A new signal for scalar top bound state production

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Abstract

We study the production and decay of a scalar ($\tilde{t}_1 \tilde{t}_1^*$) bound state $\sigma_{\tilde{t}_1}$ at hadron supercolliders, where $\tilde{t}_1$ is the lighter stop eigenstate. If $\tilde{t}_1$ has no tree-level 2–body decays, the dominant decay modes of $\sigma_{\tilde{t}_1}$ are $gg$ or, if $m_h < m_{\tilde{t}_1} \ll m_{\tilde{t}_2}$, a pair of light scalar Higgs bosons $h$. Nevertheless the branching ratio into two photons is often large enough to yield a detectable signal.

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One of the best motivated extensions of the Standard Model (SM) of particle physics is the introduction of supersymmetry (SUSY). On the one hand it solves the “hierarchy problem”, i.e. stabilizes the scale of electroweak symmetry breaking against radiative corrections which would otherwise pull it up to the unification scale $M_X$ or the Planck scale [1]. On the other hand recent precise measurements of the strength of coupling constants exclude the possibility of Grand Unification if the only light (compared to $M_X$) particles are those of the SM, while the minimal supersymmetric standard model (MSSM) [2] contains just the right degrees of freedom to allow for unification at scale $M_X \simeq 10^{16}$ GeV [3].

The MSSM predicts the existence of scalar superpartners to all known leptons and quarks. It is by now quite well known that the lighter scalar top (stop) eigenstate $\tilde{t}_1$ is expected to be lighter than the superpartners of the first two generations of quarks, and might even be the lightest strongly interacting sparticle [4, 5]. There are two reasons for this: Since the top quark is heavy, $m_t \geq 113$ GeV [6], mixing between the superpartners $\tilde{t}_L, \tilde{t}_R$ of left- and right-handed top quarks cannot be neglected, in contrast to the superpartners of light quarks. Furthermore, if we assume all squarks to have the same mass at some very high (GUT, string or Planck) scale, radiative corrections [7] will reduce the masses of $\tilde{t}_L$ and $\tilde{t}_R$ relative to those of the other squarks.

Here we study possible signals for the production of a scalar ($\tilde{t}_1\tilde{t}_1^*$) bound state $\sigma_{\tilde{t}_1}$ within the MSSM. Squark bound state production and decay were first discussed in ref. [8], for negligible mixing between the superpartners of left- and right-handed quarks. In ref. [9] it has been pointed out that this mixing can be very important for the case of $\tilde{t}_1$ bound states; in particular, the branching ratio for $\sigma_{\tilde{t}_1} \to hh$ might be large, where $h$ is the lighter scalar Higgs boson. Very recently it has been claimed [10] that $\sigma_{\tilde{t}_1} \to W^+W^-$ can have a very large branching ratio, which might give rise to interesting signals at hadron supercolliders. We computed all potentially large branching ratios of $\sigma_{\tilde{t}_1}$. We basically agree with the results of ref. [10], but were unable to reproduce those of ref. [11]. The cleanest signal for $\sigma_{\tilde{t}_1}$ production at hadron colliders arises from its $2\gamma$-decay, leaving rise to an interesting signal at hadron supercolliders. The starting point of our discussion is the stop mass matrix. Following the convention of ref. [1] we write it as (in the basis $\tilde{t}_L, \tilde{t}_R$):

$$M^2_\tilde{t} = \begin{pmatrix} m_{\tilde{t}_L}^2 + m_{\tilde{t}_R}^2 + 0.35D_Z & -m_t(A_t + \mu \cot\beta) \\ -m_t(A_t + \mu \cot\beta) & m_{\tilde{t}_R}^2 + m_{\tilde{t}_L}^2 + 0.16D_Z \end{pmatrix}. \quad (1)$$

Here, $D_Z = M_Z^2 \cos 2\beta$ with $\tan \beta = \langle H^0 \rangle / \langle H^0 \rangle$ as usual [2], $\mu$ is the supersymmetric Higgs(ino) mass parameter, $A_t$ a trilinear soft supersymmetry breaking parameter, and $m_{\tilde{t}_L,R}^2$ the non-supersymmetric contributions to the squared masses of the $\tilde{t}_L, \tilde{t}_R$ current states. The diagonalization of $M^2_\tilde{t}$ is straightforward. While the gauge interactions of $\tilde{t}_1$ only depend on the stop mixing angle $\theta_t$, the couplings of stop squarks to Higgs bosons depend on all parameters entering eq. (1); all these quantities therefore have to be specified before $\sigma_{\tilde{t}_1}$ branching ratios can be computed.

Both single stop decay and $(\tilde{t}_1\tilde{t}_1^*)$ annihilation contribute to $\sigma_{\tilde{t}_1}$ decays. In the first case either $\tilde{t}_1$ or $\tilde{t}_1^*$ decays, leaving the other squark behind. We assume that the gluino is too heavy to be produced in these decays. In general we then have to consider the following decay modes:

$$\tilde{t}_1 \to b\tilde{W}_i, \ t\tilde{Z}_j, \ c\tilde{Z}_j, \ i = 1, 2, \ j = 1, \ldots, 4, \quad (2)$$
where \( \tilde{W}_i (\tilde{Z}_j) \) denotes a generic chargino (neutralino) state. The first two decays occur at tree level and with full gauge or top Yukawa strength. It has been shown in ref.\(^[3]\) that \( c\tilde{Z}_j \) is the dominant \( \tilde{t}_1 \) decay mode if these first two modes are kinematically forbidden. However, this last decay only occurs at 1–loop level and necessitates flavor mixing; it is therefore suppressed relative to the tree–level decays by a factor \(|\epsilon|^2 \sim 10^{-7}\) \(^[3]\). We will see below that the \( c\tilde{Z}_j \) mode can therefore be neglected in the discussion of \( \sigma_{\tilde{t}_1} \) decays. The widths of (2) can be computed using the couplings of refs.\(^[2]\); the decay width of \( \sigma_{\tilde{t}_1} \) is twice that of \( \tilde{t}_1 \).

In annihilation decays \( \tilde{t}_1 \) and \( \tilde{t}_1^* \) annihilate into a flavor and color singlet final state; this kind of decay is by far dominant for the familiar \((c\bar{c})\) and \((b\bar{b})\) bound states (quarkonia). We calculated the widths for the following modes:

\[
\sigma_{\tilde{t}_1} \rightarrow gg, W^+W^-, ZZ, Z\gamma, \gamma\gamma, hh, b\bar{b}, \bar{t}t, \tilde{Z}_i\tilde{Z}_j, i,j = 1, \ldots, 4. \tag{3}
\]

Since \( \sigma_{\tilde{t}_1} \) is a scalar \((s–wave)\) state, we only need \(^[12]\) the \( \tilde{t}_1 \) velocity \( v \rightarrow 0 \) limit of the \( \tilde{t}_1\tilde{t}_1^* \) annihilation amplitudes leading to the final states of eq.(3). In this limit, the \( gg, \gamma\gamma \) and \( Z\gamma \) final states are produced via \( t–channel \) \( \tilde{t}_1 \) exchange as well as via 4–point “butterfly” interactions; \( W^+W^- \) is accessible via \( b_L \) exchange\(^[4]\), a 4–point interaction as well as scalar Higgs exchange in the \( s–channel \), while the \( ZZ \) and \( hh \) final states come from \( \tilde{t}_1 \) or \( \tilde{t}_2 \) exchange in the \( t–channel \), a 4–point coupling and Higgs exchange\(^[4]\). The production of \( QQ \) states involves \( s–channel \) scalar Higgs exchange and \( t–channel \) chargino or neutralino exchange; note that the corresponding matrix elements are proportional to the final state quark masses, so that the \( b\bar{b} \) width is very small unless \( \tan\beta \gg 1 \). Finally, neutralino pairs can be produced via \( s–channel \) Higgs exchange or \( t–channel \) top exchange.

In order to compute the decay widths for the processes (3) we have to know the “wave function at the origin” \(|\psi(0)|^2\), see ref.\(^[12]\). Recently the \( m_{\tilde{t}_1} \) dependence of this quantity has been parametrized in ref.\(^[13]\), for 4 different values of the QCD scale parameter \( \Lambda \), using a potential that reproduces the known quarkonium spectrum well; we use their fit for \( \Lambda = 0.2 \) GeV.

In fig.1 we show examples for the branching ratios of the more important processes of eq.(3) for relatively small \( m_{L,R} = 200 \) GeV. In addition to the parameters appearing in eq.(1) we have to specify the SU(2) gaugino mass \( M_2 \) (we assume the usual unification condition \( M_1 = 5/3 \tan^2 \theta_W M_2 \) and the mass \( m_P \) of the pseudoscalar Higgs boson. This then determines all relevant masses and couplings. We have included leading radiative corrections to the scalar Higgs sector involving top–stop loops\(^[14]\).

In this figure we have assumed \( M_2 = 100 \) GeV leading to a mass of about 110 GeV for the lighter chargino. For \( m_{\tilde{t}_1} > 115 \) GeV the single stop decay into \( b + \tilde{W}_1 \) (not shown) opens up and quickly dominates the total decay width. Indeed, in this region the total width of \( \sigma_{\tilde{t}_1} \) is comparable to its binding energy. Our calculations are no longer reliable in this case, since we assume that formation and decay of \( \sigma_{\tilde{t}_1} \) occur at very different time scales so that they can be treated independently; one has to use methods developed previously \(^[15]\) for \( (t\bar{t}) \) bound states instead. However, we can conclude from fig. 1 that if the tree–level single stop

\(^*\)We ignore mixing in the \( b \) sector.

\(^{+}\)\( \tilde{t}_2 \) exchange has not been included in ref.\(^[3]\), where the \( hh \) final state has been studied; this contribution is small compared to the \( \tilde{t}_1 \) exchange term for parameters leading to a sizable \( hh \) branching ratio.
decays are allowed the branching ratios for final states that might be detectable at hadron colliders (see below) are very small, less than $10^{-4}$.

In fig. 2 we have therefore varied $M_2$ along with $m_{t_1}$, so that the decays $\tilde{t}_1 \to b\tilde{\nu}_i$, $t\tilde{Z}_j$ remain closed for $m_{t_1} \leq |\mu|$. We have also chosen larger values for $m_{t_{L,R}}$ with $m_{t_1} > m_{t_R}$ as predicted by minimal supergravity models. We see that now the $Br(\sigma_{t_1} \to hh)$ shoots up very rapidly once this decay becomes kinematically allowed. The reason is that the $h\tilde{t}_1\tilde{t}_1$ coupling contains a term which, in the limit $m_{t_1}^2 \gg M_Z^2$, is exactly proportional to the $LR$ element of $\mathcal{M}_t$.

Obviously this element has to be large if $m_{t_1}$ is to be much smaller than $m_{t_{L,R}}$. As a result, the amplitude for $\tilde{t}_1\tilde{t}_1 \to hh$ is proportional to $(m_{t_{L,R}}^2/M_W m_{t_1})^2$ if $m_{t_{L,R}}^2 \gg m_{t_1}^2 \gg m_h^2/2$. The corresponding branching ratio therefore decreases with increasing $m_{t_1}$, in spite of the increasing phase space.

The branching ratios for $W^+W^-$ and $ZZ$ also increase quickly just beyond threshold. However, they do not reach the level of the $hh$ branching ratio; their amplitudes are at best $\propto (m_{t_{L,R}}/M_W)^2$, if $m_{t_{L,R}}^2 \gg m_{t_1}^2 \gg M_Z^2$. This can be understood from the equivalence theorem, which states that amplitudes involving longitudinal gauge bosons are equal to corresponding ones involving pseudoscalar Goldstone bosons $G$, if the energy of the process is $\gg m_W$. There is no diagonal $G\tilde{t}_1\tilde{t}_1$ coupling; a $G\tilde{t}_1\tilde{t}_2$ coupling with strength similar to the $h\tilde{t}_1\tilde{t}_1$ coupling does exist, but it only affects $\sigma_{t_1}$ decays via diagrams involving a heavy stop propagator. The amplitude for $\sigma_{t_1} \to Z_L Z_L$ is therefore suppressed by a factor $(m_{t_1}/m_{t_{L,R}})^2$ compared to the $hh$ amplitude. Similar arguments apply for the $W_L^+W_L^-$ amplitude. Transverse $W$ and $Z$ bosons are at best produced with ordinary (weak) gauge strength, and their couplings can even be suppressed by $\tilde{t}$ mixing. Unlike ref. we therefore never find the width into $WW$ to exceed the one into gluons. However, the authors of ref. neglected $\tilde{t}$ mixing, and assumed that $m_{t_L}$ can be varied independently of the mass of the left–handed sbottom $\tilde{b}_L$. (The same assumptions were made in ref.) Since $\tilde{t}_1$ and $\tilde{b}_L$ reside in the same $SU(2)$ doublet, this introduces a new source of explicit gauge symmetry breaking, which renders the theory nonrenormalizable.

The structure of the curves of fig. 2 around $m_{t_1}=250$ GeV occurs because for the given choice of parameters the $s$–channel heavy Higgs ($H$) exchange diagrams become resonant here ($m_H \simeq 500$ GeV), greatly enhancing the matrix elements for $tt$, $bb$, $hh$ and $\tilde{Z}\tilde{Z}_j$ final states. The enhancement of the $W^+W^-$ and $ZZ$ channels is much weaker, since the $HVV$ couplings ($V = W, Z$) are small for $m_H^2 \gg M_Z^2$. If $2m_{t_1}$ is very close to $m_H$ our treatment again breaks down; in this case the $(\tilde{t}_1\tilde{t}_1^*)$ bound states mix with $H$. The curves of figs. 1,2 also exhibit numerous minima resulting from destructive interference between different contributions to the matrix elements. Note that (at least in the limit $v \to 0$) usually only a single partial wave is accessible in $\sigma_{t_1}$ decays; if in addition only a single combination of final state helicities can be produced, destructive interference can lead to a vanishing total amplitude even far above threshold.

Clearly one could in principle learn a lot about the MSSM parameters by studying $\sigma_{t_1}$ branching ratios. In practice, however, even discovery of $\sigma_{t_1}$ may not be trivial. We focus here on $pp$ supercolliders. The total cross section for $\sigma_{t_1}$ production at the SSC ($\sqrt{s} = 40$ TeV) is shown by the solid line in fig. 3. Hadronic final states ($gg$, $bb, t\bar{t}$) will be useless for the discovery of $\sigma_{t_1}$ at hadron colliders, due to the enormous backgrounds. In ref.
the use of the $W^+W^-$ final state was advocated. However, we have seen above that $SU(2)$ gauge invariance implies a rather small rate for this final state; besides, it is not clear to us how the $W^+W^-$ invariant mass will be measured, since both $W$ bosons will have to decay leptonically in order to suppress QCD backgrounds. The $ZZ$ final state is very clean if both $Z$ bosons decay leptonically, but then the rate will be very small even at the SSC ($< 5$ events/year).

In ref.\[9\] the use of $\sigma_\tilde{t}_1 \to hh \to \tau^+\tau^-\tau^+\tau^-$ has been proposed. Since $Br(h \to \tau^+\tau^-) \simeq 8\%$ the 4 $\tau$ final state is also relatively rare ($Br < 6.4 \cdot 10^{-3}$). However, it is not clear how this final state is to be identified experimentally in a hadronic environment. A pair of isolated like–sign leptons would be a rather clean tag; assuming a 30% detection efficiency per lepton one could get more than 50 events/year for $m_\tilde{t}_1 \leq 100$ GeV if the $hh$ final state dominates. The problem is that the presence of at least 4 neutrinos in the final state makes it impossible to reconstruct the $\sigma_\tilde{t}_1$ mass. Given that there are sizable backgrounds (e.g., $\sigma(pp \to ZZX \to \tau^+\tau^-\tau^+\tau^-X) \simeq 3 \cdot 10^{-2}$ pb, leading to $\sim 300$ events/year) the identification of the 4 $\tau$ signal might be problematic.

Probably the most promising signal results from the decay $\sigma_\tilde{t}_1 \to \gamma\gamma$. Searches for intermediate mass Higgs bosons will presumably require good electromagnetic energy resolution for SSC detectors, so the reconstruction of $m(\sigma_\tilde{t}_1)$ is in principle straightforward. The signal would then be a bump in the $\gamma\gamma$ spectrum on top of the smooth background. In order to suppress QCD backgrounds (misidentified $\pi^0$'s) as well as backgrounds from photon bremsstrahlung off quarks one usually requires the signal photons to be isolated, i.e. to have little hadronic activity in a cone around the photon. In addition, the physics background from $q\bar{q}$ annihilation and $gg$ fusion is strongly peaked in forward and backward directions, due to propagator effects; it is therefore advantageous to demand the centre–of–mass scattering angle $\theta^*$ of the two photons to be large \[18\]. We impose the following cuts:

$$|\cos\theta^*| \leq 0.5; \quad (4a)$$

$$|y_\gamma| \leq 1.74. \quad (4b)$$

The cut on the rapidity $y_\gamma$ guarantees that the photons are well isolated from the beam pipes (lab scattering angle $\theta \geq 20^\circ$). It also helps to further suppress the $q\bar{q}$ background, which comes from an asymmetric initial state (valence quark on sea antiquark) and hence often undergoes a strong boost.

The dashed curves in fig. 3 show the two–photon signal for the two sets of parameters shown in figs. 1 and 2, respectively, after the cuts of eqs.(4) have been applied. Clearly there is a sizable number of signal events for almost all combinations of parameters, as long as tree–level 2–body decays of $\tilde{t}_1$ are suppressed. In order to decide whether this signal can also be seen on top of the irreducible (physics) background we have defined the minimal cross section leading to a significant signal after one year of SSC running (integrated luminosity = $10^4$ pb$^{-1}$). Following ref.\[19\] we have defined the signal to be significant if the 99% c.l. lower limit of signal + background (after cuts) is at least as large as the 99% upper limit of the background alone. The background has been integrated over an invariant mass bin of width $0.02 \cdot M_{\gamma\gamma}$, which requires good energy resolution. The resulting minimal detectable signal is shown by the dotted line in fig. 3. This line scales up with the square–root of the bin width over which the background has to be integrated, and scales down with the
square–root of the integrated luminosity. We conclude from this figure that \( \sigma_{\tilde{t}_1} \) should be quite easily detectable at the SSC if tree–level \( \tilde{t}_1 \) decays are suppressed, the \( hh \) channel is not greatly enhanced by very large off–diagonal elements of the stop mass matrix, and if one is not too close to the \( H \) pole. The long–dashed curve shows that even in the presence of a sizable branching ratio into the \( hh \) final state the signal might still be visible after several years of operation, provided the energy resolution of the calorimeters is not much worse than we assumed. We mention in passing that the superpartners of light quarks do usually have unsuppressed 2–body decays; we do therefore not expect their bound states to be detectable at the SSC.

We computed both signal and background in leading order in QCD, since no higher order calculation of the signal rate is available yet. Such calculations do exist for the background \[20\]. The result is very similar to the leading order prediction if, in addition to the cuts \( [4] \), one vetoes against the presence of high–\( p_T \) jets.

In summary, we have computed all potentially large branching ratios of a scalar stoponium bound state \( \sigma_{\tilde{t}_1} \), fully taking into account effects due to \( \tilde{t} \) mixing. The dominant decay modes are \( gg, hh \) or, if \( m(\sigma_{\tilde{t}_1}) \approx m_H, \tilde{t}\tilde{t} \). The process \( pp \rightarrow \sigma_{\tilde{t}_1}X \rightarrow \gamma\gamma X \) should be observable at the SSC, provided that \( \tilde{t}_1 \) has no unsuppressed tree–level 2–body decays.

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Figure Captions

Fig.1 Branching ratios for annihilation decays of $\sigma_i$ listed in eq. (3). The range of $m_{\tilde{t}_1}$ values shown results from varying $A_t$ between $-310$ and $-70$ GeV. The values of the other parameters are: $m_{\tilde{t}_L} = m_{\tilde{t}_R} = 200$ GeV, $m_t = -\mu = 150$ GeV, $M_2 = 100$ GeV, $m_P = 500$ GeV and $\tan \beta = 2$. The branching ratios for the $b\bar{b}$ and $t\bar{t}$ final states (not shown) are always below $10^{-3}$.

Fig.2 Branching ratios for annihilation decays of $\sigma_i$ listed in eq. (3). The range of $m_{\tilde{t}_1}$ values shown results from varying $A_t$ between $440$ and $1080$ GeV. We have increased the $SU(2)$ gaugino mass $M_2$ along with $m_{\tilde{t}_1}$ so that the tree–level single stop decays of eq. (2) remain kinematically forbidden ($M_2 = 1.5 m_{\tilde{t}_1}$). The values of the other parameters are: $m_t = 150$ GeV, $m_{\tilde{t}_L} = 400$ GeV, $m_{\tilde{t}_R} = -\mu = 300$ GeV, $m_P = 500$ GeV, and $\tan \beta = 2$. The branching ratios for the $b\bar{b}$ and $t\bar{t}$ final states (not shown) are again small, except for the vicinity of the $H$ pole ($m_{\tilde{t}_1} \simeq 250$ GeV), where the $t\bar{t}$ final state dominates.

Fig.3 Cross section for $\sigma_i$ production at the SSC. The solid line shows the total cross section multiplied with 0.01, and the dashed curves the $\gamma\gamma$ signal cross section after cuts, for the two scenarios of figs. 1 and 2. The dotted curve shows the minimal cross section giving a significant signal after one year of nominal SSC operations, as defined in the text.
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