Supersymmetric Singlet Majorons and Cosmology

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Abstract

We examine cosmological constraints on the lepton number breaking scale in supersymmetric singlet majoron models. Special attention is drawn to the model dependence arising from the particular choice of a certain majoron extension and a cosmological scenario. We find that the bounds on the symmetry breaking scale can vary substantially. Large values of this scale can be allowed if the decoupling temperature of smajoron and majorino exceeds the reheating temperature of inflation. In the opposite case an upper bound depending on the majoron model can be obtained which, however, is unlikely to be much larger than $10^{10}$ GeV.
Singlet majoron models [1] provide the simplest extension of the standard model with spontaneously broken lepton number and neutrino masses generated via the celebrated see-saw mechanism. Therefore these models are of particular interest for cosmology. The well-known cosmological constraint on the conventional singlet majoron model comes from the late-decay of heavy neutrinos into majorons ($J$) and lighter neutrinos. In the model with one singlet $S$ the amplitude for this decay has a strong suppression [2]. Then the standard cosmological constraint derived from the closure density of the universe [3] bounds the right handed neutrino (RHN) mass to be $< 10^3$ GeV taking the largest possible neutrino mass of 30 MeV and a mixing angle of order $10^{-4}$. For generic models with an arbitrary number of singlets the bound can be shifted to $\sim 10^9$ GeV. Clearly, for decreasing neutrino masses these bounds become weaker and for masses below $\sim 40$ eV no information about the RHN scale can be obtained at all.

In supersymmetric models, cosmological considerations provide independent informations due to the presence of the superpartners of the majoron (the smajoron $\sigma$ and the majorino $\psi$) as discussed first by Mohapatra and Zhang [4]. Because their coupling is suppressed by the lepton number symmetry breaking scale $V_L$ they can be expected to decouple when relativistic. In comparison to supersymmetric axion models [5] the partners of the majoron tend to couple much weaker to the standard particles due to the lack of an anomaly and the smallness of the neutrino masses. One might therefore ask if the constraints on $V_L$ from baryogenesis and overclosure will be more restrictive than corresponding constraints on axion models. We consider the range from $\sim 10^{10}$ GeV to $10^{12}$ GeV as particulary attractive for $V_L$ since several independent hints from the MSW solution of the solar neutrino problem, the hidden sector SUSY breaking and axion physics point to such scales. A question of particular interest might therefore be to what extent such large values of $V_L$ are compatible with cosmological restrictions on SUSY majoron models.

In the below we examine this matter taking two kinds of model dependences into account. On the one hand we discuss the freedom in the field theoretical parameters relevant for cosmology like masses and couplings of the smajoron and the majorino. On the other hand we analyze the influence of the cosmological scenario, namely the influence of inflation which for large $V_L$ may be able to wash out the relic densities
We start by briefly introducing the class of models we are considering here. Their superpotentials which come in addition to that of the MSSM can be written as

$$W = h_{ij} L_i H N_j + f_{ija} N_i N_j S_a + W(S_a)$$

with three families $N_i$ of RHNs and gauge singlet fields $S_a$. For most of our discussions soft breaking terms of the standard type

$$V_{soft} = m_{3/2}^2 \sum_a |\phi_a|^2 + m_{3/2} \left( \sum_a \phi_a \frac{\partial W}{\partial \phi_a} + (A - 3)W + \text{h. c.} \right)$$

are assumed. We have in mind that supergravity is broken in the hidden sector. The whole potential is invariant under $U_L(1)$ with $Q_L(L) = 1$, $Q_L(N) = -1$ and $Q_L(S_a) = q_a$, which is spontaneously broken by the VEVs $v_a$ of the gauge singlet fields $S_a$. Apart from being larger than $m_{3/2}$ the $U_L(1)$ symmetry breaking scale $V_L = (\sum |v_a|^2 q_a^2)^{1/2}$ is taken as a free parameter. The freedom for model building consists in the number of singlet fields $S_a$ and the choice of their potential $W(S_a)$. In this paper we assume that R-parity is unbroken.

For cosmological considerations we extract now the model-dependence of masses and lifetimes of smajoron and majorino. Let us first examine the interactions of the majoron supermultiplet $\Phi = ((\sigma + iJ)/\sqrt{2}, \psi)$. The interaction with the standard particles is given by the usual form

$$W = h L H N + \frac{M}{2} N N + \frac{f}{2} N N \Phi$$

with the generation indices suppressed. We also need to know higher-dimensional effective interactions. The one most important for our purpose comes from a one loop diagram built up from the interactions $L H N$ and $\Phi N N$. It leads to the effective $D=5$ operator $(\Phi + \bar{\Phi}) L \bar{L}$. Putting the fields on shell the diagram vanishes so that there is no decay mode of $\Phi$ with a $1/V_L$ suppression. However, attaching a gauge field $V$ gives the coupling

$$\frac{g h^2}{16\pi^2 V_L} (\Phi + \bar{\Phi}) L \bar{L} V.$$
Here \( h \) denotes the dominant one of Yukawa couplings \( h_{ij} \) in eq. (1). This diagram can keep the majoron supermultiplet in equilibrium after the phase transition.

To describe the interactions among the components of \( \Phi \) which generally depend on the form of \( W(S_a) \), the nonlinear approach is easier to handle since the majoron and the majorino are completely rotated away from the superpotential as well as from the soft terms. Only the Kähler part needs to be considered. Starting with the minimal Kählerian \( K = \sum_a S_a \bar{S}_a \) and writing \( S_a = v_a \exp(q_a \Phi / V_L) \), we arrive at the following Lagrangian in component fields:

\[
\mathcal{L} = (1 + \frac{\sqrt{2}x}{V_L}) \sigma \left( \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma + \frac{1}{2} \partial_\mu J \partial^\mu J + i \bar{\psi} \sigma \partial^\mu \psi \right) - \frac{x}{V_L} (\partial^\mu J) \bar{\psi} \sigma \psi + (4 \text{ fermion terms}) + O(1/v^2) \tag{5}
\]

with \( x = \sum_a \frac{|v_a|^2}{V_L^2} q_a^3 \). \tag{6}

The model dependence is now condensed in the value of \( x \) which influences the decay rates for \( \sigma \to 2J \) and \( \sigma \to 2\psi \).

The value of \( x \) can be of order one or zero at tree level. Even in the second case a nonzero value can be generated by radiative corrections. To discuss this, let us take the simplest model. It is specified by the superpotential

\[
W = h_{ij} L_{\alpha i} H_{\alpha} N_j + f_{ij} N_i N_j S + \lambda (SS' - \mu^2) Y \tag{7}
\]

which results in the following soft terms for the gauge singlet fields:

\[
V_{soft} = \sum_{\phi=n,s,s'} m_\phi^2 |\phi|^2 + m_{3/2}^2 \left[ A_n f n n s + A_s \lambda s s' y - A_y \lambda \mu^2 y + \text{h. c.} \right] . \tag{8}
\]

The boundary condition to be fulfilled at \( M_{Pl} \) is \( m_\phi = m_{3/2}, A_n = A_s = A \) and \( A_y = A - 2 \). The value of \( x \) for the above potential is \( x = (v^2 - v'^2)/(v^2 + v'^2) \). Minimization leads to

\[
x = \frac{m_s^2 - m_{s'}^2}{2\lambda^2 y^2 + m_s^2 + m_{s'}^2} \sim \frac{m_s^2 - m_{s'}^2}{4m_{3/2}^2} , \tag{9}
\]

with \( y = (A_y - A_s)m_{3/2}/2\lambda \). Obviously degenerate soft masses result in \( x = 0 \) which is not surprising since in that case a permutation symmetry \( s \leftrightarrow s' \) exists. The only way to break this symmetry is to take \( m_s \neq m_{s'} \) which can be achieved in two
ways: One can assume nonuniversal mass terms at $M_{Pl}$ thereby departing from the standard soft terms (2) or one can consider renormalization group effects. The first option clearly allows to produce any value of $x$. So let us discuss the implications of RGE starting with universal boundary masses. Due to the asymmetric coupling of $s$ and $s'$ to the RHNs their masses renormalize differently and one finds \[ x \sim \frac{f^2}{64\pi^2}(1 + A^2) \ln \left( \frac{M_{Pl}}{V_L} \right) \]

\[ = f^2(1 + A^2) \left( 2.2 \times 10^{-2} + \frac{1}{64\pi^2} \ln \left( \frac{10^{12}\text{GeV}}{V_L} \right) \right). \tag{10} \]

So it appears to be difficult to get $x > 0.1$. The lifetime of the smajoron decay $\sigma \rightarrow 2J$ is given by

\[ \tau = 2 \times 10^{-5} \left( \frac{V_L}{x 10^{12}\text{GeV}} \right) \left( \frac{10^2\text{GeV}}{m_\sigma} \right)^3 \text{sec}. \tag{11} \]

In another model discussed in ref. \[ this only one change appears. The value of $y$ is negligible, and so the value of $x$ roughly doubles.

The majoron has mass zero forgetting about possible $U_L(1)$-breaking gravity effects. Considering this effect appearing in the D=5 operators it was found that the cosmological mass density constraint gives a strong bound on the lepton number breaking scale: $V_L < 10 \text{ TeV}$ \[ this paper we neglect this effect. The smajoron cannot escape receiving a mass $m_\sigma \sim O(m_{3/2})$ from soft terms. The majorino tree level mass is model-dependent and can take values $m_{\psi,\text{tree}} \sim O((m_{3/2}/V_L)^k m_{3/2})$ with $k = 0$ or 1 typically. A small tree level majorino mass ($k > 0$) is e.g. realized in the model of ref. \[ or in the above model for nonstandard soft terms satisfying $A_s = A_y$ at $M_{Pl}$. However, in analogy to the heavy quark axion \[ , it receives a model-independent one-loop correction $m_{\psi,\text{loop}} = (1/16\pi^2)f^2 A m_{3/2}$, where $f$ is the dominant one of the couplings $f_{ija}$. Additionally, there can arise another radiative mass due to the RG evolution of the trilinear soft couplings in such models. The induced mass depends on $f$ and the completely unrestricted couplings in the $W(S_a)$-sector \[ and is therefore quit model-dependent. In any case the majorino mass

\footnote{This model can provide a light majorino with the standard universal soft terms. But it has another zero mode than the majoron, which may cause additional cosmological impacts.}
can take values far below $m_{3/2}$. To illustrate the cosmological implications we will use the above model-independent value $m_{\psi, \text{loop}}$ and examine the cosmological effects depending on the value of $f$.

When the temperature of the universe falls below the scale $V_L$, smajoron and majorino can be kept in equilibrium via effective interactions with light fields. Considering the D=5 effective interaction in eq. (4) smajoron and majorino get out of equilibrium below the temperature

$$T_L = 6.2 \times 10^{17} f_1^{-2} \sqrt{g_*/100} \left( \frac{10 \text{eV}}{m_\nu} \right)^2 \text{GeV}. \quad (12)$$

Hereafter we introduce the notation $f_n = 10^n f$. The actual decoupling temperature is the lower one between $V_L$ and $T_L$. From this we conclude that smajoron and majorino decouple when they are relativistic. Such hot relics can cause the problem of the excessive energy density if the particles are late-decaying. In regards of inflation this problem should be examined in two ways. The usual see-saw mechanism assumes a large scale $V_L \sim 10^{12} \text{GeV}$ for the lepton number breaking. It is bigger than the reheating temperature $T_R$ which should be less than about $10^{10} \text{GeV}$ in order to avoid the gravitino problem [11]. In such a case ($V_L$ and $T_L > T_R$) the relic density of majorino or smajoron can be washed away. However, there can be a significant amount of regenerated relics after the reheating. We will later consider the effects of regeneration.

It can also happen that smajoron and majorino decouple after inflation ($V_L$ or $T_L < T_R$). For this case we note that the crucial model-dependence of the cosmological problem resides, especially, on the values of $x$ and $m_\psi$. We begin with the discussion of this case. The ratio of the relic density to the entropy density is then of order $10^{-3}$. In the case that the smajoron decays into two majorons, the decay-produced majorons should be red-shifted away sufficiently in order not to upset the standard nucleosynthesis prediction. The majorons from smajoron decay should not supply an energy density larger than three tenths of the density of one neutrino species, which implies

$$\left( \frac{m_\sigma}{10^2 \text{GeV}} \right) \left( \frac{\tau_\sigma}{1 \text{sec}} \right)^{1/2} < 8.1 \times 10^{-5}. \quad (13)$$
From the lifetime given in eq. (11) one finds
\[ V_L < 1.8 \times 10^{10} x \left( \frac{m_\sigma}{10^2 \text{GeV}} \right)^{1/2} \text{GeV}. \] (14)

Thinking of models with nonuniversal scalar masses one may expect \( x = O(1) \) so that it seems difficult to push \( V_L \) higher than \( 10^{10} \) GeV. Taking the universal boundary conditions for the soft terms the value of \( x \) depends on the Yukawa coupling \( f \) between majoron superfield and RHNs. For the two explicit models we have mentioned here the situation can be summarized as
\[ V_L < 4 \times 10^8 f^2 (1 + A^2) \left( \frac{m_\sigma}{10^2 \text{GeV}} \right)^{1/2} \text{GeV}. \] (15)

When the value of \( f \) becomes extremely small, the decay channel of smajoron into two majorons is suppressed and those into two neutrinos or two sneutrinos become important. This is the region of
\[ m_\nu > 7.5 f^2 \left( \frac{m_\sigma}{10^2 \text{GeV}} \right)^{1/2} \text{eV}, \]
where the constraint in ref. [4] is applicable.

A more severe constraint comes from the decay of majorinos. Being R-parity odd, each majorino will eventually produce at least one LSP (ordinary neutralino \( \chi^0 \)) whose mass in generally exceeds the GeV-range. Then those relics of secondary LSP cause an overclosure of the universe. To avoid this problem the majorino should decay before the decoupling time of neutralinos which is about \( 10^{-6} (20 \text{ GeV}/m_{\chi^0})^2 \) sec. If the majorino is heavier than the sneutrino the decay \( \psi \to \nu \bar{\nu} \) produces a strong bound :
\[ V_L < 10^2 \left( \frac{m_\nu}{10^2 \text{eV}} \right)^{1/2} \left( \frac{m_\psi}{10^2 \text{GeV}} \right)^{1/2} \left( \frac{20 \text{ GeV}}{m_{\chi^0}} \right) \text{GeV}. \] (16)

One can also find a bound of a similiar order of magnitude if the decay channel \( \psi \to \nu \nu \tilde{H}^0 \) [4] is allowed.

Of course the majorino can be the LSP itself in which case it is a warm dark matter candidate. Such a majorino has to fulfill the overclosure bound \( m_\psi < 2 \) keV [3] which translates to a restriction on \( f \)
\[ f < 1.8 \times 10^{-3} \left( \frac{10^2 \text{GeV}}{m_{3/2}} \right)^{1/2} A^{-1/2}. \] (17)
The model has now to be such that the tree level mass is below $\sim 2$ keV. Assuming that $x$ comes entirely from RG effects eqs. (17, 15) give
\[ V_L < 1.3 \times 10^3 \left( \frac{m_\sigma}{10^2 \text{ GeV}} \right)^{1/2} \left( \frac{10^2 \text{ GeV}}{m_{3/2}} \right) \frac{1 + A^2}{A} \text{GeV} \] (18)
which is quite restrictive. It should, however, be stressed that this bound is influenced by the magnitude of RG effects and is therefore model dependent. Taking a model with a light majorino and nonuniversal soft masses it would be completely invalidated and provided $f$ satisfies eq. (17) the strongest constraint on $V_L$ would be given by eq. (14).

Let us now consider the case that smajoron and majorino decouple before or during inflation ($V_L$ and $T_L > T_R$). As already mentioned our concern is the regeneration of relics. In terms of the ratio of relic density to entropy density $Y_X = n_X/s$ the Boltzmann equation of the regenerated population for a particle $X$ becomes
\[ \frac{dY_X}{dT} = -\frac{\langle \Sigma_T v \rangle n_r^2}{sHT} \] (19)
where we neglected the decay terms. We can explicitly integrate the Boltzmann equation from the reheating temperature $T_R$ to some temperature $T$ with an initial condition $Y(T_R) = 0$. If $T_R \gg T$, we get
\[ Y_X(T) \simeq \frac{\langle \Sigma_T v \rangle n_r (T_R)^2}{H(T_R)s(T_R)} = 5.9 \times 10^{-6} \langle \Sigma_T v \rangle M_{Pl} T_R. \] (20)
We see that $Y_X$ is determined by $\langle \Sigma_T v \rangle$ and $T_R$. For our case, the D=5 interaction in eq. (4) is most important for the regeneration and $\langle \Sigma_T v \rangle \simeq g^2 h^4 / 16^3 \pi^5 V_L^2$. Then we obtain the regenerated population for smajoron or majorino:
\[ Y_{\sigma,\psi} \simeq 6.5 \times 10^{-11} f_l^2 \left( \frac{m_\psi}{10 \text{ eV}} \right)^2 \left( \frac{T_R}{10^{10} \text{ GeV}} \right). \] (21)

We first consider the relevant constraints coming from the regeneration of majorinos. If the majorino is stable or its lifetime $\tau_\psi$ is longer than the age of the universe $t_0$, the relation
\[ Y_\psi \left( \frac{m_\psi}{10^2 \text{ GeV}} \right) < 3.55 \times 10^{-11} \] (22)
has to be fulfilled in order not to overclose the universe. This puts a bound on the neutrino mass

$$m_\nu < 17 f_1^{-2} \left( \frac{10^{10} \text{GeV}}{T_R} \right)^{1/2} \left( \frac{20 \text{ GeV}}{m_\psi} \right)^{1/2} \text{eV}.$$  \hspace{1cm} (23)

If the majorino is unstable and $\tau_\psi < t_0$, the overclosure density bound is applied to the decay-produced neutralinos. For this we just replace $m_\psi$ by $m_{\chi^0}$ in the above equation.

There is another constraint coming from the high-energy neutrinos when the majorino (smajoron) is unstable. The decay of smajoron and majorino generally produces high-energy neutrinos which can be observed by experiments. For the energy range of our interest, the best bounds come from the experimental upper limit on upward-going muons from the IMB detector [12]. Here we quote the rough order of estimation

$$Y_\psi \left( \frac{m_\psi}{10^2 \text{GeV}} \right)^2 < 10^{-18} \left( \frac{t_0}{\tau_\psi} \right)^{4/3} \text{ for } \tau_\psi < t_0$$

$$< 10^{-18} \left( \frac{\tau_\psi}{t_0} \right) \text{ for } \tau_\psi > t_0.$$  \hspace{1cm} (24)

Considering the decay of the majorino into a neutrino and a sneutrino, we find for $\tau_\psi < t_0$

$$V_L < 1.6 \times 10^{10} f_1^{-3/4} \left( \frac{m_\nu}{10 \text{ eV}} \right)^{1/4} \left( \frac{10^2 \text{ GeV}}{m_\psi} \right)^{1/4} \left( \frac{10^{10} \text{ GeV}}{T_R} \right)^{3/8} \text{GeV},$$  \hspace{1cm} (25)

and for $\tau_\psi > t_0$

$$V_L > 1.1 \times 10^{17} f_1 \left( \frac{m_\nu}{10 \text{ eV}} \right)^2 \left( \frac{m_\psi}{10^2 \text{ GeV}} \right)^{3/2} \left( \frac{T_R}{10^{10} \text{ GeV}} \right)^{1/2} \text{GeV}.$$  \hspace{1cm} (26)

Here the majorino lifetime is given by

$$\tau_\psi = 1.7 \times 10^{15} \left( \frac{V_L}{10^{12} \text{ GeV}} \right)^2 \left( \frac{10 \text{ eV}}{m_\nu} \right)^2 \left( \frac{10^2 \text{ GeV}}{m_\psi} \right) \text{sec}.$$  \hspace{1cm} (27)

The above results are summarized in the figures. Fig. 1 shows the forbidden region in the $V_L - m_\nu$ space when the decay $\psi \to \nu \tilde{\nu}$ is allowed. The lines (ii) or (iii) are very sensitive to the value of $f$ and they provide no constraint unless $f$ is larger than about $10^{-4}$. As one expects, a rather strong constraint can be found for
larger values of \( f \) and \( T_R \). We note that \( V_L \) is limited between \( T_R \) and the value in eq. (25) in order to have, e.g., a 40 eV neutrino as a dark matter candidate. If the reheating temperature is pushed up to \( 10^{10} \) GeV, the line (iv) goes down below the line (a) and we find a strong bound on the neutrino mass:

\[
m_\nu < 18 f_3^{-2/3} \left( \frac{10^2 \text{GeV}}{m_\psi} \right)^{1/2} \left( \frac{10^{10} \text{GeV}}{T_R} \right)^{1/6} \text{eV}.
\]  

The size of the constrained region becomes smaller if the lifetime of the majorino is fairly long or if it is stable. When majorino is stable, only the lines (ii) and (iii) are to be taken. As an example for a long-lived majorino which can decay into the LSP, let us take the decay \( \psi \to \nu \nu \tilde{H}^0 \) with the Higgsino \( \tilde{H}^0 \) as the LSP. The lifetime can be found in ref. [4] and together with the constraint from the high-energy neutrinos it results in the lower part of Fig. 2. The constrained region becomes smaller for larger \( T_R \) as the curve moves below the line (a). For instance, there is no region forbidden if \( T_R = 10^{10} \) GeV.

Regenerated smajorons play a less important role than majorinos. With \( x \sim 1 \) one finds no constraints since then smajorons decay fast enough. In a model with the RG-induced \( x \sim 10^{-2} f^2 \), the smajoron decay into two neutrinos can be important for small values \( f < 10^{-3} \), for which the constraint from the high-energy neutrino is the same as in eqs. (25, 26). The effect increases for smaller values of \( f \) and larger values of \( T_R \). This behavior is shown in the upper curve in Fig. 2.

In conclusion we have discussed the cosmological constraints on SUSY majoron models. The majoron superpartners decouple when relativistic and we should distinguish two cases, the decoupling temperature is lower or higher than the reheating temperature after inflation.

In the first case we identified the mass of the majorino \( m_\psi \) and the quantity \( x \) which determines the strength \( x/V_L \) of the decay \( \sigma \to 2J \) as the crucial parameters influencing the cosmological constraints. From the decay of the smajoron we obtained a bound \( V_L < x 10^{10} \) GeV where in models with large RG effects or nonuniversal scalar soft masses one can think of \( x \) as being of \( O(1) \). Assuming standard soft scalar masses, however, in our examples \( x \) was determined by small RG effects leading to typical values \( x \sim 10^{-2} f^2 \) where \( f \) is the coupling of the majoron to the
RHNs. Stronger bounds were found from the majorino decay. If the majorino can
decay into the LSP, $V_L$ is constrained to be $< O(10^2 \text{ GeV})$. However, there exist
models with a negligible tree level majorino mass and a radiative mass contribu-
tion proportional to $f$. In such models the majorino can be the LSP and to avoid
overclosure $f$ has to fulfill $f < O(10^{-3})$. If at the same time $x$ is generated by RG
effects, a small $f$ suppresses the decay $\sigma \rightarrow 2J$ which results in a typical bound
$V_L < O(10^3 \text{ GeV})$. It should be stressed that this bound is avoidable: A model with
a keV majorino and $x$ of order one would be fully consistent up to $V_L \sim 10^{10} \text{ GeV}$.

For the second case we examined the effects of the regenerated relics. The con-
strained region in the $V_L - m_\nu$ plane is strongly dependent on the values of $f$ and
the reheating temperature $T_R$ as well as the decay channel of the majorino. So one
needs more informations on the various parameters to draw some definite conclu-
sions. Capitalizing the general pattern, we find a relatively strong constraints if the
majorino decays into a sneutrino and a neutrino is allowed. In this case a higher
value of $V_L$ together with a 10 eV neutrino can be allowed for a small value of the
Yukawa coupling $f < O(10^{-3})$. If the above decay channel is not allowed, we con-
clude that a higher value of $V_L$ can be easily achieved. In either case the bounds on
heavy neutrinos tend to be more restrictive. We finish by making a remark on the
recent work [13] which showed that the supersymetric singlet majoron model can
explain both chaotic inflation and baryogenesis. In this scenario the inflaton is iden-
tified with the scalar partner of a RHN. Then the primordial density perturbation
fixes the RHN mass at $M \simeq 10^{13} \text{ GeV}$. Furthermore the reheating temperature of
order $10^{10} \text{ GeV}$ could be obtained with reasonable values of parameters. Therefore
this scenario should be such that the decoupling of majorino and smajoron occurs
before the reheating and should follow the constraints we have investigated. For
instance, assuming a Yukawa coupling $f$ of order one, the above discussion shows
that the majorino should be lighter than the sneutrino.

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Figure captions

- **Fig. 1.** For the decay $\psi \rightarrow \nu \tilde{\nu}$. The curves are normalized to $T_R = 10^8$ GeV and $f = 0.1$.
  
  - The lines (a) and (b) define the region where $V_L > T_R$ and $T_L > T_R$, respectively. The region below the line (t1) corresponds to $\tau(\psi \rightarrow \nu \tilde{\nu}) < t_0$.
  
  - Below the line (i) the yukawa coupling $h$ remains in perturbative region $h < 1$, that is,
    \[
    \frac{V_L}{10^{13} \text{GeV}} \left( \frac{m_\nu}{10 \text{eV}} \right) < 6.1/f_1 .
    \]
  
  - The lines (ii) and (iii) arise from the overclosure bound for (quasi-) stable majorinos and decay-produced neutralinos, respectively. The region left to the lines is allowed. (cf. eq. (23) in the text.)
  
  - The lines (iv) and (v) show the constraints from the decay-produced energetic neutrinos. The region below the line (iv) and above the line (v) is allowed. (cf. eqs. (25, 26))

- **Fig. 2.** For the decay $\psi \rightarrow \nu \nu \tilde{H}^0$ and $\sigma \rightarrow JJ, \nu \nu$. The curves are normalized to $T_R = 10^8$ GeV and $f = 10^{-4}$.
  
  - The lines (a), (b) and (i)–(v) are as explained in Fig. 1. The lines (ii) and (iii) are now outside the experimental limit of $m_\nu$.
  
  - The lines (s1, s2) represent $\tau(\sigma \rightarrow JJ, \nu \nu) = t_0$. The line (t2) corresponds to $\tau(\psi \rightarrow \nu \nu \tilde{H}^0) = t_0$.
  
  - The lines (vi) and (vii) show the constraints from the decay-produced energetic neutrinos as in eqs. (25, 26). The region below the line (vi) is defined by
    \[
    V_L < 1.2 \times 10^8 f_3^{-7/8} \left( \frac{m_\nu}{10 \text{eV}} \right)^{1/8} \left( \frac{m_\psi}{10^2 \text{GeV}} \right)^{7/8} \left( \frac{10^8 \text{GeV}}{T_R} \right)^{3/16} \text{GeV} .
    \]
  
  and the region above the line (vii) by
    \[
    V_L > 7.6 \times 10^8 \left( \frac{m_\nu}{10 \text{eV}} \right) \left( \frac{m_\psi}{10^2 \text{GeV}} \right)^{7/4} \left( \frac{T_R}{10^8 \text{GeV}} \right)^{1/4} \text{GeV} .
    \]
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Figure 1
Figure 2