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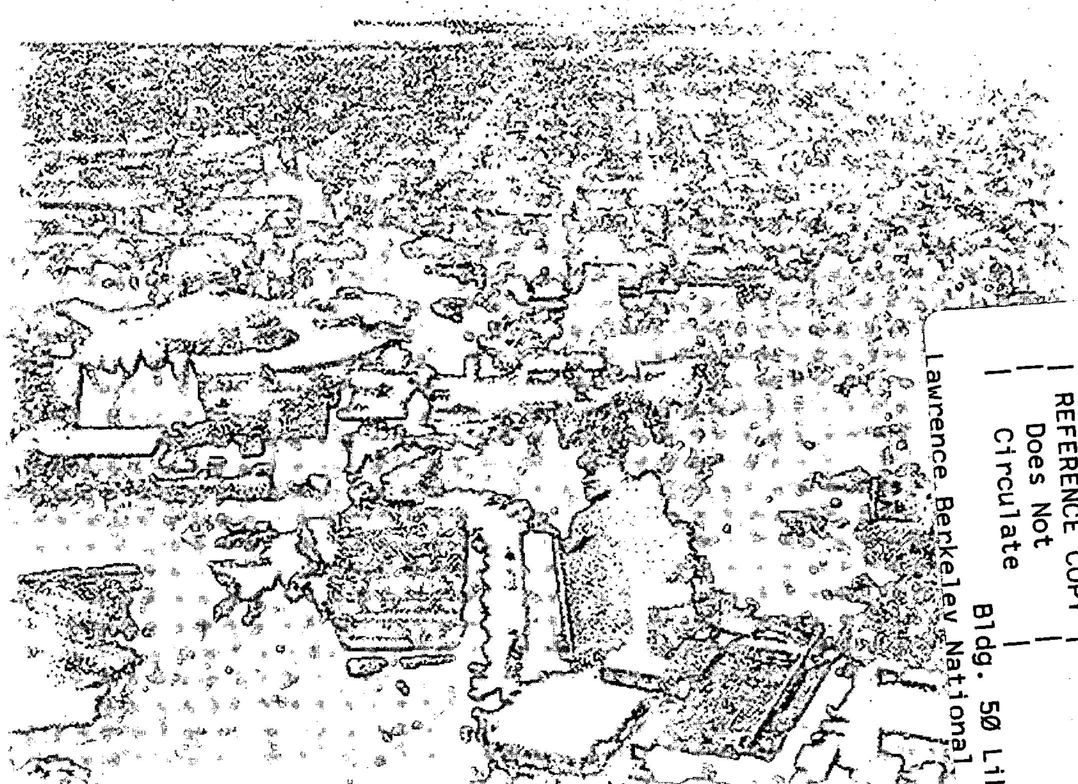

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# Sneutrino Cold Dark Matter With Lepton-Number Violation*

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**Abstract**

The tau sneutrino is proposed as a candidate for galactic halo dark matter, and as the cold dark matter (CDM) component of the universe. A lepton-number-violating sneutrino mass, $\tilde{\nu}\tilde{\nu}$, splits the tau sneutrino into two mass eigenstates: $\tilde{\nu} \rightarrow \tilde{\nu}_{\pm}$. The absence of a $Z\tilde{\nu}_{-}\tilde{\nu}_{-}$ coupling implies that the lighter mass eignestate, $\tilde{\nu}_{-}$, does not annihilate via the $s$-channel $Z$-exchange to a low cosmological abundance, and furthermore, halo sneutrinos do not scatter excessively in Ge detectors. For the majority of the relevant parameter space, the event rate in Ge detectors is $\geq 10^{-2}$ events/kg/day. The lepton number violation required for sneutrino CDM implies that the tau neutrino mass is $m_{\nu_\tau} \gtrsim 5$ MeV, large enough to be excluded by $B$ factory experiments. Events of the form $l^+l^-\not{E}$ or $jj\not{E}$, with low $m_{ll}$ or $m_{jj}$, may be observed at LEP2. A seesaw mechanism is investigated as the origin for the lepton number violation, and several other cosmological and particle physics consequences of sneutrino CDM are discussed.

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**1. Introduction.** It has been known for decades that most of the mass in the Universe is dark, *i.e.* not seen by optical methods [1]. This is deduced from studying the motion of visible objects, which is governed by the size of the gravitational force acting on them. Rotational curves of spiral galaxies and motions of galaxies in clusters are good examples, both of which indicate that most sources of gravity are not seen.

Recently, MACHOs (MAssive Compact Halo Object) were seen within the halo of the Milky Way galaxy by means of gravitational microlensing [2]. However, the determination of the MACHO mass fraction in the halo is still quite uncertain: anywhere between 10% to 100%. Moreover, the scenario of 100% MACHO fraction faces various astrophysical and cosmological difficulties (see, *e.g.*, [3]). Therefore it is quite possible that the MACHOs account for the missing dark baryons, as required by Big Bang Nucleosynthesis, but are not the dominant component of the galactic halo.

From the point of view of galaxy formation theories, and the small density fluctuations observed by COBE, the most promising candidate for the invisible source of gravity is Cold Dark Matter (CDM) [4]. Although the standard CDM model, with scale-invariant primordial density fluctuations, is not favored by the COBE data and the observed large scale structures, the small discrepancy can be accounted for by introducing a small Hot Dark Matter component [5], by "tilting" the primordial density fluctuation spectrum [6], or by introducing particles (such as $\nu_\tau$) whose decay changes the time of radiation-matter equality [7]. In all these scenarios, CDM is the dominant component of the galactic halo.

There is no CDM candidate in the standard model. On the other hand, theories of weak-scale supersymmetry are strongly motivated: they allow a symmetry description of the weak scale, they incorporate the economical description of flavor symmetry breaking by Yukawa couplings, and they successfully predict the weak mixing angle, at the percent level, from gauge coupling unification. Finally, the lightest superpartner (LSP) is a candidate for CDM.[†]

There are two obvious choices for a neutral LSP candidate for CDM: neutralinos and sneutrinos. The neutralino candidate, especially the case of the superpartner of the $U(1)_Y$ hypercharge gauge boson, the bino $\tilde{B}$, has received extensive discussion [8]. For certain choice of superpartner masses, the cosmological $\tilde{B}$ energy density can have the correct order of magnitude to be the dark matter. Its interactions in semiconductor detectors are sufficiently weak that it is an experimental challenge to directly detect this form of CDM; its detection rate can be as low as $10^{-4}$ events/kg/day. Higgsino-like [9] and mixed gaugino-Higgsino LSPs are also possible neutralino candidates for CDM.

Sneutrinos annihilate rapidly in the early universe via $s$-channel $Z$ and $t$-channel neutralino and chargino exchange. To reduce these annihilations and obtain a cosmologically significant $\Omega_{\tilde{\nu}}$, it was proposed that the sneutrinos should be light, $m_{\tilde{\nu}} \approx 2$ GeV [10]. Such light LSP sneutrinos could be obtained in minimal supergravity models, although from todays perspective such small scalar masses appear somewhat fine-tuned. This light $\tilde{\nu}$ CDM has been excluded from measurements of the $Z$ width. In supersymmetric models, a LSP

[†]We assume that its stability is guaranteed by $R$-parity.

sneutrino is expected to have a mass in the range of, say, 30—200 GeV from naturalness arguments. However, in this case the cosmological annihilation is large, leading to a low abundance. The annihilation in the early universe can be reduced by taking the sneutrino heavier, 550—2300 GeV for $0.1 \lesssim \Omega_{\tilde{\nu}} \lesssim 1$. Such a heavy sneutrino, already disfavored on theoretical grounds, is firmly excluded by the nuclear recoil direct detection searches: the $t$-channel $Z$ exchange gives a cross section four times larger than the case of a Dirac neutrino, excluding all $m_{\tilde{\nu}}$ up to 17 TeV if the sneutrino is the dominant component of the halo [11]. Sneutrino CDM is apparently firmly excluded.

This negative conclusion on sneutrino CDM is based on the implicit assumption of lepton number conservation, which implies three mass eigenstates of sneutrino, each described by a complex field. It is well known that the phenomenology of neutrinos is greatly changed by the addition of lepton number violation, and the same is true for sneutrinos — each complex field now represents two particles with different masses: $\tilde{\nu}_{\pm}$. In the minimal standard model, without right-handed neutrinos, gauge invariance and renormalizability ensure that the lepton numbers are exact symmetries. However, the standard model is surely just a low energy effective theory, and physics from high mass scales $M$ can induce lepton number violation via the operator $llhh/M$, giving Majorana neutrino masses, where $l$ and $h$ are lepton and Higgs doublets. In supersymmetric extensions of the standard model, with minimal field content and $R$-parity conservation, lepton number violation can occur at dimension four by the operator $\tilde{l}\tilde{l}hh$, which breaks supersymmetry explicitly, and gives a mass splitting between $\tilde{\nu}_{\pm}$.

**2. Phenomenology of $\tilde{\nu}$ CDM.** The purpose of this letter is to present a phenomenological analysis on the viability of the sneutrino CDM with lepton-number violation. The sneutrinos carry the same lepton numbers as their supersymmetric partners (neutrinos), and are distinguished from their anti-particles, anti-sneutrinos. They have soft supersymmetry-breaking masses which are expected to be in the range of 30—200 GeV/$c^2$. In the presence of lepton number violation, sneutrinos can mix with anti-sneutrinos because there are no other quantum numbers which forbid the mixing [12, 13]. Without loss of generality, the mass-squared matrix for a single generation of sneutrinos can be parameterized by two real parameters, $m_{\tilde{\nu}}^2$ and $\Delta m^2$:

$$\mathcal{L}_{\text{mass}} = \frac{1}{2}(\tilde{\nu}^*, \tilde{\nu}) \begin{pmatrix} m_{\tilde{\nu}}^2 & \Delta m^2/2 \\ \Delta m^2/2 & m_{\tilde{\nu}}^2 \end{pmatrix} \begin{pmatrix} \tilde{\nu} \\ \tilde{\nu}^* \end{pmatrix}, \tag{1}$$

where the positive mixing parameter $\Delta m^2$ is a consequence of the operator $\tilde{l}\tilde{l}hh$ mentioned earlier. We later identify these as the tau sneutrinos. We assume that $m_{\tilde{\nu}}^2$ is sufficiently positive that the physical mass eigenstates sneutrinos are $\tilde{\nu}_+ = (\tilde{\nu} + \tilde{\nu}^*)/\sqrt{2}$ and $\tilde{\nu}_- = i(\tilde{\nu} - \tilde{\nu}^*)/\sqrt{2}$, with eigenvalues $m_{\tilde{\nu}_\pm}^2 = m_{\tilde{\nu}}^2 \pm \Delta m^2/2$. The mass difference between $\tilde{\nu}_-$ and $\tilde{\nu}_+$ is

$$\Delta m \equiv m_{\tilde{\nu}_+} - m_{\tilde{\nu}_-} \simeq \frac{\Delta m^2}{2m_{\tilde{\nu}}} \tag{2}$$

for $\Delta m^2 \ll m_{\tilde{\nu}}^2$.

For our purpose, the most important property of the mass eigenstates $\tilde{\nu}_\pm$ is that there is no diagonal coupling to the $Z$-boson; its coupling is always off-diagonal, *i.e.*, $Z$-$\tilde{\nu}_+$-$\tilde{\nu}_-$. This result is a simple consequence of Bose symmetry, and has a crucial impact on both the cosmological sneutrino abundance and on the signal for direct detection of halo sneutrinos. With lepton number conservation, a large contribution to cosmological sneutrino annihilation comes from the $s$-channel exchange of a virtual $Z$ boson, $\tilde{\nu}\tilde{\nu}^* \to Z^* \to f\bar{f}$, where $f$ is any of the Standard Model quarks and leptons with kinematically allowed masses. Although the annihilation process is $P$-wave, the large number of allowed final states and fixed $m_Z$ makes this process important: $\sum_f \sigma(\tilde{\nu}\tilde{\nu}^* \to f\bar{f}) = 0.0072 v_{rel} m_{\tilde{\nu}}^2/(4m_{\tilde{\nu}}^2 - m_Z^2)^2$, where $v_{rel}$ is the relative velocity of the two sneutrinos. With lepton-number violation, however, this process is replaced by the *co*-annihilation $\tilde{\nu}_+\tilde{\nu}_- \to f\bar{f}$. Unless the mass splitting is too small $\Delta m \lesssim 5$ GeV (see below), the annihilation via the $s$-channel $Z$-exchange can be suppressed effectively. Moreover, the mass splitting between the sneutrino and slepton in the same multiplet is given by the $D$-term, $m_{\tilde{l}}^2 - m_{\tilde{\nu}}^2 = (1 - \sin^2\theta_W) m_Z^2 (-\cos 2\beta) > 0$, which is quite important for $m_{\tilde{\nu}_-} < m_W$ and a moderately large ratio of Higgs vevs $\tan\beta \gtrsim 2$. Then the coannihilation with the charged slepton becomes unimportant. Therefore, the dominant annihilation process is via the $t$- and $u$-channel neutralino exchange, to which we will return shortly. If $m_{\tilde{\nu}} > m_W$, however, other processes $\tilde{\nu}\tilde{\nu}^* \to W^-W^+, ZZ$ are possible via $t$-channel slepton or sneutrino exchange. In this case the mixing of the sneutrinos does not affect the annihilation process significantly and hence the earlier analyses [11] apply. Therefore, we focus on the range $m_{\tilde{\nu}_-} < m_W$. We also assume that $\tilde{\nu}_-\tilde{\nu}_- \to hh$ is not kinematically allowed.

The calculation of the cosmic abundance is standard [14]. In Fig. 1, we show the values of $\Omega_{\tilde{\nu}}$ for $H_0 = 50$ Mpc/km/s ($h_0 = 0.5$) as functions of $M_1$ with various values of $\Delta m$. We considered the most important annihilation processes as discussed above: $t$- and $u$-channel exchange of the bino $\tilde{B}$ and the neutral wino $\tilde{W}^3$, with the cross section

$$v_{rel}\sigma(\tilde{\nu}_-\tilde{\nu}_- \to \nu\nu, \bar{\nu}\bar{\nu}) = \frac{\pi}{4}\left(\frac{\alpha_Y M_1}{m_{\tilde{\nu}_-}^2 + M_1^2} + \frac{\alpha_W M_2}{m_{\tilde{\nu}_-}^2 + M_2^2}\right)^2, \tag{3}$$

and also the coannihilation effect with $s$-channel $Z$-exchange suppressed by the Boltzman factor $e^{-\Delta m/T}$. The temperature is taken at the annihilation freezeout: $T \simeq m/25$. Note that the thermal average over the initial state should include the statistical factor 1/2! to avoid double counting of states in the Boltzmann equation. The cross section depends sensitively on the relative ratio (and sign) of $M_1$ and $M_2$. The $SU(5)$ grand-unified theory predicts $M_2 = M_1 \times (3\alpha_W/5\alpha_Y)$. We vary the ratio freely for the purpose of our phenomenological analysis. There is also a contribution from the $s$-channel Higgs boson exchange into $b\bar{b}$ or $\tau^+\tau^-$, but we have checked that it is always much smaller than the neutralino exchange for the range shown in the plot. Recall that $\Omega \sim 0.03$–$0.4$ is needed for halo dark matter, while measurements at larger scales suggest somewhat larger range. Inflation predicts $\Omega = 1$. With the grand-unified gaugino mass relation, the range required for halo dark matter can be obtained with $M_1 \gtrsim 200$ GeV and $\Delta m \gtrsim 5$ GeV. With more general gaugino mass parameters, even the value preferred by inflation can be

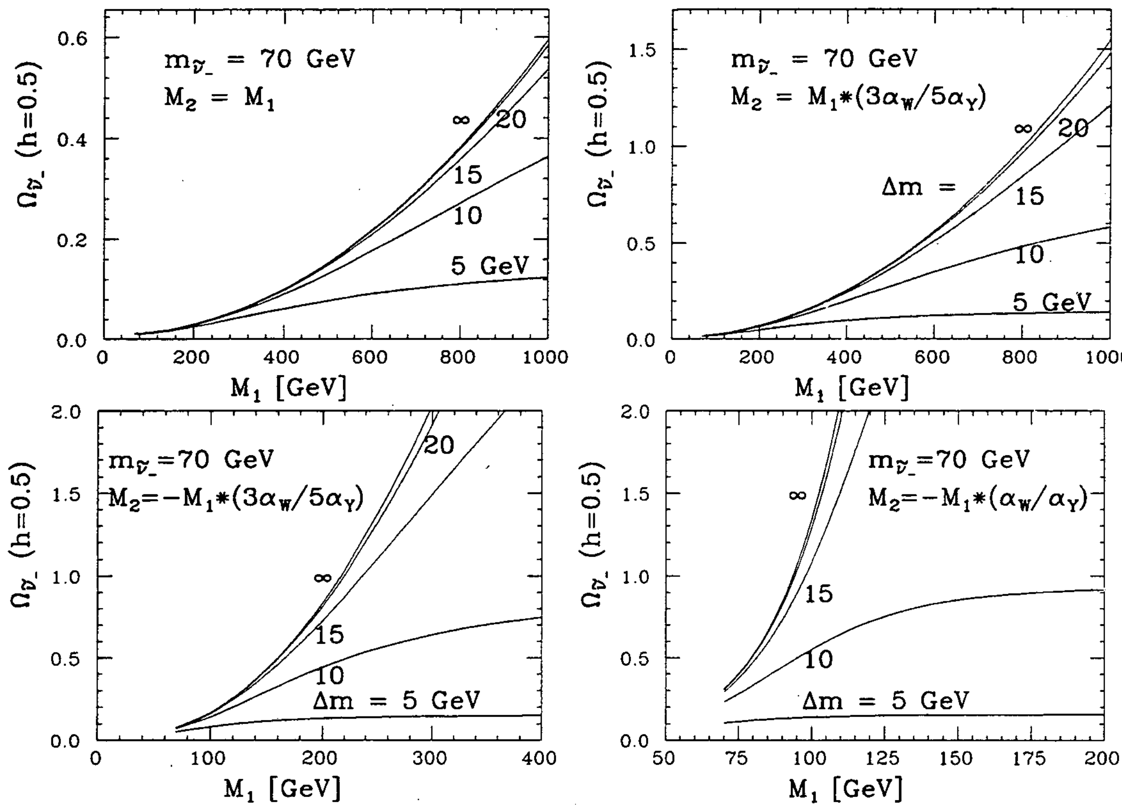


Figure 1: The present relic density of the sneutrino $\Omega_{\tilde{\nu}}$ with $H_0 = 50$ Mpc/km/s, with $m_{\tilde{\nu}_-} = 70$ GeV and $\Delta m = 5, 10, 15, 20$ GeV and $\infty$. The annihilation processes included are $t$- and $u$-channel $\tilde{B}$, $\tilde{W}^3$ exchange and coannhilation with $\tilde{\nu}_+$ via $s$-channel $Z$-exchange. Four plots assume different ratios of $M_1$ and $M_2$ as quoted.. Note the different scales in the plots.

easily obtained. Lepton-number violation allows the sneutrino to become a viable CDM candidate.

Next we consider the detection of galactic halo sneutrinos in Ge detectors. The scattering of $\tilde{\nu}_-$ cannot produce $\tilde{\nu}_+$ due to simple kinematics if $\Delta m > \beta_h^2 m_{\tilde{\nu}_-} m_A / 2(m_{\tilde{\nu}_-} + m_A) =$ 20 keV for $m_{\tilde{\nu}_-} = m_W$, $m_A = 72$ GeV for Ge, and $\beta_h = 10^{-3}$ for virialized halo particles on average. Therefore, there is no $Z$-exchange process beween the sneutrino and the nucleus, and hence the bound from the direct detection experiment described earlier does not apply.[‡] The dominant contribution to the scattering comes from the lightest Higgs boson exchange. We assume that the heavy Higgs boson, whose exchange may enhance the cross section, is sufficiently heavy such that the lightest Higgs boson has the same

[‡]The absence of the $Z\tilde{\nu}_-\tilde{\nu}_-$ coupling also implies that fewer halo sneutrinos are captured by the sun. We find that present limits on high energy neutrinos from the sun do not place a constraint on our scheme.

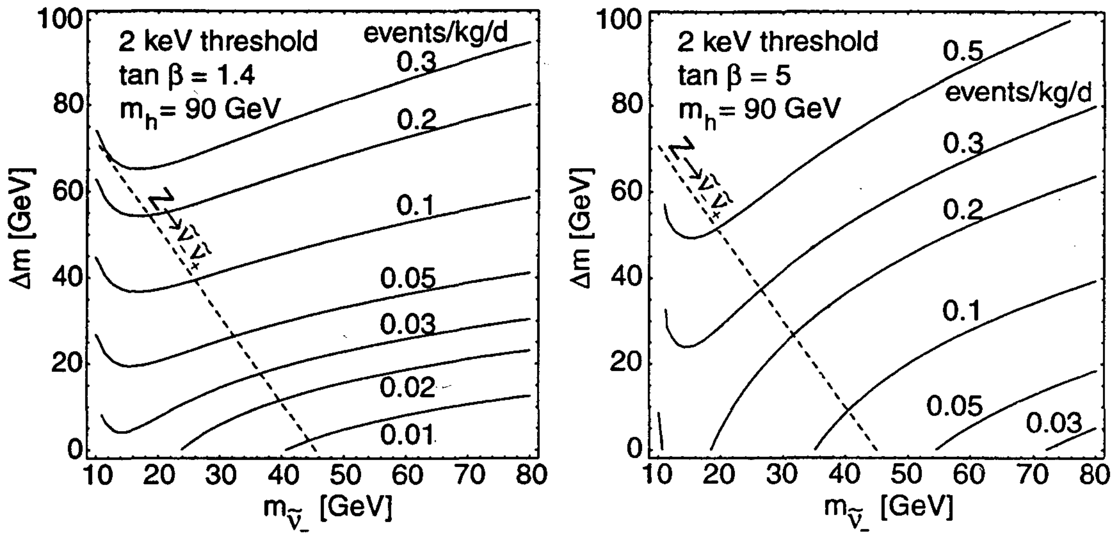


Figure 2: The detection rate of sneutrino CDM in Ge detectors, in units of events/kg/day, for a halo density $\rho = 0.3$ GeV/cm$^3$, and $m_h = 90$ GeV. The threshold energy of the detector is assumed to be 2 keV. We take $\tan\beta = 1.4$ and 5. The dotted line shows the contour for $m_{\tilde{\nu}_-} + m_{\tilde{\nu}_+} = m_Z$, and hence the region below it is excluded by LEP1.

coupling as the Standard Model Higgs boson. The coupling of the Higgs boson and the sneutrino comes from the $D$-term potential in the supersymmetric Lagrangian as well as from the SUSY breaking operator $\tilde{l}\tilde{l}hh$, which is also the origin of the sneutrino mass splitting $\Delta m^2$. For $\tan\beta > 1$, these two contributions always interfere constructively, so that the scattering cross section can be estimated as

$$\sigma = \frac{1}{81\pi(m_{\tilde{\nu}_-} + m_A)^2}\left(\frac{m_A^2(\Delta m^2 - m_Z^2\cos 2\beta)}{v^2 m_h^2}\right)^2, \tag{4}$$

with $v = 250$ GeV. Recall that the lightest Higgs boson in the Minimal Supersymmetric Standard Model has to be lighter than 130 GeV/$c^2$ and must be in a comparable range even in non-minimal extensions, if perturbativity up to the Planck scale is assumed [15].

We show the counting rate of sneutrino CDM with Ge detectors in Fig. 2. Here, we assume the local halo density $\rho = 0.3$ GeV/cm$^3$, lightest Higgs mass $m_h = 90$ GeV, and the isothermal distribution of halo particles with the average velocity $\beta_h = 10^{-3}$. The lowest value of $\tan\beta$ which keeps the top Yukawa coupling perturbative up to the GUT-scale is 1.4 and we used this value as the case with lowest possible detection rate. Another case shown is $\tan\beta = 5$. For larger $\tan\beta$ the detection rate is somewhat larger than the latter case but not much. To obtain a large enough relic density of the sneutrino, $\Delta m$ has to be larger than (5–10) GeV; otherwise coannihilation effect reduces the sneutrino abundance irrespective of the gaugino mass. In this region, the detection rate can be typically $10^{-2}$/kg/day or larger, which is within the reach of future detection of the CDM

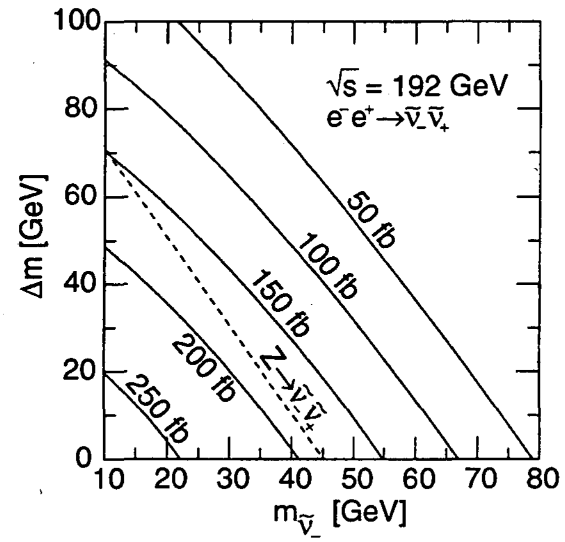


Figure 3: The production cross section for $e^+e^- \to \tilde{\nu}_+\tilde{\nu}_-$ at $\sqrt{s} = 192$ GeV (solid lines). The dotted line shows the contour for $m_{\tilde{\nu}_-} + m_{\tilde{\nu}_+} = m_Z$, and hence the region below the dotted line is excluded by LEP1.

at Ge detector.[§] If $m_h$ is increased to its maximum value of 135 GeV, these rates are decreased by about a factor of four. Nevertheless, direct detection searches are able to probe a larger fraction of the relevant parameter space for sneutrino CDM than for neutralino CDM.

It is an important question what part of the $(m_{\tilde{\nu}_-}, \Delta m)$ parameter space is allowed by current collider experiments. For $\Delta m$ large enough to give a significant cosmological abundance, any $\tilde{\nu}_+$ produced at colliders will decay into $\tilde{\nu}_- jj$ or $\tilde{\nu}_- l^+ l^-$ inside the detector. The signature at LEP1 and LEP2 results from the pair production $e^+e^- \to \tilde{\nu}_-\tilde{\nu}_+$ via $s$-channel $Z$-exchange, with a subsequent decay $\tilde{\nu}_+ \to \tilde{\nu}_- jj$ or $\tilde{\nu}_- l^+ l^-$. This is similar to the signature of the higgsino-like neutralino $\tilde{\chi}^0_1\tilde{\chi}^0_2$ production and the subsequent decay of $\tilde{\chi}^0_2$. The LEP1 constraint is basically $m_{\tilde{\nu}_-} + m_{\tilde{\nu}_+} = 2m_{\tilde{\nu}_-} + \Delta m < m_Z$ [25]. This limit is shown in Fig. 2.

At $\sqrt{s} = 172$ GeV, we find the cross section $\sigma(e^+e^- \to \tilde{\nu}_+\tilde{\nu}_-)$ to be smaller than $\sim 300$ fb, while the current upper bound on the neutralino production cross section is about 800 fb or larger [26]. Hence the sneutrino LSP is not constrained by this bound. However, in the near future, with the full luminosity of LEP2, the cross section will be constrained to be below 100–200 fb [27]. In Fig. 3 the cross section $\sigma(e^+e^- \to \tilde{\nu}_+\tilde{\nu}_-)$ is shown for $\sqrt{s} = 192$ GeV: LEP2 may probe a significant portion of the interesting parameter space.

§We also estimated the nuclear form factor suppression using the formula in [16]. The suppression factor is always less than a factor of two for this light range of CDM mass.

So far we have phenomenologically parametrized the sneutrino mass matrix by varying $m_{\tilde{\nu}}^2$ and $\Delta m^2$ freely. One cannot, however, make $\Delta m^2$ arbitrarily large because the lepton number violation in the sneutrino mass matrix induces a Majorana mass for its partner neutrino from one-loop diagrams. The authors of Ref. [13] analyzed this question and found¶

$$m_\nu \gtrsim \frac{\Delta m}{2 \times 10^3}. \tag{5}$$

Given the value of the $\Delta m^2$ necessary to keep a large enough cosmic abdundance of sneutrino, we conclude that the sneutrino CDM must be the tau-sneutrino. The current limit on the tau neutrino mass is $m_{\nu_\tau} < 18.2$ MeV/$c^2$ [17], so there is room for $\Delta m^2$ in the cosmologically interesting range. An important consequence of sneutrino CDM is that the $\nu_\tau$ mass is in the region of 10 MeV. The asymmetric $B$-factory experiments at SLAC and KEK, BABAR and BELLE, will be able to exclude the finite $\nu_\tau$ mass down to 2 MeV/$c^2$ range [18]. It will be particularly interesting if both the direct search experiments for CDM see a signal and BABAR or BELLE measure a finite $\nu_\tau$ mass. It would be possible to study the consistency of the two results to determine the underlying parameter set.

**3. A model with right-handed neutrinos.** What underlying theory of lepton number violation could lead to $\Delta m \approx 10$ GeV? Since $\tilde{\nu}_-$ must be stable, we seek an $R$-parity conserving origin for the $\Delta L = 2$ operator $O_1 = [llhhz]_F/\Lambda$ where $h$ is the up-type Higgs doublet and $z$ is a dimensionless spurion field $z = A\theta^2$.‖ This operator gives $\Delta m = (A/m_{\tilde{\nu}})(v^2 \sin^2\beta/\Lambda)$. In general one expects $O_1$ to be accompanied by $O_2 = [llhh]_F/\Lambda$, which leads to $m_\nu = v^2 \sin^2\beta/\Lambda$ and the relation $\Delta m = (A/m_{\tilde{\nu}})m_\nu$. In theories with supersymmetry broken in a hidden sector of supergravity, one finds $(A/m_{\tilde{\nu}}) \approx 1$, giving $\Delta m \approx m_\nu \leq 18.2$ MeV/$c^2$, which is a factor $10^3$ too small.

Suppose that operators $O_{1,2}$ arise on integrating out a heavy right-handed neutrino, $\nu_R$, which has the interactions $[\frac{1}{2}\lambda S \nu_R \nu_R + h_\nu l \nu_R h]_F$. Heavy particles may be coupled to large supersymmetry breaking without upsetting the gauge hierarchy, so that we consider $\langle S \rangle = V + \theta^2 F$, a mass for the right-handed neutrino is generated $M = \lambda V$. The effective operators $O_1$, $O_2$ are obtained upon integrating out the right-handed neutrino, with $\Lambda = 2\lambda V/h_\nu^2$, $A = F/V$ and hence the relation

$$\Delta m = \frac{F}{V m_{\tilde{\nu}}} m_\nu. \tag{6}$$

Hence, if $\nu_R$ is coupled to a field $S$ which has $F/V \approx 10^3 m_{\tilde{\nu}}$, then $\Delta m$ is sufficient to allow $\tilde{\nu}_-$ to be CDM. Since $\nu_R$ is vector-like with respect to the standard model gauge group, it is not surprising that it is coupled to larger symmetry breakings than the light matter — this is the motivation for the seesaw mechanism itself — however, we have no convincing argument for the magnitude of $F/V$.

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¶This result assumes no accidental cancellations between tree level and various one-loop contributions to the neutrino mass.

‖Here and below, $[\ldots]_F$ refers to the $F$-component of the chiral superfield in square brackets.

There is an important constraint on the scale $M$. The sum of the tree-level neutrino mass and the one-loop induced term should not be larger than the experimental limit of 18.2 MeV/$c^2$, and, barring a possible cancellation, we require $m_\nu^{\rm tree} = (h_\nu v \sin\beta)^2/2M \lesssim$ 20 MeV which bounds $M$ from below. On the other hand, a sufficient $\Omega_{\tilde\nu}$ requires $\Delta m^2 = 2Am_\nu^{\rm tree} \gtrsim 500$ GeV$^2$. This requires $A \gtrsim .0075 \cdot M/h_\nu^2$. Finally, the large supersymmetry breaking $A$ in the right-handed neutrino generates corrections to $m_{\tilde l}^2$ and $m_{\tilde h}^2$ via two-loop diagrams. This can be calculated using the method of Giudice and Rattazzi to leading order in $A$ [19], and we find

$$\delta m_{\tilde l}^2 = \frac{1}{(16\pi^2)^2} A^2 h_\nu^2 (4h_\nu^2 + 3h_t^2 - 3g^2 - g'^2), \tag{7}$$

$$\delta m_{\tilde h}^2 = \frac{1}{(16\pi^2)^2} A^2 h_\nu^2 (4h_\nu^2 - 3g^2 - g'^2), \tag{8}$$

where $h_t$ is the top quark Yukawa coupling, and $g$, $g'$ are $SU(2) \times U(1)$ gauge coupling constants. On naturalness grounds, these corrections should not be larger than about (100 GeV)$^2$. Combined with the lower bound on $A$ found above, we find an upper bound on $M$; in fact for $\delta m_{\tilde h,\tilde l}^2 < (100\text{GeV})^2$ we find that the allowed region is given by $M = 100$–700 TeV and $h_\nu = 0.2$–$0.7$.** It is interesting that the scale $M$ is comparable to the one found in the simplest theories of gauge mediation, so that $S$ can be identified as the singlet field which gives gauge mediated supersymmetry breaking [20]. However, for $\tilde\nu_-$ to be the LSP, it is necessary that the gravitino mass be larger than $m_{\tilde\nu_-}$, which requires the existence of a larger primordial supersymmetry breaking in the theory $F_P \geq 10^{10}$GeV $\gg F$. This happens when the messenger U(1) gauge coupling is somewhat small. The gravitino heavier than the sneutrino is actually cosmologically favorable because the gravitino LSP is rather problematic [21].

**4. Conclusions** In the minimal supersymmetric standard model, the sneutrino is firmly excluded as a CDM candidate. The $Z\tilde\nu^\dagger\tilde\nu$ coupling leads to rapid cosmological annihilation, and therefore low values of $\Omega_{\tilde\nu}$, unless $m_{\tilde\nu}$ is very large, in which case the same coupling leads to a large and excluded event rate in Ge detectors of halo CDM particles. In this letter we have shown that sneutrino CDM is allowed in supersymmetric theories with lepton number violation. A lepton number violating sneutrino mass implies that each flavor of sneutrino has two distinct mass states $\tilde\nu_\pm$. In this case there is only an off-diagonal $Z$ coupling, $Z\tilde\nu_+\tilde\nu_-$, so that if the mass splitting of these two states is larger than about 5 GeV, and if the lightest sneutrino has a mass in the range of about 40—80 GeV, $\Omega_{\tilde\nu}$ in the interesting range of 0.1 to 1 can result. We have shown that the seesaw mechanism, which gives small neutrino masses from integrating out heavy right-handed neutrinos, can also lead to the required lepton number violation in the sneutrino mass matrix.

There are three important, pre-LHC/LC tests for sneutrino CDM:

**The constraints from $\delta m_{\tilde l}^2$, $\delta m_{\tilde h}^2$ are somewhat subjective. Also a different model of lepton number violation (such as a weak-triplet lepton exchange generating $O_1$, $O_2$) leads to very different results for $\delta m_{\tilde l}^2$, $\delta m_{\tilde h}^2$ and hence the constraints here are model-dependent.

- Galactic halo sneutrinos will scatter in Ge detectors with an event rate $\geq 10^{-2}$ events/kg/day, for most of the relevant parameter range (see Fig. 2).

- $m_{\nu_\tau} \gtrsim 5$ MeV, unless different contributions to the neutrino mass are fined tuned to cancel. This mass range of $\nu_\tau$ can be excluded by the $B$-factory experiments.

- Events of the form $l^+l^- \not{E}$ or $jj\not{E}$, with low $m_{ll}$ or $m_{jj}$, may be observed at LEP2, with $\sqrt{s} = 192$ GeV (see Fig. 3). They result from $\tilde{\nu}_+\tilde{\nu}_-$ pair production, followed by $\tilde{\nu}_+$ decay.

There are further important consequences of sneutrino CDM:

(1) There are important new collider signatures of supersymmetry. Squark and gluino production at hadron colliders leads to events with substantial missing transverse energy, which is carried away by the undetected $\tilde{\nu}_-$. However, a large fraction of these events have $\tilde{\nu}_+$ in the decay chain, and when these decay to $\tilde{\nu}_-$ they can produce lepton pairs with small invariant mass. This decay, $\tilde{\nu}_+ \to \tilde{\nu}_- l^+l^-$, becomes an important characteristic feature of many supersymmetric signals. Also, the lightest Higgs boson may decay dominantly to sneutrinos, $h \to \tilde{\nu}_+\tilde{\nu}_+, \tilde{\nu}_-\tilde{\nu}_-$. It is possible that only the invisible $\tilde{\nu}_-\tilde{\nu}_-$ channel is kinematically allowed.

(2) The $\nu_\tau$, with its mass in the expected (5–20) MeV range, would overclose the universe if it is stable. A visible decay, such as $\nu_\tau \to \nu_{e,\mu}\gamma$ or $\nu_{e,\mu}e^+e^-$, and the invisible $3\nu$ mode, are disfavored for a variety of reasons. The $\nu_\tau$ should decay into a massless boson $\nu_\tau \to \nu_{e,\mu} f$, with $f$ a Majoron or familon, whose phenomenology was discussed recently in detail [23]. The existence of such a massless boson is natural if the lepton number is broken spontaneously at the mass scale of right-handed neutrino [22]. It is interesting to note that a $\nu_\tau$ in this mass range, and with lifetime $10^{-2}$ sec $\lesssim \tau_\nu \lesssim 1$ sec, improves the situation with Big-Bang Nucleosynthesis [24].

(3) For a critical universe with $\Omega_{\tilde{\nu}} = 1$, unification of the gaugino mass parameters is strongly disfavored.

(4) In the early universe, the large lepton number violation in the neutrino sector, together with high temperature $B+L$ sphaleron transitions, may wash out the cosmological baryon asymmetry [28]. Hence the baryon asymmetry should either be generated at low temperatures, beneath the electroweak phase transition, or protected from sphaleron washout by condensates [29] or by other exact symmetries.

(5) The sneutrinos must be coupled more strongly to supersymmetry breaking than occurs in the simplest supergravity models. Such mediation of supersymmetry breaking can readily occur via the gauge singlet right-handed neutrino.

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