Small neutrino masses from supersymmetry breaking

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(Received 11 October 2000; published 12 November 2001)

An alternative to the conventional seesaw mechanism is proposed to explain the origin of small neutrino masses in supersymmetric theories. The masses and couplings of the right-handed neutrino field are suppressed by supersymmetry breaking, in a way similar to the suppression of the Higgs doublet mass μ. New mechanisms for light Majorana and Dirac neutrinos arise, depending on the degree of suppression. Superpartner phenomenology is greatly altered by the presence of weak scale right-handed sneutrinos, which may have a coupling to a Higgs boson and a left-handed sneutrino. The sneutrino spectrum and couplings are quite unlike the conventional case—the lightest sneutrino can be the dark matter and predictions are given for event rates at upcoming halo dark matter direct detection experiments. Higgs boson decays and search strategies are changed. Copious Higgs boson production at hadron colliders can result from cascade decays of squarks and gluinos.

DOI: 10.1103/PhysRevD.64.115011 PACS number(s): 12.60.Jv

I. INTRODUCTION

Why are neutrinos much lighter than charged leptons, but not absolutely massless? It is universally recognized that this can be simply and elegantly understood in an $SU(2)\times U(1)$ effective theory. The most general, gauge-invariant interactions of dimension less than 6, which can lead to masses for the known leptons from the vacuum expectation value (VEV) of a Higgs doublet, are

$$\mathcal{L}_{eff}=\lambda LEH+\frac{\lambda'}{M}LLHH, \tag{1}$$

where L and E are the lepton doublet and singlet fields, and H is the Higgs doublet. The dimensionless matrix of Yukawa couplings, λ, has a hierarchy of eigenvalues to describe the masses of the charged leptons. Such a hierarchy could result by promoting the couplings to fields, $\lambda\rightarrow\phi/M$, which acquire VEV's to sequentially, spontaneously break the flavor symmetry G_F. Such a flavor symmetry will also result in a certain structure for the neutrino mass matrix via λ'. The mass scale M is the cutoff of the low-energy effective theory. The crucial point is that if this cutoff is very large, for example the Planck or gauge coupling unification scale, then the neutrino masses are very small. While the charged lepton masses are linear in the Higgs VEV v, the neutrino masses are quadratic:

$$m_\nu\approx\frac{v^2}{M}. \tag{2}$$

The power of this effective field theory approach is that no assumption need be made about the full theory at or above the scale M. The only assumption is that the low-energy theory is the most general allowed by its symmetries. Nevertheless, there is a very simple theory which does lead to the dimension 5 operator of (1). Right-handed neutrino fields N are introduced, with Majorana masses M and couplings to the lepton doublets:

$$\mathcal{L}=\frac{M}{2}NN+\xi LNH, \tag{3}$$

where M and ξ are mass matrices. On integrating out the heavy neutrinos, the well-known seesaw mechanism gives λ'/M in Eq. (1) by $\xi^T M^{-1}\xi$ [1].

In supersymmetric theories, there is a very important reason for questioning this simple view of neutrino masses: the low-energy effective theory must contain more fields than the leptons, the Higgs boson, and their superpartners. In particular, there are two Higgs doublet superfields H_u and H_d, and there is another sector of the theory which spontaneously breaks supersymmetry (SUSY) and triggers electroweak symmetry breaking. At first sight these additions seem irrelevant to the question of neutrino masses, but closer inspection reveals new opportunities. An important objection to the minimal supersymmetric standard model is that it is not the most general low-energy effective theory consistent with $SU(3)\times SU(2)\times U(1)$ gauge symmetry. The gauge symmetry allows a bare mass term for the Higgs fields, $[\mu H_u H_d]_F$, giving the expectation $\mu\approx M$, which removes the Higgs fields from the low-energy theory. This is known as the "μ problem" in supersymmetric theories.

In this paper we take supersymmetry to be broken in a hidden sector, at the intermediate scale m_I, via fields Z: $\langle F_Z\rangle\approx m_I^2\approx vM_{Pl}$, where M_{Pl} is the Planck mass. The supersymmetry breaking is communicated to the standard model by supergravitational interactions, so that the cutoff for the low-energy effective field theory is M_{Pl}. This gives rise to superpartner masses at the weak scale in the usual way. In this case, the above "μ problem" is easily solved by introducing a further global symmetry, G, and dividing the light matter superfields into two types. There are fields which are chiral with respect to the gauge interactions, such as L

and E, which are guaranteed to be massless until the gauge symmetry is broken, and there are fields which are vectorlike, such as H_u+H_d, which are kept massless only via G. Furthermore, the fields ϕ which break the flavor symmetry do not break G—the vectorlike fields acquire mass only from supersymmetry breaking. There is a subset of the Z fields, which we call X, that transform nontrivially under G. This ensures that μ is of the order of the supersymmetry breaking scale, as it must be for electroweak symmetry breaking to occur successfully [2,3]. In particular, the operator

$$\frac{1}{M_{Pl}}[X^\dagger H_u H_d]_D \tag{4}$$

leads to $\mu\approx F/M_{Pl}\approx m_I^2/M_{Pl}\approx v$, where F is the VEV of the highest component of X.

The right-handed neutrino fields N, like H_u+H_d, are vectorlike with respect to the gauge interactions; hence the crucial question becomes how they transform under G. If they are also vectorlike with respect to G, they will acquire a large mass, from which the seesaw proceeds via Eq. (3) as usual. Alternatively, they may be protected from acquiring a large mass by G, in which case they will appear in the low-energy effective theory. The question of neutrino masses now becomes much richer, since it is necessary to study all possible interactions of L and N in the low-energy theory. In particular, mass terms can be induced via interactions with Higgs VEV's and with X VEV's.

By analogy with Eq. (4), the right-handed neutrinos may acquire mass via the operator

$$\frac{1}{M_{Pl}}[X^\dagger NN]_D, \tag{5}$$

giving the right-handed neutrinos a Majorana mass of order of the weak scale, v. The seesaw mechanism can still be operative if the Yukawa couplings are suppressed. For example, if one of the X fields acquires an A component VEV at the intermediate scale, this occurs via

$$\frac{1}{M_{Pl}}[XLNH_u]_F. \tag{6}$$

In general, the coefficients of these higher dimensional operators are understood to depend on flavor symmetry breaking, and are functions of ϕ/M_{Pl}. The seesaw now gives light Majorana masses $m_\nu\approx(m_I v/M_{Pl})^2/(m_I^2/M_{Pl})\approx v^2/M_{Pl}$—the usual result.

Even simpler possibilities occur: the G quantum numbers may prevent a right-handed Majorana mass to very high order, so that the dominant mass term is Dirac. For example the operator

$$\frac{1}{M_{Pl}^2}[X^\dagger LNH_u]_D \tag{7}$$

dominates either if the operator (6) is forbidden by G or if X does not have an A component VEV, and gives light Dirac neutrinos of mass $m_\nu\approx Fv/M_{Pl}\approx v^2/M_{Pl}$.

In Sec. II, we study the general low-energy effective theory for the interactions of L and N with H_u and X, with a view to studying the interesting forms for the neutrino mass matrices. The above Majorana and Dirac cases are considered further.

What are the consequences of our proposal that neutrino masses are suppressed relative to charged lepton masses by supersymmetry breaking factors? While the most immediate consequence is that it opens up a new class of models for neutrino masses, the most important consequence may be that supersymmetric phenomenology can be drastically altered. This is largely due to the possibility that the right-handed sneutrino, $\tilde{n}$, can now be at the weak scale and couple via a new A-type interaction:[1]

$$\frac{1}{M_{Pl}}[XLNH_u]_F\supset v\tilde{l}\tilde{n}h_u. \tag{8}$$

The structure of this interaction is investigated in Sec. III. Its consequences for the sneutrino mass spectrum is studied in Sec. IV, and its consequences for the lightest sneutrino as the cosmological dark matter in Sec. V. We find two interesting cases where this occurs, and each case predicts a characteristic signal in upcoming experiments to directly detect halo dark matter.

The A-term interaction $\tilde{l}\tilde{n}h_u$ can significantly alter the decay branching ratios for charged and neutral Higgs bosons in supersymmetric theories. This is studied in Sec. VI, where we find that, for certain ranges of parameters, the decays $h\to\tilde{\nu}\tilde{n}$ and $H^\pm\to\tilde{l}^\pm\tilde{n}$ can be the dominant decay modes. Further consequences for collider phenomenology, arising from the A-term changing the decay chain of $\tilde{\nu}$, are discussed in Sec. VII. Finally, in Sec. VIII we study the rare lepton flavor violation implied by our mechanism for neutrino mass generation.

Here we mention a few other interesting alternatives to the seesaw mechanism that were considered previously. One extensively explored possibility is that small neutrino masses arise from R-parity violation in supersymmetric theories [4]. Also, a small Yukawa coupling can, in principle, at least, be understood with the Froggatt-Nielsen mechanism [5], through the ratio v/M of the VEV of a supersymmetry conserving spurion over some higher mass scale. Such a VEV could arise from dimensional transmutation, or, as recently discussed in Ref. [6], from radiative symmetry breaking, which was employed in Ref. [7] to give small masses to sterile states which mix with ordinary neutrinos. The possibility of using supersymmetry breaking operators to generate light sterile states was initially explored in Ref. [8], and the possibility of using supersymmetry breaking operators to radiatively generate Majorana masses for the left-handed neutrinos was considered in Ref. [9]. Unlike the theories of Ref.

[1]If we wish to combine the unsuppressed A terms arising from Eq. (8) with the light Dirac neutrinos whose masses are generated by Eq. (7), we must require that the A component VEV of X be small: $[X]_A\leqslant v$.

[7] and [8], the models we construct give just three light neutrino mass eigenstates (although simple extensions lead to additional light states), and unlike the work of Ref. [9], the neutrino masses in our framework arise at tree level. Most important, as discussed above, the present class of models features a natural mechanism for generating weak scale A terms involving the N states, which has not been previously discussed.

II. SMALL NEUTRINO MASSES FROM F TERM SUSY BREAKING

We begin by considering the low-energy effective theory which describes the interactions of the leptons L and N with the Higgs doublet H_u and the fields X which spontaneously break both supersymmetry and the global symmetry group G. We impose R parity, which changes the sign of the L and N superfields, as well as the superspace coordinate θ, but leaves H_u and X unchanged. Expanding in powers of $1/M$,

$$\mathcal{L}_{eff}=[c_{4,1}XNN+c_{4,2}LNH_u]_F+\frac{1}{M}([c_{5,1}XLNH_u+c_{5,2}(LH_u)^2+c_{5,3}N^4+c_{5,4}(XN)^2]_F+[c_{5,5}X^\dagger NN]_D)+\frac{1}{M^2}([c_{6,1}X^\dagger LNH_u+c_{6,2}XX^\dagger NN]_D+[c_{6,3}X^3N^2+c_{6,4}X(LH_u)^2]_F+\cdots). \quad (9)$$

Here and below, the energy scale M is the ultraviolet cutoff of the low-energy theory, such as the Planck scale or the grand unified theory (GUT) scale.

We have included all possible operators even though many may be excluded by the global symmetry G, depending on the model. The bare mass $[NN]_F$ is always forbidden by G and is not shown. Likewise, the mass terms XN are forbidden by R parity. The flavor structure is not shown explicitly—there are three L fields and one or more N field, and in general the coefficients $c_{i,j}$ are power series in flavor symmetry breaking parameters ϕ/M. If X acquires an F component VEV, the operator $[X(LH_u)^2]_F$ leads to radiatively generated Majorana neutrino masses, even without the existence of N fields [9]. In the explicit models constructed below, this operator is forbidden, so that $c_{6,4}$ vanishes. There are other dimension six operators, such as $[N^\dagger NL^\dagger L]_D$, but these will not affect the structure of the model nor the phenomenology, so we will not discuss them. It is possible that even higher dimension operators are important for generating neutrino masses, but in the explicit models discussed below the above expansion will be sufficient. However, in Sec. V we will consider important consequences that lepton-number violating dimension seven operators can have in the context of dark matter.

Before we consider particular symmetries, let us explore which combinations of operators are phenomenologically interesting. In later sections, we will exhibit particular symmetries which realize these scenarios. Since N must appear in combination with some of the X_i superfields, $c_{4,2}=c_{5,3}=0$. If we allow $c_{4,1}\neq 0$, when X_i obtain F component VEV's, this term will generate an intermediate scale mass for the N scalar $\tilde{n}$, and this case is thus of less phenomenological interest. Furthermore, precisely the same term must be omitted in the minimal supersymmetric standard model (MSSM) (XH_uH_d), and its omission here seems very natural. For these reasons we will take $c_{4,1}=0$.[2]

The structure of the theory can vary, depending on a few elements. In particular, if X develops an A component VEV, the size of the Yukawa couplings will be different. There are three different natural values for the A component: M, $\sqrt{F}$ and zero. If it is M, then the Yukawa couplings from the $c_{5,1}$ term in Eq. (10) are of order 1, which is phenomenologically unacceptable. We consider the other two alternatives in greater detail. We will begin by considering situations with only one generation.

A. One light Dirac neutrino

If the A component of X is zero, but X does gain an F component, we generate Yukawas

$$\left[\frac{X^\dagger}{M}LH_uN\right]_D=\frac{F}{M}[LH_uN]_F. \quad (10)$$

If we assume that G_F sets $c_{5,5}=0$ in Eq. (10), so that there is no Majorana mass for the right-handed neutrino, then we have generated a Yukawa of order M_{SUSY}/M. When the Higgs field takes on a VEV, we then have a mass for the neutrino $O(v^2/M)$. This is astonishing, because we now have a naturally light *Dirac* neutrino, with a mass of the correct size to explain the observed phenomena associated with neutrino mass.[3] If this is correct, then experiments studying neutrino mass have *already* begun to probe the structure of supersymmetry breaking.

As discussed in Sec. I, the present class of models features a natural mechanism for generating weak scale A terms through the operator[4]

$$\left[\frac{X}{M}LH_uN\right]_F=\left[\frac{F}{M}LH_uN\right]_A. \quad (11)$$

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[2]If the A-component VEV of X is zero, one might also think that $c_{4,1}$ must be zero for another reason. If one allowed such a large supersymmetry breaking mass for N, such that the fermion was present in the weak-scale theory, but the scalar was not, one might worry that loop effects would destabilize the v/M hierarchy. However, all dangerous diagrams that appear are suppressed by small Yukawa couplings and are harmless.

[3]Whether these are precisely the right size is not of particular concern. This is an effective theory and M could easily be M_{GUT}, or some other scale, in which case the Yukawas are larger.

[4]In fact, if $\langle X\rangle=\theta^2F$ is generated in the global supersymmetry limit, supergravity effects will generate a small shift in the A component, giving $\langle X\rangle\sim F/M_{Pl}+\theta^2F$ [10]. Thus the large A terms and the small Dirac neutrino masses can be generated from the operator of Eq. (11) alone [11,12].

Much of the rest of the paper will be spent working out some of the consequences of these A terms.

Operators like Eqs. (10) and (11) could be selected by a symmetry $U(1)_L \otimes U(1)_N$, with fields X and $\overline{X}$, with charges (1,1), and $(-1,-1)$, respectively. E, L and N have charges $(-1,0)$, (1,0) and (0,1), respectively. With these charges, the only operators of dimension 6 or less allowed are

$$\left[\frac{c_{6,1}}{M^2}X^\dagger LNH_u\right]_D + \left[\frac{c_{5,1}}{M}\overline{X}LNH\right]_F. \tag{12}$$

The superfields X and $\overline{X}$ could acquire F component VEVs, but no A component VEV's if embedded in an O'Raifeartaigh model. Given a superpotential

$$W = S(Y\overline{Y} - \mu^2) + Y^2\overline{X} + \overline{Y}^2 X, \tag{13}$$

the minimum of the scalar potential will occur with $\langle y\rangle = \langle \overline{y}\rangle = \mu/\sqrt{3}$. Here Y and $\overline{Y}$ have charges (1/2,1/2) and $(-1/2,-1/2)$, respectively. Two linear combinations of s, x and $\overline{x}$ will have positive masses at tree level, while the third independent combination will obtain its mass in the one-loop effective potential, stabilizing $\langle x\rangle = \langle \overline{x}\rangle = 0$. Note that the presence of the superpotential term $[(\overline{Y}^2/M^2)LNH_u]_F$ generates a contribution to the Yukawas of the same order of magnitude as that from $[(X^\dagger/M^2)LNH_u]_D$. Also note that this breaks $U(1)_N \otimes U(1)_L$, but preserves $U(1)_{L-N}$, which is the ordinary lepton number symmetry. We will refer to this scenario, in which right-handed neutrinos couple with suppressed Yukawas, but have no Majorana masses, as ''Dirac masses from supersymmetry breaking,'' or ''sDirac'' for short.

B. One light Majorana neutrino

An alternative to generating light Dirac neutrino masses from supersymmetry breaking is to instead generate light Majorana masses. We begin by considering the operators

$$\left[\frac{X^\dagger}{M}NN\right]_D + \left[\frac{X}{M}LNH_u\right]_F. \tag{14}$$

Such terms could be justified by an R symmetry, where N has an R charge 2/3, L and H both have an R charge 0, and X has an R charge 4/3. If X takes on an A component VEV $\langle X\rangle|_{\theta=0} = \sqrt{F}$, as well as an F component VEV, $F \approx (10^{11}\,\mathrm{GeV})^2 \approx vM_{Pl}$,[5] then the second term in Eq. (14) generates a weak scale A term, but now generates a Yukawa coupling to the Higgs boson as well, roughly of the size $\sqrt{F}/M \approx 10^{-8}$. If this were the end of the story, then we would have a (Dirac) neutrino mass ~ 1 keV. However, the first term of (14) will now generate a Majorana mass for N of the order $F/M \approx 100$ GeV, yielding a LR neutrino mass matrix

$$\begin{pmatrix} 0 & v\sqrt{F}/M \\ v\sqrt{F}/M & F/M \end{pmatrix}. \tag{15}$$

After integrating out the heavy N fermion, we are left with a Majorana mass for the neutrino with a size $m_\nu \approx v^2/M$, again reproducing the seesaw result.[6]

We will refer to this scenario, in which the right-handed neutrinos have Yukawas $O(\sqrt{F}/M)$ and weak scale Majorana masses, as ''Majorana mass from supersymmetry breaking,'' or sMajorana for short.

III. FLAVOR STRUCTURES

In Sec. II, we concerned ourselves simply with the origin of the neutrino mass itself, but did not address the additional question of what determines the structures of these masses when we include additional generations. As discussed in Sec. I, we are considering a scenario in which the global symmetry of the theory is $G_F \otimes G \otimes SUSY$. G must include some symmetry to keep the Higgs doublets and right-handed neutrinos light, and G_F may include symmetries such as $U(3)^6$ which relate the different generations.

The key feature of our model is that the supersymmetry breaking spurions also contain charges under G. When these spurions acquire F, and possibly A component VEV's, they break G. Of course, they need not be charged merely under G, but potentially under some larger group H, where $G_F \otimes G \supset H \supset G$. In the most minimal framework, $H = G$ and would contain only those symmetries which are necessary to suppress the μ term and the right-handed neutrino masses, for instance $U(1)_L \otimes U(1)_N$ in Sec. II A, or the R symmetry in Sec. II B.

With such an assumption, the textures of the neutrino mass matrices and the A terms would be determined by supersymmetry preserving elements of the theory. For instance, in the sDirac scenario, the couplings would be given by

$$[\lambda_{ij}H_u L^i N^j]_F = [X^\dagger \Lambda_{ij} H_u L^i N^j]_D, \tag{16}$$

$$A_{ij}h_u \tilde{l}^i \tilde{n}^j = [\overline{X}\Lambda'_{ij}H_u L^i N^j]_F, \tag{17}$$

where Λ_{ij} and Λ'_{ij} are supersymmetry *preserving* but flavor breaking spurions. As such, an explanation of structure of the Yukawas and A terms is beyond the scope of our scenarios. However, because of this, in the presence of a flavor symmetry, we are able to easily relate the structure of A_{ij} and λ_{ij}.

[5]For instance, in the SUSY breaking model of Ref. [13], a chiral superfield naturally develops A and F components of the same order of magnitude under the dynamical assumption that a constant appearing in the Kahler potential is negative [11].

[6]A weak scale Majorana mass for N could have been generated with nonzero $c_{5,4}$ as well. However, the scalars would then have a supersymmetry breaking mass squared $O(\sqrt{F}M_W)$. Thus, for the same reasons we took $c_{4,1}=0$, we do not consider the case where G_F allows nonzero $c_{5,4}$.

However, in the sMajorana case, the couplings will be given by

$$A_{ij}h_u\tilde{l}^i\tilde{n}^j=[X_F\Lambda_{ij}H_uL^iN^j]_F, \tag{18}$$

$$[\lambda_{ij}H_uL^iN^j]_F=[X_A\Lambda_{ij}H_uL^iN^j]_F. \tag{19}$$

Here the Yukawas are precisely the same as the A terms, but due to the potential mixing of the N's, it is impossible to necessarily relate the A term matrix to the neutrino mass matrix obtained by the seesaw mechanism.

We consider this to be the minimal scenario, in which H is as small as possible, so to speak. However, H can easily be much larger. Indeed, people have investigated the possible effects of baryon and lepton number violation operators from supersymmetry breaking [9]. *A priori*, there is no reason why we should reject the possibility that $H=G_F\otimes G$. If that were the case, we would write the sDirac couplings as

$$[\lambda_{ij}H_uL^iN^j]_F=[X_{ij}^\dagger H_uL^iN^j]_D, \tag{20}$$

$$A_{ij}h_u\tilde{l}^i\tilde{n}^j=[\overline{X}_{ij}H_uL^iN^j]_F. \tag{21}$$

Now the spurions X and $\overline{X}$ carry generation indices themselves. We assume that charged fermion Yukawas are generated in a supersymmetry preserving sector of the theory. Since the flavor structure of X and $\overline{X}$ is entirely determined in the supersymmetry breaking sector, they need not be aligned with those of the supersymmetry preserving Yukawas. Consequently, a large mixing between ν_μ and ν_τ is natural. Of course, the small mixing between ν_e and this heavy state ($\theta_{e3}<0.16$ as required by CHOOZ [14]), must be viewed as somewhat of an accident, but not necessarily a fine tuning. This is similar to the anarchy proposal of Ref. [15], except that here we need not relate the Yukawas of ν and e, and a hierarchy of eigenvalues could still possibly occur in X. Consequently, even for sDirac neutrinos, anarchic aspects of the theory are reasonable. Whether this is compatible with supersymmetric flavor changing constraints is an important question. We will address this further in Sec. VIII.

It is important to note that we do not need three right-handed neutrinos for the cases of Secs. II A and II B. The presence of just two N states is enough to generate either two massive Dirac or two massive Majorana neutrinos. The remaining neutrino is simply a massless Weyl neutrino. In a certain sense, this is more minimal than with three N's, but the phenomenology is largely unchanged.

One limit of this could be that there is, in fact, only one N. Here, one might generate Majorana masses for the neutrinos through an ordinary GUT seesaw, while a sDirac N would contribute a fourth mass eigenstate, resulting in four Majorana neutrinos. Given appropriate G_F charges, other, more exotic possibilities may exist, such as combinations of sDirac and sMajorana neutrinos.

IV. SNEUTRINO MASS MATRIX

In the MSSM, the sneutrino and charged slepton masses are intimately related:

$$m_{\tilde{\nu}}^2=m_L^2+\frac{1}{2}m_Z^2\cos 2\beta,$$

$$m_{\tilde{l}_L}^2=m_L^2+\left(\sin^2\theta_W-\frac{1}{2}\right)m_Z^2\cos 2\beta, \tag{22}$$

where m_L is the soft scalar mass for the left-handed sleptons. For $\tan\beta>1$, $\cos 2\beta<0$ and the D-term splitting pushes the sneutrino mass down and the charged slepton mass up. The present experimental bound $m_{\tilde{l}_L}>70$ GeV still allows for light sneutrinos due to this splitting. However, if in the future it becomes established that $m_{\tilde{l}_L}$ is very large, much of the phenomenology associated with light sneutrinos will be ruled out in the MSSM. In our model, with light right-handed sneutrinos, the story is quite different, both because the A terms provide an additional source of splitting between the sneutrino and charged slepton masses, and because the right-handed sneutrino mass is not linked to slepton masses by gauge invariance. Thus, even if m_L is quite large it is still possible to have significant change in phenomenology that would otherwise be absent.

To better understand the spectrum, we consider a single sneutrino generation with mass-squared matrix:

$$m_{\tilde{\nu}}^2=\begin{pmatrix} m_L^2+\frac{1}{2}m_Z^2\cos 2\beta & Av\sin\beta \\ Av\sin\beta & m_R^2\end{pmatrix}. \tag{23}$$

Given that m_L, m_R and A are independent parameters, this matrix can have very different eigenvalues. We plot the mass spectra for various choices of m_L and m_R as a function of A in Fig. 1.

An independent lower bound on the sneutrino mass in the MSSM, $m_{\tilde{\nu}}>44$ GeV, comes from the measurement of the invisible width of the Z, and is also altered by the addition of right-handed sneutrinos. The lightest sneutrino in our model is a superposition of left- and right-handed states:

$$\tilde{\nu}_1=-\tilde{\nu}_L\sin\theta+\tilde{\nu}_R\cos\theta. \tag{24}$$

If this state is light enough to be produced in Z decays, its contribution to the Z width is given by

$$\delta\Gamma=\frac{\sin^4\theta}{2}\left(1-\left(\frac{2m_{\tilde{\nu}_1}}{m_Z}\right)^2\right)^{3/2}\Gamma_\nu \tag{25}$$

where $\Gamma_\nu=167$ MeV is the Z width to ordinary neutrinos. If we take the current LEP limit of 2 MeV [16], the $\sin^4\theta$ factor allows us to evade the bounds regardless of mass provided $\sin\theta<0.39$, which is a very mild constraint.

Although m_L, m_R and A are independent parameters, we can gain some intuition into their sizes from their renormalization group running. In fact, it is somewhat natural to have $m_R<m_L$ in the low-energy theory. The runnings of m_L and m_R are governed by

$$\frac{dm_L^2}{dt}=-\frac{3}{16\pi^2}g^2M_2^2-\frac{3}{80\pi^2}g_Y^2M_1^2+\frac{1}{16\pi^2}A^2, \tag{26}$$

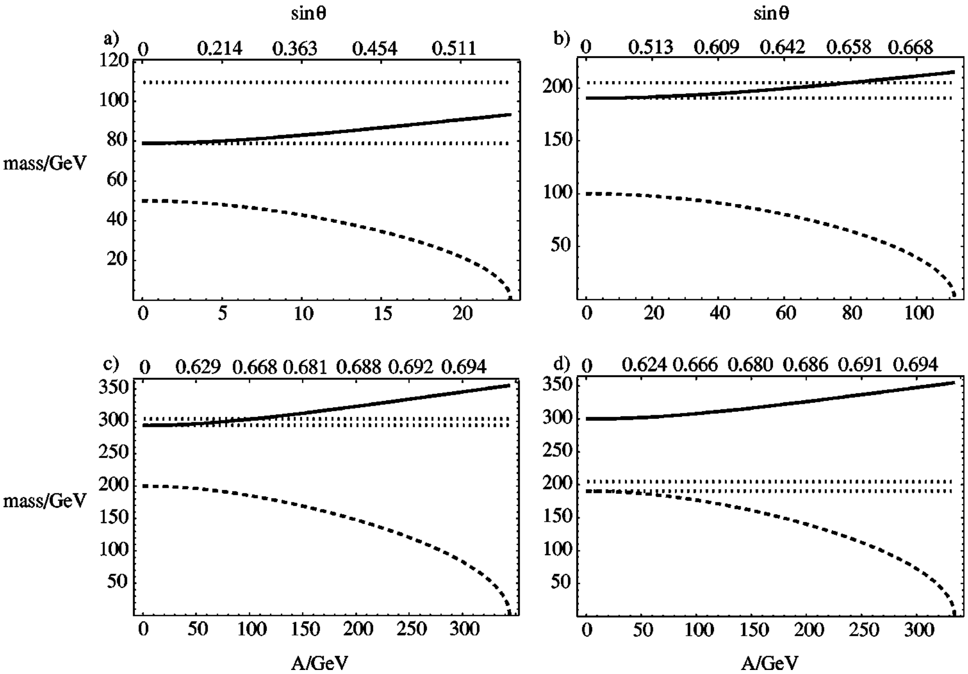

FIG. 1. Slepton mass spectra as a function of A for (a) $m_L = 100$ GeV, $m_R = 50$ GeV, (b) $m_L = 200$ GeV, $m_R = 100$ GeV, (c) $m_L = 300$ GeV, $m_R = 200$ GeV, and (d) $m_L = 200$ GeV, $m_R = 300$ GeV. The solid line is the mass of the heavier sneutrino, the dashed line that of the lighter sneutrino. The dotted lines are the masses of the sneutrino (lower dotted) and charged slepton (higher dotted) in the MSSM. Curves are drawn for $\tan\beta = 5$, and are relatively insensitive to $\tan\beta$.

$$\frac{dm_R^2}{dt} = \frac{2}{16\pi^2} A^2, \tag{27}$$

where $t = \ln(\mu^2/\mu_0^2)$. Since $\tilde{n}$ is a standard model singlet, there are no gaugino loops to drive its mass upward as we run the energy scale down from M_{Pl} to M_W. Likewise, there are new, sizable loop diagrams arising from the A terms (Fig. 2), which push the soft masses down. However, two states ($\tilde{l}$ and $\tilde{\nu}$) can propagate in the loop contributing to m_R, while only one ($\tilde{n}$) can propagate in the loop contributing to m_L, pushing $m_{\tilde{n}}$ down faster than $m_{\tilde{\nu}}$.

In summary, the mass matrix can have two *very* different mass eigenstates, and there can be particles that couple very much like $\tilde{\nu}$, but with suppressed couplings and masses unrelated to our expectations from the MSSM. The lightest is likely to be dominantly $\tilde{n}$, such that its mass is not restricted by Z decay data.

V. SNEUTRINO DARK MATTER

One of the appealing features of R-parity conserving supersymmetric theories with gravity-mediated supersymmetry breaking is that the lightest superpartner (LSP) is a good candidate for dark matter. Searches for superheavy hydrogen have ruled out a charged LSP, leaving the neutralino and the sneutrino as candidates for dark matter.

A number of direct searches for dark matter have been carried out [18–20] which essentially excluded sneutrino dark matter in the MSSM unless $m_{\tilde{\nu}} < 10$ GeV. However, as we have already discussed, measurements of the invisible width of the Z exclude such a light sneutrino. Within our framework, the Z width provides only a mild constraint, and we are free to explore the possibility of a light sneutrino dark matter candidate. A second, equally important point is that the $\sin^2\theta$ suppression of the light sneutrino coupling to the Z boson greatly reduces the sneutrino-nucleon cross section, making it possible even for a heavier $\tilde{\nu}$ to evade direct detection. Finally, if we include lepton number violation, the lightest sneutrino cannot scatter elastically via Z exchange [21], further diminishing the limits from direct searches. Sneutrino dark matter requires the presence of the A term of Eq. (8), and hence is linked to the other phenomenology of the interaction.

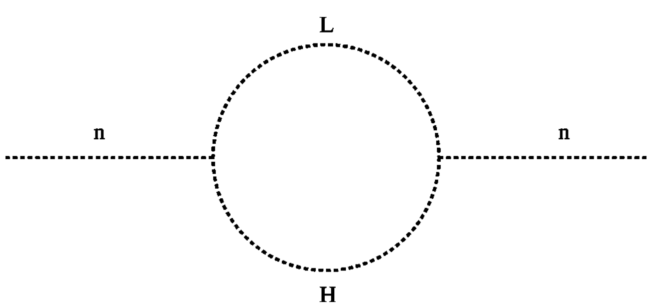

FIG. 2. Loops contributing to the running of $m_{\tilde{n}}$.

A. Dark matter without lepton number violation

1. Light sneutrinos

We can determine the relic density of light sneutrinos ($m_{\tilde{\nu}} < 10$ GeV) through standard methods [22]. The dominant annihilation process is through t-channel neutralino exchange. If the neutral W-ino exchange dominates, we find

$$\Omega_{\tilde{\nu}_1} h^2 \approx \left(\frac{M_{\tilde{W}}}{100\ \ \mathrm{GeV}}\right)^2 \left(\frac{0.19}{\sin\theta}\right)^4, \tag{28}$$

where h is the normalized Hubble parameter. We find it highly significant that for W-ino masses of roughly 100 GeV and angles roughly half of the limit from the invisible Z

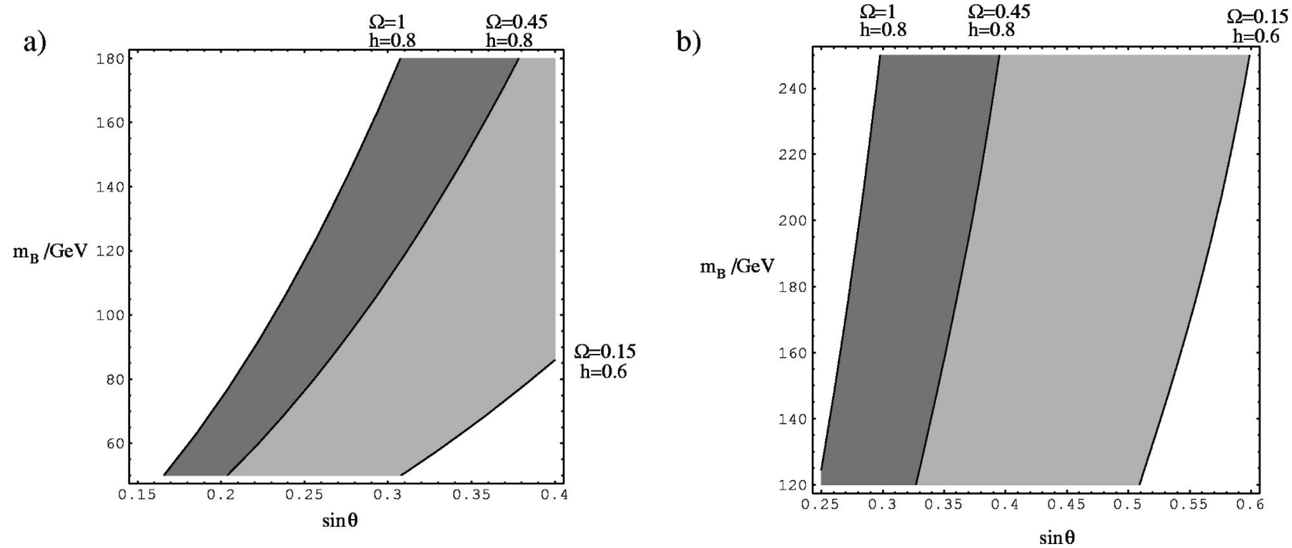

FIG. 3. Contours of Ωh^2, where h is the normalized Hubble parameter, as a function of the B-ino mass m_B (assuming GUT unification of gaugino masses) and $\sin\theta$. Both shaded regions yield relic densities below overclosure, with the lighter shaded region corresponding to values of Ω preferred by supernovae data [17]. In (a) we take $m_{\tilde{\nu}}=10$ GeV, and in (b) we take $m_{\tilde{\nu}}=100$ GeV, $A=20$ GeV, $\tan\beta=50$, and $m_h=115$ GeV. For (b), direct detection bounds are evaded only in the lepton-number-violating case.

width, we have a cosmologically interesting amount of sneutrino dark matter. In Fig. 3(a) we show the relic density of sneutrino dark matter considering all annihilation processes. If this scenario is correct, and $\tilde{\nu}_1$ is the dark matter, then the mixing angle must be near the limit from the invisible Z width measurements, making a future detection possible. In particular, such a sneutrino would almost certainly be seen in the upcoming CRESST experiment [23].

Relic sneutrinos captured by the sun will annihilate into neutrinos that can induce upward-going-muon events on earth [24]. The flux of these muons is constrained to be less than $10^{-14}\,\mathrm{cm}^{-2}\,\mathrm{s}^{-1}$ [25]. For a 10-GeV sneutrino with $\sin\theta=0.2$, we calculate a flux, using the formulas of Ref. [22], that is roughly three times this, assuming that all neutrinos produced are muon type when they reach the earth[7] (we find a much smaller flux due to capture by the earth itself). The actual muon flux could be suppressed depending on the flavor of the LSP sneutrino and on neutrino oscillation parameters; for instance, for an electron-type sneutrino and the small angle Mikheyev-Smirnov-Wolfenstein (MSW) solution to the solar neutrino problem, the flux would be smaller by roughly a factor of 1000.

2. Heavier sneutrinos

For heavier sneutrinos, the strongest current direct detection limit comes from CDMS [18], which, under the assumption of A^2 scaling, constrains the nucleon-relic cross section to be less than $(2-3)\times10^{-42}\,\mathrm{cm}^2$ for relic masses of $O(100$ GeV). The cross section for ordinary sneutrino-nucleus scattering is

$$\sigma=\frac{G_F^2}{2\pi}\mu^2((A-Z)-(1-4\sin^2\theta_W)Z)^2, \qquad (29)$$

where μ is the sneutrino-nucleus reduced mass. In our framework this cross section comes with an additional $\sin^4\theta$ suppression, implying that for $m_{\tilde{\nu}}$ much larger than the nucleon mass m_N, the CDMS constraint can be evaded by requiring

$$\sin^4\theta<2\pi\left(\frac{A}{(A-Z)-(1-4\sin^2\theta_W)Z}\right)^2\left(\frac{2\times10^{-42}\,\mathrm{cm}^2}{G_F^2 m_N^2}\right). \qquad (30)$$

Taking $A=73$ and $Z=32$ for Ge^{73} gives $\sin\theta<0.17$. The DAMA Collaboration [19] reported a positive signal corresponding to a relic-nucleon cross section of roughly $(2-10)\times10^{-42}\,\mathrm{cm}^2$ and a relic mass $\sim$30–100 GeV. In our framework this range in cross section corresponds approximately to $0.17<\sin\theta<0.25$.

In Fig. 4, we plot contours of Ωh^2 for $m_{\tilde{\nu}}=100$ GeV and a B-ino mass of 200 GeV. For this choice of parameters, and for large enough A, the dominant annihilation processes in the early universe are s-channel Higgs exchange into W^+W^- and Z pairs, which have cross sections proportional to $A^2\sin^2 2\theta$ rather than $\sin^4\theta$. These are also the dominant annihilation processes for sneutrinos trapped in the sun. This is important because the alternative process is annihilation directly into neutrinos via t-channel neutralino exchange, which would likely lead to a much larger signal at indirect detection experiments. Assuming that 1/3 of all neutrinos produced in the sun are muon type upon reaching the earth, we find that indirect detection constrains $\sin\theta<0.18$ for the parameters we have chosen, comparable to the CDMS constraint. The interesting relic abundances indicated in Fig. 4 lead us to conclude that the prospects for sneutrino dark matter with $m_{\tilde{\nu}}\sim100$ GeV are quite interesting in our model.

[7]For light sneutrinos and large A, the rate for sneutrino capture by the sun is dominated by Higgs exchange and the flux can be much larger. Here we assume $A\simeq10$ GeV, so that the capture rate is dominated by Z exchange.

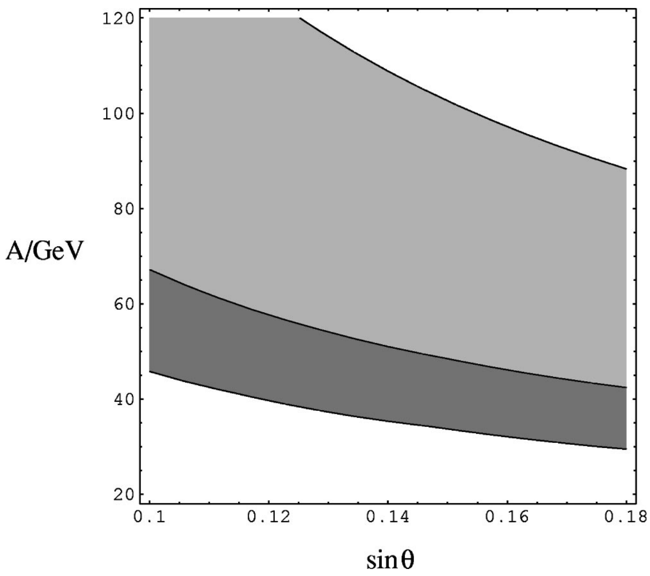

FIG. 4. Contours of Ωh^2 for $m_{\tilde{\nu}}=100$ GeV as a function of $\sin\theta$ and A. The meanings of the lighter and darker shaded regions are the same as in Figs. 3(a) and 3(b). We take $\tan\beta=50$, $m_h=115$ GeV, and a bino mass of 200 GeV, and we assume a GUT unification of gaugino masses.

B. Dark matter with lepton number violation

As previously explored [21], the presence of lepton number violation changes the limits from direct searches for dark matter significantly. In our model, lepton number violation can reside in the $m_{nn}^2\tilde{n}\tilde{n}$ term in the Lagrangian. Such a term could easily arise from dimension seven operators in the Lagrangian, such as

$$\left[\frac{X^\dagger X X^\dagger}{M^3}N^2\right]_D. \tag{31}$$

The presence of this lepton number violation splits the CP-even and -odd states $\tilde{\nu}_+$ and $\tilde{\nu}_-$. However, the coupling to the Z is off diagonal, i.e., $Z\tilde{\nu}_+\tilde{\nu}_-$. Consequently, for large enough $\Delta m=|m_{\tilde{\nu}_+}-m_{\tilde{\nu}_-}|$, the LSP sneutrino cannot scatter off nuclei via Z-exchange, eliminating constraints arising from CDMS, DAMA, and the Heidleberg-Moscow Ge experiment [20]. More precisely, the scattering is kinematically forbidden if $\Delta m>\beta_h^2 m_{\tilde{\nu}}m_A/2(m_{\tilde{\nu}}+m_A)$, where m_A is the mass of the target nucleus, and $\beta_h=10^{-3}$ for virialized halo particles on average. For example, taking $m_{\tilde{\nu}}=100$ GeV and a Ge target, we require $\Delta m>20$ keV. Since $\Delta m=m_{nn}^2/m_{\tilde{\nu}}$, this corresponds to $m_{nn}^2>(45\,\mathrm{MeV})^2$, which is of the order of what we expect from Eq. (31).[8]

The effects of lepton number violation in dark matter have been previously explored [21]. However, in the model previously proposed, there were no singlet sneutrinos, so the mass splitting Δm was required to be adequately large so as to suppress coannihilation between $\tilde{\nu}_+$ and $\tilde{\nu}_-$ via s-channel Z exchange. In our model, this process is further suppressed by $\sin^4\theta$ in the cross section, so that even with small mass splittings from dimension seven or higher operators, we can still generate a cosmologically interesting abundance.

Unlike Ref. [21], we now have the $A\tilde{\nu}\tilde{n}h$ coupling, which yields an extra contribution to the scattering of the lightest sneutrino off of nuclei via Higgs exchange. The coupling of Higgs bosons to nucleons is larger than just that from scattering off of valence quarks [26], but it is still quite small. Using the numerical value for the Higgs boson–nucleon coupling from Ref. [27], we find that the sneutrino-nucleon cross section obtained from Higgs exchange alone is[9]

$$\sigma=\left(\frac{A\sin\beta\sin 2\theta-(\sqrt{2}M_Z^2/v)\cos 2\beta\sin^2\theta}{100\,\mathrm{GeV}}\right)^2 \times\left(\frac{100\,\mathrm{GeV}}{m_{\tilde{\nu}}}\right)^2\left(\frac{115\,\mathrm{GeV}}{m_h}\right)^4(3\times 10^{-43}\,\mathrm{cm}^2). \tag{32}$$

For broad ranges of parameters this cross section falls well below the current direct detection limits. For instance, it is quite reasonable to consider $m_{\tilde{\nu}}\sim 100$ GeV even for values of $\sin\theta$ larger than those that allow one to evade CDMS in the lepton-number conserving case. Future experiments [28,29] should be able to probe an additional three orders of magnitude down from the present constraint, giving a significant probe of sneutrino dark matter over a considerable range of parameters, for both the lepton-number conserving and lepton-number violating cases.

Since the dominant annihilation process in the early universe is s wave, there is little dependence of the relic abundance on the sneutrino mass given that the sneutrino is relatively light ($\lesssim 30$ GeV). For these relatively light sneutrinos, Fig. 3(a) is still qualitatively applicable. For larger sneutrino masses, Z and Higgs pole effects or production of W and Z pairs can be relevant. The sneutrino relic density is shown in Fig. 3(b) for $m_{\tilde{\nu}}=100$ GeV and $A=20$ GeV.

The lepton number violating mass m_{nn}^2 can induce radiative corrections to the neutrino mass through neutralino loops. If the splitting Δm is too large, the possibility exists of generating neutrino masses radiatively which are large enough to affect the overall analysis. Such a possibility will be explored elsewhere [30]. For our purposes here, we limit ourselves to the case where the mixing between $\tilde{n}$ and $\tilde{l}$ is small enough to suppress this radiatively generated mass [which is why a relatively small value for A is taken in Fig. 3(b)].

We conclude that the possibility of evading direct detection through lepton number violation leads to another interesting version of sneutrino dark matter in our framework. Moreover, the elastic scattering of sneutrinos from nuclei via

[8]Somewhat higher values of Δm may be required to prevent indirect detection due to sneutrino capture and annihilation in the sun. Here we simply assume that Δm is large enough to evade indirect detection as well.

[9]We take the decoupling limit for the Higgs sector.

Higgs exchange is just below the current limits, and potentially detectable at upcoming dark matter searches.

VI. HIGGS DECAYS

The unsuppressed $A\tilde{l}\tilde{n}h_u$ coupling in our scenario can lead to interesting collider phenomena. If kinematically allowed, $\tilde{\nu}_1\tilde{\nu}_1^*$ is typically the dominant decay mode for the light Higgs boson. There is a similar situation in the MSSM [31]: provided the sneutrinos are sufficiently light and that $\tan\beta$ is not too close to 1, the $\tilde{\nu}\tilde{\nu}^*h$ coupling proportional to $M_Z\cos 2\beta$ leads to a partial width into sneutrinos that is larger than that into $b\bar{b}$ by two orders of magnitude. Assuming that the sneutrinos decay invisibly into $\chi_1^0\nu$ (or that the sneutrino itself is the LSP), a light Higgs boson that decays dominantly into sneutrinos would be very difficult to discover at the CERN Large Hadron Collider (LHC), leaving the Next Linear Collider (NLC) the opportunity for discovery through the process $e^+e^-\to Z^*\to Zh$.

In the MSSM, the Z width measurement, the theoretical bound $m_H\lesssim 130$ GeV, and the relation between the $\tilde{\nu}$ and $\tilde{l}_L$ masses constrain the region of parameter space in which the light Higgs bosons can decay into sneutrinos. For example, if in the future it becomes established that $m_{\tilde{l}_L}\gtrsim 105$ GeV, the decay $h\to\tilde{\nu}\tilde{\nu}^*$ will be ruled out in the MSSM.

In the present scenario, however, the sneutrino mass spectrum is expected to be quite different from that in the MSSM, as discussed in Sec. IV. Even if m_L (and therefore $m_{\tilde{l}_L}$) is quite large, it is still possible for the light Higgs decay into sneutrinos to be kinematically allowed. To explore this possibility quantitatively, we consider a single generation of sneutrinos with the mass matrix of Eq. (23), whose four free parameters are m_L, m_R, $\tan\beta$, and A. For simplicity we consider the case in which the splitting between m_L^2 and m_R^2 is generated by RG running from the GUT scale to the weak scale, and adopt $m_R^2=m_L^2-0.4A^2-0.5m_{1/2}^2$, with $m_{1/2}^2$, the universal gaugino mass, set to 100 GeV. The region of (m_L^2,A) parameter space in which the Z width constraint is met and $\Gamma(h\to\tilde{\nu}\tilde{\nu})>\Gamma(h\to b\bar{b})$ holds are displayed in Fig. 5 for $m_h=130$ GeV and $\tan\beta=2$ (the plot is very similar for high $\tan\beta$). As expected, there is a band in parameter space that yields invisible Higgs decays: for the region shown, the window for A is roughly 10 GeV at fixed m_L. The band persists for large m_L, with the window for A scaling as $\sim 1/m_L$.

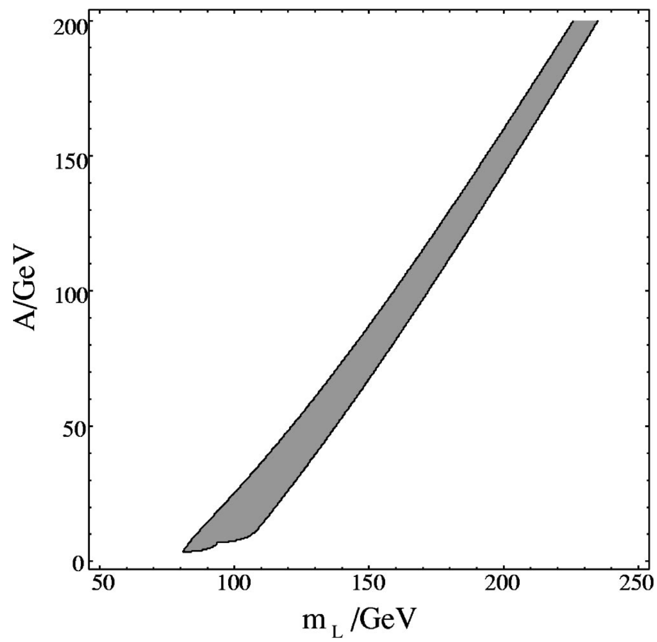

FIG. 5. The region of parameter space in which the Z width constraint is met and $\Gamma(h\to\tilde{\nu}\tilde{\nu})>\Gamma(h\to b\bar{b})$ holds, for $\tan\beta=2$ and $m_h=130$ GeV. We take the splitting between m_L and m_R to be generated by RG running from the GUT scale, as discussed in the text.

The $A\tilde{l}\tilde{n}h_u$ coupling can also alter the details of charged Higgs decays. If $m_{H^\pm}<m_t$, one can look for the charged Higgs bosons through the process $p\bar{p}\to t\bar{t}$, with one or both of the top quarks decaying into H^+b $(H^-\bar{b})$. The standard analysis exploits the fact that the charged Higgs boson is coupled most strongly to $\tau\nu$ (in contrast to the universally coupled W), and so if produced should lead to a surplus of τ's. This analysis has been applied at the Tevatron to establish lower bounds on the charged Higgs mass for $\tan\beta\lesssim 1$ and $\tan\beta\gtrsim 35$ [32]. The region of intermediate $\tan\beta$ will be only partially accessible to Run II of the Tevatron, but should be covered entirely at the LHC [33]. Similarly, the Higgs search at the LHC for high $\tan\beta$ employs the decay of heavy neutral Higgs states $H^0,A^0\to\tau^+\tau^-$, which can be suppressed due to the decay modes into sneutrinos.

It has already been pointed out that if the charged Higgs decays into SUSY particles, the analysis changes drastically [34]. If the decay into a charged slepton and a sneutrino is kinematically allowed, the MSSM Lagrangian term $-(g/\sqrt{2})M_W\sin 2\beta H^+\tilde{\nu}^*\tilde{l}_L$ ensures that $H^\pm\to\tilde{\nu}\tilde{l}_L$ dominates over the Yukawa-coupling-induced decay into $\tau\nu$ for small $\tan\beta$. Even for large $\tan\beta$, it is still possible for the decay into $\tilde{\nu}\tilde{\tau}$ to dominate due to the coupling $-(g/\sqrt{2})\times(m_\tau/M_W)(\mu+A_\tau\tan\beta)H^+\tilde{\nu}^*\tilde{\tau}_R$ (of course, μ and A_τ must not take on values that push the lightest charged slepton mass below the experimental bound). In this case, the excess τ's produced will have lower energy than when produced directly via $H^\pm\to\tau^+\nu$.

An unsuppressed $A\tilde{l}\tilde{n}h_u$ coupling introduces the added possibility of $H^\pm\to\tilde{l}_L\tilde{n}$ decays. If kinematically allowed, this is another process that can dominate over the direct decay into fermions for small $\tan\beta$. Once again, kinematical considerations for this decay are modified from the MSSM decays both because of the additional mass splitting between the sneutrinos and charged sleptons, and because the Z-width constraint does not apply to a sneutrino mass eigenstate that is chiefly right-handed. Thus it is conceivable that $H^\pm\to\tilde{l}_L\tilde{n}$ is the only SUSY decay mode allowed. Another important difference is that if one supposes that the flavor structure of the A coupling is similar to that of the neutrino masses, then one expects the charged Higgs to decay into $\tilde{\mu}\tilde{n}$

and $\tilde{\tau}\tilde{n}$ with similar probabilities, leading to an excess of both μ's and τ's over e's.[10]

Even if $m_{H^\pm} > m_t$, it is still possible for $H^\pm \to \tilde{l}_L \tilde{n}$ to be a dominant decay, for small to intermediate $\tan\beta$. The width is proportional to $(A\cos\beta)^2/m_H$, to be compared with $g^2 m_H/m_W^2[(m_t\cot\beta)^2+(m_b\tan\beta)^2]$ for $H^+ \to t\bar{b}$. To obtain a competitive rate requires A/m_H to be sizable, which is most easily achieved kinematically when m_H itself is large. In this regime, decays into other SUSY particles are also likely to be important.

VII. OTHER COLLIDER PHENOMENOLOGY FROM *A* TERMS

The unsuppressed A terms can have other interesting consequences for collider phenomenology, due both to their effect on the particle spectrum and because of the trilinear scalar vertex itself. Here we briefly consider a few examples, first for a visibly decaying $\tilde{\nu}$ and second for an invisible $\tilde{\nu}$. We have already seen that it is natural in our scenario to have large mass splittings among the various sneutrino states, so it is easily conceivable that there will be sneutrinos in both categories.

A. Visibly decaying sneutrinos

As in the standard framework of the MSSM, sneutrino decays into $\chi_2^0\nu$ and $\chi_1^\pm l^\mp$, if kinematically allowed, lead to final states with, e.g., $2l\not{E}_T$, $ljj\not{E}_T$, or $jj\not{E}_T$. For example, possible decay chains include[11] $\chi_2^0 \to \chi_1^0 Z^{(*)}/l\tilde{l}^{(*)}$ and $\chi_1^\pm \to \chi_1^0 W^{\pm(*)}/\nu\tilde{l}^{(*)}$, followed by $Z^{(*)} \to ll/jj$, $W^{\pm(*)} \to l\nu/jj$, and $\tilde{l}^{(*)} \to l\chi_1^0$. A possible signal for sneutrino pair production at the NLC is thus $4l\not{E}_T$. Sneutrino pair production should be distinguishable from neutralino pair production due to the different angular distributions and the different β dependences at threshold. As far as this signal is concerned, the distinctive feature of our model is the admixture of the gauge singlet $\tilde{n}$ in the sneutrino mass eigenstate. Decomposing the mass eigenstate as $\tilde{\nu}\sin\theta + \tilde{n}\cos\theta$, the sneutrino pair production rate will be suppressed by a factor $\sin^4\theta$ relative to its MSSM value. By performing a scan in energy one would be able to see the $4l\not{E}_T$ signal turn on for an isolated sneutrino mass eigenstate. Then, knowing the masses and mixings of the charginos and neutralinos, one could infer from the measured rate the value of $\sin\theta$, and demonstrate that the sneutrino produced is only partly left handed.

If a heavier $\tilde{\nu}_2$ state is sufficiently split from a lighter $\tilde{\nu}_1$, the unsuppressed A term induces the decay $\tilde{\nu}_2 \to \tilde{\nu}_1 h$, providing an interesting new way to produce Higgs particles. One might wonder how efficient this method of producing Higgs bosons would be at the LHC, through cascade decays such as $\tilde{q} \to q\chi^\pm$, $\chi^\pm \to l\tilde{\nu}_2$, $\tilde{\nu}_2 \to \tilde{\nu}_1 h$. The production cross section of squarks and gluinos at the LHC depends sensitively on their masses. Taking $m_{\tilde{q}} = 1.2 m_{\tilde{g}} = 300$ GeV, the cross section is ~ 2 nb at $\sqrt{s} = 14$ TeV [35], roughly 50 times larger than the cross section for gluon fusion Higgs boson production at that energy for $m_h \sim 100-130$ GeV in the decoupling limit [36], the regime we consider here. If $m_{\tilde{q}} = 1.2 m_{\tilde{g}} = 700$ GeV, the cross sections are comparable.

Assuming that A is sizable and that $m_{\tilde{\nu}_2} - m_{\tilde{\nu}_1} > m_h$, the question of whether or not there is an appreciable branching fraction for the cascades to produce $\tilde{\nu}_1 h$ depends largely on $m_{\tilde{\nu}_2}$. For simplicity consider the case of a B-ino-like χ_1^0 and W-ino-like χ_2^0 and $\chi_1^\pm$, and take $\tilde{\nu}_2$ and $\tilde{\nu}_1$ to be essentially left and right handed, respectively. If $m_{\tilde{\nu}_2} > m_{\chi_2^0}, m_{\chi_1^\pm}$, then $\tilde{\nu}_2$ will never be produced in the cascades, because the colored particles decay via $(\tilde{g} \to)\tilde{q} \to \chi_1^0 q/\chi_2^0 q/\chi_1^\pm q$. If $m_{\tilde{\nu}_2} < m_{\chi_2^0}, m_{\chi_1^\pm}$, then the branching fraction for producing $\tilde{\nu}_1 h$ is the product of three probabilities: first, the probability of the gluino or squark decaying into a neutralino or chargino heavier than $\tilde{\nu}_2$; second, the probability of that gaugino decaying into $\tilde{\nu}_2\nu/\tilde{\nu}_2 l$ rather than into a lighter gaugino or $\tilde{l}\nu/\tilde{l}l$; and third, the probability that $\tilde{\nu}_2$ decays into $\tilde{\nu}_1 h$ rather than $\chi_1^0\nu$.[12]

None of these probabilities is likely to be smaller than $\sim 1/\text{few}$ if $m_{\tilde{\nu}_2} < m_{\chi_2^0}, m_{\chi_1^\pm}$, so in this case the rate for Higgs boson production from $\tilde{\nu}_2$ decay at the LHC could easily be comparable to or even much larger than that due to $gg \to h$. Moreover, the cascade products (additional jets, and an energetic lepton from $\chi_1^\pm \to \tilde{\nu}_2 l$ decay, for instance), allow for detection via the $h \to b\bar{b}$ mode, so that the signal is further enhanced relative to $gg \to h \to \gamma\gamma$ by a factor of a 1000.[13]

For example, suppose that $m_{\chi_1^0} < m_{\tilde{\nu}_1} < m_{\tilde{\nu}_2} < m_{\tilde{l}} < m_{\chi_2^0}, m_{\chi_1^\pm} < m_{\tilde{q}} = 1.2 m_{\tilde{g}} = 300$ GeV.[14] Then the first probability is $\sim 1/2$, because $\tilde{q}_R$ couples to the B-ino, while $\tilde{q}_L$ "prefers" W-inos. The second probability is also $\sim 1/2$, because χ_0^2 and $\chi_1^\pm$ are equally likely to produce $\tilde{l}$ and $\tilde{\nu}_2$ and do not couple to χ_1^0 in the limit that it is pure B-ino. The third probability is determined by $\Gamma(\tilde{\nu}_2 \to \tilde{\nu}_1 h)/\Gamma(\tilde{\nu}_2 \to \chi^0\nu) \sim A^2/(m_{\tilde{\nu}_2}^2 g_1^2)$, and can be as large as $\sim 1/2$. So in this case, the branching fraction for producing $\tilde{\nu}_1 h$ could be $\sim 1/10$,

[10]Here we assume that one neutrino mass is hierarchically heavier than the others, and that $H^\pm \to \tilde{l}\tilde{n}$ is kinematically allowed for all flavors.

[11]For now we ignore the role an additional, lighter $\tilde{\nu}$ might play in these decay chains. For instance, χ_2^0 might decay invisibly into $\tilde{\nu}\nu$, as discussed below.

[12]Both $\text{BR}(\chi \to \tilde{\nu}_1\nu/\tilde{\nu}_1 l)$ and $\text{BR}(\tilde{\nu}_2 \to Z\tilde{\nu}_1)$ are suppressed if $\tilde{\nu}_1$ is essentially right-handed.

[13]Thanks to Ian Hinchcliffe for pointing this out.

[14]If gaugino mass unification is imposed for this mass ordering, then $m_{\tilde{g}}$ is forced to be much heavier, >700 GeV, in order for $m_h < m_{\tilde{\nu}_2} < m_{\chi_2^0}, m_{\chi_1^\pm}$ to be satisfied. In this case the squark and gluino production cross section is comparable, at best, to the $gg \to h$ cross section.

leading to a Higgs production rate about ten times larger than that from gluon fusion (not just five times larger, since each strong production event gives two sparticles that can potentially produce a Higgs boson). In fact, in this case the rate of Higgs *pair* production through $\tilde{\nu}_2$ decay is as large as the rate for $gg\to h$. These pair production events would lead to striking final states $b\bar{b}b\bar{b}llX$, with the invariant masses of both $b\bar{b}$ pairs equal to m_h. Note that even if $m_{\tilde{q}}\sim m_{\tilde{g}}\sim 700$ GeV, the rate for cascade Higgs production is only down from the $gg\to h$ rate by a factor of ~ 10, and would still likely allow for discovery because the signal is not suppressed by the $h\to\gamma\gamma$ branching ratio.

Even more striking is the possibility that $\tilde{\nu}_2$ and $\tilde{\nu}_1$ are the two lightest supersymmetric particles, with $m_{\tilde{\nu}_2}>m_{\tilde{\nu}_1}+m_h$. In this case, every squark and gluino produced yields a Higgs particle in its cascade. So for $m_{\tilde{q}}=1.2m_{\tilde{g}}=300$ GeV, gluon fusion would account for only one in every ~ 100 Higgs particles produced at the LHC.

Production of $\tilde{\nu}_2$'s and their subsequent decay could also be an interesting source of Higgs particles at the NLC. The rate of Higgs production through $e^+e^-\to\tilde{\nu}_2\tilde{\nu}_2$ is typically at least an order of magnitude lower than that due to $e^+e^-\to Zh$ and WW fusion for $\sqrt{s}=500$ GeV [37,38]. However, $e^+e^-\to\tilde{\nu}_2\tilde{\nu}_2$ could lead to sizable Higgs *pair* production. For example, for $\sqrt{s}=500$ GeV and $m_{\tilde{\nu}}=200$ GeV,

$$\sigma(e^+e^-\to hh\tilde{\nu}_1\tilde{\nu}_1)\simeq 6 \ \ \text{fb}\cos^4\theta(\text{BR}(\tilde{\nu}_2\to\tilde{\nu}_1 h))^2, \quad (33)$$

compared to a cross section of ~ 0.3–0.5 fb for the double Higgs-strahlung process $e^+e^-\to Zhh$ in the decoupling regime, for $m_h\sim 100$–130 GeV [39]. Especially if $\tilde{\nu}_2$ is lighter than all gauginos except for a B-ino-like state, it is easy to choose m_L, m_R, and A so that $\cos^4\theta(\text{BR}(\tilde{\nu}_2\to\tilde{\nu}_1 h))^2$ is not more than an order of magnitude suppression [there is even the possibility, as mentioned above, that $\text{BR}(\tilde{\nu}_2\to\tilde{\nu}_1 h)=1$]. Thus a possible signature of our scenario is an excess of events with missing energy plus two $b\bar{b}$ pairs whose invariant masses are equal to the Higgs boson mass, beyond the number expected from double Higgs-strahlung followed by $Z\to\nu\bar{\nu}$.

B. Invisible sneutrinos

The motivation for considering this case in our scenario is that the A terms suppress the masses of the lighter sneutrinos, making it more likely than in standard schemes that some $\tilde{\nu}$'s can either only decay invisibly, or not decay at all. Let us assume that χ_1^0 and $\tilde{\nu}$ are the two lightest supersymmetric particles. One immediate question is whether the clean trilepton signal from $\chi_1^{\pm}\chi_2^0$ production at hadron machines remains, since χ_2^0 has access to the invisible two body decay mode $\tilde{\nu}\nu$. However, provided $m_{\chi_2^0}>m_{\tilde{l}}$, χ_2^0 can also decay through the visible two body mode $\tilde{l}l$. If the branching ratio for this decay is not too small, the trilepton signal survives, because $\chi_1^{\pm}$ decays into $\tilde{\nu}l$ and possibly $\tilde{l}\nu$, if the latter is kinematically accessible. Especially interesting is the particular case where only the lightest $\tilde{\nu}$ mass eigenstate is lighter than $\chi_1^{\pm}$, and $m_{\chi_1^{\pm}}<m_{\tilde{l}}<m_{\chi_2^0}$. The flavor of the lepton produced in $\chi_1^{\pm}\to\tilde{\nu}l$ is correlated with the flavor of the light $\tilde{\nu}$. Moreover, it is reasonable to expect the lightest sneutrino to be coupled to the largest A term, and so, assuming that the flavor structure for the A terms resembles that of the neutrino masses (the connection between the two is most immediate in the sDirac case), the lightest $\tilde{\nu}$ is likely to be mixture of $\tilde{\nu}_\mu$ and $\tilde{\nu}_\tau$. In this case, $\sim 1/2$ of the leptons produced in the $\chi_1^{\pm}$ decays are μ's, while very few e's are produced, leading to roughly seven μ's for every four e's in the trilepton signal, assuming that the χ_2^0 decays produce equal numbers of each lepton flavor.

An invisible $\tilde{\nu}_1$ state can be produced along with a heavier, visibly decaying $\tilde{\nu}_2$ at e^+e^- colliders through s-channel Z exchange (and t-channel chargino exchange for $\tilde{\nu}_e$), provided that $\sin 2\theta$ is not too small. The decays of the heavier $\tilde{\nu}$ would lead to the signal $2l+\not{E}_T$ at the NLC. If the masses of heavier, visibly decaying sneutrinos have already been established through pair production, it should be possible to measure the mass of a light, invisible $\tilde{\nu}$ using the energy spectrum endpoints for the leptons produced in this process.

As mentioned above, if $m_{\chi_2^0}<m_{\tilde{l}}$, then the only two-body decay for χ_2^0 is into $\tilde{\nu}\nu$, so that both $\tilde{\nu}$ and χ_2^0 decay invisibly. The process $e^+e^-\to\gamma+\not{E}_T$ has been shown to be a feasible means for detecting the presence of these extra carriers of $\not{E}_T$ at the NLC [40].

Finally, suppose that $\tilde{\nu}$ and $\tilde{l}$, rather than χ_1^0 and $\tilde{\nu}$, are the lightest supersymmetric particles. Thus $m_{\tilde{l}}<m_{\chi_1^0}$, and χ_1^0 decays visibly into $\tilde{l}l$. Meanwhile, the NLSP $\tilde{l}$ has only three-body decays, into $\tilde{\nu}l'\nu'$ and $\tilde{\nu}jj$. In this case, a signal for slepton pair production is $ljj\not{E}_T$, a characteristic signature for chargino pair production (although the two cases are distinguishable by their angular distributions, for instance) [41].

VIII. FLAVOR CHANGING SIGNALS

In our framework, lepton flavor violating contributions to m_L^2 can arise at tree level due to the same spurion(s) responsible for the Dirac neutrino masses and A terms. This possibility exists in both of the scenarios discussed in Sec. II, but the connection between the flavor structure of m_L^2 and that of the neutrino masses is most direct for the sDirac case, which we therefore consider here for simplicity.

Suppose that X, a chiral superfield with $\langle X\rangle=\theta^2 F_X$, has the appropriate flavor structure to induce Dirac neutrino masses via

$$\frac{1}{M^2}[LX^{\dagger}NH_u]_D. \quad (34)$$

Then one can also write down the Lagrangian term

$$\frac{1}{M^2}[(L^\dagger X)(X^\dagger L)]_D, \tag{35}$$

giving potentially large lepton-flavor violating contributions to m_L^2. As discussed in Sec. III, X could alternatively be a product of multiple spurions, some of which conserve flavor and break supersymmetry, and others which do the opposite. There are additional contributions to m_L^2 from the spurion that generates the A terms. We will assume that the flavor structure of this spurion is identical to that of X: this is especially likely in the case that the supersymmetry-breaking piece of X is flavor conserving.

The contributions of Eq. (35) might lead, for example, to slepton oscillation signals at the NLC [42] or, as considered here, to rare lepton decays. Of course, similar contributions arise in more standard schemes as well: a flavor breaking spurion $\langle\phi\rangle/M=\lambda$ that generates Majorana neutrino masses through $(1/M)[\lambda_{ij}L_iHL_jH]_F$ can also appear in $(1/M^2)[L^\dagger\lambda^\dagger\lambda LZ^\dagger Z]_D$, where Z breaks supersymmetry but not flavor. If the lepton-flavor violating contributions to m_L^2 are not screened by much larger universal contributions, then a generic flavor structure for F_X in Eq. (35), or for λ in the standard case, leads to unacceptably large flavor violating signals. On the other hand, not every structure for F_X or λ leads to a phenomenologically acceptable neutrino mass matrix. In light of this, we briefly consider forms for F_X motivated by neutrino phenomenology, and estimate the flavor changing signals they induce. We will not have specific flavor symmetries in mind that motivate the textures we will consider. Moreover we will estimate only the lepton flavor violating signals due to the nonuniversal contributions to m_L^2, and make the simplifying assumption that all other potential sources of flavor violation (for instance, an A_e matrix that is not aligned with λ_e) vanish.

For our discussion we will assume that one Dirac neutrino is hierarchically heavier than the others, and we will also take there to be three N states, although this is not an important assumption. Using our freedom to choose a basis for the N's, we consider the leading order flavor structure:

$$F_X\sim\begin{pmatrix} 0 & 0 & 0\\ 0 & 0 & 0\\ 0 & \alpha & \beta\end{pmatrix}, \tag{36}$$

with α and β comparable. That is, there is large $\nu_\mu-\nu_\tau$ mixing as indicated by SuperKamiokande, but ν_e has only a small component in the heavy state, to satisfy the CHOOZ bound [14].

Since we assume that Eq. (35) gives contributions to m_L^2 as large as the universal ones, the form taken for F_X suggests that the 23 entry of m_L^2 will be comparable in size to the diagonal entries. In this case the branching ratio for the process $\tau\to\mu\gamma$ is near the current experimental limit for slepton masses near 100 GeV [43,44].

Another potential signal is $\mu\to e\gamma$. The size of the branching ratio depends on the higher order contributions to F_X and is highly model dependent. In the Abelian flavor symmetry models considered in Ref. [44], with right handed neutrinos integrated out above the flavor scale, both large angle MSW and vacuum oscillation solutions to the solar neutrino problem generally lead to too large a rate for $\mu\to e\gamma$, and even models compatible with the small angle MSW solution force the slepton masses above $\sim$500 GeV. Here we do not construct flavor models for light Dirac neutrinos but instead simply consider the texture

$$F_X\sim\begin{pmatrix} \epsilon & \epsilon & \epsilon\\ \epsilon & \epsilon & \epsilon\\ \epsilon & 1 & 1\end{pmatrix}, \tag{37}$$

with only the order of magnitude of each entry, and not its precise value, indicated. One finds in this case that $F_X^\dagger F_X$ has eigenvalues $\sim\epsilon^2$, ϵ^2, and 1, and mixing angles $\theta_{23}\sim 1$, $\theta_{13}\sim\epsilon$, and $\theta_{12}\sim 1$. Thus this case is most likely to correspond to either large angle MSW or vacuum oscillation solutions to the solar neutrino problem. For vacuum oscillations, we take $\epsilon^2\sim\Delta m_\odot^2/\Delta m_{atm}^2\sim 10^{-7}$. Since the 12 and 13 entries of $F_X^\dagger F_X$ are both of order ϵ, we obtain the order of magnitude relation

$$\frac{B(\mu\to e\gamma)}{1.2\times 10^{-11}}\sim 0.03\left(\frac{100\ \mathrm{GeV}}{\tilde{m}}\right)^4, \tag{38}$$

where $\tilde{m}$ is a typical slepton mass, the lightest neutralino is taken to be photinolike, and $m_{\tilde{\gamma}}^2/\tilde{m}=0.3$ [43,44]. This case is thus safe as far as $\mu\to e\gamma$ is concerned. Depending on the assortment of assumed order unity factors we have ignored, it is still possible, for light sparticle masses, that the branching ratio will be accessible to future experiments.

On the other hand, for the large angle MSW solution we should take $\epsilon^2\sim 10^{-2}$, leading to

$$\frac{B(\mu\to e\gamma)}{1.2\times 10^{-11}}\sim\left(\frac{700\ \mathrm{GeV}}{\tilde{m}}\right)^4, \tag{39}$$

where we again take $m_{\tilde{\gamma}}^2/\tilde{m}=0.3$. Thus $\mu\to e\gamma$ forces the sparticle masses to be heavy. One should keep in mind that these estimates have been obtained ignoring other possible sources of lepton flavor violation, and for a particular texture for F_X.

One choice for the higher order entries in F_X that is compatible with the small angle MSW solution to the solar neutrino problem is

$$F_X\sim\begin{pmatrix} \epsilon^2 & \epsilon & \epsilon\\ \epsilon^2 & \epsilon & \epsilon\\ \epsilon^2 & 1 & 1\end{pmatrix}, \tag{40}$$

leading to eigenvalues for $F_X^\dagger F_X\sim\epsilon^4$, ϵ^2, and 1, and mixing angles $\theta_{23}\sim 1$, $\theta_{13}\sim\epsilon^2$, and $\theta_{12}\sim\epsilon$. Choosing $\epsilon\sim 1/30$ leads to acceptable masses and mixings. Since the 12 and 13 entries of $F_X^\dagger F_X$ are order ϵ^2, $B(\mu\to e\gamma)$ is suppressed by roughly a factor of 10^4 relative to the large angle MSW case.

TABLE I. Possible scenarios which achieve $m_\nu = v^2/M$. We only allow the various couplings to take values in powers of the intermediate scale $m_I \approx 10^{11}$ GeV, as would occur if the couplings were generated in the supersymmetry breaking sector. m_N is the Majorana mass for the right-handed neutrino, and m_D is the Dirac mass coupling the left- and right-handed neutrinos. m_{LL} is the left-handed neutrino Majorana mass. $\times$ indicates that m_{LL} cannot occur at v^2 for the given m_N.

	Seesaw theories				Non-seesaw theories	
Mass scale ($M_{Pl}=1$)	Conventional seesaw		sMajorana		Conventional EFT	sDirac
1	m_N					
m_I		m_N				
$m_I^2 = v$	m_D		m_N			
m_I^3			m_D	m_N		
$m_I^4 = v^2$	m_{LL}	$\times$	m_{LL}	$\times$	m_{LL}	m_D

If large universal contributions to m_L^2 are present, they will suppress the effects of the flavor violating contributions induced by F_X at tree level. However, even if we ignore Eq. (35) entirely and take a universal form for m_L^2 at, say, the GUT scale, we still obtain potentially interesting flavor violating signals. This is because the A terms generate nonuniversal contributions to m_L^2 radiatively:

$$\delta m_L^2 = \frac{1}{8\pi^2} A^\dagger A \log(M_{GUT}/M_{SUSY}). \tag{41}$$

This effect was studied in Refs. [45,46] in the context of seesaw theories with a high scale for the right-handed neutrinos. In these models, the Dirac neutrino Yukawa couplings are sizable and generate additional non-universal contributions, leading to

$$\delta m_L^2 = \frac{1}{8\pi^2}(A^\dagger A + 3\lambda^\dagger \lambda m_{3/2}^2)\log(M_{GUT}/M_N) \tag{42}$$

for a universal scalar mass $m_{3/2}$ and a right handed neutrino scale M_N. Note that while in our framework only the A term contributions are present and not those from the Yukawa couplings, the logarithm is larger than in the seesaw case because the right handed neutrinos remain in the effective theory down to low energies.

To calculate rates for flavor changing processes due to Eq. (41) for a given set of MSSM parameters, one needs to know the A matrix. This ambiguity is at the same level as in Ref. [46], where the Dirac neutrino Yukawa couplings are unknown: the size of the largest coupling is fixed by atmospheric neutrino data only once the scale of the right handed neutrinos is specified (the A terms are chosen proportional to the Yukawas, with the scale set by the universal gaugino mass). Setting $M_N = 10^{13}$ GeV, the authors of Ref. [46] considered textures for λ based on Abelian symmetries and found promising signals for $\mu \to e\gamma$ and $\tau \to \mu\gamma$ for significant portions of parameter space. If we take $A = \lambda m_{3/2}$ in Eq. (41), then our framework yields very similar signals to those of Ref. [46] for a given choice of SUSY parameters and a given form for λ.

IX. CONCLUSIONS

The relationship $m_\nu \approx v^2/M$, where v is of the order of weak scale and M an ultraviolet cutoff, has been tremendously successful in describing the small mass of the neutrino. Whether arising from Planck-scale suppressed operators in an effective theory, or from a particular realization such as the seesaw mechanism, this has been generally interpreted to signify the presence of lepton-number violating physics at the scale M, well above the reach of laboratory high-energy physics. This belief is predicated upon the idea that the coupling of the neutrino to the lepton-number violating sector of physics is order one, as might be expected in a GUT seesaw, for instance.

The likelihood of a light Dirac neutrino has been discounted for decades. Given the observed mass scales for neutrinos from solar and atmospheric neutrino data, we would need a Yukawa coupling at $O(10^{-12})$ or smaller, which appears difficult to understand when the smallest known Yukawa λ_e is $O(10^{-5})$. There is a possibility to explain the needed small Yukawa coupling as a consequence of a new flavor symmetry G_F broken only very slightly. In previous efforts, the factorization of the symmetry group into $G_F \otimes SUSY$ has been extended to the factorization of the symmetry breaking itself: VEV's which break G_F are supersymmetry preserving, while VEV's breaking SUSY ($F \sim m_I^2 \sim v M_{Pl}$) are G_F conserving. If this is not the case, however, the Yukawa coupling can have an additional suppression factors in powers of $m_I/M_{Pl} \sim 10^{-8}$.

The lightness of the Higgs doublets in supersymmetric theories ($\mu \sim v$) suggests that such a factorization of symmetry breaking is inadequate. There should exist an additional symmetry group G which is broken in the supersymmetry breaking sector. Given that the Higgs boson is kept light by G, we may ask whether there might be other particles such as right-handed neutrinos, singlet under the standard model, also kept light by G.

One immediate consequence is that the physics responsible for the relationship $m_\nu \approx v^2/M$ is not occurring at the scale M, and, in particular, that the mass of the standard model singlet state may be much lighter than M—even as light as m_ν itself. The numerous new possibilities, employing only VEV's in integer powers of $\sqrt{F}=m_I$, are summarized in Table I. In particular, there are interesting possibilities of sDirac (a light Dirac neutrino) and sMajorana (a light Majorana neutrino with a weak-scale right-handed neutrino) scenarios in a single generation.

The idea that G, which protects m_N, is broken in the supersymmetry breaking sector is by no means purely philosophical. In theories in which G is broken by a supersymmetry preserving VEV, we typically expect *all* couplings of N supermultiplet to be highly suppressed. On the other hand, the case G is broken in the supersymmetry breaking sector immediately invites the possibility of unsuppressed A terms for right-handed scalar neutrinos, which radically alter the phenomenology of this scenario relative to previous ones.

For instance, Higgs physics can be modified drastically, both in production and decay. The mass spectrum of sleptons is dramatically altered and the presence of a light $m_{\tilde{\nu}}<45$ GeV sneutrino is permitted. Collider signatures can be changed dramatically. With or without lepton number violation, the sneutrino is revived as a dark matter candidate. All of these things are easily realized if G is broken by supersymmetry breaking VEV's.

Furthermore, if the supersymmetry breaking sector breaks flavor symmetries that are also broken in a separate supersymmetry conserving sector, we have a potentially new understanding of the large mixing in the neutrino sector. Since the VEV's of the fields X which break supersymmetry need not be aligned with flavor violating VEV's that preserve supersymmetry, there is no reason to expect the mass eigenstates of neutrinos to be aligned with those of charged fermions, although they may still have a hierarchical structure.

When viewed from the perspective of the μ problem, such a scenario is exceedingly natural. The presence of unsuppressed A terms provides not only exciting phenomenology, but also the promise that this scenario can be tested in the near future. While experiments will provide the ultimate test of these ideas, this framework provides exciting possibilities for connections between what previously have seemed separate elements of supersymmetric theories.

ACKNOWLEDGMENTS

This work was supported in part by the Director, Office of Science, Office of High Energy and Nuclear Physics, Division of High Energy Physics of the U.S. Department of Energy under Contract DE-AC03-76SF00098 and in part by the National Science Foundation under grant PHY-95-14797.

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