Impact of squark pair decay modes on search for supersymmetric neutral Higgs bosons

A. Bartl,² H. Eberl,² K. Hidaka,³ T. Kon,⁴ W. Majerotto² and Y. Yamada⁵

Abstract

We present a detailed study of the decays of the neutral Higgs bosons \( H^0 \) and \( A^0 \) within the Minimal Supersymmetric Standard Model. We find that the supersymmetric modes \( \tilde{t}\tilde{t} \) and \( \tilde{b}\tilde{b} \) can dominate the \( H^0 \) and \( A^0 \) decays in a wide range of the model parameters due to large Yukawa couplings and mixings of \( \tilde{t} \) and \( \tilde{b} \). Compared to the conventional modes \( t\bar{t} \) and \( b\bar{b} \), these modes have very distinctive signatures. This could have a decisive impact on the neutral Higgs boson searches at future colliders.
The existence of two or more Higgs bosons would be a clear indication that
the Standard Model must be extended. For example, the Minimal Supersymmetric
Standard Model (MSSM) \[1\] with two Higgs doublets predicts the existence of five
physical Higgs bosons $h^0$, $H^0$, $A^0$, and $H^{\pm}$ \[2, 3\]. In searching for the Higgs bosons
it is necessary to study not only the production mechanism, but also all their poss-
sible decay modes. Their decays to supersymmetric (SUSY) particles could be very
important if they are kinematically allowed. This is indeed the case for the charged
Higgs boson $H^+$. If all SUSY particles are very heavy, the $H^+$ decays dominantly
into $t\bar{b}$; the decays $H^+ \to \tau^+\nu$ and/or $H^+ \to W^+h^0$ are dominant below the $t\bar{b}$
threshold \[2, 4\]. In refs.\[5, 6\] all decay modes of $H^+$ including the SUSY-particle
modes were studied in the case that the SUSY-particles are relatively light; it
was shown that the SUSY decay mode $H^+ \to t\bar{b}$ can be dominant in a large region of
the MSSM parameter space due to large $t$ and $b$ quark Yukawa couplings and large $\tilde{t}$-
and $\tilde{b}$-mixings, and that this could have a decisive impact on $H^+$ searches at future
colliders. In this paper we extend this study to the decay of the neutral Higgs bosons
$H^0$ and $A^0$.

The lighter stop $\tilde{t}_1$ can be much lighter than the other squarks and even lighter
than the $t$ quark due to large $\tilde{t}_L$-$\tilde{t}_R$ mixing being proportional to the large top Yukawa
coupling $h_t$ and the $\tilde{t}$-mixing parameters $A_t$ and $\mu$ \[7\]. Similarly, the lighter sbottom
$\tilde{b}_1$ can also be much lighter than the other squarks \[4\]. In the case of large $\tilde{t}$- and
$\tilde{b}$- mixings one also expects the couplings of $H^0\tilde{t}\tilde{t}$, $A^0\tilde{t}\tilde{t}$, $H^0\tilde{b}\tilde{b}$ and $A^0\tilde{b}\tilde{b}$ to be large.
Here we show explicitly that the modes $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ can indeed dominate the $H^0$ and
$A^0$ decays in a wide range of the MSSM parameters.

First we summarize current experimental limits on the squark masses. The D∅ group at FNAL obtained mass bounds for squarks and gluinos from the negative search for them. However, they obtained no mass bound for the mass-degenerate five flavors of squarks ($\tilde{q}$) (other than the stop ($\tilde{t}$)) within the MSSM in case the lightest neutralino $\tilde{\chi}_1^0$ (assumed to be the lightest SUSY particle (LSP)) is heavier than $\sim 70\text{GeV}$; this is mainly due to the massive LSP effect. On the other hand, the LEP-I experiments set the general bound on the $\tilde{q}$ mass $m_{\tilde{q}} \gtrsim 45\text{GeV}$ [11, 12]. Experiments dedicated to $\tilde{t}$ search at LEP1.5 and Tevatron have set some limits on the lighter stop ($\tilde{t}_1$) mass for relatively light $\tilde{\chi}_1^0$, e.g., typically $m_{\tilde{t}_1} \gtrsim 55\text{GeV}$ (for no $\tilde{t}$-mixing) [11] and $m_{\tilde{t}_1} \gtrsim 100\text{GeV}$ [12], respectively. However, they set no limit on $m_{\tilde{t}_1}$ for $m_{\tilde{\chi}_1^0} \gtrsim 50\text{GeV}$ [11, 12].

In the MSSM the properties of the charginos $\tilde{\chi}_i^\pm$ ($i = 1, 2$) and neutralinos $\tilde{\chi}_j^0$ ($j = 1, \ldots, 4$) are completely determined by the parameters $M$, $\mu$ and $\tan \beta = v_2/v_1$, assuming $M' = (5/3) \tan^2 \theta_W M$. Here $M$ ($M'$) is the SU(2) (U(1)) gaugino mass, $\mu$ is the higgsino mass parameter, and $v_1$ ($v_2$) is the vacuum expectation value of the Higgs $H_1^0$ ($H_2^0$) [1]. Here $m_{\tilde{\chi}_1^+} < m_{\tilde{\chi}_2^+}$ and $m_{\tilde{\chi}_1^0} < \cdots < m_{\tilde{\chi}_4^0}$. To specify the squark sector additional (soft SUSY breaking) parameters $M_Q$, $M_U$, $M_D$ (for each generation) and $A$ (for each flavor) are necessary. The mass matrix for stops in the base of ($\tilde{t}_L$, $\tilde{t}_R$) reads [4, 5]

$$M_{\tilde{t}}^2 = \begin{pmatrix} m_{\tilde{t}_L}^2 & a_{\tilde{t}} m_t \\ a_{\tilde{t}} m_t & m_{\tilde{t}_R}^2 \end{pmatrix}$$
with

\[ m_{t_L}^2 = M_{\tilde{Q}}^2 + m_t^2 + m_Z^2 \cos 2\beta (I_1^3 - c_t \sin^2 \theta_W), \quad (2) \]

\[ m_{t_R}^2 = M_{\tilde{U}}^2 + m_t^2 + m_Z^2 \cos 2\beta e_t \sin^2 \theta_W, \quad \text{and} \]

\[ a_t m_t = m_t (A_t - \mu \cot \beta). \quad (4) \]

For the \( \tilde{b} \) system analogous formulae hold but with \( M_{\tilde{U}}^2 \) replaced by \( M_{\tilde{D}}^2 \) in eq.(3), and instead of eq.(4), \( a_b m_b = m_b (A_b - \mu \tan \beta) \). \( \tilde{b}_L - \tilde{b}_R \) mixing may be important for large \( A_b \), \( \mu \) and \( \tan \beta \). The slepton sector is fixed by adding (soft SUSY breaking) parameters \( M_{\tilde{L}} \) and \( M_{\tilde{E}} \).

The masses and couplings of the Higgs bosons \( H^\pm, H^0, h^0 \) and \( A^0 \), including leading Yukawa corrections, are fixed by \( m_{A^0}, \tan \beta, m_t, M_{\tilde{Q}}, M_{\tilde{U}}, M_{\tilde{D}}, A_t, A_b \) and \( \mu \). \( H^0 \) (\( h^0 \)) and \( A^0 \) are the heavier (lighter) CP-even and CP-odd neutral Higgs bosons, respectively. For the Yukawa corrections to the \( h^0 \) and \( H^0 \) masses and their mixing angle \( \alpha \) we use the formulae of ref.[13]. We include also the potentially large Yukawa corrections to the triple Higgs vertices of \( H^0 h^0 h^0 \), \( H^0 A^0 A^0 \) and \( h^0 A^0 A^0 \) using the formulae of ref.[14].

In the following, we take for simplicity \( M_{\tilde{Q}} = M_{\tilde{U}} = M_{\tilde{D}} \) (for the third generation), \( M_{\tilde{L}} = M_{\tilde{E}} = M_{\tilde{Q}} \), and \( A_t = A_b = A_\tau \equiv A \). Thus we have as free parameters

\( m_{A^0}, \mu, \tan \beta, M_{\tilde{Q}} \) and \( A \).

We calculate the widths of all possibly important modes of \( H^0 \) and \( A^0 \) decays:

(i) \( H^0 \to t\bar{t}, b\bar{b}, c\bar{c}, \tau^+\tau^-, W^+W^-, Z^0Z^0, h^0h^0, A^0A^0, W^\pm H^\mp, Z^0A^0, t_i\bar{t}_j, b_i\bar{b}_j, \tilde{\ell}_i\tilde{\ell}_j^* \), \( \tilde{\nu}_\ell \tilde{\nu}_\ell \) (\( \ell = e, \mu, \tau \)), \( \tilde{\chi}_i^+\tilde{\chi}_j^- \), \( \tilde{\chi}_k\tilde{\chi}_l^0 \), and (ii) \( A^0 \to t\bar{t}, b\bar{b}, c\bar{c}, \tau^+\tau^-, Z^0h^0, t_1\bar{t}_2, t_2\bar{t}_1, b_1\bar{b}_2, \) \( h^0 \).

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\(\tilde{b}_2\tilde{b}_1, \tilde{\tau}_2^\pm, \tilde{\tau}_2^- \tilde{\tau}_1^+, \tilde{\chi}_k^0\tilde{\chi}_l^0\). Formulae for these widths are found in ref. [2]. In principle, also the decays \(H^0 \rightarrow \tilde{q}_\alpha \tilde{q}_\alpha\) \((q = u, d, c, s\) and \(\alpha = L, R\)) could contribute via their gauge couplings. As the squarks of the first two generations are supposed to be heavy, these decays will be strongly phase-space suppressed. Even if they were kinematically allowed, they would have a rate at most comparable to that of \(H^0 \rightarrow \tilde{\ell}_i^+ \tilde{\ell}_j^-\) and \(\tilde{\nu}_i \tilde{\nu}_\ell\) (see fig. 2 below). We neglect loop induced decay modes (such as \(H^0 \rightarrow gj, \gamma\gamma\)) and three-body decay modes [15].

In order not to vary many parameters, in the following we fix \(m_t = 180\text{GeV}\) and \(\mu = 300\text{GeV}\), and take the values of \(M\) and \(\tan\beta\) such that \(m_{\tilde{\chi}_1^0} \simeq 70\text{GeV}\) for which the DØ bounds on \(m_{\tilde{q}}\) and \(m_{\tilde{t}_1}\) [8, 12] and the LEP1.5 bound on \(m_{\tilde{t}_1}\) [11] disappear.

In fig.1 the contour lines for the branching ratios \(B(H \rightarrow \tilde{t}_i \tilde{t}_j, \tilde{b}_i \tilde{b}_j) \equiv \sum_{i,j=1,2}(B(H \rightarrow \tilde{t}_i \tilde{t}_j)+B(H \rightarrow \tilde{b}_i \tilde{b}_j))\) \((H = H^0\) or \(A^0\)) are plotted in the \(A - M_{\tilde{Q}}\) plane for \((m_{A^0}\text{GeV}), M\text{GeV}, \tan\beta\) = (450, 160, 2)(a), (500, 160, 2)(b), (450, 146, 12)(c) and (500, 146, 12)(d), for which \(m_{\tilde{\chi}_1^+}\text{GeV} = 128, 128, 131\) and 131, respectively. In the plots we have required \(m_{\tilde{t}_i, \tilde{b}_i} > m_{\tilde{\chi}_1^0}\) \((\simeq 70\text{GeV})\). In the allowed regions of fig.1 we find \(m_{H^0} \simeq m_{A^0}\) and \(m_{H^0} > 55\text{GeV},\) which satisfies the LEP limit \(m_{H^0} > 44\text{GeV}\) [16]. We see that the branching ratios \(B(H^0 \rightarrow \tilde{t}_i \tilde{t}_j, \tilde{b}_i \tilde{b}_j)\) and \(B(A^0 \rightarrow \tilde{t}_i \tilde{t}_j, \tilde{b}_i \tilde{b}_j)\) can be larger than 50% in a sizable region. In this region, these SUSY decay modes dominate over the conventional modes.

In fig.2 we show the \(m_{A^0}\) dependence (in the \(m_{A^0}\) range for \(m_{H^0} > 50\text{GeV}\)) of the important branching ratios of \(H^0\) and \(A^0\) decays for \((M_{\tilde{Q}}\text{GeV}), A\text{GeV}\),
$M$(GeV), $\tan \beta = (102, 325, 160, 2)$ (a, b), $(218, 405, 145, 30)$ (c), and $(145, 274, 146, 12)$ (d). In these three cases we have $(m_{\tilde{t}_1}, m_{\tilde{t}_2}, m_{\tilde{b}_1}, m_{\tilde{b}_2}, m_{\chi^+}) = (100, 270, 100, 115, 128)$ (a, b), $(80, 386, 80, 304, 132)$ GeV (c), and $(80, 310, 80, 199, 131)$ GeV (d). We see that in these cases the sum of the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ modes dominates the $H^0$ and $A^0$ decays in a wide range of $m_{A^0}$. Here note that $m_{H^0} \simeq m_{A^0}$ in the $m_{A^0}$ range shown here, and that the $A^0$ does not couple to $\tilde{t}_i\tilde{t}_i$, $\tilde{b}_i\tilde{b}_i$ and $\tilde{\tau}^+\tilde{\tau}^-$ ($i = 1, 2$).

As for $h^0$ decay, we have found that the decay $h^0 \to \tilde{t}_1\tilde{t}_1$ is kinematically allowed only in a very limited region of the MSSM parameter space.

The reason for the dominance of the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ modes in the $H^0$ decay is as follows: The modes $t\bar{t}$, $b\bar{b}$, $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ (whose couplings to $H^0$ are essentially $\sim h_t \sin \alpha$, $\sim h_b \cos \alpha$, $\sim (A - \mu \cot \alpha)h_t \sin \alpha$ and $\sim (A - \mu \tan \alpha)h_b \cos \alpha$, respectively) can be strongly enhanced relative to the other modes due to the large Yukawa couplings $h_{t,b}$. In addition, the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ modes can dominate over the $t\bar{t}$ and $b\bar{b}$ modes, respectively, in the case the $\tilde{q}$-mixing parameters $A$ and $\mu$ are large. Furthermore, in this case $\tilde{t}_1$ and $\tilde{b}_1$ tend to be light due to a large mass-splitting. Here note that the effects of the bottom Yukawa coupling $h_b$ and $\tilde{b}$-mixing play a very important role for large $\tan \beta$ (see fig.2c). The reason for the dominance of the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ modes in the $A^0$ decay is similar to that in the $H^0$ decay.

Quite generally, $B(H \to \tilde{t}\tilde{t}, \tilde{b}\tilde{b})$ depends on the parameters $m_{A^0}$, $M_Q$, $A$, $\mu$, $\tan \beta$ and more weakly on $M$. For a given $m_{A^0}$ the strongest dependence is that on $M_Q$ to which $m_{\tilde{t}}$ and $m_{\tilde{b}}$ are sensitive (see fig.1). $B(H \to \tilde{t}\tilde{t}, \tilde{b}\tilde{b})$ can be quite large in a substantial part of the parameter region kinematically allowed for the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$.
modes. We find that the dominance of the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ modes is fairly insensitive to the assumption $M_\tilde{Q} = M_\tilde{U} = M_\tilde{D} = M_\tilde{L} = M_\tilde{E}$. As seen in fig.1 the dependence on $A$ is also strong. Concerning the assumption $A_t = A_b = A_\tau$, we have found no significant change of $B(H \to \tilde{t}\tilde{t}, \tilde{b}\tilde{b})$ ($H = H^0, A^0$) as compared to fig.2, when we take $A_{b,\tau}/A_t = \pm 0.5, \pm 1, \pm 2$ keeping $A_t = A$. As $|\mu|$ increases, the dominance of the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ modes becomes more pronounced because of the increase of the $(H^0, A^0)$ couplings to $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$. We also find that $B(H \to \tilde{t}\tilde{t}, \tilde{b}\tilde{b})$ ($H = H^0, A^0$) is nearly invariant under $(\mu, A) \to (-\mu, -A)$.

Here we have not included QCD radiative corrections to the hadronic modes. It is shown in ref.\cite{17} that both the standard QCD correction (due to gluon-quark loop) and SUSY QCD correction (due to gluino-squark loop) to the widths of $H^0, A^0 \to \tilde{t}\tilde{t}, \tilde{b}\tilde{b}$ can be large, but that the two corrections can partly or even totally cancel each other. The QCD corrections to $H^0, A^0 \to \tilde{t}\tilde{t}, \tilde{b}\tilde{b}$ are not known. It is shown in ref.\cite{6} that the QCD corrections to the closely related process $H^+ \to \tilde{t}\tilde{b}$ are significant, but that they do not invalidate the tree-level conclusion of ref.\cite{5} on the dominance of the $\tilde{t}\tilde{b}$ mode in a wide parameter region. As the structures of the couplings of $H^0$ and $A^0$ to $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ ($\tilde{t}\tilde{b}$ and $\tilde{b}\tilde{b}$) are similar to those of $H^+$ to $\tilde{t}\tilde{b}$ ($tb$), we can expect that the QCD corrections would not invalidate our tree-level conclusion on the dominance of the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ modes in a large parameter region.

As for the signatures of the $H^0$ and $A^0$ decays, typical $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ signals are shown in table 1. They have to be compared with the conventional $tt$ and $bb$ signals, respectively, $(H^0, A^0) \to \tilde{t}\tilde{t} \to (W^+b)(W^-\bar{b}) \to (f\bar{f}')(b\bar{f}')\bar{b}$ (i.e., 6 jets (j’s), 4j’s + 1
isolated charged lepton ($\ell^\pm$) + missing energy-momentum ($\not{p}$), or 2j’s + $\ell^+\ell^− + \not{p}$) and $(H^0, A^0) \rightarrow b\bar{b}$ (i.e., 2j’s). Note that $B(\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0) \simeq 1$ if $m_{\tilde{t}_1} < m_{\tilde{c}_1^+} + \not{p}_t$ and $m_{\tilde{c}_1^−} < m_{\tilde{t}_1} < m_{b} + m_{\tilde{\chi}_1^0}$, $m_{\tilde{c}_1^−} + m_{\tilde{w}}$ (in cases (a), (b) and (d) of table 1), and $B(\tilde{t}_1 \rightarrow c\tilde{\chi}_1^0) \simeq 0$ otherwise (in case (c)) [18]. (In principle there is a region $m_{\tilde{b}} + m_{\tilde{w}} + m_{\tilde{\chi}_1^0} < m_{\tilde{t}_1}$ where also the decay $\tilde{t}_1 \rightarrow b\tilde{W}^+\tilde{\chi}_1^0$ plays a role [14]. However, this decay does not occur in the cases considered here.) As seen in table 1, the $\tilde{t}\bar{\tilde{t}}$ ($\tilde{b}\bar{\tilde{b}}$) signals have general features which distinguish them from the $t\bar{t}$ ($b\bar{b}$) signals: (i) more $\not{p}$ due to the emission of two LSP’s and hence less energy-momentum of jets and the isolated-charged-lepton in case of a short decay chain, or (ii) more jets and/or more isolated-charged-leptons in case of a larger decay chain. Moreover, depending on the values of the MSSM parameters, the $\tilde{t}\bar{\tilde{t}}$ and $\tilde{b}\bar{\tilde{b}}$ signals could have the following remarkable features: (i) production of same-sign dileptons $\ell^±\ell^±$ (e.g., in case (g)), which could yield same-sign isolated dilepton events such as $e^+e^- \rightarrow H^0Z^0 \rightarrow (\ell^+\ell^+ \text{ or } \ell^−\ell^−) + j's + \not{p}$; (ii) less bottom-jet activity (e.g., in cases (a), (b), (d) with $(h^0, Z^0) \rightarrow (\ell^+\ell^− \text{ or } \nu\bar{\nu})$) or more bottom-jet activity (e.g., in cases (c), (d), (f) with $(h^0, Z^{(*)}) \rightarrow b\bar{b}$); and (iii) emission of a real $Z^0$ or $h^0$ (e.g., in cases (b), (c), (d), (f)). The identification of the sign of charged leptons and the tagging of $b$- and $c$-quark jets, $h^0$, $Z^0$ and $W^±$ would be very useful [2] in discriminating the $\tilde{t}\bar{\tilde{t}}$ and $\tilde{b}\bar{\tilde{b}}$ signals from the $t\bar{t}$ and $b\bar{b}$ signals as well as in suppressing the background.

The suitable places for $H^0$ and $A^0$ search would be $e^+e^−$ colliders [20] and hadron supercolliders [4]. If the $\tilde{t}\bar{\tilde{t}}$ and $\tilde{b}\bar{\tilde{b}}$ modes dominate the $H^0$ and $A^0$ decays, it decisively influences the signatures of $H^0$ and $A^0$. For example, it can strongly
suppress the conventional discovery modes of $H^0$ and $A^0$ at the hadron supercollider
such as the ”gold-plated” $H^0 \rightarrow Z^0Z^0$ mode (viable for small $\tan\beta$ and $2m_Z \lesssim m_{H^0} \lesssim 2m_t$) (see fig.2a) and $H^0$, $A^0 \rightarrow \tau^-\tau^+$ mode (viable for large $\tan\beta$) (see figs.2c, d). Clearly it would be necessary to perform a detailed Monte-Carlo study to separate the signals from the background. Such a study is, however, beyond the scope of this article.

We have shown that the SUSY modes $\tilde{t}\bar{\tilde{t}}$ and $\tilde{b}\bar{\tilde{b}}$ can dominate the $H^0$ and $A^0$ decays in a large allowed region of the MSSM parameter space due to large $t$ and $b$ quark Yukawa couplings and large $\tilde{t}$ - and $\tilde{b}$ -mixings. The $\tilde{t}\bar{\tilde{t}}$ and $\tilde{b}\bar{\tilde{b}}$ modes have very distinctive signatures as compared to the conventional modes $t\bar{t}$, $\bar{b}b$ and $\tau^-\tau^+$. This could decisively influence the $H^0$ and $A^0$ search at future $e^+e^-$ and hadron colliders.

While preparing this manuscript a paper by Djouadi et al. [21] appeared dealing with a similar subject. They also point out the possible importance of the squark pair modes in $H^0$ and $A^0$ decays. They, however, studied them in a strongly constrained supergravity model in contrast to our work, which is in the general framework of the MSSM.

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Figure Captions

Fig.1 Contour lines of $B(H^0 \to \tilde{t}\tilde{t}, \tilde{b}\tilde{b})$ (a, c) and $B(A^0 \to \tilde{t}\tilde{t}, \tilde{b}\tilde{b})$ (b, d) in the $A - M_{\tilde{Q}}$ plane for $(m_{A^0} \text{GeV}, M \text{GeV}, \mu \text{GeV}, \tan \beta) = (450, 160, 300, 2)$ (a), $(500, 160, 300, 2)$ (b), $(450, 146, 300, 12)$ (c) and $(500, 146, 300, 12)$ (d). The contour of $B = 0$ (dashed line) represents the kinematical boundary; the lines (iii), (iv) and (v) are the contour lines of $2m_{\tilde{t}_1} = m_{H^0} \ (\sim 450 \text{ GeV})$, $2m_{\tilde{b}_1} = m_{H^0} \ (\sim 450 \text{ GeV})$ and $m_{\tilde{b}_1} + m_{\tilde{b}_2} = m_{A^0} \ (= 500 \text{ GeV})$, respectively. The shaded area is excluded by the requirement $m_{\tilde{t}_1,\tilde{b}_i} > m_{\tilde{\chi}_1^0} \ (\sim 70 \text{ GeV})$; the lines (i) and (ii) are the contour lines of $m_{\tilde{t}_1}$ and $m_{\tilde{b}_1} = 70 \text{ GeV}$, respectively, above which $m_{\tilde{t}_1}$ and $m_{\tilde{b}_1} > 70 \text{ GeV}$.

Fig.2 The $m_{A^0}$ dependence of important branching ratios of the $H^0$ and $A^0$ decays for $(M_{\tilde{Q}} \text{GeV}, A \text{GeV}, M \text{GeV}, \mu \text{GeV}, \tan \beta) = (102, 325, 160, 300, 2)$ (a, b), $(218, 405, 145, 300, 30)$ (c) and $(145, 274, 146, 300, 12)$ (d). The sum over all mass eigenstates and/or flavors is taken for $B(\tilde{t}\tilde{t})$, $B(\tilde{b}\tilde{b})$, $B(\tilde{\chi}_1^+\tilde{\chi}_1^-)$, $B(\tilde{\chi}_1^0\tilde{\chi}_1^0)$, $B(\tilde{\tau}^+\tilde{\tau}^-)$ and $B(\tilde{\nu}\tilde{\nu})$. The shoulders in the curves for the $\tilde{t}\tilde{t}$ and $\tilde{b}\tilde{b}$ modes correspond to opening of the $\tilde{t}_2\tilde{t}_2$ and $\tilde{b}_2\tilde{b}_2$ channels, respectively.
Figure 1

(a) $H^0$ decay
$tan\beta=2$
$m_{A^0}=450\text{GeV}$

(b) $A^0$ decay
$tan\beta=2$
$m_{A^0}=500\text{GeV}$

Fig. 1
\[ A^0 \text{ decay} \]
\[ \tan \beta = 12 \]
\[ m_{A^0} = 450 \text{GeV} \]

\[ H^0 \text{ decay} \]
\[ \tan \beta = 12 \]
\[ m_{A^0} = 500 \text{GeV} \]

Fig. 1
Fig. 2
Fig. 2
Table 1: The typical $\tilde{t}\tilde{t}$ and $b\bar{b}$ signals of the $(H^0, A^0)$ decays in comparison to the conventional $t\bar{t}$ and $b\bar{b}$ signals, respectively. The decay chains (a), (b) and (d) are possible if $m_{\tilde{t}^1_i} < m_{\chi_1^\pm} + m_{\tilde{b}_i}$ and $m_{\tilde{b}_i^0} < m_{\chi_1^0} < m_{\chi_2^0} + m_{\tilde{b}_i}$. $p_j, \ell^\pm, Z^{(*)}$ and $f$ denote missing energy-momentum, jet, isolated charged lepton, real (or virtual) $Z^0$ boson, and $(q, \ell^-, \nu)$, respectively. $A^0$ does not couple to $\tilde{q}_1 \bar{q}_1$ and $\tilde{q}_2 \bar{q}_2$.

| Typical decay chains | Signatures |
|-----------------------|------------|
| (a) $H^0 \rightarrow$ | $2j's + \bar{p}$ |
| $\bar{t}_1 \rightarrow e\chi_1^0$ | (more $p$, less $b$ activity) |
| $\bar{t}_1 \rightarrow e\chi_1^0$ | |
| (b) $H^0, A^0 \rightarrow$ | $4j's + \bar{p}, 2j's + \bar{p}$ |
| $\tilde{t}_1 \rightarrow e\chi_1^0$ | |
| $\tilde{t}_2 \rightarrow [\tilde{t}_1 Z^0 \rightarrow (e\chi_1^0)(f\bar{f})]$ | $2j's + e^+e^- + \bar{p}$ |
| or | $(Z^0$ or $h^0$ emission ; |
| $\tilde{t}_1 h^0 \rightarrow (e\chi_1^0)(\bar{b}b$ or $\tau^+\tau^-)$ | less $b$ activity if $Z^0$ or $h^0$ |
| (c) $H^0, A^0 \rightarrow$ | $8j's + \bar{p}, 6j's + e^+e^- + \bar{p},$ |
| $\tilde{t}_2 \rightarrow e\chi_1^0$ | $4j's + e^+e^- + \bar{p},$ |
| $\tilde{t}_2 \rightarrow [\tilde{t}_1 Z^0 \rightarrow (e\chi_1^0)(f\bar{f})]$ | $2j's + e^+e^- + \bar{p},$ |
| or | $(Z^0$ or $h^0$ emission ; |
| $\tilde{t}_1 h^0 \rightarrow (e\chi_1^0)(\bar{b}b$ or $\tau^+\tau^-)$ | more $b$ activity if $Z^0$ or $h^0$ |
| (d) $H^0 \rightarrow$ | $6j's + \bar{p}, 4j's + e^+e^- + \bar{p},$ |
| $\tilde{t}_2 \rightarrow e\chi_1^0$ | $4j's + e^+e^- + \bar{p},$ |
| $\tilde{t}_2 \rightarrow [\tilde{t}_1 Z^0 \rightarrow (e\chi_1^0)(f\bar{f})]$ | $2j's + e^+e^- + \bar{p},$ |
| or | $(Z^0$ or $h^0$ emission ; |
| $\tilde{t}_1 h^0 \rightarrow (e\chi_1^0)(\bar{b}b$ or $\tau^+\tau^-)$ | more $b$ activity if $Z^0$ or $h^0$ |
| (e) $H^0, A^0 \rightarrow$ | $2j's + \bar{p}$ |
| $\tilde{b}_{1,2} \rightarrow b\chi_1^0$ | (more $p$) |
| $\tilde{b}_{1,2} \rightarrow b\chi_1^0$ | |
| (f) $H^0, A^0 \rightarrow$ | $8j's + \bar{p}, 6j's + e^+e^- + \bar{p},$ |
| $\tilde{b}_1 \rightarrow b\chi_1^0$ | $4j's + e^+e^- + \bar{p},$ |
| $\tilde{b}_2 \rightarrow [\tilde{b}_1 Z^0 \rightarrow (e\chi_1^0)(f\bar{f})]$ | $6j's + e^+e^- + \bar{p},$ |
| or | $2j's + e^+e^- + \bar{p},$ |
| $\tilde{b}_1 h^0 \rightarrow (e\chi_1^0)(\bar{b}b$ or $\tau^+\tau^-)$ | $(Z^0$ or $h^0$ emission ; |
| (g) $H^0 \rightarrow$ | $8j's + \bar{p}, 6j's + e^+e^- + \bar{p},$ |
| $\tilde{b}_2 \rightarrow b\chi_1^0$ | $4j's + e^+e^- + \bar{p},$ |
| $\tilde{b}_2 \rightarrow [\tilde{b}_1 Z^0 \rightarrow (e\chi_1^0)(f\bar{f})]$ | $2j's + e^+e^- + \bar{p},$ |
| $\tilde{b}_1 h^0 \rightarrow (e\chi_1^0)(\bar{b}b$ or $\tau^+\tau^-)$ | (same-sign dilepton ; |
| (h) $H^0, A^0 \rightarrow$ | trilepton) |