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Charged Higgs production from SUSY particle cascade decays at the LHC

Aseshkrishna DATTA, Abdelhak DJOUADI,
Monoranjan GUCHAIT and Yann MAMBRINI

[1] Laboratoire de Physique Mathématique et Théorique, UMR5825–CNRS,
Université de Montpellier II, F–34095 Montpellier Cedex 5, France.

Abstract

We analyze the cascade decays of the scalar quarks and gluinos of the Minimal Supersymmetric extension of the Standard Model, which are abundantly produced at the Large Hadron Collider, into heavier charginos and neutralinos which then decay into the lighter ones and charged Higgs particles, and show that they can have substantial branching fractions. The production rates of these Higgs bosons can be much larger than those from the direct production mechanisms, in particular for intermediate values of the parameter $\tan\beta$, and could therefore allow for the detection of these particles. We also discuss charged Higgs boson production from direct two–body top and bottom squark decays as well as from two– and three–body gluino decays.

1. Introduction

The most distinctive signature of an extended Higgs sector, compared to the Standard Model (SM) where only one scalar doublet is needed to break the electroweak symmetry leading to a single neutral Higgs particle, is the discovery of charged Higgs bosons. For instance, in the Minimal Supersymmetric (SUSY) extension of the Standard Model (MSSM) [1], two Higgs doublets are present, leading to the existence of a quintet of scalar particles: two CP-even neutral Higgs bosons h and H, a CP–odd neutral Higgs boson A and two charged Higgs particles $H^\pm$ [2]. SUSY constraints on the Higgs spectrum impose that the charged Higgs boson mass is related to the pseudoscalar Higgs mass, $M_{H^\pm}^2 = M_A^2 + M_W^2$. With the present experimental limit on the A boson mass, $M_A \gtrsim 93.5$ GeV [3], this leads to a bound $M_{H^\pm} \gtrsim 120$ GeV. In fact, in the popular minimal Supergravity models (mSUGRA) with universal boundary conditions at the Grand Unification scale and where the electroweak symmetry breaking is induced radiatively [4], the $H^\pm$ boson, as well as the H and A neutral Higgs particles, tend to be rather heavy, with masses of the order of a few hundred GeV [5]. The states are therefore kinematically accessible only at the LHC [6], at future e^+e^- linear colliders [7] or muon colliders [8].

The discovery of $H^\pm$ bosons at the LHC through the standard processes is rather difficult, if not impossible in some areas of the MSSM parameter space [9]. This is mainly due to the fact that the production rates are controlled by the charged Higgs boson Yukawa couplings to up– and down–type fermions. Using the notation of the first generation, the latter are given by [2]:

$$\frac{gV_{ij}}{\sqrt{2}M_W} H^+ \left[\cot\beta \, m_u \, \bar{u}_i d_{jL} + \tan\beta \, m_d \, \bar{u}_i d_{jR}\right], \tag{1}$$

where V_{ij} is the CKM matrix with $u_i = u, c, t$, $d_i = d, s, b$ and $\tan\beta = v_2/v_1$ is the ratio of the vacuum expectation values of the two Higgs doublets needed to break the electroweak symmetry in the MSSM. For values $\tan\beta > 1$, as is the case in the MSSM, the couplings to down–type (up–type) fermions are enhanced (suppressed). Only the couplings to the top and bottom isodoublet quarks are therefore important, in particular for small ~ 1 and large $\sim m_t/m_b$ values[1].

A light charged Higgs particle can be searched for at the Tevatron in top decays through the process $p\bar{p} \to t\bar{t}$, with at least one of the top quarks decaying via $t \to H^+ b$, leading to a surplus of τ leptons due to the $H^\pm \to \tau^\pm \nu$ decay, an apparent breakdown of τ–μ–e universality. For small and large values of $\tan\beta$, the branching ratio BR$(t \to H^+ b)$ is large and would allow for the detection of the signal at the Tevatron [10]. The situation for intermediate values of $\tan\beta$, where the $H^- t\bar{b}$ coupling is small, leading to a tiny $t \to H^+ b$ branching ratio, is rather difficult and should await for the LHC. Indeed, detailed analyses of the ATLAS and CMS collaborations have shown that the entire range of $\tan\beta$ values should be covered for $M_{H^\pm} \lesssim m_t$ using this process [6].

[1]Interestingly, these two regions of $\tan\beta$ are favored by Yukawa coupling unification. However, the experimental bound on the lightest h boson at LEP2, $M_h \gtrsim 113.5$ GeV [3] in the decoupling regime where the $H^\pm$ bosons are heavy, rules out the low $\tan\beta$ scenario.

For charged Higgs bosons with masses $M_{H^\pm} > m_t$, the two production mechanisms which potentially have sizeable cross sections at the LHC are [11, 12]:

$$
\begin{aligned}
pp \to &\quad gb(g\bar{b}) \to tH^- \ (\bar{t}H^+) \\
pp \to &\quad gg/q\bar{q} \to tH^-\bar{b} + \bar{t}H^+ b
\end{aligned}
\tag{2}
$$

The signal cross section from the $2 \to 2$ mechanism $gb \to tH^-$, where the b quark is obtained from the proton, is 2–3 times larger than the $2 \to 3$ process $gg/q\bar{q} \to t\bar{b}H^-$, where the H^- boson is radiated from a heavy quark line. When the decays $H^+ \to t\bar{b}$ and $t \to Wb$ take place, the first process gives rise to 3 b–quarks in the final state while the second one gives 4 b–quarks; both processes contribute to the inclusive production where at most 3 final b–quarks are tagged[2]. However, the cross sections are rather small: even for the extreme values $\tan\beta = 2$ and 40, they hardly reach the level of a picobarn for a charged Higgs boson mass $M_{H^\pm} = 200$ GeV. For intermediate values of $\tan\beta$ and/or larger $H^\pm$ masses, the cross sections are too small for these processes to be useful. For instance, for the value $\tan\beta = 10$, the cross section is below the level of a few femtobarn for $M_{H^\pm} \gtrsim 250$ GeV.

Other mechanisms for $H^\pm$ production at hadron colliders are the Drell–Yan type process for pair production through γ and Z boson exchange, $q\bar{q} \to H^+H^-$ [13], the gluon–gluon fusion process for pair production, $gg \to H^+H^-$ [14], and the associated production process with W bosons in gg fusion and $q\bar{q}$ annihilation, $q\bar{q}, gg \to H^\pm W^\mp$ [15]. However, the production rates are rather small at the LHC: for the quark–antiquark processes because of the low quark luminosities at high energies and for the gluon–gluon fusion processes because they are induced by loops of heavy quarks (and squarks) and are thus suppressed by additional coupling factors. The cross sections are at the femtobarn level for large enough charged Higgs boson masses, and are therefore too low to be easily useful in the complicated and hostile environment of hadron colliders.

In this paper, we show that there is a potentially large source of the $H^\pm$ bosons of the MSSM at the LHC: the cascade decays of squarks and gluinos, which are abundantly produced in pp collisions, thanks to their strong interactions[3]. Squarks and gluinos can decay into the heavy charginos and neutralinos, $\chi_2^\pm, \chi_3^0$ and χ_4^0, and if enough phase space is available, the latter could decay into the lighter charginos/neutralinos, $\chi_1^\pm, \chi_1^0$ and χ_2^0, and charged Higgs bosons[4]:

$$
\begin{aligned}
pp \to \tilde{g}\tilde{g}, \tilde{q}\tilde{q}, \tilde{q}\tilde{g} \ &\to \ \chi_2^\pm, \chi_3^0, \chi_4^0 + X \\
&\to \ \chi_1^\pm, \chi_2^0, \chi_1^0 + H^\pm \ + X
\end{aligned}
\tag{3}
$$

[2]However, the two processes have to be properly combined to avoid double counting of the contribution where a gluon gives rise to a $b\bar{b}$ pair that is collinear to the initial proton. The cross section of the inclusive process in this case is mid–way between those of the two production mechanisms eqs. (2) [9].

[3]These decays have been discussed in the past [16] for charged Higgs bosons with masses below ~ 150 GeV and which can thus also be produced in top quark decays.

[4]We will not specifically consider here the similar production of neutral Higgs particles, $\Phi = h, H, A$, since the later can be produced abundantly in standard processes such as as gluon–gluon fusion, $gg \to \Phi$ and associated production with heavy quarks $gg, q\bar{q} \to b\bar{b}\Phi, t\bar{t}\Phi$ [6].

These processes are similar to the ones with cascade decays of strongly interacting SUSY particles into the next–to–lightest neutralino χ_2^0 which then decays into the lightest h boson and the lightest neutralino [which is expected to be the lightest SUSY particle in the MSSM], a process which has been discussed in the literature; see for instance Ref. [17].

Charged Higgs bosons could also be searched for, if kinematically possible, in the direct decays of heavy third generation squarks into their lighter partners,

$$\tilde{Q} \to \tilde{Q}' H^{\pm} \;\; \text{with} \;\; \tilde{Q}, \tilde{Q}' = \tilde{t}, \tilde{b} \tag{4}$$

or in direct gluino three–body decays into heavy quarks, their partners squarks and $H^{\pm}$ bosons,

$$\tilde{g} \to Q' \tilde{Q} H^{\pm} \;\; \text{with} \;\; \tilde{Q} = \tilde{t}, \tilde{b} \tag{5}$$

The remainder of the paper is organised as follows. In the next section, we discuss the production cross sections of squarks and gluinos at the LHC as well as their main decay modes. In section 3, we analyse the decays of the heavy charginos and neutralinos into the lighter charginos/neutralinos and Higgs or gauge bosons and estimate the production cross sections for $H^{\pm}$ bosons at the LHC in some specific scenarii. In section 4, we discuss the direct production of $H^{\pm}$ bosons in the two–body decays of top and bottom squarks and three–body decays of gluinos. A short conclusion is given in section 5.

2. Production and decay modes of Squarks and Gluinos

2.1 Production cross sections

In proton–proton collisions, gluino pairs are produced through $q\bar{q}$ annihilation and gluon–gluon fusion, $q\bar{q}, gg \to \tilde{g}\tilde{g}$. Squark pairs can be produced through t–channel exchange, $qq \to \tilde{q}\tilde{q}$, while squark–antisquark pairs are produced in both $q\bar{q}$ annihilation and gluon–gluon fusion. Finally, mixed squark–gluino production proceeds through s– and t– channel gq annihilation. The production cross sections at the tree–level are given in Ref. [18]. They have however, to be supplemented by the next–to–leading order QCD radiative corrections, which stabilize the theoretical predictions. The K–factors can be rather large, ranging from $K = \sigma_{\rm NLO}/\sigma_{\rm LO} \sim 1$ to 2 [19], depending on the ratio of the squark to gluino masses. Taking for the renormalisation and factorization scale the average mass m of the two produced sparticles results in a conservative estimate of the total production cross section. We will therefore not include the K–factors in our analysis and set the scale at which the cross section is evaluated to be the average mass of the final particles.

The cross sections for pair production and associated production of squarks and gluinos are shown in Fig. 1 as functions of the squark and gluino masses. The CTEQ3L [20] parameterization of the parton densities has been used. In Fig. 1a, we choose $m_{\tilde{g}} = m_{\tilde{q}}$ while in Figs. 1b and 1c, we take respectively, $m_{\tilde{q}} = 1.2 m_{\tilde{g}}$ and $m_{\tilde{g}} = 1.2 m_{\tilde{q}}$. We assume the left– and right–handed scalar partners of the light quarks, including the b–quark, to be degenerate in mass and we sum the individual cross sections. We show separately in Fig. 1a, the total cross section for the pair production of the lightest top squark, $\sigma_{q\bar{q}}, gg \to \tilde{t}_1 \tilde{t}_1^*)$, which depends only on $m_{\tilde{t}_1}$ at the tree–level.

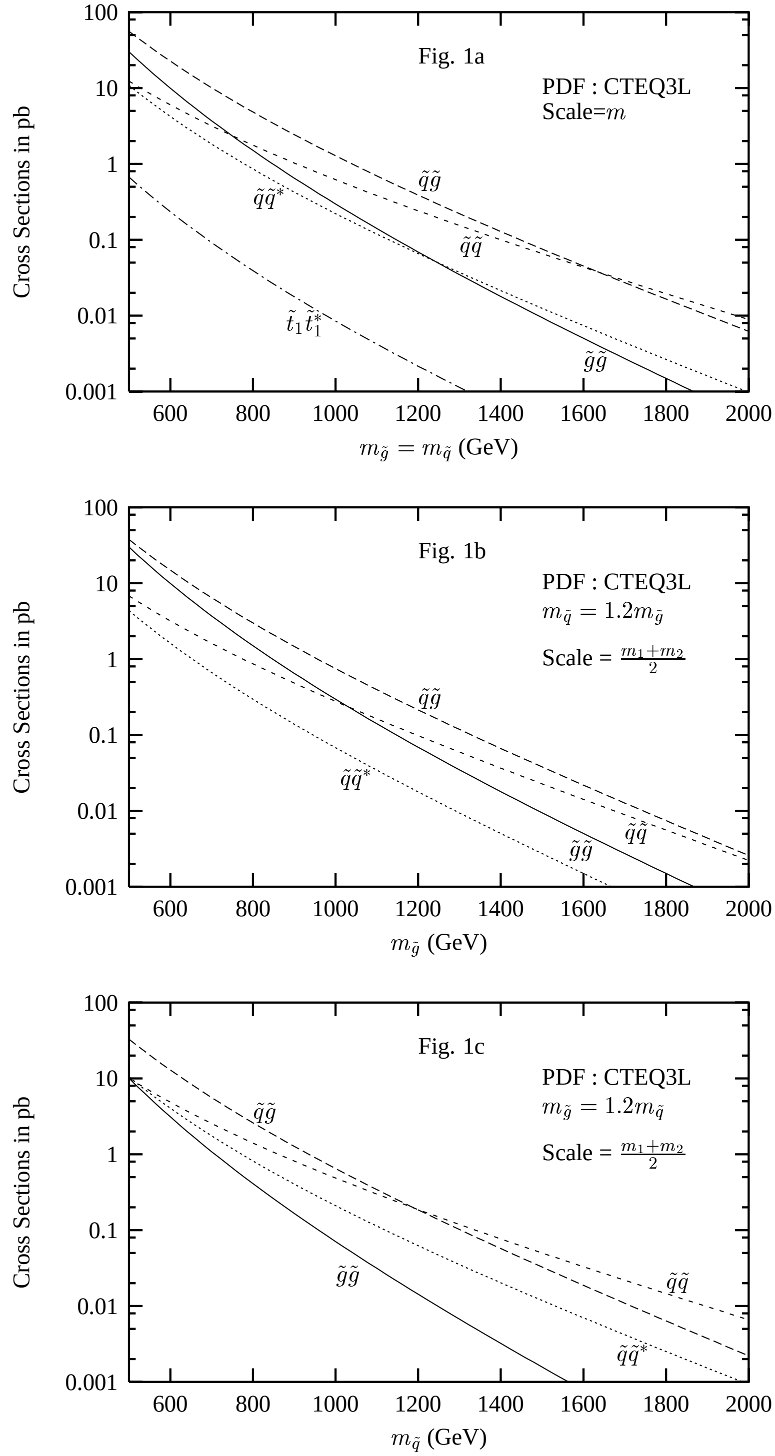

Figure 1: Pair and associated production cross sections at the lowest order for the strongly interacting SUSY particles at LHC with $\sqrt{s} = 14$ TeV as a function of the final state masses for representative choices of squark and gluino masses. $\sigma_{\tilde{q}\tilde{q}}$ and $\sigma_{\tilde{q}\tilde{g}}$ include the contributions from $\sigma_{\tilde{q}\tilde{q}^*}$ and $\sigma_{\tilde{g}\tilde{g}}$ respectively; $\sigma_{\tilde{t}_1\tilde{t}_1^*}$ is presented only in Fig. 1a.

We see that the largest cross section is due to squark and gluino associated production, in which squarks and anti–squarks of the five light flavors have been added. It ranges from $\sigma_{\tilde{q}\tilde{g}} \sim 50$ pb for a gluino mass $m_{\tilde{g}} \sim 500$ GeV to $\mathcal{O}$ (1 pb) for $m_{\tilde{g}} \sim 1$ TeV. In the low mass range, it is followed by the cross section for gluino (squark) pair production if the gluino is lighter (heavier) than squarks. The total cross sections decrease quickly with increasing final state masses. In the case of top squarks, the total production cross section is an order of magnitude smaller than the one for squark–antisquark pairs, $\sigma_{\tilde{t}_1\tilde{t}_1^*} \sim \frac{1}{12}\sigma_{\tilde{q}\tilde{q}^*}$ if all squarks, including $\tilde{t}_1$, have the same mass; the cross section for the production of $\tilde{t}_2$, which is expected to be heavier than $\tilde{t}_1$ because of the generally large mixing in the stop sector, is smaller than $\sigma_{\tilde{t}_1\tilde{t}_1^*}$.

Summing up all cross sections, one obtains $\sigma_{\tilde{q}+\tilde{g}} \sim 110$ (3) pb for $m_{\tilde{g}} \sim m_{\tilde{q}} \sim 500$ (1000) GeV. This means that with the expected integrated luminosity at the LHC, $\int \mathcal{L} \sim 300$ fb^{-1}, a total of $3 \cdot 10^7$ to 10^5 events can be collected by the ATLAS and CMS collaborations in a course of few years. It is therefore tempting to use this very large sample of events to look for particularly interesting decay channels of squarks and gluinos which, even if they have branching ratios below the percent level, would still lead to a rather large number of final state events that could be studied in detail.

2.2 Squark Decays

The main decay modes of squarks, if they are lighter than the gluino, will be into their partner quarks and neutralinos, $\tilde{q}_i \rightarrow q\chi_j^0$ [j=1–4], as well as quarks and charginos, $\tilde{q}_i \rightarrow q'\chi_j^\pm$ [j=1–2]. Taking into account the mass of the final quark, which would be appropriate in the case of top squark decays, the partial decay widths are given at the tree–level by [the QCD corrections to these decay modes have been calculated in Ref. [21]]:

$$
\begin{aligned}
\Gamma(\tilde{q}_i \rightarrow q\chi_j^0) &= \frac{\alpha\lambda^{\frac{1}{2}}(\mu_q^2, \mu_{\chi_j^0}^2)}{4} m_{\tilde{q}_i}\left[({a_{ij}^{\tilde{q}}}^2 + {b_{ij}^{\tilde{q}}}^2)(1-\mu_q^2-\mu_{\chi_j^0}^2) - 4a_{ij}^{\tilde{q}}b_{ij}^{\tilde{q}}\mu_q\mu_{\chi_j^0}\epsilon_{\chi_j}\right] \\
\Gamma(\tilde{q}_i \rightarrow q'\chi_j^\pm) &= \frac{\alpha\lambda^{\frac{1}{2}}(\mu_{q'}^2, \mu_{\chi_j^+}^2)}{4} m_{\tilde{q}_i}\left[({a_{ij}^{\tilde{q}}}^2 + {b_{ij}^{\tilde{q}}}^2)(1-\mu_{q'}^2-\mu_{\chi_j^+}^2) - 4a_{ij}^{\tilde{q}}b_{ij}^{\tilde{q}}\mu_{q'}\mu_{\chi_j^+}\right]
\end{aligned} \tag{6}
$$

where $\lambda(x,y) = 1+x^2+y^2-2(xy+x+y)$ is the usual two–body phase space function with the reduced masses $\mu_X = m_X/m_{\tilde{q}_i}$ and ϵ_{χ_j} is the sign of the eigenvalue of the neutralino χ_j^0. In terms of $s_W^2 = 1-c_W^2 \equiv \sin^2\theta_W$, the squark electric charge, weak isospin and the mixing angle θ_q which turns the left– and right–handed states into the mass eigenstates $\tilde{q}_1 = c_{\theta_q}\tilde{q}_L + s_{\theta_q}\tilde{q}_R$ and $\tilde{q}_2 = -s_{\theta_q}\tilde{q}_L + c_{\theta_q}\tilde{q}_R$, one has for the couplings among neutralinos, quarks and squarks:

$$
\begin{aligned}
\left\{ \begin{array}{c} a_{j1}^{\tilde{q}} \\ a_{j2}^{\tilde{q}} \end{array} \right\} &= -\frac{m_q r_q}{\sqrt{2}M_W s_W} \left\{ \begin{array}{c} s_{\theta_q} \\ c_{\theta_q} \end{array} \right\} - e_{Lj}^q \left\{ \begin{array}{c} c_{\theta_q} \\ -s_{\theta_q} \end{array} \right\} \\
\left\{ \begin{array}{c} b_{j1}^{\tilde{q}} \\ b_{j2}^{\tilde{q}} \end{array} \right\} &= -\frac{m_q r_q}{\sqrt{2}M_W s_W} \left\{ \begin{array}{c} c_{\theta_q} \\ -s_{\theta_q} \end{array} \right\} - e_{Rj}^q \left\{ \begin{array}{c} s_{\theta_q} \\ c_{\theta_q} \end{array} \right\}
\end{aligned} \tag{7}
$$

with $r_u = Z_{j4}/\sin\beta$ and $r_d = Z_{j3}/\cos\beta$ for up and down–type fermions, and

$$e^q_{Lj} = \sqrt{2}\left[e_q\, Z'_{j1} + \left(I^3_q - e_q\, s^2_W\right)\frac{1}{c_W s_W}\, Z'_{j2}\right] \ , \ e^q_{Rj} = -\sqrt{2}\, e_q\left[Z'_{j1} - \frac{s_W}{c_W}\, Z'_{j2}\right] \tag{8}$$

while for the couplings among charginos, fermions and sfermions, $\tilde{q}_i - q' - \chi^+_j$, one has for up–type and down–type sfermions:

$$\begin{aligned}
\left\{\begin{array}{c} a^{\tilde{u}}_{j1} \\ a^{\tilde{u}}_{j2}\end{array}\right\} &= \frac{V_{j1}}{s_W}\left\{\begin{array}{c} -c_{\theta_u} \\ s_{\theta_u}\end{array}\right\} + \frac{m_u\, V_{j2}}{\sqrt{2}\, M_W s_W\, s_\beta}\left\{\begin{array}{c} s_{\theta_u} \\ c_{\theta_u}\end{array}\right\} \\
\left\{\begin{array}{c} b^{\tilde{u}}_{j1} \\ b^{\tilde{u}}_{j2}\end{array}\right\} &= \frac{m_d\, U_{j2}}{\sqrt{2}\, M_W s_W\, c_\beta}\left\{\begin{array}{c} c_{\theta_u} \\ -s_{\theta_u}\end{array}\right\}
\end{aligned} \tag{9}$$

$$\begin{aligned}
\left\{\begin{array}{c} a^{\tilde{d}}_{j1} \\ a^{\tilde{d}}_{j2}\end{array}\right\} &= \frac{U_{j1}}{s_W}\left\{\begin{array}{c} -c_{\theta_d} \\ s_{\theta_d}\end{array}\right\} + \frac{m_d\, U_{j2}}{\sqrt{2}\, M_W s_W\, c_\beta}\left\{\begin{array}{c} s_{\theta_d} \\ c_{\theta_d}\end{array}\right\} \\
\left\{\begin{array}{c} b^{\tilde{d}}_{j1} \\ b^{\tilde{d}}_{j2}\end{array}\right\} &= \frac{m_u\, V_{j2}}{\sqrt{2}\, M_W s_W\, s_\beta}\left\{\begin{array}{c} c_{\theta_d} \\ -s_{\theta_d}\end{array}\right\}
\end{aligned} \tag{10}$$

In these expressions, U, V and Z are the (real) diagonalizing matrices for the chargino and neutralino states [22] with:

$$Z'_{i1} = Z_{i1}c_W + Z_{i2}s_W \ , \ Z'_{i2} = -Z_{i1}s_W + Z_{i2}c_W \ , \ Z'_{i3} = Z_{i3} \ , \ Z'_{i4} = Z_{i4} \tag{11}$$

If squarks are heavier than the gluino, they can also decay into gluino–quark final states, for which the partial decay width is given by:

$$\Gamma(\tilde{q}_i \to q\tilde{g}) \ = \ \frac{2\alpha_s \lambda^{\frac{1}{2}}(\mu^2_q, \mu^2_{\tilde{g}})}{3} m_{\tilde{q}_i}\left[1 - \mu^2_q - \mu^2_{\tilde{g}} - 4a^{\tilde{q}}_{i\tilde{g}} b^{\tilde{q}}_{i\tilde{g}} \mu_q \mu_{\tilde{g}}\right] \tag{12}$$

with the same notation as previously and the squark–quark–gluino coupling are:

$$a^q_{1\tilde{g}} = b^q_{2\tilde{g}} = \sin\theta_q \quad , \quad a^q_{2\tilde{g}} = -b^q_{1\tilde{g}} = \cos\theta_q \tag{13}$$

We will now discuss some scenarii for these decay modes, starting with the case of the scalar partners of light quarks and continuing with the special case of top squarks.

2.2.1 The case of the scalar partners of light quarks

If squarks are heavier than the gluino, they will decay most of the time into quark plus gluino final states. This is essentially due to the fact that these are strong interaction decays, compared to their weak interaction decays into charginos and neutralinos. The large value of the strong coupling constant, $\alpha_s/\alpha \sim 10$, makes the decays into gluinos an order of magnitude larger. In the opposite case, $m_{\tilde{q}} < m_{\tilde{g}}$, the right–handed squarks will decay [for small quark masses] only into quarks and neutralinos, while left–handed squarks decay into both charginos and neutralinos. Two scenarii are possible, depending on the chargino/neutralino textures:

(i) Gaugino–limit: If the lighter chargino and neutralinos are gaugino–like, that is if the higgsino mass parameter is much larger than the wino and bino mass parameters, $|\mu| \gg M_{1,2}$ with M_1 and M_2 related by the GUT constraint $M_2 \sim 2M_1$, the masses are such that $m_{\chi_2^0} \sim m_{\chi_1^\pm} \sim 2m_{\chi_1^0} \sim M_2$ while the heavy chargino and neutralinos have masses, $m_{\chi_3^0} \sim m_{\chi_4^0} \sim m_{\chi_2^\pm} \sim |\mu|$. Squarks will have then the tendency to decay into the lighter ino states not only because of a more favorable phase space, but also because for the partners of the light quarks, the higgsino component of the quark–squark–ino coupling [which is proportional to m_q, see the previous formulae for the couplings] is very small. Therefore squarks will decay dominantly into the lighter charginos and neutralinos.

(ii) Higgsino–limit: If, the lighter chargino and neutralinos are higgsino–like, that is $M_{1,2} \gg |\mu|$, the trend is reversed and one would have the mass hierarchies, $m_{\chi_2^0} \sim m_{\chi_1^\pm} \sim m_{\chi_1^0} \sim |\mu|$ and $2m_{\chi_3^0} \sim m_{\chi_2^\pm} \sim m_{\chi_4^0} \sim M_2$. If allowed by phase space, i.e. for $m_{\tilde{q}} > M_2$, (left–handed) squarks will decay into all possible neutralino and chargino combinations in principle. However, because the higgsinos couple proportionally to the quark masses, the partial decay widths into the lighter chargino and neutralinos are tiny, and squarks will dominantly decay into the heavier chargino and neutralinos, $\tilde{q} \to q'\chi_2^\pm, q\chi_3^0$ and $q\chi_4^0$. [In the mixed region, $M_2 \sim |\mu|$ decays into all charginos and neutralinos are possible if phase space allowed.]

The two scenarii (i) and (ii) are exemplified in the upper and lower panels of Fig. 2, respectively, where we display the "average" squark branching ratios into gluinos, charginos and neutralinos as a function of the squark mass for the value $\tan\beta = 10$. This "average" branching ratio is defined as the branching fraction of the decay of any squark [except for top squarks] into a given neutralino or chargino [or gluino] final state; this means that we sum all possibilities for left– and right–handed squarks as well as for up–type and down–type squarks keeping proper track of flavors and chirality to reach a given final state. [Note that here we neglect the small effect of the b–quark Yukawa coupling which is not enough enhanced for the value of $\tan\beta$ we are using, and treat the $\tilde{b}$ squark in the same footing as the first and second generation squarks]. The gluino mass is fixed to $m_{\tilde{g}} = 3M_2$, which is typically the case in models with unified gaugino masses at the GUT scale which lead to the tree–level relation $m_{\tilde{g}} \sim M_3 \sim 3M_2$ at the low–energy scale.

We see that in scenario (i), where the wino and higgsino mass parameters are set to $M_2 = 150$ GeV and $\mu = 400$ GeV, squarks will decay mainly into the lighter chargino $\chi_1^\pm$ and neutralinos $\chi_{1,2}^0$ in the low mass range, i.e. $m_{\tilde{q}} \lesssim 450$ GeV, with branching ratios of the order of 50% for the charged and neutral decay channels, while the strong interaction $\tilde{q} \to q\tilde{g}$ decay channel becomes by far dominating above the gluino mass threshold, with a branching ratio reaching 90% at large squark masses. In scenario (ii), where the higgsino and wino mass parameters are set to $\mu = 150$ GeV and $M_2 = 400$ GeV, the gluino is much heavier than the squarks, $m_{\tilde{g}} \gtrsim 1.2$ TeV, and squarks will decay almost 100% of the time into the heavier chargino and neutralinos, with a dominance at large squark masses, of the charged decay mode, $\tilde{q} \to q'\chi_2^\pm$, which reaches a branching fraction of 50% due to the usual dominance of the charged currents over the neutral currents.

Thus, there exist situations in which heavy squarks can decay into the heavier chargino and neutralinos with significant rates.

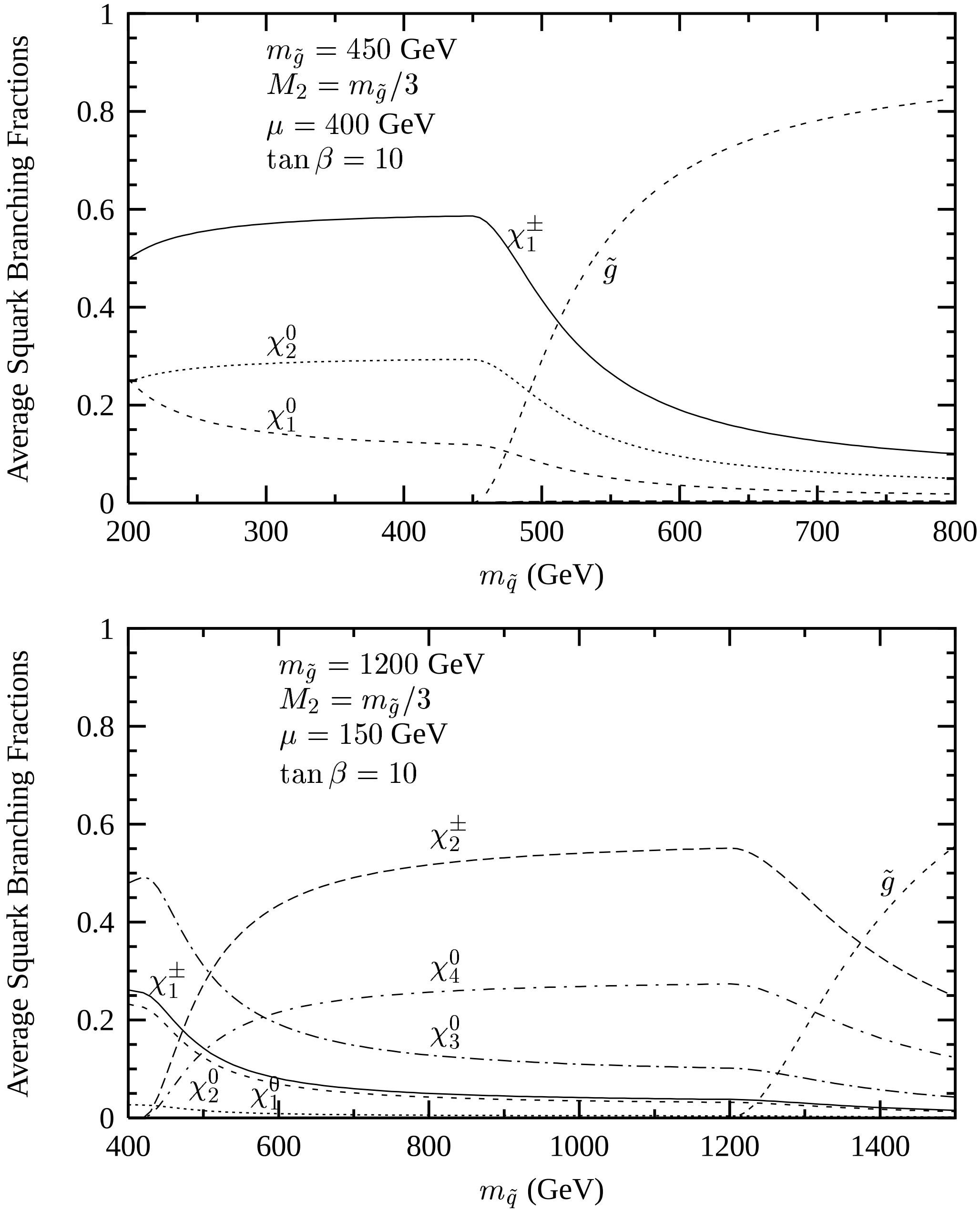

Figure 2: Branching ratios of squarks decaying into gluinos, charginos and neutralinos as a function of their mass for $\tan\beta = 10$ and $m_{\tilde{g}} = 450$ GeV (top) 1.2 TeV (bottom). The gaugino mass is fixed to $M_2 = m_{\tilde{g}}/3$ while the value of the parameter μ is 400 GeV and 150 GeV in scenarii (top) and (bottom) respectively.

2.2.2 The case of the top squarks

The case of top squarks is special: because of the large value of top quark Yukawa coupling, there is a sizeable splitting between the two states $\tilde{t}_2$ and $\tilde{t}_1$, and the latter is in general much lighter than all other squarks. In some cases, gluinos can be lighter than the scalar partners of light squarks, but heavier than $\tilde{t}_1$. In particular, one can have the mass hierarchy $m_{\tilde{q}} \geq m_{\tilde{g}} \geq m_{\tilde{t}_1} + m_t$ and the squarks $\tilde{q}$ will decay almost exclusively into gluinos and quarks as discussed previously, and the gluinos will decay into the only two–body decay mode which is allowed, i.e. $\tilde{g} \to t\tilde{t}_1$. This means that all strongly interacting particles produced at the LHC could decay dominantly into top and stop final states.

The branching ratios for $\tilde{t}_1$ decays into $\chi^+ b$ and $\chi^0 t$ final states are shown in Fig. 3 for a mass of $m_{\tilde{t}_1} = 600$ GeV and $\tan\beta = 10$. In the upper (lower) panel, they are shown as functions of $\mu(M_2)$ for a fixed value of $M_2 = 150$ GeV ($\mu = 400$ GeV). As can be seen, there are large regions where the decay $\tilde{t}_1 \to b\chi_2^+$ is dominating. In particular, one can see that for small values of μ (~ 150 GeV) and large values of M_2 (~ 400 GeV), i.e. in the higgsino region, BR($\tilde{t}_1 \to b\chi_2^+$) can reach the level of 50%. But in contrast to the scalar partners of light quarks, the branching ratio can also be large in the gaugino region, $M_2 = 150$ GeV and $\mu = 400$ GeV, since now the couplings of top/stop states to higgsinos are enhanced by m_t. The decays into neutralinos, $\tilde{t}_1 \to t\chi^0_{3,4}$ are also sizeable in the gaugino region, but the phase space is limited because of the large value of m_t.

2.3 Gluino decays

If gluinos are heavier than squarks, their only relevant decay channel will be into quark plus squark final states, $\tilde{g} \to q\tilde{q}$. The partial decay width, including the quark mass and the squark mixing angle to take into account the possibility of decays into top squarks, using the notation introduced previously, is given by [here, $\mu_X = m_X/m_{\tilde{g}}$]:

$$\Gamma(\tilde{g} \to q\tilde{q}_i) \;=\; \frac{\alpha_s \lambda^{\frac{1}{2}}(\mu_q^2, \mu_{\tilde{q}_i}^2)}{8} m_{\tilde{g}} \Big[1 + \mu_q^2 - \mu_{\tilde{q}_i}^2 + 4 a^{\tilde{q}}_{i\tilde{g}} b^{\tilde{q}}_{i\tilde{g}} \mu_q \mu_{\tilde{g}} \Big] \qquad (14)$$

In the case, where the gluino is lighter than the squarks, it will mainly decay into final states involving a quark–antiquark pair and charginos or neutralinos. The Dalitz density for the decay, taking into account the masses of the final fermions, is given by [23]:

$$\frac{\mathrm{d}\Gamma}{\mathrm{d}x_u \mathrm{d}x_d}(\tilde{g} \to \chi_j u\bar{d}) \;=\; \frac{\alpha\alpha_s}{32\pi} m_{\tilde{g}} \sum_{k,l=1}^{2} \Big[\mathrm{d}\Gamma^u_{kl} + \mathrm{d}\Gamma^d_{kl} + \mathrm{d}\Gamma^{ud}_{kl} \Big] \qquad (15)$$

$$\begin{aligned}
\mathrm{d}\Gamma^{\tilde{d}}_{kl} \;=\;& \frac{1}{(-\mu_u - \mu_{\tilde{d}_k} + \hat{u})(-\mu_u - \mu_{\tilde{d}_l} + \hat{u})} \Big\{ (a^d_{jk} a^d_{jl} + b^d_{jk} b^d_{jl}) [(a^d_{k\tilde{g}} a^d_{l\tilde{g}} + b^d_{k\tilde{g}} b^d_{l\tilde{g}})(-\hat{u}^2 + \hat{u} \\
\times\;& (1 + \mu_\chi + \mu_d + \mu_u) - (\mu_\chi + \mu_d)(1 + \mu_u)) - 2\sqrt{\mu_u}(a^d_{k\tilde{g}} b^d_{l\tilde{g}} + a^d_{l\tilde{g}} b^d_{k\tilde{g}})(\mu_d + \mu_\chi - \hat{u})] \\
+\;& 2(a^d_{jk} b^d_{jl} + a^d_{jl} b^d_{jk}) \sqrt{\mu_\chi \mu_d} [(a^d_{k\tilde{g}} a^d_{l\tilde{g}} + b^d_{k\tilde{g}} b^d_{l\tilde{g}})(\hat{u} - 1 - \mu_u) - 2\sqrt{\mu_u}(a^d_{k\tilde{g}} b^d_{l\tilde{g}} + a^d_{l\tilde{g}} b^d_{k\tilde{g}})] \Big\}
\end{aligned}$$

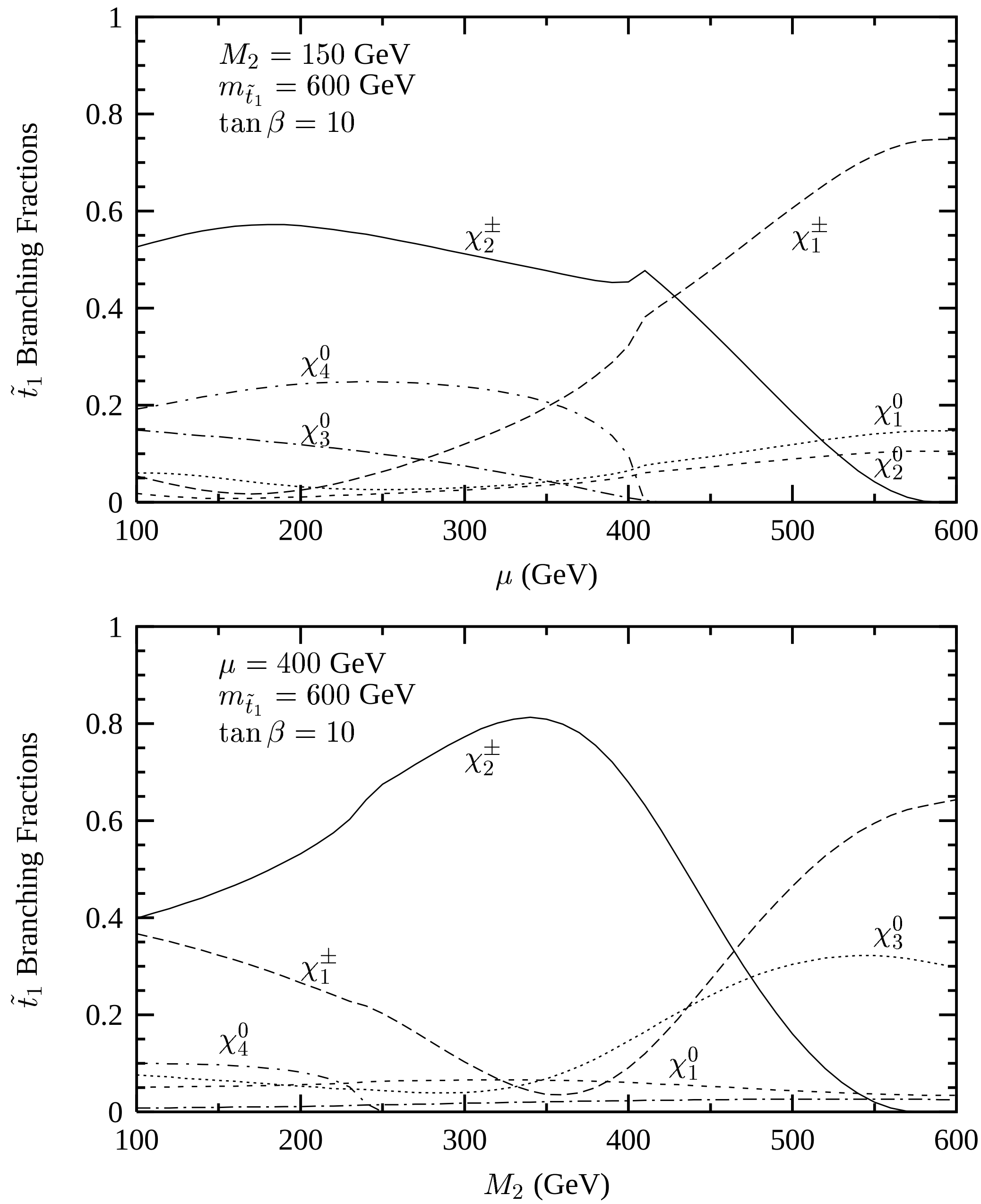

Figure 3: Branching ratios for top squark decaying into charginos and neutralinos as a function of μ with $M_2 = 150$ GeV (top) and M_2 with $\mu = 400$ GeV (bottom). We have fixed the other parameters to $\tan\beta = 10$ and $m_{\tilde{t}_1} = 600$ GeV.

$$\begin{aligned}
\mathrm{d}\Gamma^{\tilde u}_{kl} &= \frac{1}{(-\mu_d-\mu_{\tilde u_k}+\hat t)(-\mu_d-\mu_{\tilde u_l}+\hat t)}\Big\{(a^u_{jk}a^u_{jl}+b^u_{jk}b^u_{jl})[(a^u_{k\tilde g}a^u_{l\tilde g}+b^u_{k\tilde g}b^u_{l\tilde g})(-\hat t^2+\hat t \\
&\times\ (1+\mu_\chi+\mu_u+\mu_d)-(\mu_\chi+\mu_u)(1+\mu_d))-2\sqrt{\mu_d}(a^u_{k\tilde g}b^u_{l\tilde g}+a^u_{l\tilde g}b^u_{k\tilde g})(\mu_u+\mu_\chi-\hat t)] \\
&+\ 2(a^u_{jk}b^u_{jl}+a^u_{jl}b^u_{jk})\sqrt{\mu_\chi\mu_u}[(a^u_{k\tilde g}a^u_{l\tilde g}+b^u_{k\tilde g}b^u_{l\tilde g})(\hat t-1-\mu_d)-2\sqrt{\mu_d}(a^u_{k\tilde g}b^u_{l\tilde g}+a^u_{l\tilde g}b^u_{k\tilde g})]\Big\} \\
\mathrm{d}\Gamma^{\tilde u\tilde d}_{kl} &= \sum_{k,l=1}^{2}\frac{-2}{(-\mu_u-\mu_{\tilde d_l}+\hat u)(-\mu_d-\mu_{\tilde u_k}+\hat t)}\Big\{[a^u_{k\tilde g}a^u_{jk}b^d_{l\tilde g}b^d_{jl}+a^d_{l\tilde g}a^d_{jl}b^u_{k\tilde g}b^u_{jk}]\sqrt{\mu_u\mu_d} \\
&\quad(\hat u+\hat t-\mu_u-\mu_d)+[a^u_{k\tilde g}a^u_{jk}a^d_{l\tilde g}b^d_{jl}+b^u_{k\tilde g}b^u_{jk}b^d_{l\tilde g}a^d_{jl}]\sqrt{\mu_d}(\hat t-\mu_\chi-\mu_u) \\
&\quad+[a^u_{jk}a^d_{l\tilde g}a^d_{jl}b^u_{k\tilde g}+b^u_{jk}b^d_{l\tilde g}b^d_{jl}a^u_{k\tilde g}]\sqrt{\mu_\chi\mu_d}(\hat u-\mu_u-1)-2[a^u_{jk}a^d_{jl}b^u_{k\tilde g}b^d_{l\tilde g}+a^u_{k\tilde g}a^d_{l\tilde g}b^u_{jk}b^d_{jl}] \\
&\quad\sqrt{\mu_\chi}\sqrt{\mu_u\mu_d}+[a^u_{jk}b^u_{k\tilde g}b^d_{l\tilde g}b^d_{jl}+a^u_{k\tilde g}a^d_{l\tilde g}a^d_{jl}b^u_{jk}]\sqrt{\mu_u}(\hat u-\mu_\chi-\mu_d) \\
&\quad+[a^u_{jk}a^d_{l\tilde g}b^u_{k\tilde g}b^d_{jl}+a^u_{k\tilde g}a^d_{jl}b^u_{jk}b^d_{l\tilde g}](\hat u\hat t-\mu_\chi-\mu_u\mu_d)+[a^u_{k\tilde g}a^u_{jk}a^d_{l\tilde g}a^d_{jl}+b^u_{k\tilde g}b^u_{jk}b^d_{l\tilde g}b^d_{jl}]\sqrt{\mu_\chi} \\
&\quad(\hat u+\hat t-\mu_\chi-1)+[a^u_{k\tilde g}a^u_{jk}a^d_{jl}b^d_{l\tilde g}+b^u_{k\tilde g}b^u_{jk}b^d_{jl}a^d_{l\tilde g}]\sqrt{\mu_\chi\mu_u}(\hat t-\mu_d-1)\Big\} \qquad (16)
\end{aligned}$$

where $x_u=2E_u/m_{\tilde g}, x_d=2E_d/m_{\tilde g}$ are the reduced energies of the final quarks, $\mu_X=m_X^2/m_{\tilde g}^2$ the reduced masses and $\hat t=(p_{\tilde g}-p_u)^2/m_{\tilde g}^2=1-x_u+\mu_u$, $\hat u=(p_{\tilde g}-p_d)^2/m_{\tilde g}^2=1-x_d+\mu_d$. The squark-quark couplings to charginos and neutralinos, a^q_{jl} and b^q_{jl} and the couplings to gluinos $a^q_{l\tilde g}$ and $b^q_{l\tilde g}$ have been given previously. The fully integrated partial decay width, in the case of massless final state quarks, can be found in Ref. [24].

The various gluino branching fractions are shown in Fig. 4 as a function of μ for a gluino mass $m_{\tilde g}=800$ GeV and for squark masses fixed to $m_{\tilde q}=1.2m_{\tilde t_1}=1.2m_{\tilde g}$. Two different scenarii are exemplified: in the upper panel of Fig. 4, the approximate GUT relation $M_2=m_{\tilde g}/3$ is adopted while in the lower part of Fig. 4, this GUT relation is relaxed and we set $M_2=150$ GeV, i.e. with gaugino–like $\chi_1^\pm,\chi_{1,2}^0$ for large μ values. [We however, keep the equality of the wino and bino mass parameters at the GUT scale, leading to $M_2\simeq 2M_1$ at low energies]. In both scenarii, the top squark is lighter compared to other squarks [which is generally the case, as discussed previously], $m_{\tilde t_1}=m_{\tilde g}$, and its virtuality in the decay of the gluinos is smaller. Since, in addition, the top squark has large couplings to the higgsino states, the decay modes $\tilde g\to t\bar b\chi_i^-$ and $\tilde g\to t\bar t\chi_i^0$ [and the charge conjugate states] are more important, the former being in general dominant because of the larger phase–space and the stronger charged current couplings.

In the universal scenario with $M_2\sim m_{\tilde g}/3\sim 270$ GeV, the branching ratio for the decay $\tilde g\to t\bar b\chi_2^-$ is at the level of $\sim 30\%$ for small values of $\mu\lesssim 300$ GeV, where $\chi_{1,2}^\pm$ are mixtures of gauginos and higgsinos, while the branching ratios of the decays into top quarks and $\chi_{3,4}^0$ final states are smaller by a factor of 3. The branching ratios decrease with increasing μ since $\chi_{3,4}^0$ and $\chi_2^\pm$ become heavier and the phase space for the decay is reduced. However, even for values of μ around 400 GeV, BR($\tilde g\to t\bar b\chi_2^-$) stays above the level of 10%. In the non–universal scenario with $M_2=150$ GeV, $\chi_2^\pm$ are almost higgsino–like for relatively large μ values, $\mu\gtrsim 300$ GeV. BR($\tilde g\to bt\chi_2^-$) is larger than what was previously for small μ values while for large μ values, it is similar to the previous case.

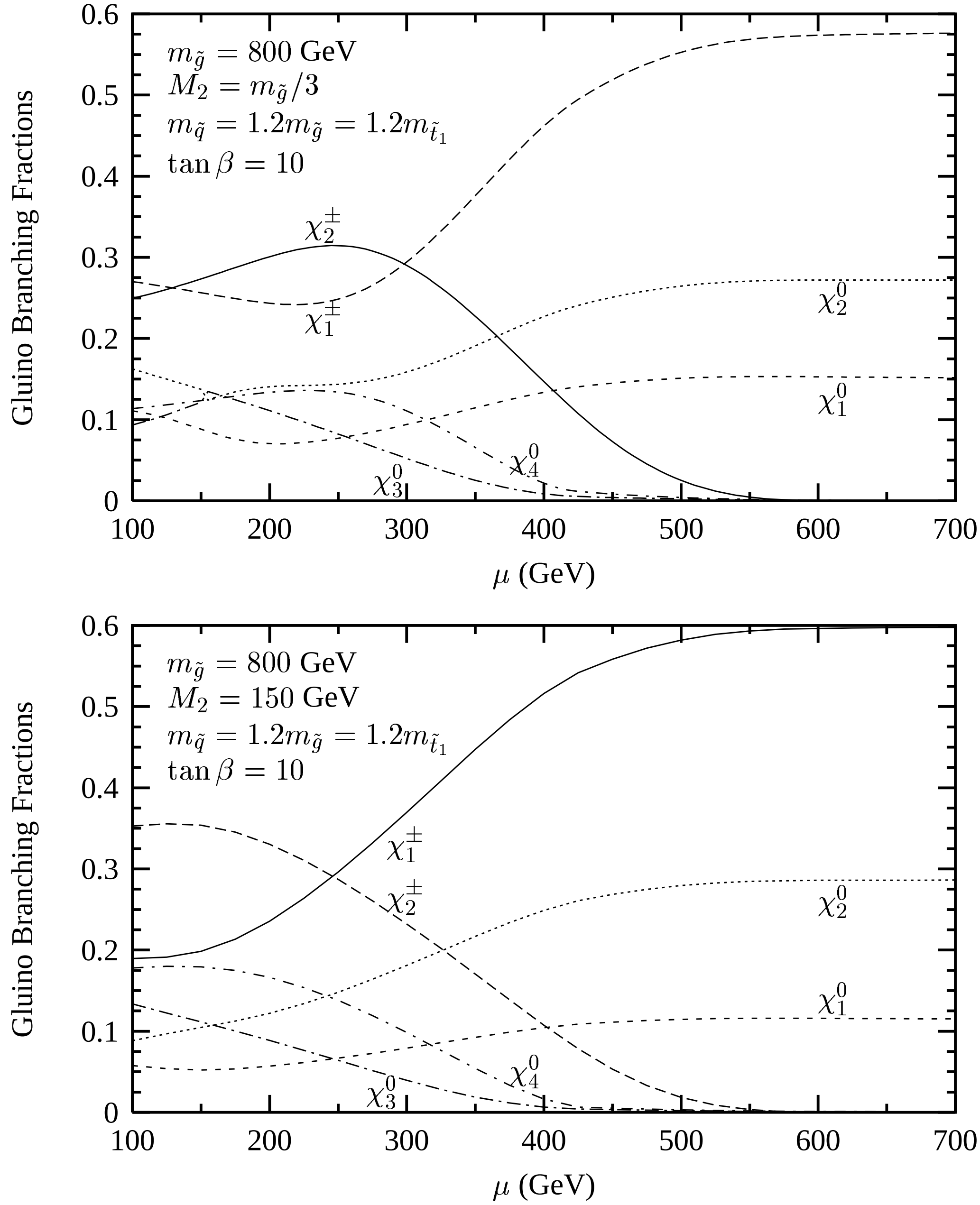

Figure 4: Branching ratios of gluinos into electroweak gauginos as functions of μ for $\tan\beta = 10$. The gluino mass is fixed to $m_{\tilde{g}} = 800$ GeV while the squark masses are $m_{\tilde{q}} = 1.2 m_{\tilde{t}_1} = 1.2 m_{\tilde{g}}$. The wino mass parameter is fixed to $M_2 = 2M_1 = m_{\tilde{g}}/3$ (upper curve) and $M_2 = 2M_1 = 150$ GeV (lower curve).

3. Decays into Charged Higgs Bosons

3.1 Two–body decays of heavy charginos and neutralinos

The heavier chargino and neutralinos will mainly decay into the lighter chargino and neutralino states and, if enough phase space is available, into sfermion–fermion pairs. The partial decay widths of the two–body decays, [including the possibility of massive fermions in the last case], are given by:

$$\Gamma(\chi_i \to f\tilde{f}_j) = \frac{\alpha N_c}{8} m_{\chi_i} \left[\left((a^f_{ij})^2 + (b^f_{ij})^2 \right) (1 - \mu_{\tilde{f}_j} + \mu_f) + 4\sqrt{\mu_f} a^f_{ij} b^f_{ij} \right] \lambda^{\frac{1}{2}}(\mu_f, \mu_{\tilde{f}_j}) \quad (17)$$

$$\Gamma(\chi_i \to \chi_j V) = \frac{\alpha}{8} m_{\chi_i} \lambda^{\frac{1}{2}}(\mu_{\chi_j}, \mu_V) \Big\{ -12\sqrt{\mu_{\chi_j}} G^L_{jiV} G^R_{jiV} + \left[(G^L_{jiV})^2 + (G^R_{jiV})^2 \right] (1 + \mu_{\chi_j} - \mu_V) + (1 - \mu_{\chi_j} + \mu_V)(1 - \mu_{\chi_j} - \mu_V)\mu_V^{-1} \Big\} \quad (18)$$

$$\Gamma(\chi_i \to \chi_j H_k) = \frac{\alpha}{8} m_{\chi_i} \lambda^{\frac{1}{2}}(\mu_{\chi_j}, \mu_{H_k}) \Big\{ \left[(G^L_{ijk})^2 + (G^R_{ijk})^2 \right] (1 + \mu_{\chi_j} - \mu_{H_k}) + 4\sqrt{\mu_{\chi_j}}\, G^L_{ijk} G^R_{ijk} \Big\} \quad (19)$$

The couplings among charginos, neutralinos and fermions/sfermions have been given previously, while the chargino and neutralino couplings to the Higgs and the gauge bosons are given by [$H_k = h, H, A$ and $H^\pm$ for $k = 1, 2, 3, 4$]:

$$G^{L,R}_{\chi^0_i \chi^+_j W^+} = G^{L,R}_{ijW} \quad \text{with} \quad \begin{array}{l} G^L_{ijW} = \frac{1}{\sqrt{2}s_W}[-Z_{i4}V_{j2} + \sqrt{2}Z_{i2}V_{j1}] \\ G^R_{ijW} = \frac{1}{\sqrt{2}s_W}[Z_{i3}U_{j2} + \sqrt{2}Z_{i2}U_{j1}] \end{array} \quad (20)$$

$$G^{L,R}_{\chi^-_i \chi^+_j Z} = G^{L,R}_{ijZ} \quad \text{with} \quad \begin{array}{l} G^L_{ijZ} = \frac{1}{c_W s_W} \left[-\frac{1}{2} V_{i2}V_{j2} - V_{i1}V_{j1} + \delta_{ij} s_W^2 \right] \\ G^R_{ijZ} = \frac{1}{c_W s_W} \left[-\frac{1}{2} U_{i2}U_{j2} - U_{i1}U_{j1} + \delta_{ij} s_W^2 \right] \end{array} \quad (21)$$

$$G^{L,R}_{\chi^0_i \chi^0_j Z} = G^{L,R}_{ijZ} \quad \text{with} \quad \begin{array}{l} G^L_{ijZ} = -\frac{1}{2 s_W c_W}[Z_{i3}Z_{j3} - Z_{i4}Z_{j4}] \\ G^R_{ijZ} = +\frac{1}{2 s_W c_W}[Z_{i3}Z_{j3} - Z_{i4}Z_{j4}] \end{array} \quad (22)$$

$$G^{L,R}_{\chi^0_i \chi^+_j H^+} = G^{L,R}_{ij4} \quad \text{with} \quad \begin{array}{l} G^L_{ij4} = \frac{c_\beta}{s_W} \left[Z_{j4}V_{i1} + \frac{1}{\sqrt{2}} \left(Z_{j2} + \tan\theta_W Z_{j1} \right) V_{i2} \right] \\ G^R_{ij4} = \frac{s_\beta}{s_W} \left[Z_{j3}U_{i1} - \frac{1}{\sqrt{2}} \left(Z_{j2} + \tan\theta_W Z_{j1} \right) U_{i2} \right] \end{array}$$

$$G^{L,R}_{\chi^-_i \chi^+_j H^0_k} = G^{L,R}_{ijk} \quad \text{with} \quad \begin{array}{l} G^L_{ijk} = \frac{1}{\sqrt{2}s_W} \left[e_k V_{j1}U_{i2} - d_k V_{j2}U_{i1} \right] \\ G^R_{ijk} = \frac{1}{\sqrt{2}s_W} \left[e_k V_{i1}U_{j2} - d_k V_{i2}U_{j1} \right] \epsilon_k \end{array} \quad (23)$$

$$G^{L,R}_{\chi^0_i \chi^0_j H_k} = G^{L,R}_{ijk} \quad \text{with} \quad \begin{array}{l} G^L_{ijk} = \frac{1}{2s_W} \left(Z_{j2} - \tan\theta_W Z_{j1} \right) \left(e_k Z_{i3} + d_k Z_{i4} \right) \; + \; i \leftrightarrow j \\ G^R_{ijk} = \frac{1}{2s_W} \left(Z_{j2} - \tan\theta_W Z_{j1} \right) \left(e_k Z_{i3} + d_k Z_{i4} \right) \epsilon_k \; + \; i \leftrightarrow j \end{array} \quad (24)$$

where $\epsilon_{1,2} = -\epsilon_3 = 1$ and the coefficients e_k and d_k read

$$e_1/d_1 = c_\alpha/-s_\alpha \ , \ e_2/d_2 = -s_\alpha/-c_\alpha \ , \ e_3/d_3 = -s_\beta/c_\beta \tag{25}$$

The branching ratios of the heavier chargino $\chi_2^\pm$ and $\chi_{3,4}^0$ into the lighter ones $\chi_1^\pm$ and $\chi_{1,2}^0$ and gauge and Higgs bosons are shown in Figs. 5 and 6 for, respectively, the two scenarii (i) gaugino–limit and (ii) higgsino–limit discussed previously. Squarks and sleptons are assumed to be too heavy to play a role here[5]. To analyze them, it is useful to discuss first the couplings of charginos and neutralinos to the Higgs and gauge bosons.

The Higgs bosons couples preferentially to mixtures of gauginos and higgsinos. This means that in the gaugino–like or higgsino like regions, the couplings of the Higgs bosons which involve heavy and light chargino/neutralino states are maximal, while the couplings involving only heavy or light ino states are suppressed by powers of M_2/μ for $|\mu| \gg M_2$ or powers of μ/M_2 for $|\mu| \ll M_2$. To the contrary, the gauge boson couplings to charginos and neutralinos are important only for higgsino–like states. Thus, in principle, the (higgsino or gaugino–like) heavier chargino and neutralinos $\chi_2^\pm$ and $\chi_{3,4}^0$, will dominantly decay, if phase space allowed, into Higgs bosons and the lighter χ states.

However, in the asymptotic limit where the heavier chargino and neutralino masses are very large, $m_{\chi_i} \gg m_{\chi_j}, M_{H_k}, M_V$, the decay widths into Higgs bosons grow as m_{χ_i},

$$\Gamma(\chi_i \to \chi_j H_k) \sim \frac{1}{8}\alpha m_{\chi_i}[(G^L_{ijk})^2 + (G^R_{ijk})^2] \tag{26}$$

while the decay widths into gauge bosons grow as $m^3_{\chi_i}$

$$\Gamma(\chi_i \to \chi_j V) \sim \frac{\alpha m^3_{\chi_i}}{8M_V^2}[(G^L_{ijV})^2 + (G^R_{ijV})^2] \tag{27}$$

This is due to the longitudinal component of the gauge boson propagators which introduce extra powers of the χ_i four–momentum in the decay amplitudes. The suppression of the $(G^{L,R}_{ijV})^2$ squared coupling by powers of $(\mu/M_2)^2$ or $(M_2/\mu)^2$ depending on whether we are in the gaugino or higgsino region will be compensated by the power m_χ^2/M_Z^2 from the matrix element squared. Therefore, the branching ratios for the decays of heavy χ particles into lighter ones and Higgs or gauge bosons will have the same order of magnitude. Of course, as usual, the charged current decay modes will be more important than the neutral current decay modes.

This is exemplified in the figures. In both higgsino and gaugino regions, Fig. 5 and Fig. 6 respectively, the decays of charginos $\chi_2^\pm$ and neutralinos $\chi_{3,4}^0$ into lighter charginos and neutralinos and Higgs bosons are not the dominant ones. Still, decays into $H^\pm$ bosons will have substantial branching fractions of the order of 20 to 30%.

[5]We assume in our analysis that the heavier charginos and neutralinos are coming from the decays of heavier squarks, including top squarks which is dealt as a special case. Decays of charginos and neutralinos into sleptons, which can be lighter than squarks, are relevant only if the former particles are gaugino–like since the higgsino–slepton–lepton couplings are rather tiny, unless $\tan\beta$ is very large in which case the decays into $\tilde{\tau}$'s could play a role. However, in the large $\tan\beta$ case, $H^\pm$ particles can be produced directly with large cross sections.

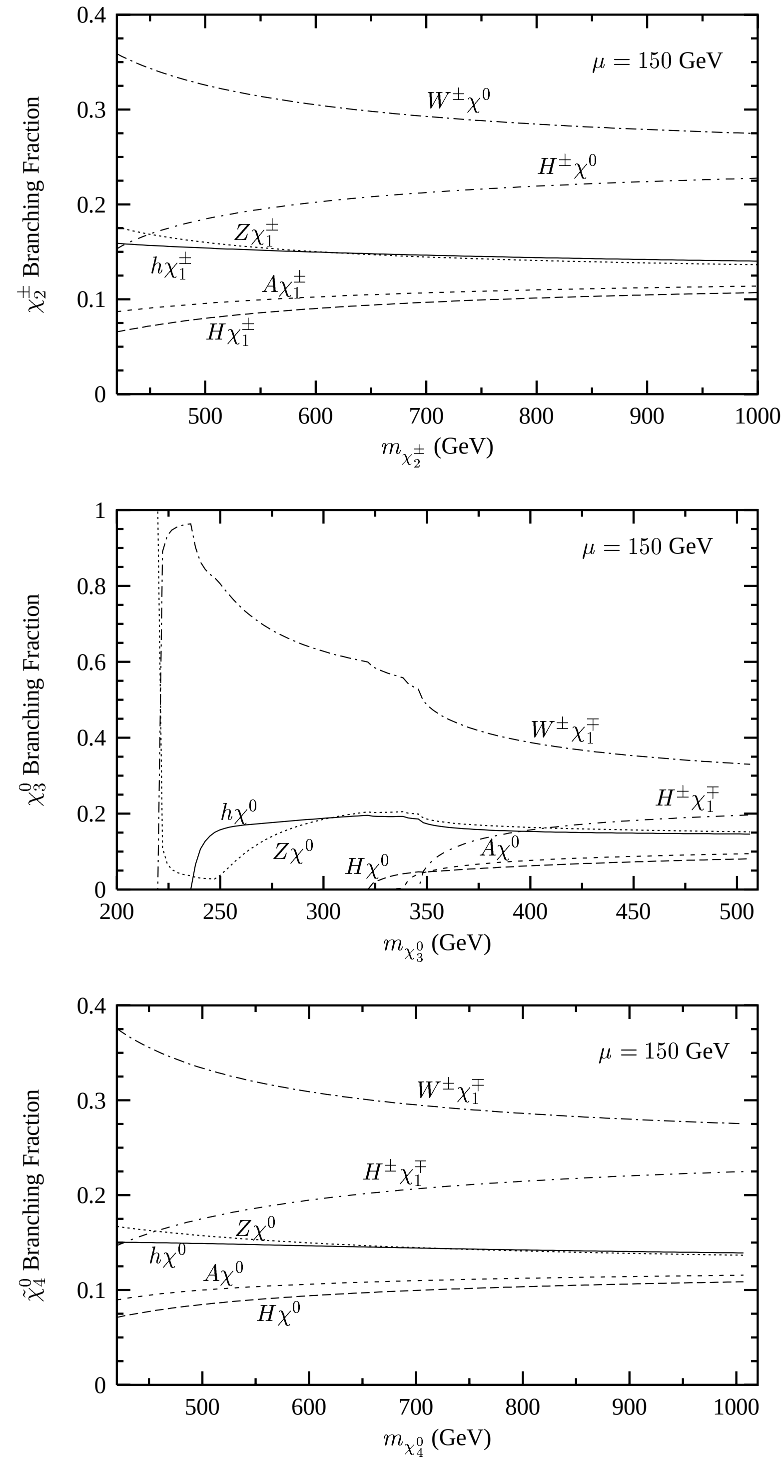

Figure 5: Branching ratios of heavier chargino and neutralinos into the lighter ones and gauge/Higgs bosons as functions of their masses for $\tan\beta = 10$. The charged Higgs boson mass is $M_{H^\pm} = 200$ GeV, μ is fixed to 150 GeV while M_2 varies with the heavy ino mass; χ^0 represents the lighter χ_1^0 and χ_2^0 neutralinos for which the branching ratios are added.

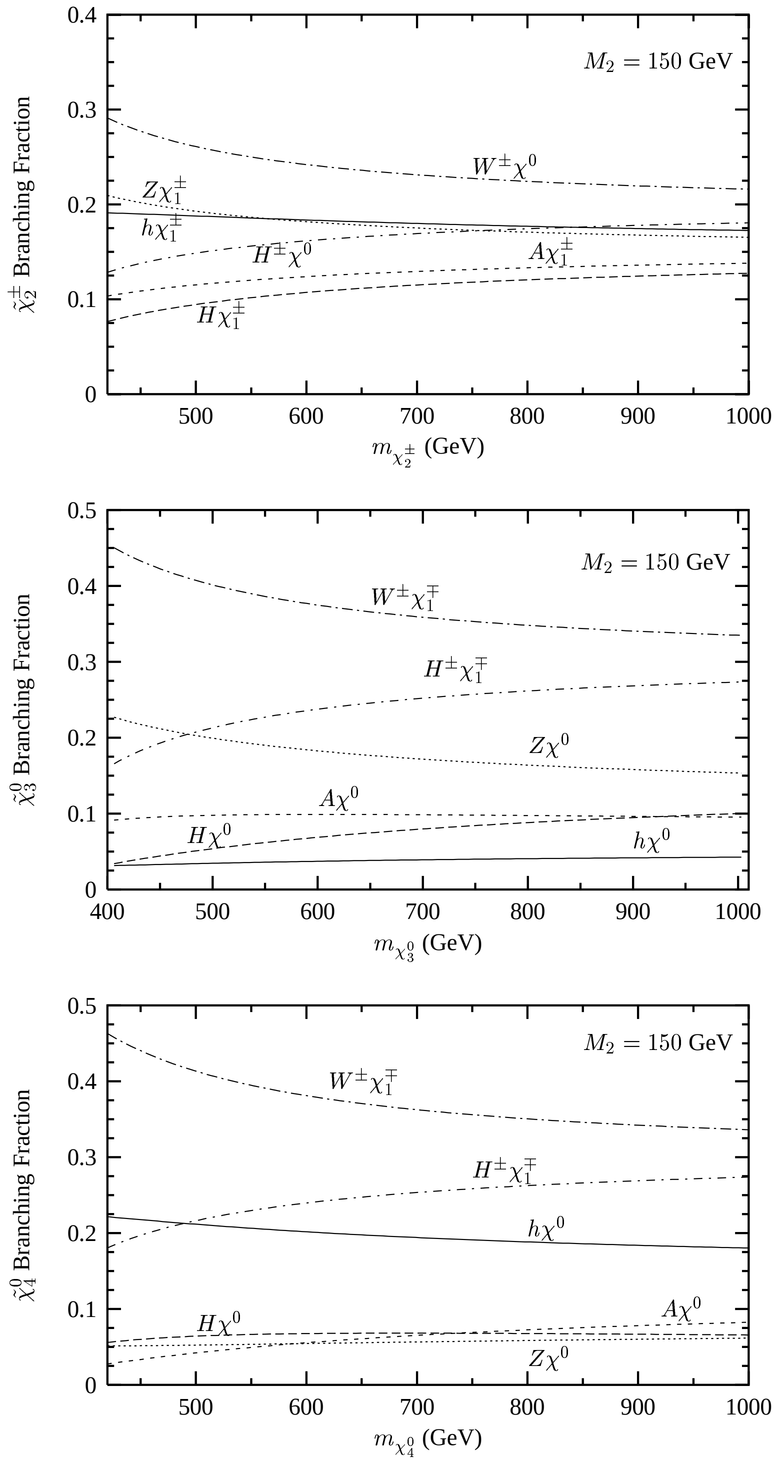

Figure 6: Branching ratios of heavier chargino and neutralinos into the lighter ones and gauge/Higgs bosons as functions of their masses for $\tan\beta = 10$. The charged Higgs boson mass is $M_{H^\pm} = 200$ GeV, M_2 is fixed to 150 GeV while μ varies with the heavy ino mass; χ^0 represents the lighter χ_1^0 and χ_2^0 neutralinos for which the branching ratios are added.

3.2 Cascade decays of squarks and gluinos into $H^\pm$ bosons

3.2.1 The case of squarks

We first focus on the scenario where gluinos are heavier than the scalar partners of the light quarks and in this case, all gluinos produced in the processes $pp \to \tilde{g}\tilde{g}$ and $pp \to \tilde{q}\tilde{g}$ will decay into squarks and light quarks. The squarks will subsequently decay into their partner quarks and all types of charginos or neutralinos. The heavier charginos and neutralinos will then decay into gauge bosons and Higgs bosons and the lighter chargino and neutralino states. Fig. 7 shows the final cross sections times branching ratio to obtain *one* charged Higgs boson in the final state, $\sigma \times \mathrm{BR}(\to H^\pm)$. For the total production cross sections, we use the same set of parton densities and the same scale as in section 2.1.

The particularly favorable case that we will discuss here is as follows. We choose a common squark mass of $m_{\tilde{q}} = 800$ GeV and fix $\tan\beta = 10$. We also fix $\mu = 150$ GeV and vary the wino mass parameter M_2 in such a way that for $M_2 \gtrsim 300$ GeV we have higgsino–like lighter chargino and neutralinos. For the gluino mass, we will assume in the first step that the gaugino masses are unified at the GUT scale so that $m_{\tilde{g}} = 3M_2$, and in the second step, we relax this universality assumption and fix the gluino mass to a constant value of $m_{\tilde{g}} = 900$ GeV, while still varying M_2. In the latter scenario, the gluino is not allowed kinematically to decay into top squarks if $m_{\tilde{t}_1} \simeq m_{\tilde{g}}$, but in the universal case since $m_{\tilde{g}} \sim 3M_2 \gtrsim m_{\tilde{t}_1} + m_t$, gluino decays into stops have to be taken into account. In the analysis, we will treat all squarks on equal footing and assume that the branching fraction for top squarks decaying into the heavier charginos and neutralinos are the same as for the scalar partners of light squarks. [This is a rather conservative approach since we have seen in the previous section that the branching ratios for stop decays into bottom quarks and the heavier chargino states can be larger.] The charged Higgs bosons masses are chosen to be $M_{H^\pm} = 180$, 200 and 300 GeV.

In the universal scenario of Fig. 7 [upper panel], we see that the cross sections times branching ratios for $H^\pm$ final states exceed the level of 0.1 pb in most of the displayed M_2 range. This means that with a luminosity of $\int \mathcal{L} = 300$ fb^{-1} which is expected to be collected at the LHC, around ~ 30.000 charged Higgs bosons can be produced through the cascade decays in this scenario. For small M_2 values, the states $\chi^0_{3,4}$ and χ^+_2 are not heavy enough for the decays into $H^\pm$ bosons to occur, in particular for large $M_{H^\pm}$. When these decays are allowed, $\sigma \times \mathrm{BR}(\to H^\pm)$ values of the order of 1 pb for $M_{H^\pm} \sim 180$ GeV and 0.3 pb for $M_{H^\pm} \sim 300$ GeV can be reached. For increasing values of M_2, the gluino mass increases and the cross section for associated squark and gluino production, which is the largest in this case, as well as for gluino pair production drop and $\sigma \times \mathrm{BR}(\to H^\pm)$ decreases accordingly; at some stage, only the cross section for squark production survives [since $m_{\tilde{q}}$ is fixed]. The decrease of $\sigma \times \mathrm{BR}(\to H^\pm)$ with increasing M_2 is also due to the more suppressed phase space for $\tilde{q} \to q'\chi_2^\pm, q\chi_4^0$ since for large M_2, $m_{\chi_4^0}, m_{\chi_2^\pm} \sim M_2$. For even larger M_2 values, $M_2 \gtrsim 650$ GeV, the channel $\chi_3^0 \to H^\pm\chi_1^\mp$ opens up, and since the phase space is more favorable, because $m_{\chi_3^0} \sim M_2/2$, $\sigma \times \mathrm{BR}(\to H^\pm)$ increases again.

In the non–universal case [lower panel of Fig. 7], with a gluino mass fixed to $m_{\tilde{g}} = 900$ GeV, the cross sections times branching ratios for $H^\pm$ final states can be much larger. In the low M_2 region, the situation is similar to the previous case in which the gluino was rather light and contributed substantially to the production cross sections. However, for increasing M_2 values, $\sigma_{\tilde{g}\tilde{g}}$ and $\sigma_{\tilde{g}\tilde{q}}$ stay constant contrary to the universal case, while the phase space for the decays of the heavier charginos and neutralinos into charged Higgs bosons increases. $\sigma \times \mathrm{BR}(\to H^\pm)$ can then reach the picobarn level even for relatively large charged Higgs boson masses, $M_{H^\pm} \sim 300$ GeV, leading to samples containing several hundred thousand charged Higgs bosons events with the high luminosity $\int \mathcal{L} = 300$ fb^{-1}.

3.2.2 The case of gluinos

The production cross sections times branching ratios for squarks decaying into gluinos, with the subsequent three–body decays of the gluinos into the heavier charginos and neutralinos which then decay into the lighter χ states and charged Higgs bosons are shown in Fig. 8. Again we set $\tan\beta = 10$ and fix the gluino and squark masses to $m_{\tilde{q}} = 1.2 m_{\tilde{g}}, m_{\tilde{g}} = 800$ GeV; the $H^\pm$ boson masses are taken to be $M_{H^\pm} = 180, 200$ and 300 GeV. Then we vary the higgsino parameter μ and take the universal scenario where M_2 and M_3 are related, as well as the non–universal scenario where $M_2 = 150$ GeV so that the heavier χ states are higgsino–like for large enough values of $|\mu|$. In both cases the cross sections for squark and gluino production are constant and the variation of $\sigma \times \mathrm{BR}(H^\pm)$ is only due to the variation of the branching ratios $\mathrm{BR}(\tilde{g} \to \chi^0_{3,4} qq, \chi_2^\pm qq')$ and $\mathrm{BR}(\chi^0_{3,4}, \chi_2^\pm \to \chi_1^\pm H^\mp, \chi^0_{1,2} H^\pm)$.

In the universal scenario, with $m_{\tilde{g}} \sim M_3 \sim 3M_2$, $\sigma \times \mathrm{BR}(H^\pm)$ is relatively large for small values of μ and $M_{H^\pm}$, when the gaugino–like heavy χ states are light enough for the decays $\chi^0_4 \to \chi_1^\pm H^\mp$ and $\chi_2^\pm \to \chi^0_{1,2} H^\pm$ to occur. In the mixed region, $\mu \sim M_2$, the mass difference between the heavy and light χ states are too small to allow for decays in $H^\pm$ bosons. For large values of μ, $\sigma \times \mathrm{BR}(H^\pm)$ increases to reach values of the order of ~ 0.1 pb for $M_{H^\pm} \sim 200$ GeV [in particular when the additional channels $\chi^0_3 \to \chi_1^\pm H^\mp$ open up] before it drops out because of the gradually closing phase space for the decays $\tilde{g} \to q\bar{q}\chi^0_{3,4}, qq'\chi_2^\pm$.

In the non–universal case with $M_2 = 150$ GeV, the region with $\mu \lesssim 200$ GeV is cut–away since the mass difference between the heavier and lighter χ states is too small for the decays into $H^\pm$ bosons to occur. For larger μ, the situation becomes similar to the universal case and $\sigma \times \mathrm{BR}(H^\pm)$ values of the order of 0.3 pb can be reached for not too large $H^\pm$ masses and favorable regions of parameter space, leading to a sample of a hundred thousand charged Higgs boson events for the luminosity expected at the LHC.

3.2.3 The case of top squarks

The cross section for top squark production at the LHC can be rather large when $m_{\tilde{q}} \geq m_{\tilde{g}} \geq m_{\tilde{t}_1} + m_t$; in this case all squarks and gluinos will decay into $\tilde{t}_1$ final states. The stops will decay into heavier χ states, with rates shown in Fig. 3, and the latter will possibly decay into lighter χ particles and $H^\pm$ bosons as shown in Figs. 5 and 6. The situation will then be almost exactly similar to gluino decays discussed in the previous subsection. The reason is that gluino decays into the heavy χ particle occurred mainly

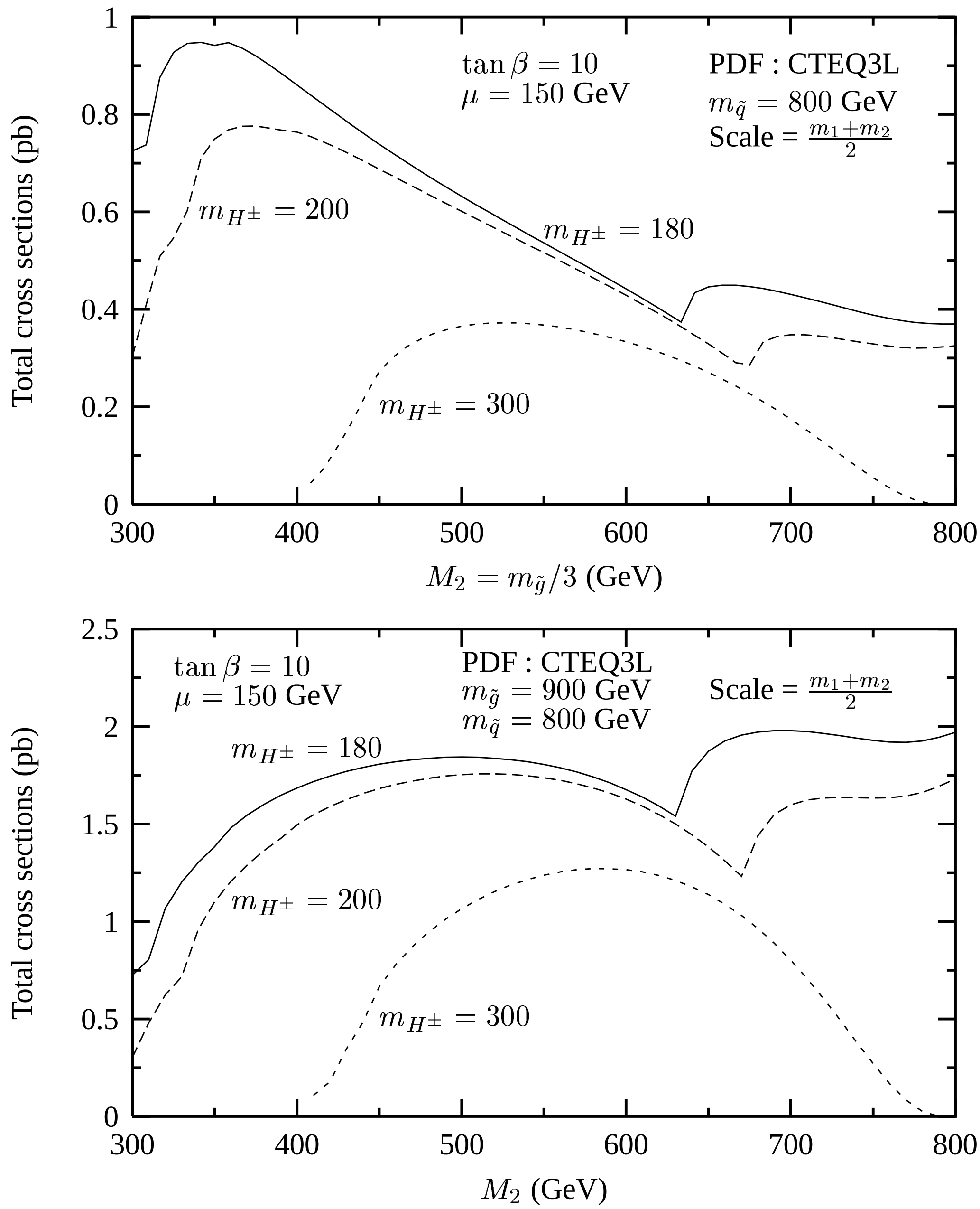

Figure 7: Cross sections times branching ratios for gluinos decaying into squarks and squarks decaying through cascades into charged Higgs bosons with masses $M_{H^\pm} = 180, 200$ and 300 GeV. They are shown as functions of M_2 with $\mu = 150$ GeV and $\tan\beta = 10$. The squark mass is fixed to $m_{\tilde q} = 800$ GeV, while the gluino mass is $m_{\tilde g} = 3M_2$ (upper curve) and $m_{\tilde g} = 900$ GeV (lower curve).

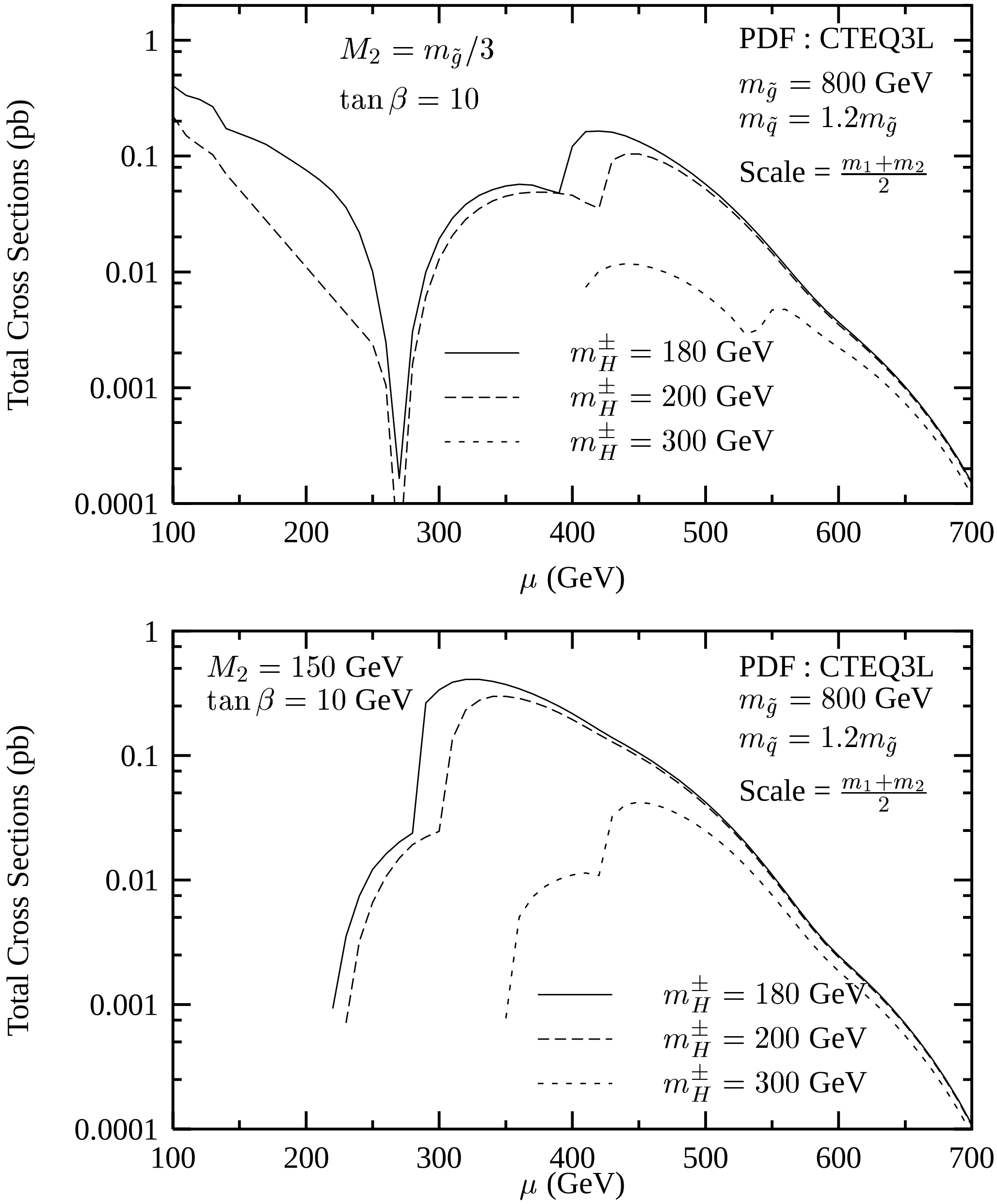

Figure 8: Cross sections times branching ratios for squarks decaying into gluinos with the gluinos decaying through cascades into charged Higgs bosons with masses $M_{H^\pm} = 180, 200$ and 300 GeV. They are shown as functions of μ for $\tan\beta = 10$. The gluino mass is fixed to $m_{\tilde{g}} = 800$ GeV and the squark mass to $m_{\tilde{q}} = 1.2 m_{\tilde{g}}$. In the upper figure, $M_2 = m_{\tilde{g}}/3$ while $M_2 = 150$ GeV in the lower figure.

through the virtual exchange of $\tilde{t}_1$ and the branching ratios were controlled by the $\tilde{t}_1 b \chi_2^{\pm}$ and $\tilde{t}_1 t \chi_{3,4}^0$ vertices. The only difference with the previous discussion will be due to the smaller value of $m_{\tilde{t}_1}$ [since now $\tilde{t}_1$ is not virtual anymore] which will suppress the phase space for $\tilde{t}_1 \to b\chi_2^+, t\chi_{3,4}^0$ decays.

Note that an additional contribution to the cross section will come from the direct production of stop quarks $gg/q\bar{q} \to \tilde{t}_1 \tilde{t}_1^*$, which for $m_{\tilde{t}_1} \sim 500$ GeV can reach the picobarn level. This may substantially increase the number of charged Higgs bosons in the final state through the cascade decay $\tilde{t}_1 \to b\chi_2^+ \to b\chi_{1,2}^0 H^+$ for instance.

4. Direct H^+ decays of squarks and gluinos

4.1 Two–body decays of squarks into Higgs bosons

If the mass splitting between two squarks of the same generation is large enough, as is generally the case of the $(\tilde{t}, \tilde{b})$ iso–doublet, the heavier squark can decay into a lighter one plus a gauge boson $V = W, Z$ or a Higgs boson $\Phi = h, H, A, H^{\pm}$. The partial decay widths are given at the tree–level by [the QCD corrections to these decay modes have also been calculated and can be found in Ref. [25]]:

$$\begin{aligned}
\Gamma(\tilde{q}_i \to \tilde{q}'_j V) &= \frac{\alpha}{4M_V^2} m_{\tilde{q}_i}\, g^2_{\tilde{q}_i \tilde{q}'_j V}\, \lambda^{3/2}(\mu_V^2, \mu_{\tilde{q}'_j}^2) && (28) \\
\Gamma(\tilde{q}_i \to \tilde{q}'_j \Phi) &= \frac{\alpha}{4} m_{\tilde{q}_i}\, g^2_{\tilde{q}_i \tilde{q}'_j \Phi}\, \lambda^{1/2}(\mu_\Phi^2, \mu_{\tilde{q}'_j}^2) && (29)
\end{aligned}$$

In these equations, the couplings of the Higgs bosons to squarks, $g_{\tilde{q}_i \tilde{q}'_j \Phi}$, read in the case of neutral Higgs bosons:

$$\begin{aligned}
g_{\tilde{q}_1 \tilde{q}_2 h} &= \frac{1}{4 s_W M_W} \left[M_Z^2 s_{2\theta_q} (2I_q^3 - 4e_q s_W^2) \sin(\alpha+\beta) + 2m_q c_{2\theta_q} (A_q r_2^q + 2I_q^3\, \mu\, r_1^q) \right] \\
g_{\tilde{q}_1 \tilde{q}_2 H} &= \frac{1}{4 s_W M_W} \left[- M_Z^2 s_{2\theta_q} (2I_q^3 - 4e_q s_W^2) \cos(\alpha+\beta) + 2m_q c_{2\theta_q} (A_q r_1^q - 2I_q^3\, \mu\, r_2^q) \right] \\
g_{\tilde{q}_1 \tilde{q}_2 A} &= -g_{\tilde{q}_2 \tilde{q}_1 A} = \frac{-m_q}{2 s_W M_W} \left[\mu + A_q (\tan\beta)^{-2I_q^3} \right] && (30)
\end{aligned}$$

with the coefficients $r_{1,2}^q$ as [α is a mixing angle in the CP–even Higgs sector of the MSSM, and at the tree–level, can be expressed only in terms of M_A and $\tan\beta$]

$$r_1^t = \frac{\sin\alpha}{\sin\beta} \;\; , \;\; r_2^t = \frac{\cos\alpha}{\sin\beta} \;\; , \;\; r_1^b = \frac{\cos\alpha}{\cos\beta} \;\; , \;\; r_2^b = -\frac{\sin\alpha}{\cos\beta} \; . \qquad (31)$$

In the case of the charged Higgs boson, the couplings to squarks are given, in terms of the squark mixing matrices $R^{\tilde{q}}$, by

$$g_{\tilde{q}_i \tilde{q}'_j H^\pm} = \frac{1}{2 s_W M_W} \sum_{k,l=1}^{2} \left(R^{\tilde{q}} \right)_{ik} C^{kl}_{\tilde{q}\tilde{q}' H^\pm} \left(R^{\tilde{q}'} \right)^{\mathrm{T}}_{lj} \qquad (32)$$

with the matrix $C_{\tilde{q}\tilde{q}'H^\pm}$ summarizing the couplings of the $H^\pm$ bosons to the squark current eigenstates and it is given by

$$C_{\tilde{t}\tilde{b}H^\pm} \;=\; \sqrt{2}\begin{pmatrix} m_b^2 \tan\beta + m_t^2/\tan\beta - M_W^2 \sin 2\beta & m_b\,(A_b \tan\beta + \mu) \\ m_t\,(A_t/\tan\beta + \mu) & 2\,m_t\,m_b/\sin 2\beta \end{pmatrix} \tag{33}$$

For the couplings of squarks to the W and Z gauge bosons, one has

$$\begin{aligned} g_{\tilde{q}_1\tilde{q}_2 Z} &= g_{\tilde{q}_2\tilde{q}_1 Z} = \frac{2I_q^3 s_{2\theta_q}}{4 s_W c_W} \\ g_{\tilde{q}_i\tilde{q}'_j W} &= \frac{1}{\sqrt{2}s_W}\begin{pmatrix} c_{\theta_q}c_{\theta'_q} & -c_{\theta_q}s_{\theta'_q} \\ -s_{\theta_q}c_{\theta'_q} & s_{\theta_q}s_{\theta'_q} \end{pmatrix} \end{aligned} \tag{34}$$

In Fig. 9, we display the branching ratio for the decays of a bottom squark into the lightest top squark and a charged Higgs boson, $\tilde{b}_1 \to \tilde{t}_1 H^-, \tilde{b}_1^* \to \tilde{t}_1^* H^+$ as a function of the parameter μ with three values of $M_2 = 200$, 300 and 400 GeV. We have fixed $\tan\beta = 10$ and the sbottom, stop and charged Higgs boson masses to $m_{\tilde{b}_1} = 832$ GeV, $m_{\tilde{t}_1} = 430$ GeV and $M_{H^\pm} = 200(300)$ GeV in the upper (lower) panel. The other competing decays of sbottoms are decays into bottom quarks and neutralinos and top quarks and charginos as well as bottom and gluino final states. We have assumed universal gaugino masses so that the gluino mass is given by $m_{\tilde{g}} \sim 3M_2$; the strong interaction decay channel $\tilde{b}_1 \to b\tilde{g}$ is therefore only open for $M_2 = 200$ GeV and is dominant in this case.

As can be seen, for $M_2 \geq 300$ GeV [i.e when there is no phase space for the decay channel $\tilde{b}_1 \to b\tilde{g}$ to occur], BR($\tilde{b}_1 \to \tilde{t}_1 H^-$) can be substantial for large μ values, $\mu \gtrsim 700$ GeV, possibly exceeding the level of 50%. The reason for this feature, besides the fact that for $\mu \gtrsim 800$ GeV, the $\tilde{b}_1$ decays into the heavier chargino and neutralinos are kinematically closed, is that the sbottom–stop–$H^\pm$ coupling is strongly enhanced and becomes larger than the sbottom–bottom–gaugino coupling which controls the sbottom decays into the lighter chargino and neutralinos. For smaller values of M_2, as pointed out earlier, the decay $\tilde{b}_1 \to b\tilde{g}$ becomes accessible and would be the dominant decay channel.

Another possibility would be the decays of the heaviest top squark into a sbottom quark and charged Higgs bosons, $\tilde{t}_2 \to \tilde{b}_{1,2}H^+$: since the $\tilde{t}_2 - \tilde{t}_1$ mass splitting can be large, there might be enough phase space for this process to occur. However, the cross sections for $\tilde{t}_2$ from direct production are not large, and if the gluinos are heavier than squarks, they will decay into $\tilde{t}_2 t$ final states only less than $\sim 10\%$ of the time. This decay channel is therefore less favored than the previous one.

Note that for light top squarks, $m_{\tilde{t}_1} \lesssim m_t + m_{\chi_1^0}$ and $m_b + m_{\chi_1^\pm}$, and light charged Higgs bosons $M_{H^\pm} \lesssim m_{\tilde{t}_1} - m_{\chi_1^0}$, the three body decay $\tilde{t}_1 \to bH^+\chi_1^0$ [which is mediated by virtual chargino, top or sbottom exchange] is accessible and would compete with the other possible three–body decay mode $\tilde{t}_1 \to bW^+\chi_1^0$ and the loop induced and flavor changing decay $\tilde{t}_1 \to c\chi_1^0$. In some areas of the parameter space, the stop decay into charged Higgs bosons can be dominant; see the discussions in Ref. [26].

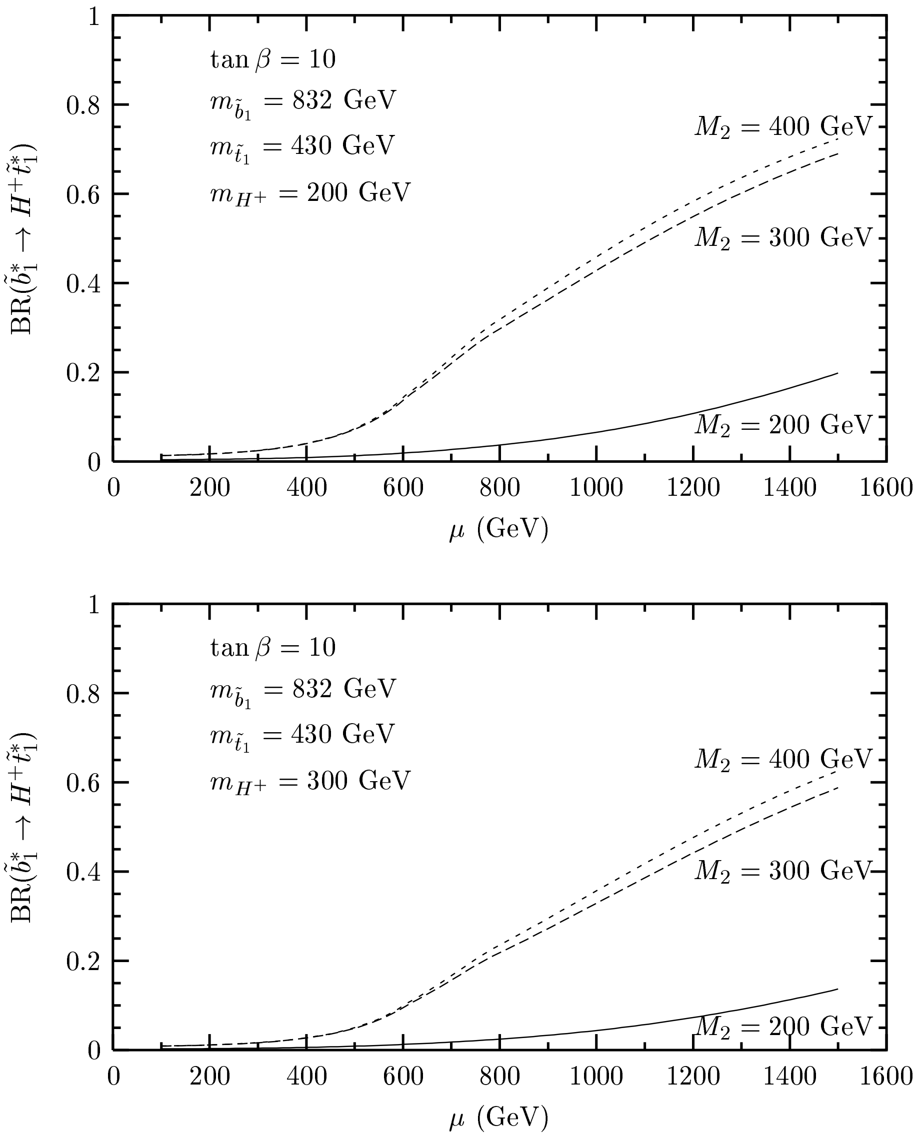

Figure 9: The branching ratios for bottom squarks decaying into top squarks and charged Higgs bosons as a function of μ for $\tan\beta = 10$ and $M_2 = 200, 300$ and 400 GeV. The charged Higgs boson mass is taken to be $M_{H^\pm} = 200$ and 300 GeV, in the upper and lower panels, respectively. The two squark masses are taken to be $m_{\tilde{b}_1} = 832$ GeV and $m_{\tilde{t}_1} = 430$ GeV.

4.2 Three body decays of gluinos into charged Higgs bosons

Finally, we discuss the direct decays of gluinos into top squarks, bottom quarks and charged Higgs bosons, mediated by virtual top quark or bottom squark exchanges; Fig. 10.

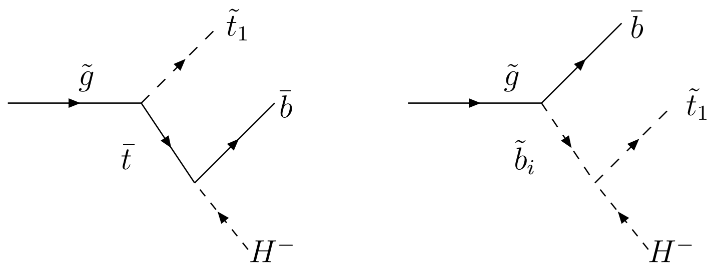

Figure 10: The Feynman diagrams contributing to the three–body decay $\tilde{g} \to \tilde{t}_1 \bar{b} H^-$.

The Dalitz density for this decay mode, taking into account all the masses of the final state particles, including the bottom quark, is given by:

$$\frac{\mathrm{d}\Gamma}{\mathrm{d}x_1 \mathrm{d}x_2}(\tilde{g} \to H^- \bar{b} \tilde{t}_1) \;=\; \frac{\alpha \alpha_s}{64\pi} m_{\tilde{g}} \left[\mathrm{d}\Gamma_t + \mathrm{d}\Gamma_{\tilde{b}} + 2\mathrm{d}\Gamma_{t\tilde{b}} \right] \tag{35}$$

In terms of $x_1 = 2E_{H^\pm}/m_{\tilde{g}}$, $x_2 = 2E_b/m_{\tilde{g}}$ and the reduced masses $\mu_X = m_X/m_{\tilde{g}}$, the squared t, $\tilde{b}$ contributions and the $t\tilde{b}$ interference are given by

$$\begin{aligned}
\mathrm{d}\Gamma_t \;&=\; \frac{1}{(x_1 + x_2 - 1 + \mu_{\tilde{t}_1} - \mu_t)^2} \Big\{ \mu_t x_2 \left[y_t^2 c_\beta^2 (b_{1\tilde{g}}^{\tilde{t}})^2 + y_b^2 s_\beta^2 (a_{1\tilde{g}}^{\tilde{t}})^2 \right] \\
&\quad + \left[y_t^2 c_\beta^2 (a_{1\tilde{g}}^{\tilde{t}})^2 + y_b^2 s_\beta^2 (b_{1\tilde{g}}^{\tilde{t}})^2 \right] \left[x_1 (1 - \mu_b + \mu_{H^+} - \mu_{\tilde{t}_1} - x_1 - x_2) + x_2 (\mu_{H^+} - \mu_b) \right] \Big\} \\
\mathrm{d}\Gamma_{\tilde{b}} \;&=\; \sum_{i,j=1}^{2} \frac{g_{\tilde{t}_1 \tilde{b}_i H^+} g_{\tilde{t}_1 \tilde{b}_j H^+} [a_{i\tilde{g}}^{\tilde{b}} a_{j\tilde{g}}^{\tilde{b}} + b_{i\tilde{g}}^{\tilde{b}} b_{j\tilde{g}}^{\tilde{b}}] x_2}{(1 + \mu_b - x_2 - \mu_{\tilde{b}_i})(1 + \mu_b - x_2 - \mu_{\tilde{b}_j})} \\
\mathrm{d}\Gamma_{t\tilde{b}_1} \;&=\; \sum_{i=1}^{2} \frac{g_{\tilde{t}_1 \tilde{b}_H^+}}{(1 + \mu_b - x_2 - \mu_{\tilde{b}_i})(x_1 + x_2 - 1 + \mu_{\tilde{t}_1} - \mu_t)} \Big\{ -\sqrt{\mu_t} x_2 \left[y_t c_\beta b_{i\tilde{g}}^{\tilde{b}} b_{1\tilde{g}}^{\tilde{t}} + y_b s_\beta a_{i\tilde{g}}^{\tilde{b}} a_{1\tilde{g}}^{\tilde{t}} \right] \\
&\quad + [y_t c_\beta b_{i\tilde{g}}^{\tilde{b}} a_{1\tilde{g}}^{\tilde{t}} + y_b s_\beta a_{i\tilde{g}}^{\tilde{b}} b_{1\tilde{g}}^{\tilde{t}}] (x_1 + x_2 - 1 - \mu_b - \mu_{H^+} + \mu_{\tilde{t}_1}) \Big\}
\end{aligned} \tag{36}$$

The various couplings have been given in the previous sections and the Yukawa couplings of top and bottom quarks are given by: $y_t = m_t/(\sqrt{2} s_W M_W s_\beta)$ and $y_b = m_b/(\sqrt{2} s_W M_W c_\beta)$.

To obtain the partial decay width, one has to integrate over x_1 and x_2 with boundary conditions:

$$2\sqrt{\mu_{H^+}} \;\leq\; x_1 \;\leq 1 + [\mu_{H^+} - (\sqrt{\mu_b} + \sqrt{\mu_{\tilde{t}_1}})^2] \tag{37}$$

$$s_{\min} \;\leq x_2 \;\leq s_{\max} \tag{38}$$

with

$$
\begin{aligned}
s_{\min} &= \frac{1}{2}\frac{(x_1-2)(x_1-1-\mu_b+\mu_{\tilde{t}_1}-\mu_{H^+})-\sqrt{\Delta}}{1-x_1+\mu_{H^+}} \\
s_{\max} &= \frac{1}{2}\frac{(x_1-2)(x_1-1-\mu_b+\mu_{\tilde{t}_1}-\mu_{H^+})+\sqrt{\Delta}}{1-x_1+\mu_{H^+}}
\end{aligned} \tag{39}
$$

and

$$
\Delta = (\mu_{H^+}-x_1^2)\left[\frac{1}{4}\mu_b\mu_{\tilde{t}_1}-(x_1-1+\mu_b+\mu_{\tilde{t}_1}-\mu_{H^+})^2\right] \tag{40}
$$

The branching fraction for the three–body decay, $\mathrm{BR}(\tilde{g}\to\tilde{t}_1\bar{b}H^-+\tilde{t}_1^*bH^+)$ is illustrated in Fig. 11 as a function of μ for $\tan\beta=10$. We have chosen squark masses of $m_{\tilde{q}}=m_{\tilde{b}_i}=1$ TeV, a gluino mass slightly lower, $m_{\tilde{g}}=900$ GeV, and the lighter stop mass to be $m_{\tilde{t}_1}=433$ GeV; for the charged Higgs boson mass we take three values: $M_{H^\pm}=190, 230$ and 310 GeV. Thus in this scenario, all squarks [including bottom squarks] will decay into almost massless quarks and gluinos and the latter will dominantly decay into the lighter top squarks and top quarks. The three–body decays $\tilde{g}\to\tilde{t}_1\bar{b}H^-$ and $\tilde{g}\to\tilde{t}_1^*bH^+$ have therefore to compete with a strong interaction two–body decay, which has a large phase space in this case. This is the reason why the branching ratio hardly exceeds the one percent level, which occurs for large μ values when the $\tilde{t}\tilde{b}H^\pm$ couplings are enhanced. Note that the smallness of the branching ratio is also due to the smallness of the tbH^+ coupling for the chosen value of $\tan\beta$; for larger or smaller values of $\tan\beta$, the branching ratio can be significantly larger [e.g. an order of magnitude larger than in Fig. 11 for $\tan\beta=30$].

In spite of the small branching ratio, the number of $H^\pm$ final states can be rather large at the LHC in the kinematical configuration shown in Fig. 11. Indeed, for the chosen squark and gluino masses, the cross section for gluinos coming from direct production and from squark decays is at the level of ~ 5 picobarn. With branching fractions of the order of a few percent, this means that the cross section for $H^\pm$ production via gluino decays can reach the level of a fraction of a picobarn leading, in the favorable case, to a sample of a few ten thousand H^+ events with an integrated luminosity of $\int\mathcal{L}=300$ fb^{-1}.

Finally, let us note that for charged Higgs bosons lighter than top quarks, the decay $t\to bH^+$ will occur at the two–body level and the branching fraction can be rather large [of the order of 10 percent for $\tan\beta=10$]. In the kinematical configuration discussed often in this analysis, i.e. squarks heavier than gluinos which are heavier than $m_{\tilde{t}_1}+m_t$, all strongly interacting particles will cascade into stop and top quarks, with the latter decaying into charged Higgs bosons and bottom quarks. The number of H^+ boson can be therefore substantial; see also the discussion given in [16]. These processes will complement [and possibly, also compete with] searches of $H^\pm$ bosons from top quark decays with these particles produced directly in pp collisions.

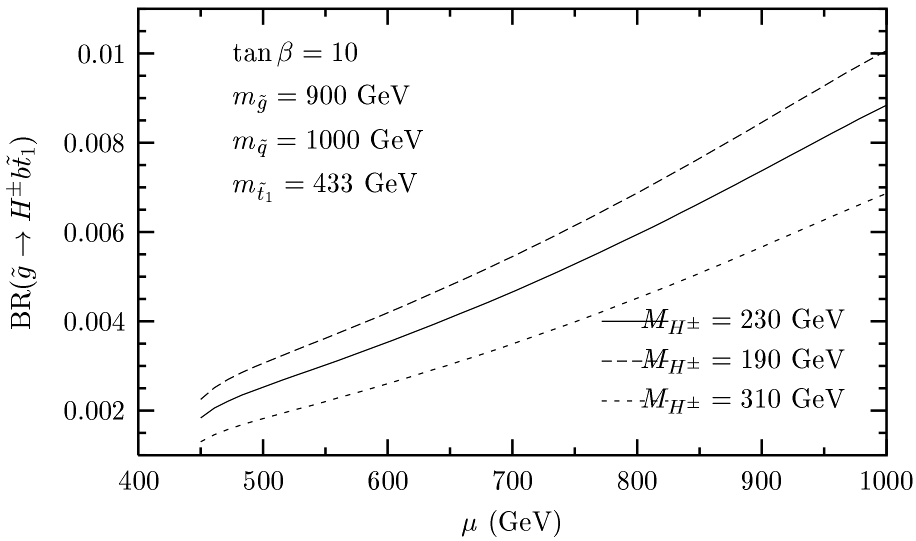

Figure 11: The branching ratios for direct decays of gluinos into bottom quarks, top squarks and charged Higgs bosons as a function of μ for $\tan\beta = 10$ and $M_{H^\pm} = 190, 230$ and 310 GeV. The squark masses are taken to be $m_{\tilde{q}} = 1$ TeV and $m_{\tilde{t}_1} = 433$ GeV, while the gluino mass is $m_{\tilde{g}} = 900$ GeV.

5. Conclusions

We have analyzed the production of the charged Higgs particles of the unconstrained Minimal Supersymmetric Standard Model from cascade as well as direct decays of squarks and gluinos, which have large production cross sections at the LHC due to their strong couplings.

Squarks and gluinos can decay into the heavier chargino and neutralinos, which subsequently decay into the lighter chargino and neutralino states and charged Higgs bosons. We have shown that the branching fractions for these decay modes can be rather large in some regions of the MSSM parameter space, and that the cross sections times the branching ratios for $H^\pm$ production can be significantly large. This would allow the possibility to produce large samples of these states at the LHC, in particular in situations [$H^\pm$ masses larger than the top quark mass or intermediate values of $\tan\beta$, for instance] where their production cross sections in the direct processes [associated production with bottom quarks or with tb final states] are rather small.

We have also discussed the possibility of producing charged Higgs bosons either through direct decays of the scalar partners of heavy quarks, as is the case in the process $\tilde{b}_{1,2} \to H^-\tilde{t}_1$, or in direct three–body decays of gluinos, $\tilde{g} \to \tilde{t}_1\bar{b}H^-$ and $\tilde{t}_1^* bH^+$. In some favorable regions of the parameter space, the branching ratios of these decay channels, in particular in the former case, can be large enough to allow for the detection of the charged Higgs bosons. For gluino decays, the situation becomes even more favorable for light $H^\pm$ bosons when the two-body decays $t \to bH^+$ are kinematically possible, leading to a surplus of events compared to the case $pp \to \bar{t}t \to H^\pm + X$.

Thus, there is an additional source of charged Higgs particles at the LHC with interesting signals[6] in most cases, since the final states involve missing energy, multi–leptons [from the decay cascades] and heavy flavors [b and t quarks]. These signals would help in detecting these particles in the difficult environment of the LHC. In our parton level analysis of these process, we did not attempt to scan the entire MSSM parameter space, discuss the final state topologies and the possible backgrounds, etc., although for the cross sections and branching ratios, we tried to be rather conservative [for instance by not including K–factors, etc.]. A detailed analysis taking into account all backgrounds, selection and detection efficiencies in a realistic way, which is beyond the scope of the present paper, is required to assess in which part of the MSSM parameter space these final states can be isolated experimentally.

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[6]We thank Filip Moortgat from the CMS collaboration for a discussion on this point.

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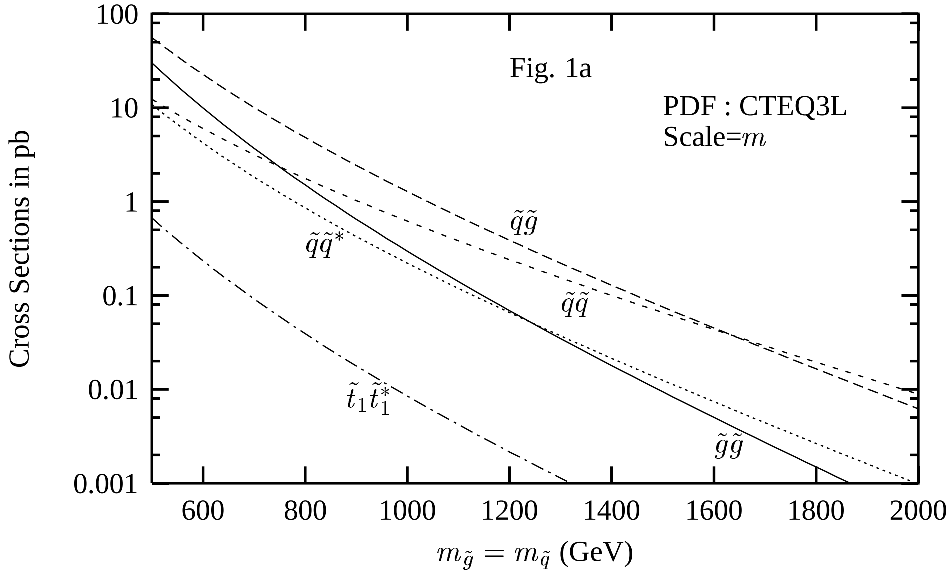
Fig. 1a
PDF : CTEQ3L
Scale=m
Cross Sections in pb
100
10
1
0.1
0.01
0.001
600
800
1000
1200
1400
1600
1800
2000
$m_{\tilde{g}} = m_{\tilde{q}}$ (GeV)
$\tilde{q}\tilde{g}$
$\tilde{q}\tilde{q}^*$
$\tilde{q}\tilde{q}$
$\tilde{t}_1\tilde{t}_1^*$
$\tilde{g}\tilde{g}$

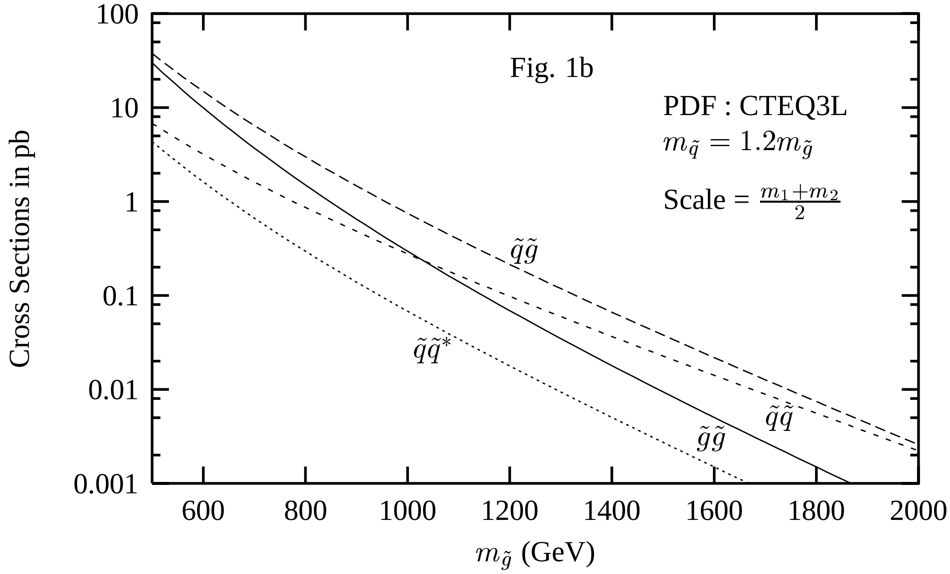

Fig. 1b
PDF : CTEQ3L
$m_{\tilde{q}} = 1.2 m_{\tilde{g}}$
Scale = $\frac{m_1+m_2}{2}$
$\tilde{q}\tilde{g}$
$\tilde{q}\tilde{q}^*$
$\tilde{g}\tilde{g}$
$\tilde{q}\tilde{q}$
Cross Sections in pb
100
10
1
0.1
0.01
0.001
600
800
1000
1200
1400
1600
1800
2000
$m_{\tilde{g}}$ (GeV)

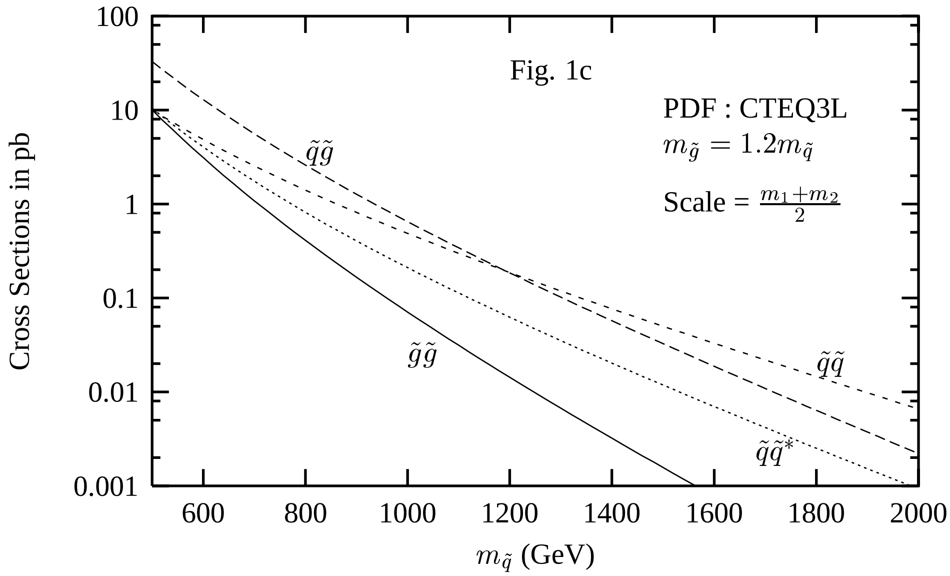
Fig. 1c
PDF : CTEQ3L
$m_{\tilde{g}} = 1.2 m_{\tilde{q}}$
Scale = $\frac{m_1+m_2}{2}$
$\tilde{q}\tilde{g}$
$\tilde{g}\tilde{g}$
$\tilde{q}\tilde{q}$
$\tilde{q}\tilde{q}^*$
Cross Sections in pb
100
10
1
0.1
0.01
0.001
600
800
1000
1200
1400
1600
1800
2000
$m_{\tilde{q}}$ (GeV)