mu \) plane, especially for large \(|A_{\tau}| \) and/or \(|\mu| \), as seen in Fig. 1. We also see that the branching ratios of the bosonic decays are almost unchanged by \( A_{\tau} \) increase with \( \tan \beta \). tan \( \beta \) tan \( \beta \) tan \( \beta \) tan \( \beta \) tan \( \beta \) tan \( \beta \). In Fig. 2 we show the individual branching ratios of the \( \tilde{\tau}_2 \) and \( \tilde{\nu}_\tau \) decays as a function of \( \tan \beta \) for \( m_{\tilde{\tau}_1} = 250 \) GeV, \( m_{\tilde{\tau}_2} = 500 \) GeV, \( A_{\tau} = 800 \) GeV and \( \mu = 1000 \) GeV, for the two cases \( m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R} \) and \( m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R} \). We see that the branching ratios of the boson modes increase with \( \tan \beta \) and become important, and even dominant if \( m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R} \), for large \( \tan \beta \) (\( \gtrsim 15 \)). As already explained, this comes from the increase of \( h_\tau \) and \( a_{\tau} m_{\tau} \) with \( \tan \beta \). We have checked that \( \Gamma(\tilde{\tau}_2 \to \tilde{\tau}_1 H^0) \) itself increases with \( \tan \beta \).

In Fig. 3 we show the \( m_{\tilde{\tau}_2} \) dependence of the \( \tilde{\tau}_2 \) and \( \tilde{\nu}_\tau \) decay branching ratios for \( m_{\tilde{\tau}_1} = 250 \) GeV, \( A_{\tau} = 800 \) GeV, \( \mu = 1000 \) GeV and \( \tan \beta = 30 \). We see that the branching ratios of the bosonic decays decrease with increasing \( m_{\tilde{\tau}_2} \). This behavior comes from the fact that in the large \( m_{\tilde{\tau}_2} \) limit the decay widths of the bosonic and fermionic modes are proportional to \( m_{\tilde{\tau}_2}^{-1} \) and \( m_{\tilde{\tau}_2} \), respectively.

In all figures one can see that the total branching ratio of the bosonic decays of \( \tilde{\tau}_2 \) is substantially larger for \( m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R} \) than for \( m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R} \). The reason is the following: First, as already explained, the decays of \( \tilde{\tau}_2 \) into \( \tilde{\nu}_\tau \) are kinematically allowed only for \( m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R} \), as seen in Figs. 2 and 3. Second, for the parameter set used in Figs. 2 and 3, \( \tilde{\chi}_1 \sim W^- \). Hence the decay \( \tilde{\tau}_2 \rightarrow \nu_\tau \tilde{\chi}_1 \) is strongly suppressed for \( m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R} \) (i.e. \( \tilde{\tau}_2 \sim \tilde{\tau}_R \)), but not for \( m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R} \), as seen in Figs. 2 and 3; the \( \tilde{\tau}_2 \) decay into \( \tilde{\chi}_2 \sim H^- \) is kinematically forbidden. This results in a rather strong enhancement of the bosonic modes for the case \( m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R} \).

We find that the importance of the bosonic modes is fairly insensitive to the choice of the values of \( M_A \), \( M_{Q,U,D} \), and \( A_{t,b} \). The decays to \( H^0, A^0 \) and \( H^\pm \) are kinematically suppressed for large \( M_A \). However, the remaining \( h^0 \) and gauge boson modes can still be important and even dominant. For example, for \( m_{\tilde{\tau}_1} = 250 \) GeV, \( m_{\tilde{\tau}_2} = 500 \) GeV, \( A_{\tau} = 800 \) GeV, \( \mu = 1000 \) GeV, \( \tan \beta = 30 \) and \( m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R} \), the \( \tilde{\tau}_2 \) decays into \( (H^0, A^0, H^\pm) \) are forbidden for \( M_A > 260 \) GeV. Nevertheless, in this case we have \( \text{BR}(\tilde{\tau}_2 \to \tilde{\tau}_1 h^0) = 13\%\) and \( \text{BR}(\tilde{\tau}_2 \to \tilde{\ell}V) = 52\%\). We have also checked that our results do not change significantly for smaller values of \( M_A \) in the range where the condition (i) is satisfied.

Now we discuss the signatures of the \( \tilde{\tau}_2 \) and \( \tilde{\nu}_\tau \) decays. We compare the signals of the decays into bosons (Eqs. (4, 5)) with those of the decays into fermions (Eq. (6)). The bosonic decays always produce cascade decays. In principle, the final states of the bosonic decays can also be generated from fermionic decays. For example, the final particles of the decay chain

\[
\tilde{\tau}_2^- \rightarrow \tilde{\tau}_1^- + (h^0, H^0, A^0 \text{ or } Z^0) \rightarrow (\tau^- \tilde{\chi}_1^0) + (b\bar{b}) \tag{11}
\]

are the same as those of

\[
\tilde{\tau}_2^- \rightarrow \tau^- + \tilde{\chi}_{2,3,4}^0 \rightarrow \tau^- + ((h^0, H^0, A^0 \text{ or } Z^0) + \tilde{\chi}_1^0) \rightarrow \tau^- + (b\bar{b}\tilde{\chi}_1^0). \tag{12}
\]
Likewise,

\[
\tilde{\nu}_\tau \to \tilde{\tau}_1^- + (H^+ \text{ or } W^+) \to (\tau^- \tilde{\chi}_1^0) + (q\bar{q}')
\]

has the same final particles as

\[
\tilde{\nu}_\tau \to \tau^- + \tilde{\chi}^+_1 \to \tau^- + ((H^+ \text{ or } W^+) + \tilde{\chi}_1^0) \to \tau^- + (q\bar{q}\tilde{\chi}_1^0).
\]

Nevertheless, the decay distributions of the two processes (11) and (12) ((13) and (14)) are in general different from each other due to the different intermediate states. For example, the \(\tau^-\) in the chains (11) and (13) tends to be softer than the \(\tau^-\) in (12) and (14), respectively. A similar argument holds for the quark pairs in the decay chains. Moreover, the distribution of the missing energy-momentum carried by \(\tilde{\chi}_1^0\) could be significantly different in (11) and (12) ((13) and (14)) since it is emitted from a different sparticle. Detailed Monte Carlo simulations are necessary to investigate the experimental consequences of the bosonic decays.

In conclusion, we have shown that the decays of \(\tilde{\tau}_2\) and \(\tilde{\nu}_\tau\) into Higgs or gauge bosons, such as \(\tilde{\tau}_2^- \to \tilde{\tau}_1^- + (h^0, H^0, A^0 \text{ or } Z^0)\) and \(\tilde{\tau}_2^- \to \tilde{\nu}_\tau + (H^- \text{ or } W^-)\), can be very important in a fairly wide MSSM parameter region with large mass splitting between \(\tilde{\tau}_1\) and \(\tilde{\tau}_2\), large \(\tan\beta\), and large \(|A_{\tau}|\) and/or \(|\mu|\). Compared to the fermionic decay modes, these bosonic decay modes could have significantly different decay distributions. This could have important implications for the searches of \(\tilde{\tau}_2\) and \(\tilde{\nu}_\tau\) and the determination of the MSSM parameters at future colliders.

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Figure Captions

**Figure 1**: Branching ratios of $\tilde{\tau}_2$ and $\tilde{\nu}_\tau$ decays in the $A_{\tau}\mu$ plane for $m_{\tilde{\tau}_1} = 250$ GeV, $m_{\tilde{\tau}_2} = 500$ GeV and $\tan \beta = 30$ in the cases of $m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R}$ (a – c) and $m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R}$ (d – i); (a, d) $\sum$ BR[$\tilde{\tau}_2^- \rightarrow \tilde{\tau}_1^+ + (h^0, H^0, A^0)$, $\tilde{\nu}_\tau + H^-$], (b, e) $\sum$ BR[$\tilde{\tau}_2^- \rightarrow \tilde{\tau}_1^+ + Z^0$, $\tilde{\nu}_\tau + W^-$], (c, f) $\sum$ BR[$\tilde{\tau}_2^- \rightarrow \tilde{\tau}_1^- + (h^0, H^0, A^0, Z^0)$, $\tilde{\nu}_\tau + (H^-, W^-)$], (g) BR[$\tilde{\nu}_\tau \rightarrow \tilde{\tau}_1^- + H^+$], (h) BR[$\tilde{\nu}_\tau \rightarrow \tilde{\tau}_1^- + W^+$], and (i) $\sum$ BR[$\tilde{\nu}_\tau \rightarrow \tilde{\tau}_1^- + (H^+, W^+)$]. The gray areas are excluded by the conditions (i) to (iii) given in the text.

**Figure 2**: $\tan \beta$ dependence of $\tilde{\tau}_2$ (a, b) and $\tilde{\nu}_\tau$ (c) decay branching ratios for $m_{\tilde{\tau}_1} = 250$ GeV, $m_{\tilde{\tau}_2} = 500$ GeV, $A_{\tau} = 800$ GeV and $\mu = 1000$ GeV in the cases of $m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R}$ (a) and $m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R}$ (b, c). ”Gauge/Higgs + X” refers to the sum of the gauge and Higgs boson modes. The gray areas are excluded by the conditions (i) to (iii) given in the text.

**Figure 3**: $m_{\tilde{\tau}_2}$ dependence of $\tilde{\tau}_2$ (a, b) and $\tilde{\nu}_\tau$ (c) decay branching ratios for $m_{\tilde{\tau}_1} = 250$ GeV, $A_{\tau} = 800$ GeV, $\mu = 1000$ GeV, and $\tan \beta = 30$ in the cases of $m_{\tilde{\tau}_L} < m_{\tilde{\tau}_R}$ (a) and $m_{\tilde{\tau}_L} \geq m_{\tilde{\tau}_R}$ (b, c). ”Gauge/Higgs + X” refers to the sum of the gauge and Higgs boson modes. The gray areas are excluded by the conditions (i) to (iii) given in the text. Note that $m_{\tilde{\nu}_\tau} \sim m_{\tilde{\tau}_2}$ in (c).
Fig. 1
Fig. 2
Fig. 3