Light Scalar Top Quarks and Supersymmetric Dark Matter

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Abstract

A stable neutralino $\tilde{\chi}_1^0$, assumed to be the lightest supersymmetric particle, is a favored particle physics candidate for the cosmological Dark Matter. We study co–annihilation of the lightest neutralino with the lighter scalar top quark $\tilde{t}_1$. We show that for natural values of the neutralino mass, $\lesssim 300$ GeV, the $\tilde{\chi}_1^0-\tilde{t}_1$ mass difference has to exceed $\sim 10$ to $30$ GeV if $\tilde{\chi}_1^0$ is to contribute significantly to the Dark Matter. Scenarios with smaller mass splitting, where $\tilde{t}_1$ is quite difficult to detect at collider experiments, are thus cosmologically disfavored. On the other hand, for small $\tilde{t}_1-\tilde{\chi}_1^0$ mass splitting, we show that co–annihilation allows very large neutralino masses, $m_{\tilde{\chi}_1^0} \sim 5$ TeV, without “overclosing” the Universe.
There is convincing evidence [1] that most matter in the Universe is dark (non–
luminous),

$$0.2 \lesssim \Omega_{DM} \leq 1$$  \hspace{1cm} (1)

where $\Omega_{DM}$ is the Dark Matter (DM) density in units of the critical density, so that $\Omega = 1$

is the mass. On the other hand, analyses of Big Bang nucleosynthesis [2] imply that most DM is non–baryonic (although dark baryons probably exist as well).

One of the favorite particle physics candidates for DM is the lightest neutralino $\tilde{\chi}^0_1$, assumed to be the lightest supersymmetric particle (LSP). It is stable if R-parity is conserved [4]; this is also a sufficient (although not necessary) condition for avoiding very fast nucleon decay in supersymmetric theories. The LSP makes an attractive DM candidate since the primary motivation for its introduction comes from particle physics arguments [5]: supersymmetry stabilizes the huge hierarchy between the weak and Grand Unification scales against radiative corrections, and if it is broken at a sufficiently high scale, it allows to understand the origin of the hierarchy in terms of radiative breaking of the Standard Model (SM) electroweak $SU_L(2) \times U(1)_Y$ gauge symmetry; furthermore it allows for a consistent unification of the gauge couplings.

Supersymmetric contributions to DM then come as extra bonus, and for wide regions of parameter space, the LSP relic density falls in the preferred range eq.(1). This is true in particular if the LSP is mostly a superpartner of the U(1)$_Y$ gauge boson, i.e. bino–like, and if both $m_{\tilde{\chi}^0_1}$ and the masses of SU(2) singlet scalar leptons fall in the natural range below a few hundred GeV [6] (but above [7] the mass range excluded by the LEP experiments).

The previous statement assumes that $\tilde{\chi}^0_1\tilde{\chi}^0_1$ annihilation reactions are the only processes that change the number of superparticles at temperatures around $T_F \simeq m_{\tilde{\chi}^0_1}/20$, where the neutralino $\tilde{\chi}^0_1$ decouples from the plasma of SM particles. It has been known for some time [8] that this is not true if the mass splitting between the LSP and the next–to–lightest supersymmetric particle $\tilde{P}$ is small. In this case, reactions of the type

$$\tilde{\chi}^0_1 + X \leftrightarrow \tilde{P} + Y,$$  \hspace{1cm} (2)

where $X, Y$ are SM particles, occur much more frequently at a temperature $T \sim T_F$ than $\tilde{\chi}^0_1\tilde{\chi}^0_1$ annihilation reactions do. The rate of the latter kind of process is proportional to two powers of the Boltzmann factor $\exp(-m_{\tilde{\chi}^0_1}/T_F) \approx \exp(-20)$, whereas for $m_{\tilde{\chi}^0_1} \simeq m_{\tilde{P}}$ the rate for reaction (2) is linear in this factor. These reactions will therefore maintain relative equilibrium between the states $\tilde{\chi}^0_1$ and $\tilde{P}$ until long after all superparticles decouple from the Standard Model plasma.

The total number of superparticles can then not only be changed by $\tilde{\chi}^0_1\tilde{\chi}^0_1$ annihilation, but also by the “co–annihilation” processes

$$\tilde{\chi}^0_1 + \tilde{P} \leftrightarrow X + Y \text{ and } \tilde{P} + \tilde{P}^{(*)} \leftrightarrow X + Y.$$  \hspace{1cm} (3)

Eventually all particles $\tilde{P}$ and $\tilde{P}^*$ will decay into $\tilde{\chi}^0_1$ (plus SM particles). In order to compute today’s LSP relic density, we therefore only have to solve the Boltzmann equation for the sum $n_{SUSY}$ of densities $n_i$ of all relevant species of superparticles. In this sum
contributions from reactions (2) cancel, since they do not change the total number of superparticles. One thus has

$$\frac{dn_{\text{SUSY}}}{dt} = -3Hn_{\text{SUSY}} - \sum_{i,j} \langle \sigma_{ij} v \rangle \left( n_i n_j - n_{eq}^i n_{eq}^j \right)$$

$$= -3Hn_{\text{SUSY}} - \langle \sigma_{\text{eff}} v \rangle \left( n_{SUSY}^2 - n_{SUSY}^{eq2} \right). \quad (4)$$

Here, $H$ is the Hubble parameter, $\langle \ldots \rangle$ denotes thermal averaging, $v$ is the relative velocity between the two annihilating superparticles in their center–of–mass frame, and the superscript “eq” indicates the equilibrium density. In the second step we made use of the fact that, as argued above, all relevant heavier superparticles maintain relative equilibrium to the neutralino LSP until long after the temperature $T_F$. This allowed us to sum all superparticle annihilation processes into an “effective” cross section; schematically

$$\sigma_{\text{eff}} \propto g_{\tilde{\chi}^0} \sigma(\tilde{\chi}^0_{\tilde{\chi}^0}) + g_{\tilde{\chi}^0} B_{\tilde{P}} \sigma(\tilde{\chi}^0_1 \tilde{P}) + g_{\tilde{P} \tilde{P}} (B_{\tilde{P}})^2 \sigma(\tilde{P} \tilde{P}^{(*)}). \quad (5)$$

where the $g_{ij}$ are multiplicity factors, and

$$B_{\tilde{P}} = (m_{\tilde{P}}/m_{\tilde{\chi}^0})^{3/2} e^{-(m_{\tilde{P}}-m_{\tilde{\chi}^0})/T} \quad (6)$$

is the temperature dependent relative Boltzmann factor between the $\tilde{P}$ and $\tilde{\chi}^0_1$ densities. The final LSP relic density $\Omega h^2$, where $h = 0.65 \pm 0.15$ is the scaled Hubble constant, is then essentially inversely proportional to $\langle \sigma_{\text{eff}} v \rangle$ at $T_F \simeq m_{\tilde{\chi}^0_1}/20$. Co–annihilation can therefore reduce the LSP relic density by a large factor, if $\delta m \equiv m_{\tilde{P}} - m_{\tilde{\chi}^0_1} \ll m_{\tilde{\chi}^0_1}$ and $\sigma(\tilde{\chi}^0_1 \tilde{P}) + \sigma(\tilde{P} \tilde{P}^{(*)}) \gg \sigma(\tilde{\chi}^0_{\tilde{\chi}^0})$. This is true in particular if $\tilde{\chi}^0_1$ is a light, $m_{\tilde{\chi}^0_1} < M_W$, higgsino or SU(2)–gaugino. More recently it has been pointed out that co–annihilation with light sleptons can reduce the relic density of a bino–like LSP by about one order of magnitude.

In this letter we study co–annihilation of neutralinos with the lighter scalar top (stop) eigenstate $\tilde{t}_1$. Compared to the other squarks, $m_{\tilde{t}_1}$ is reduced by contributions of the large top quark Yukawa coupling to the relevant renormalisation group equations, as well as by mixing between SU(2) doublet and singlet stops. While we do not know of any model that predicts $m_{\tilde{t}_1} \simeq m_{\tilde{\chi}^0_1}$, a close mass degeneracy is possible in many models, e.g. in the popular minimal Supergravity (mSUGRA) model. Moreover, scenarios with small $\tilde{t}_1$–$\tilde{\chi}^0_1$ mass splitting are of great concern for experimenters, since $\tilde{t}_1$ decays then release little visible energy, making $\tilde{t}_1$ production very difficult to detect at both $e^+e^-$ and hadron colliders.

In contrast to the cases mentioned earlier, for $\tilde{P} = \tilde{t}_1$ it is not entirely obvious that reactions of the type (2) will indeed be much faster than $\tilde{\chi}^0_1 \tilde{\chi}^0_1$ annihilation processes. In the absence of flavor mixing, one would have to chose $X = W, Y = b$ or vice versa. However, for a temperature $T < M_W$, the $W$–density is itself quite small, so reaction (2) would be much faster than $\tilde{\chi}^0_1 \tilde{\chi}^0_1$ annihilation only for $m_{\tilde{\chi}^0_1}$ significantly above $M_W$. On the other hand, most supersymmetric models predict some amount of flavor mixing in...
the squark sector, even if it is absent at some high energy scale. As a result, for small $\delta m$ the dominant $\tilde{t}_1$ decay mode is usually its flavor changing 2–body decay into $\tilde{\chi}^0 + c$ \cite{11,12,13}. For $\tilde{t}_1$ masses of current experimental interest the dominant contribution to \(\tilde{t}_1 \tilde{t}_1\) therefore comes from $X = c, Y = \text{nothing}$, i.e. (inverse) $\tilde{t}_1$ decay. If the effective $c \tilde{t}_1 \tilde{\chi}^0$ coupling is suppressed by a small mixing angle $\epsilon$, the condition that \(\tilde{t}_1 \tilde{t}_1\) is much faster than $\tilde{\chi}^0 \tilde{\chi}^0$ annihilation reads

$$\epsilon^2 e^{-\delta m/T_F} \gg \alpha e^{-m_{\tilde{\chi}^0}/T_F}, \quad (7)$$

where the extra factor of $\alpha \sim 0.01$ occurs since we are comparing $2 \leftrightarrow 1$ reactions with $2 \leftrightarrow 2$ processes. For $\delta m \lesssim T_F \sim m_{\tilde{\chi}^0}/20$ we then only need $\epsilon > e^{-10} \approx 5 \cdot 10^{-5}$. In what follows we will assume that this is true, or that $\tilde{\chi}^0$ is sufficiently heavy that $\tilde{\chi}^0 + W^+ \leftrightarrow \tilde{t}_1 + \tilde{b}$ is fast.

Another property of the stop is that it has strong interactions. A leading order calculation of $\sigma(\tilde{\chi}^0 \tilde{\chi}^0)$ and $\sigma(\tilde{t}_1 \tilde{t}_1)$ will therefore not be very reliable. Unfortunately a full higher order calculation is highly nontrivial, since one would need to include finite temperature effects (e.g. in order to cancel Coulomb singularities in the non–relativistic limit). We expect these unknown higher order QCD corrections to be more important than the contributions of higher partial waves. In the calculation of the cross sections $\sigma(\tilde{\chi}^0 \tilde{\chi}^0)$ and $\sigma(\tilde{t}_1 \tilde{t}_1)$ we therefore only include the leading, $S$–wave contribution; however, the $P$–wave contributions to $\tilde{\chi}^0 \tilde{\chi}^0$ annihilation process \cite{11,12,13} are included. Our co–annihilation cross sections will thus only be accurate to a factor of 2 or so. Due to the exponential dependence of $\sigma_{\text{eff}}$ on $\delta m$, see eqs. (13,14), the bounds on the $\tilde{t}_1 - \tilde{\chi}^0$ mass splitting that will be inferred from upper or lower bounds on $\Omega \tilde{\chi} b^2$ should nevertheless be fairly accurate.

The existence of unknown, but probably large, higher order corrections also means that we can ignore all $\tilde{t}_1$ annihilation reactions that involve more than the minimal required number of electroweak gauge couplings. However, we treat the top and bottom quarks Yukawa couplings on the same footing as the strong coupling (the latter Yukawa coupling will be large only for $\tan \beta \sim m_t/m_b$). Altogether we therefore computed the cross sections for the following processes:

$$\begin{align*}
\tilde{\chi}^0 \tilde{t}_1 & \rightarrow t \ g, \ t \ H^0_1, \ \ b \ H^+; \\
\tilde{t}_1 \tilde{t}_1 & \rightarrow t \ t; \\
\tilde{\chi}^* \tilde{t}_1 & \rightarrow g \ g, \ H^0_1 H^0_j, \ H^+ H^-, \ \ b \ \tilde{b}, \ \ t \ \tilde{t}, \quad (8)
\end{align*}$$

where $H^0 \equiv h, H, A$ is one of the three neutral Higgs bosons of the minimal supersymmetric Standard Model (MSSM) \cite{16}. The cross sections for $\tilde{\chi}^0 \tilde{t}_1$ and $\tilde{t}_1 \tilde{t}_1$ annihilation are identical to those in the first and second lines of eq.(8), respectively. We have performed two independent calculations of these cross sections. One calculation was based on trace techniques and the usual polarization sum for external gluons; here the non–relativistic limit (to extract the $S$–wave contribution \cite{11,12,13}) was only taken at the end. The second method uses helicity amplitudes \cite{13}; in this case the non–relativistic limit can already be taken at the beginning of the calculation. [Note that the cross sections for $\tilde{t}_1 \tilde{t}_1^* \rightarrow H^0_1 g$ vanish in this limit.] Explicit expressions for these cross sections will be published elsewhere. [Note that in their analytical expressions, the authors of Refs.\cite{11,12,13} do not keep...
the mass of the relevant SM fermion $f$; these terms are irrelevant for their case $f = \tau$, whereas we have to keep a finite value for the top quark mass, $m_t \neq 0$. Ref.\cite{11} also did not include $f_L - f_R$ mixing, which in our case is crucial for obtaining a light $\tilde{t}_1$. In the relevant limit we agree with Ref.\cite{11}. We disagree with Ref.\cite{12} for the $tg$ (or $\tau\gamma$) final state, by a factor $m_{\tilde{\chi}_0^0}/\sqrt{\tau}$.

Our calculation of the relic density closely follows ref.\cite{8}. In particular, we keep the temperature dependence \begin{equation} \rho_{\text{ann}}(T) \end{equation} of $\sigma_{\text{ann}}$ when computing the “annihilation integral” (essentially the integral of the Boltzmann equation for $T > T_F$).

We use a variant of the minimal Supergravity model \cite{5} for our numerical analysis. In particular, we assume a common gaugino mass, a common sfermion mass $m_0$, and a common trilinear soft breaking parameter $A_0$ at the Grand Unification scale $M_X = 2 \cdot 10^{16}$ GeV. However, we allow the soft breaking masses of the two Higgs doublets to differ from $m_0$. In practice, this means that we keep the higgsino mass parameter $\mu$ and the mass $m_A$ of the CP–odd Higgs boson as free parameters at the weak scale. The final free parameter is the ratio $\tan\beta$ of vacuum expectation values of the two Higgs fields.

For illustration, we take $\mu = -2M_2$, where $M_2 \simeq 2m_{\tilde{\chi}_1^0}$ is the SU(2) gaugino mass. This implies that the LSP is bino–like, which is the most natural choice for this type of model \cite{17}. It is also conservative, since a higgsino–like LSP will have larger couplings to the (s)top, and hence even larger co–annihilation cross sections eq.(8). We also chose a large sfermion mass, $m_0 = 2M_2$. In the absence of co–annihilation this choice is usually incompatible \cite{8} with the upper bound on the LSP relic density, which we conservatively take as $\Omega_{\tilde{\chi}_1^0} h^2 \lesssim 0.5$.

In Fig. 1 we show contours of constant $\Omega_{\tilde{\chi}_1^0} h^2$ in the $(m_{\tilde{\chi}_1^0}, \delta m)$ plane for a scenario with moderate $\tan\beta$ and a very heavy Higgs spectrum, $m_A = 5M_2$. This latter choice implies that the only Higgs boson relevant for the calculation of the LSP relic density is the light CP–even scalar $h$, with mass $m_h \leq 130$ GeV. This is a conservative scenario in the sense that it minimizes the number of final states contributing in eqs.(8), and also leads to a small $\tilde{\chi}_1^0\tilde{\chi}_1^0$ annihilation cross section. We see that scenarios with very large $\delta m$ are indeed excluded by the upper bound on $\Omega_{\tilde{\chi}_1^0} h^2$. The peak in the contour $\Omega_{\tilde{\chi}_1^0} h^2 = 0.5$ at $m_{\tilde{\chi}_1^0} \simeq m_t$ is due to $\tilde{\chi}_1^0 \tilde{\chi}_1^0 \rightarrow \tau\tilde{\tau}$, which has a sizable $S$–wave cross section if $\tilde{t}_1$ is not too heavy and $m_{\tilde{\chi}_1^0}$ is not much above $m_t$. The much smaller bumps at $m_{\tilde{\chi}_1^0} \simeq 130$ GeV are due to $hh$ final states becoming accessible.

On the other hand, for very small values of $\delta m$ and $m_{\tilde{\chi}_1^0}$ in the range indicated by naturalness arguments ($\lesssim 0.3$ TeV, corresponding to a gluino mass $m_{\tilde{g}} \lesssim 2$ TeV), we find that the LSP cannot contribute significantly to the solution of the Dark Matter puzzle, since its relic density is too small. In particular, one needs $\Omega_{\tilde{\chi}_1^0} h^2 > 0.025$ for $\tilde{\chi}_1^0$ to form galactic haloes. We see that even for the present very conservative choice of parameters one needs a $\tilde{t}_1\tilde{t}_1$ mass splitting of at least 9 to 19 GeV (6 to 10%) to satisfy this lower bound on $\Omega_{\tilde{\chi}_1^0} h^2$. This mass splitting is large enough for standard $\tilde{t}_1$ search methods at $e^+e^-$ colliders \cite{13,19} to have reasonably high efficiency. If we require that $\Omega_{\tilde{\chi}_1^0} h^2$ lies in the currently favored “best fit” range between about 0.1 and 0.2, $\delta m$ has to be between 11 and 33 GeV. Unfortunately this is still not high enough for current $\tilde{t}_1$ search strategies at the Tevatron \cite{14} to be sensitive.
Figure 1: Contours of constant $\Omega \tilde{\chi}^2 = 0.5$ (solid), 0.1 (dashed) and 0.025 (dotted) in the $(m_{\tilde{\chi}_1^0}, \delta m)$ plane, where $\delta m = m_{\tilde{t}_1} - m_{\tilde{\chi}_1^0}$. We took $\mu$, $m_0$ and $m_A$ to be fixed multiples of $M_2 \simeq 2m_{\tilde{\chi}_1^0}$, as indicated, whereas $\tan \beta = 10$ has been kept fixed. The parameter $A_0$ varies between about 2.5$m_0$ and 3.2$m_0$, with larger $A_0$ values corresponding to smaller values of $\delta m$.

So far we have focused on LSP masses in the range favored by naturalness arguments. It is sometimes claimed [18] that the upper bound on $\Omega \tilde{\chi}^2$ implies that LHC experiments must find superparticles if the MSSM is correct and $\tilde{\chi}_1^0$ is bino–like. Unfortunately this is not true; for $\delta m \to 0$ an LSP mass up to 4 TeV, corresponding to a gluino mass in excess of 20 TeV, cannot be excluded from this cosmological argument. [As noted above, our estimates for $\tilde{t}_1$ annihilation cross sections are not very reliable. However, even if we over–estimated them by a factor of 2, the bound on $m_{\tilde{\chi}_1^0}$ would only be reduced by a factor of $\sqrt{2}$, and would thus still allow sparticle masses far above the range to be covered by the LHC.]

In Fig. 2 we show analogous results for a light spectrum of Higgs bosons and large $\tan \beta$, where the bottom Yukawa coupling is sizable; the choice $m_A = 0.35M_2 \simeq 0.7m_{\tilde{\chi}_1^0}$ ensures that all Higgs pair final states will be accessible for $m_{\tilde{\chi}_1^0} > 100$ GeV. However, we keep the previous (large) values for $|\mu|$ and $m_0$. Nevertheless we see that for natural values of $m_{\tilde{\chi}_1^0}$, requiring $\Omega \tilde{\chi}^2 > 0.025$ now implies $\delta m > 20$ GeV. Moreover, the LSP makes a good DM candidate, i.e. $\Omega \tilde{\chi}^2 \sim 0.1$, only for $\delta m \gtrsim 40$ GeV; this is sufficiently large to permit $\tilde{t}_1$ searches at the Tevatron [14, 20]. Finally, for $\delta m \to 0$, cosmology now allows an LSP mass up to 6 TeV, corresponding to a gluino mass of about 30 TeV! Obviously the upper bound on $\delta m$ that follows from $\Omega \tilde{\chi}^2 > 0.025$ for natural values of $m_{\tilde{\chi}_1^0}$, as well
as the absolute upper bound on $m_{\tilde{\chi}_1^0}$ that follows from $\Omega_{\chi} h^2 \leq 0.5$, are even higher if we chose smaller values for $m_0$ and/or $|\mu|$.

Figure 2: As in Fig. 1, except that we took a large value of $\tan\beta$ and a light Higgs boson spectrum. The regions below and to the left of the heavy dotted lines are excluded by Higgs boson searches at LEP.

In conclusion, we have shown that scenarios with very small $\tilde{t}_1 - \tilde{\chi}_1^0$ mass splitting would permit an LSP mass of several TeV without “overclosing” the Universe. This shows once again [6] that the upper bound on the LSP relic density does not guarantee that LHC experiments will detect superparticles, even if the MSSM is correct; of course, (third generation) superparticles with masses out of the reach of the LHC can hardly be argued to be “natural”. On the other hand, for $\tilde{\chi}_1^0$ and $\tilde{t}_1$ masses of present experimental interest, and indeed for the entire natural range of these masses, $\tilde{\chi}_1^0$ cannot contribute significantly to the Dark Matter in the Universe unless the $\tilde{t}_1 - \tilde{\chi}_1^0$ mass difference is large enough for conventional $\tilde{t}_1$ search strategies at $e^+e^-$ colliders to be effective. This does not imply that collider searches for $\tilde{t}_1$ nearly degenerate with $\tilde{\chi}_1^0$ [21] should not be continued; a positive signal would definitely exclude $\tilde{\chi}_1^0$ as DM candidate, which is not easy to accomplish with cosmological Dark Matter searches. However, since Dark Matter is known to exist, for natural values of $m_{\tilde{\chi}_1^0}$ a very small $\tilde{t}_1 - \tilde{\chi}_1^0$ mass splitting would require physics beyond the MSSM.

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