Undetected Higgs decays and neutrino masses in gauge mediated, lepton number violating models

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ABSTRACT: We discuss SUSY models in which renormalizable lepton number violating couplings hide the decay of the Higgs through $h \to \chi_1^0 \chi_1^0$ followed by $\chi_1^0 \to \tau jj$ or $\chi_1^0 \to \nu_\tau jj$ and also explain neutrino masses. This mechanism can be made compatible with gauge mediated SUSY breaking.

KEYWORDS: Higgs Physics, Neutrino Physics, Supersymmetric Standard Model, Supersymmetry Phenomenology.

Contents

1. The little hierarchy problem and its solutions

The minimal supersymmetric extension of the standard model (MSSM) predicts a light Higgs boson. While theory predicts a tree level Higgs mass which is at most the mass of the Z boson, the current experimental lower bound from LEP [1] is 114.4 GeV. Evading the experimental lower bound requires significant one loop corrections which can be achieved only by fine tuning of parameters [2]. This *little hierarchy problem*, while nowhere near as severe as the original gauge hierarchy problem, has excited a lot of theoretical interest. A variety of solutions has been proposed [3]. Some of them introduce new degrees of freedom to enhance the contributions to the Higgs mass, while others allow for non-standard decays of the Higgs, which would have been missed at LEP. The latter can greatly alter the experimental search strategy for the Higgs and supersymmetry (SUSY) at the LHC.

In this paper, we will pursue the suggestion of [4] and [5], that the Higgs can decay into light gauginos, which in turn decay via renormalizable lepton number violating couplings, into jets plus neutrinos. This decay would have been missed at LEP if the Higgs is between $85 - 100\,\text{GeV}$, and the gauginos are less than half the Higgs mass. We will discuss the FermiLab constraints on this scenario in this paper, as well as constraints on like sign dilepton decays of the Higgs, which inevitably accompany the decays with neutrinos. We find that there are plausible models in which the branching ratios for like sign dileptons are small enough to evade the strong, model independent, bounds from FNAL.

Our purpose is to go beyond the work of [5] in two ways. First of all, we incorporate the L violating mechanism for hiding the Higgs into gauge mediated SUSY breaking models. Secondly, we also exploit the lepton number violating operators to generate the neutrino masses. The seesaw mechanism for generating neutrino masses, requires one to introduce a new mass scale, an order of magnitude or so below the unification scale $M_U \sim 2 \times 10^{16}\,\text{GeV}$.[1] Renormalizable lepton number violating operators in SUSY can provide a natural alternative [7]. Our aim is to see whether this can be combined with gauge mediation and simultaneously hide the Higgs.

We will find that certain restrictions must be placed on the L violating operators in order to achieve all of these goals. Most of our considerations are quite general, but we will specialize to the Pentagon model [8] in order to investigate whether an appropriate discrete symmetry can be found, which automatically implies these restrictions. We will also assume, as in the Pentagon model, that we have a singlet chiral field S, with an SH_uH_d coupling. This changes the tree level prediction for the lightest Higgs, and allows us to have a Higgs obeying the model independent OPAL lower bound of 82 Gev for the Higgs mass, even when $\tan\beta \sim 1$. This is important, because we find that we need such a value of $\tan\beta$ in order for our model to predict a like sign dilepton signal compatible with the model independent bounds for the Tevatron. For $\tan\beta \sim 1$ it is natural for the lightest neutralino to decouple approximately from charged leptons. We will discuss this in detail below.

Our attitude toward the magnitude of the possible L violating operators is influenced by our knowledge of the Yukawa couplings in the standard model. Many of these are surprisingly small. Given the strong constraints on flavor changing neutral currents, we think that the most plausible explanation of Yukawa textures is the Froggatt-Nielsen mechanism [9] operating near the unification scale. It then seems clear that the flavor structure of L violating operators will be similarly constrained. Rather than trying to formulate a full high energy theory of these textures, we merely take away the lesson that dimensionless L violating couplings might be anomalously small, and that one of them might be much larger than all the others.

The MSSM also contains dimension two L violating operators, analogous to the μ term, with H_d replaced with a linear combination of L_i. Clearly, an explanation of the magnitude of the dimension two parameter is necessary to a complete low energy theory. We will adopt the philosophy of the NMSSM, in which this parameter is the vacuum expectation value (VEV) of a low energy singlet, and the bare dimension two couplings are forbidden by a

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[1]It has been suggested that this scale arises naturally, as $M_{\text{seesaw}} = \frac{M_U^2}{m_P}$ [6].

discrete symmetry.

Other recent analyses of hard to find Higgs decays can be found in [10].

2. Constraints on $h^0 \to \chi^0\chi^0 : \chi^0 \to (\tau, \nu_\tau)jj$

We want to investigate the LEP bound on the Higgs mass in the MSSM where the lightest Higgs boson is produced by Higgs-strahlung of the Z boson (or maybe through Z or W-fusion processes). The cascade decay we are interested in consists of the decay of the lightest Higgs to two next-to-lightest SUSY particle (NLSP) neutralinos followed by an R-parity violating (RPV) decay of each neutralino to one third generation lepton plus two quarks. We do the computation in the narrow-width limit where the cascade is divided into a two-body decay and two three-body decays.

2.1 BR of the lightest Higgs to two neutralinos

The partial decay width $\Gamma(h^0 \to \chi_i^0\chi_j^0)$ is given by

$$\Gamma(h^0 \to \chi_i^0\chi_j^0) = \frac{\lambda^{1/2}(m_{h^0}^2, m_{\chi_i^0}^2, m_{\chi_j^0}^2)}{16\pi m_{h^0}^3 \times 2^{\delta_{ij}}} \left(2|Y^{ij}|^2(m_{h^0}^2 - m_{\chi_i^0}^2 - m_{\chi_j^0}^2)\right.$$
$$\left. -2[(Y^{ij})^2 + (Y^{ij*})^2]m_{\chi_i^0}m_{\chi_j^0}\right) \qquad (2.1)$$

where

$$\lambda(x, y, z) = x^2 + y^2 + z^2 - 2xy - 2xz - 2yz \qquad (2.2)$$

$$Y^{ij} = \frac{1}{2}(-N_{i3}^* \sin\alpha - N_{i4}^* \cos\alpha)(gN_{j2}^* - g'N_{j1}^*) + \{i \leftrightarrow j\}. \qquad (2.3)$$

Here N diagonalizes the neutralino mass matrix M_{χ^0} which can be written at tree level as

$$M_{\chi^0} = \begin{pmatrix} M_1 & 0 & -g'v_d/\sqrt{2} & g'v_u/\sqrt{2} \\ 0 & M_2 & gv_d/\sqrt{2} & -gv_u/\sqrt{2} \\ -g'v_d/\sqrt{2} & gv_d/\sqrt{2} & 0 & -\mu \\ g'v_u/\sqrt{2} & -gv_u/\sqrt{2} & -\mu & 0 \end{pmatrix} \qquad (2.4)$$

where $N^* M_{\chi^0} N^{-1} = \text{diag}(m_{\chi_1^0}, m_{\chi_2^0}, m_{\chi_3^0}, m_{\chi_4^0})$ with $|m_{\chi_1^0}| < |m_{\chi_2^0}| < |m_{\chi_3^0}| < |m_{\chi_4^0}|$.

The total decay width is expected to be dominated by decays of the lightest Higgs to neutralinos (when kinematically allowed), thus the branching ratio can be approximated as

$$\text{BR}(h^0 \to \chi_i^0\chi_j^0) = \frac{\Gamma(h^0 \to \chi_i^0\chi_j^0)}{\Gamma(h^0 \to \text{all})} \sim \frac{\Gamma(h^0 \to \chi_i^0\chi_j^0)}{\Gamma(h^0 \to \text{SM}) + \Gamma(h^0 \to \text{neutralinos})}. \qquad (2.5)$$

2.2 BR of the neutralino to one lepton plus two quarks

The decay of the neutralino to one lepton plus two quarks occurs through the R-parity violating vertex $\lambda'_{ijk}\epsilon_{ab}L_i^a Q_j^b \bar{D}_k \subset \mathcal{W}$. Since squarks are assumed much heavier than sleptons, decays with off-shell squarks are sub-dominant contributions to the partial decay widths. Moreover, assuming there is no mixing in the sfermion sector $\tilde{f}_L$ and $\tilde{f}_R$ are mass

particle	Y^H	T_3	Q
ℓ	$-\frac{1}{2}$	$-\frac{1}{2}$	-1
ν	$-\frac{1}{2}$	$\frac{1}{2}$	0

Table 1: Hypercharges for leptons and neutrinos.

eigenstates. This reduces the number of Feynman diagrams since only left-handed sleptons and sneutrinos are relevant to the R-parity violating vertex $\lambda'_{ijk}\epsilon_{ab}L_i^a Q_j^b \bar{D}_k$. Thus only decays with off-shell $\tilde{\ell}_{Li}$ or $\tilde{\nu}_i$ are possible. Finally, since the kinematically allowed final state standard model fermions (for a NSLP neutralino with mass $m_{\chi_1^0} \sim 30\,\text{GeV}$ only the top quark is excluded as a final state fermion) are much lighter than any sparticles, one can compute the partial decay widths in the limit of vanishing fermion masses. This introduces a maximal error of the order $\mathcal{O}\left(\frac{m_b}{m_{\chi_1^0}}\right) \sim 0.13$ for a NLSP neutralino with mass $m_{\chi_1^0} \sim 30\,\text{GeV}$. With these assumptions, the partial decay width computations simplify greatly and one can get analytical results.

Thus, with these assumptions, the partial decay widths $\Gamma(\chi_l^0 \to \ell_i u_j \bar{d}_k) = \Gamma(\chi_l^0 \to \bar{\ell}_i \bar{u}_j d_k)$ and $\Gamma(\chi_l^0 \to \nu_i d_j \bar{d}_k) = \Gamma(\chi_l^0 \to \bar{\nu}_i \bar{d}_j d_k)$ are

$$\Gamma(\chi_l^0 \to \ell_i u_j \bar{d}_k) = \frac{N_c m_{\chi_l^0}(|c_1|^2 + |c_2|^2)}{1024\pi^3}\left[6\rho - 5 + 2(\rho-1)(3\rho-1)\ln\left(\frac{\rho-1}{\rho}\right)\right] \quad (2.6)$$

$$\Gamma(\chi_l^0 \to \nu_i d_j \bar{d}_k) = \frac{N_c m_{\chi_l^0}|c_1|^2}{1024\pi^3}\left[6\rho - 5 + 2(\rho-1)(3\rho-1)\ln\left(\frac{\rho-1}{\rho}\right)\right] \quad (2.7)$$

where

$$c_1 = \sqrt{2}\lambda'_{ijk}(gT_3^{f_i}N_{l2}^* + g'Y_{f_i}^H N_{l1}^*) \quad (2.8)$$

$$c_2 = \lambda'_{ijk}\frac{m_{\ell_i}}{v_d}N_{l3} \quad (2.9)$$

$$\rho = \left(\frac{m_{\tilde{f}_i}}{m_{\chi_l^0}}\right)^2. \quad (2.10)$$

Here $N_c = 3$ is the number of colors and Martin's notation is used for the hypercharge, i.e. $Q = T_3 + Y^H$ [11]. Moreover, the first term in c_1 represents the fermion/sfermion coupling to the wino, the second term in c_1 represents the fermion/sfermion coupling to the bino and c_2 represents the fermion/sfermion coupling to the higgsino. Table 1 reviews the needed hypercharges.

Since the dominant decay of the lightest Higgs is expected to be $h^0 \to \chi_1^0 \chi_1^0$, the NLSP neutralino decay is the one of interest. The total decay width for the NLSP neutralino is dominated by the kinematically allowed R-parity violating vertex processes discussed above,

$$\Gamma(\chi_1^0 \to \text{all}) \sim \sum_{i,j,k}\left[\Gamma(\chi_1^0 \to \ell_i u_j \bar{d}_k) + \Gamma(\chi_1^0 \to \nu_i d_j \bar{d}_k) + \text{antiparticles}\right] \quad (2.11)$$

$$= 2\sum_{i,j,k}\left[\Gamma(\chi_1^0 \to \ell_i u_j \bar{d}_k) + \Gamma(\chi_1^0 \to \nu_i d_j \bar{d}_k)\right] \quad (2.12)$$

particle	χ_1^0	χ_2^0	χ_3^0	χ_4^0	$\chi_1^\pm$	$\chi_2^\pm$
mass (in GeV)	30	125	150	300	105	295

Table 2: Neutralino and chargino masses (in GeV) for $M_1 = 50$ GeV, $M_2 = 250$ GeV, $\mu = +150$ GeV and $\tan\beta = 1$.

where the sums over u_j is limited to $j = \{1, 2\}$ since the top quark is not kinematically allowed. Thus the relevant branching ratios are

$$\mathrm{BR}(\chi_1^0 \to \ell_i u_j \bar{d}_k) = \frac{\Gamma(\chi_l^0 \to \ell_i u_j \bar{d}_k)}{\Gamma(\chi_1^0 \to \mathrm{all})} \tag{2.13}$$

$$\mathrm{BR}(\chi_1^0 \to \nu_i d_j \bar{d}_k) = \frac{\Gamma(\chi_l^0 \to \nu_i d_j \bar{d}_k)}{\Gamma(\chi_1^0 \to \mathrm{all})} \tag{2.14}$$

with the same results for final state antiparticles. Notice that the neutralino decay to the gravitino (mass $m_{3/2}$),

$$\Gamma(\chi_1^0 \to X\tilde{G}) = \frac{m_{\chi_1^0}^5}{48\pi m_P^2 m_{3/2}^2}\left(1 - \frac{m_X^2}{m_{\chi_1^0}^2}\right)^4, \tag{2.15}$$

is sub-dominant in our direct gauge mediation model. This will be quantified in the next section.

2.3 Total branching ratio in the regime $\tan\beta \sim 1$: numerical example

The total branching ratio for one particular final state in the cascade decay of interest is simply the product of the appropriate branching ratios (in the narrow-width limit). For example, for the cascade decay $h^0 \to \chi_1^0 \chi_1^0 \to \ell_{i_1} u_{j_1} \bar{d}_{k_1} \ell_{i_2} u_{j_2} \bar{d}_{k_2}$ the total branching ratio is

$$\mathrm{BR}(h^0 \to \ell_{i_1} u_{j_1} \bar{d}_{k_1} \ell_{i_2} u_{j_2} \bar{d}_{k_2}) = \mathrm{BR}(h^0 \to \chi_1^0 \chi_1^0)\mathrm{BR}(\chi_1^0 \to \ell_{i_1} u_{j_1} \bar{d}_{k_1})\mathrm{BR}(\chi_1^0 \to \ell_{i_2} u_{j_2} \bar{d}_{k_2}). \tag{2.16}$$

Our interest lies in decays with final state tau leptons and tau neutrinos since these processes have not been studied extensively at LEP. We will evaluate these branching ratios in the limit $\tan\beta \sim 1$, as is predicted in the Pentagon model. We do this here because it turns out that in this regime, the like sign dilepton contribution to Higgs decay is naturally suppressed. The FNAL bounds on this process will be more difficult to satisfy for larger values of $\tan\beta$.

Following [4] we want the lightest neutralino χ_1^0 to be mostly bino. This can be achieved with $M_1 = 50$ GeV, $M_2 = 250$ GeV and $\mu = +150$ GeV which lead to the masses shown in table 2 and the following mixing matrix, and mixing matrix

$$N = \begin{pmatrix} -0.89 & 0.15 & -0.30 & 0.30 \\ -0.44 & -0.48 & 0.54 & -0.54 \\ 0 & 0 & 0.71 & 0.71 \\ -0.09 & 0.86 & 0.35 & -0.35 \end{pmatrix}. \tag{2.17}$$
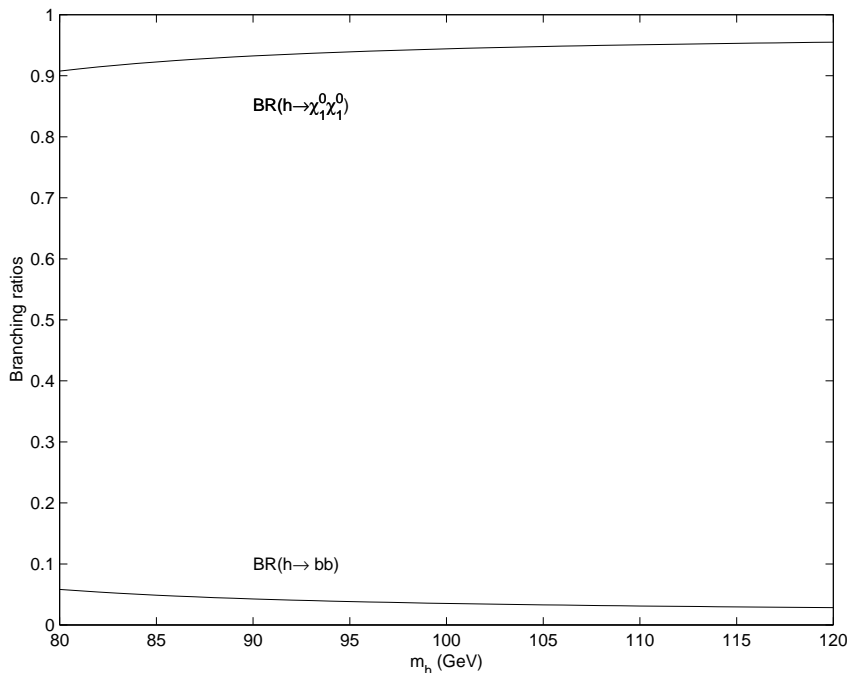
Figure 1: Branching ratio of $h^0 \to \chi_1^0 \chi_1^0$ (upper line) and $h^0 \to b\bar{b}$ (lower line) as a function of the Higgs mass (in GeV) for $M_1 = 50\,\text{GeV}$, $M_2 = 250\,\text{GeV}$, $\mu = +150\,\text{GeV}$, $\tan\beta = 1$ and $\alpha = -\frac{\pi}{8}$.

particle	$\tilde{\ell}_1 = \tilde{e}$	$\tilde{\ell}_2 = \tilde{\mu}$	$\tilde{\ell}_3 = \tilde{\tau}$	$\tilde{\nu}_1 = \tilde{\nu}_e$	$\tilde{\nu}_2 = \tilde{\nu}_\mu$	$\tilde{\nu}_3 = \tilde{\nu}_\tau$
m (in GeV)	73	94	81.9	75	75	75

Table 3: PDG bounds on the slepton masses and appropriate sneutrino masses.

Close to the Higgs decoupling limit this leads to $\text{BR}(h^0 \to \chi_1^0 \chi_1^0) \sim 0.9$ for m_{h^0} between $85 - 100\,\text{GeV}$. Figure 1 shows the branching ratios $\text{BR}(h^0 \to \chi_1^0 \chi_1^0)$ and $\text{BR}(h^0 \to b\bar{b})$ as a function of the Higgs mass for a mixing angle $\alpha = -\frac{\pi}{8}$. Here, since $\tan\beta = 1$, we go slightly away from the Higgs decoupling limit in order to satisfy the experimental bound on $\xi^2 \text{BR}(h^0 \to b\bar{b})$ with $\xi = \sin(\beta - \alpha)$ for m_h as low as $85\,\text{GeV}$ [1].

Using PDG bounds [12] on the masses of the sleptons and sneutrinos[2] (see table 3) and appropriate values for the R-parity violating coupling and for the R-parity violating coupling [7]

$$\lambda'_{ijk} = 0 \;\forall\; i = \{1,2\}, \; j,k = \{1,2,3\} \tag{2.18}$$

$$(\lambda'_{3jk}) = \begin{pmatrix} 0.001 & 0.001 & 0 \\ 0.001 & 0.001 & 0 \\ 0 & 0 & 0 \end{pmatrix} \tag{2.19}$$

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[2]For $\tan\beta = 1$, theory forces the sneutrinos to be almost degenerate with the sleptons. However from the non-SM invisible width of the Z-boson sneutrinos could be as light as 45 GeV.

one obtains $\text{BR}(\chi_1^0 \to \tau u_j \bar{d}_k) = 0.019 \ \forall \ j,k = \{1,2\}$, $\text{BR}(\chi_1^0 \to \nu_\tau d_j \bar{d}_k) = 0.11 \ \forall \ j,k = \{1,2\}$ and zero otherwise. The difference between the branching ratios comes from the mixing matrix N and the quantum numbers T_3 and Y^H. Indeed, since the sneutrino is almost degenerate with the stau one can forget about the ρ-dependent part of the decay width and focus on the c_1-bino and wino contributions to the decay width. The c_1-bino contributions to the branching ratio for τ and ν_τ are the same *but* the c_1-wino contributions to the branching ratio for τ and ν_τ have opposite sign due to the weak isospin charges. From the mixing matrix one can see that the contributions partly cancel for τ and add up for ν_τ. Though smaller, the c_1-wino contributions are about a third of the c_1-bino contributions which leads to a suppression

$$\frac{\text{BR}(\chi_1^0 \to \tau u_j \bar{d}_k)}{\text{BR}(\chi_1^0 \to \nu_\tau d_j \bar{d}_k)} \sim \frac{|c_{1,\tau}|^2}{|c_{1,\nu_\tau}|^2} \sim \frac{(-1+3)^2}{(1+3)^2} \sim \frac{1}{4}. \tag{2.20}$$

This kind of suppression is generic in the region of the parameter space where the lightest neutralino is light and mostly bino and the lightest chargino satisfies the lower bound on its mass since the mixing matrix does not change by much.

Notice also that the c_2-higgsino contribution to the charged lepton decay width is small due to the lepton mass suppression factor in the limit $\tan\beta \sim 1$ and can thus be neglected. This leads to an even bigger decoupling of the neutralinos to the charged leptons and more easily evades the FNAL bounds.

As shown before the branching ratios for antiparticles in the final state are the same. Thus the total branching ratios for the following cascade decays are

$$\text{BR}(h^0 \to \tau\tau + 4\,\text{jets}) \sim 0.0055 \tag{2.21}$$

$$\text{BR}(h^0 \to \tau\nu_\tau + 4\,\text{jets}) \sim 0.03 \tag{2.22}$$

$$\text{BR}(h^0 \to \nu_\tau\nu_\tau + 4\,\text{jets}) \sim 0.17 \tag{2.23}$$

with $m_{h^0} \sim 90\,\text{GeV}$ and $\alpha = -\frac{\pi}{8}$. Here any group of particles associated to the neutralino decay products can be changed to its antiparticle counterpart without changing the branching ratios. Thus one obtains like sign ditau (or di-antitau) events with the same branching ratio then tau-antitau events. The total decay width of the Higgs and the NLSP neutralino are

$$\Gamma(h^0 \to \text{all}) \sim 0.02\,\text{GeV} \tag{2.24}$$

$$\Gamma(\chi_1^0 \to \text{all}) \sim 0.11 \cdot 10^{-10}\,\text{GeV} \tag{2.25}$$

and this may lead to displaced vertices since $c\tau_{\chi_1^0 \to \text{all}} \sim 19\,\mu\text{m}$.

As mentioned previously, the neutralino to gravitino decay is sub-dominant, even for small R-parity violating couplings. Indeed the decay widths ratio is parametrically

$$\frac{\Gamma(\chi_1^0 \to \gamma\tilde{G})}{\Gamma(\chi_1^0 \to (\tau,\nu_\tau)jj)} \sim \frac{64\pi^2 m_{\chi_1^0}^4 |N_{11}|^2}{3N_c m_P^2 m_{3/2}^2 (g\lambda')^2} \left[6\rho - 5 + 2(\rho-1)(3\rho-1)\ln\left(\frac{\rho-1}{\rho}\right)\right]^{-1}. \tag{2.26}$$

This is of order 10^{-2} and leads to a neutralino to gravitino branching ratio of order 10^{-3} and thus can be safely neglected. Here the decay product is mostly photon and the gravitino mass is assumed to be $m_{3/2} \sim 1\,\text{eV}$. Cosmological constraints on light gravitinos restrict their mass to be $< 10 - 20\,\text{eV}$. Within the Pentagon model, the additional hypothesis of Cosmological SUSY Breaking gives $m_{3/2} \sim \Lambda^{1/4} \sim 10^{-3}\,\text{eV}$. For such super-light gravitinos, the gravitino decay channel of the gaugino, which is ruled out experimentally for light gauginos, would dominate, and our scenario would not be viable.

2.4 $\tan\beta \sim 1$ and the little hierarchy problem

The reader familiar with the supersymmetric standard model may be a bit confused at this point. We have invoked R-parity violating couplings to weaken the experimental bound on the Higgs mass, in order to avoid the little hierarchy problem. However, our discussion of like sign dilepton constraints used the Pentagon model relation $\tan\beta \sim 1$ to motivate the decoupling of light neutralinos from charged leptons. In the MSSM, the lightest Higgs mass vanishes at tree level when $\tan\beta = 1$. Fortunately, our model is not even approximately the MSSM. If we compute the Higgs sector potential neglecting Pentagon model corrections to the Kahler potential of S (which is plausible if g_S is small), we find the potential for the scalar components of neutral chiral superfields

$$V = \frac{g_1^2 + g_2^2}{8}(|H_u|^2 - |H_d|^2)^2 + |g_\mu S|^2(|H_u|^2 + |H_d|^2) + |kg_S\Lambda_5^2 + g_\mu H_u H_d + g_T S^2|^2, \quad (2.27)$$

where k is a strong Pentagon interaction correction, of order 1. This potential has a hypercharge Goldstone boson if the Higgs fields have nonzero VEVs, but all the other scalars are massive, with mass determined by a combination of the couplings and the large scale $\Lambda_5 \sim 1\,\text{TeV}$. To obtain Higgs VEVs of the right order of magnitude, we must take g_S small, but the other Yukawa couplings are bounded only by perturbative unification.[3] The latter constraint probably does not allow one to make the Higgs mass large enough to evade the conventional bounds, but there is no problem in evading the model independent bounds of $\sim 82\,\text{GeV}$ on the Higgs mass. R-parity violating Higgs decays then make the conventional bound irrelevant.

2.5 Gaugino mass relationships and the chargino mass bound

There is a strict lower mass bound of $102.7\,\text{GeV}$ set on charginos that decay through RPV operators. We note that if we assume the minimal gauge mediated prediction for the gaugino mass ratio, a chargino of $102.7\,\text{GeV}$ would make it impossible to have neutralinos less than $50\,\text{GeV}$. In order to allow neutralinos of less than half the Higgs mass and satisfy the chargino mass bound, we must alter the minimal gauge mediation predictions for gaugino masses. We present here a simple method proposed by [13] of adding to the hidden sector multiple scalar fields which get supersymmetric and non-supersymmetric masses, and coupling them with different strengths to parts of a single 10 and $\overline{10}$ of messengers. The

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[3]In fact, the Pentagon model, or any other model with 5 messengers, and a TeV scale threshold, has problems with perturbative unification, irrespective of the size of these Yukawas.

Decay channel	Experiment	Bound	Search
–	Tevatron	< 10 events for $1\,\mathrm{fb}^{-1}$	$p\overline{p}$ collisions
$bbqq+\not{E}_T$	Tevatron	$\sigma < 3$ pb	$t\overline{t}$ production cross-section
$bbbb+\not{E}_T$	LEP	$m_{h^0} < 84.5\,\mathrm{GeV}$ (worst case)	ZH_2H_2 cascade decay
$q\overline{q}'q\overline{q}'+\not{E}_T$	LEP	$m_{h^0} > 105\,\mathrm{GeV}$	WW^* with invisible Z decay
–	LEP	$\sigma < 1$ pb	direct neutralino search

Table 4: Higgs mass and production cross-section bounds for various searches.

superpotential is thus

$$W = r_i X_i u\overline{u} + \gamma_i X_i q\overline{q} + \lambda_i X_i l\overline{l} \qquad (2.28)$$

for

$$X_i = x_i + \theta^2 F_i \qquad (2.29)$$

The resulting gaugino eigenstates are determined by three mass parameters instead of a single mass parameter in the minimal case. The parameters are,

$$\Lambda_l = \frac{\lambda_i F_i}{\lambda_i x_i}, \quad \Lambda_q = \frac{\gamma_i F_i}{\gamma_i x_i}, \quad \Lambda_u = \frac{r_i F_i}{r_i x_i} \qquad (2.30)$$

and the resulting gaugino mass parameters are

$$M_1 = \frac{1}{2}\frac{\alpha_1}{4\pi}\left(\frac{4}{3}\Lambda_q + 2\Lambda_l + \frac{8}{3}\Lambda_u\right), \quad M_2 = \frac{3}{2}\frac{\alpha_2}{4\pi}\Lambda_q, \quad M_3 = \frac{1}{2}\frac{\alpha_3}{4\pi}(\Lambda_u + 2\Lambda_q) \qquad (2.31)$$

Here we have changed the ordinary gaugino mass ratio with minimal extra structure and it is easy to get a M_1 much lighter than M_2.

We note also that within the Pentagon model, we do not expect minimal gauge mediation predictions for gaugino masses to be valid.

3. Bounds on jets $+ \not{E}_T$

We will now assess the effects of current Tevatron and LEP searches on the viability of our Higgs decays. In particular there are five searches with relevant final states: Tevatron's inclusive search with like sign dilepton events, Tevatron's measurement of the ditop cross-section, LEP's search for cascade decays of the Higgs, LEP's Higgs to WW^* search and direct searches for neutralinos with R-parity violation.

3.1 Like sign dileptons

The Tevatron inclusive search for new physics with like sign dilepton events [14] is a strong model independent bound that must be evaded. The CDF collaboration looked at a data sample of $1\,\mathrm{fb}^{-1}$ and observed no significant excess in an inclusive selection or in a SUSY-optimized selection. The SUSY-optimized selection requires large missing transverse energy and is thus irrelevant in our model since we get like sign dilepton events without missing transverse energy. The expected number of events in the inclusive selection is 33.2 ± 4.7

while the observed number of events is 44. Our model could explain this slight excess. Indeed, light Higgs production cross-section at the Tevatron is around 1 pb which would lead to an extra number of like sign dilepton events of about 10 (without taking into account any cuts). However, one must remember that the Tevatron search focused on first and second generation leptons while our like sign dilepton events are ditau events which are more subtle to study.

3.2 Tevatron $t\bar{t}$ search

The Tevatron $t\bar{t}$ [15] search in principle looks very similar to our signal; it places limits on the $t\bar{t}$ production cross-section by looking for multi-jet events with two b quarks in the final state and with missing energy, which is not found in the b quark direction. There are several reasons to expect this search to be insensitive to our decay. First, since the Higgs is produced close to threshold at the Tevatron, the decay products are all relatively soft. The most missing energy we expect is from final states with two neutrinos of about 15 GeV. Since the decay products are isotropically distributed, it is unlikely that the two neutrinos will be found in the same hemisphere, and the total missing energy vector will not be large enough to pass missing energy cuts in this search. However, the most important remark is that the Higgs production cross-section in the regime of interest is a few picobarns [16]. The same can be said for the direct neutralino production cross-section. The two sigma error bars for the $t\bar{t}$ production cross-section at Tevatron are about 3 pb. Thus, the jet plus missing energy events found by Tevatron, are all consistent with coming from top and cannot put strong bounds on $LQ\bar{D}$ couplings, even if a significant number of events were to pass the missing energy cut. Finally, there is a question about how many b quarks are in the final state. The Fermilab search insisted on two b tagged jets as one of their primary cuts. Since the flavor structure of the $LQ\bar{D}$ couplings is undetermined, we could easily construct models, which further suppressed b jets. Thus gaugino decays through lepton number violation could have easily escaped detection. However, since the missing energy threshold is small, if there are too many tagged b's in the event it will be discarded as QCD background. Therefore events with four b's in the final state are less likely to be picked up by this search.

3.3 Higgs cascade decays

LEP conducted three searches which have bearing on these decays. One is the cascade decay $e^+e^- \to H_2 Z \to H_1 H_1 Z$. In the case of invisible Z decay and $H_1 \to b\bar{b}$ there is a four jet plus missing energy signal. This search only applies to final states with a maximum number of b quarks. Details may be found in [4]. This channel, as detailed in figure 12, puts a bound on the Higgs mass of less than 85 GeV. Reconstruction of the missing Z may further reduce the sensitivity of this search down to OPAL's inclusive limit, which is 82 GeV [19]. LEP also conducted a search in the channel $hZ \to WW^*Z$ where the Z decays invisibly and the W's decay to qq' [20]. This search is not applicable to final states where all four quarks are the same flavor, or to final states where the quarks are all up type or down type. For example, this search is insensitive to our 4b plus missing energy final state. When the search is applicable to our final states, it can only put a bound of

105 GeV on the Higgs mass. We detail the Higgs mass limits and production cross-section limits in the table above.

Having convinced ourselves that current searches do not rule out our scenario, we note that the kinematics of the gaugino decays are sufficiently different from the top decays, and that one could imagine finding them in a dedicated search. One could attempt a search without b tagging, but simple searches for jets plus significant missing energy are very difficult. In the standard 4b search the Higgs, which decays to $b\bar{b}$, is produced in association with one or two b quarks [18]. This search required at least 3 b-tags, but overall it was not very sensitive since it was sensitive to the 3b background. The best chance might be to assume that the Higgs decays have the maximal number of b's in the final state, and to modify the Tevatron 4b Higgs search to be sensitive to missing energy [17]. One may imagine a similar search which requires multiple b tags and a missing energy cut. This scenario faces the same problem that we detailed in the ditop measurement above, the smallness of the missing energy vector. The missing energy at the Tevatron must be more than what we would expect from semi-leptonic b decay, and more than what would result from the mis-measurement of a 4 b event with no missing energy. Passing these cuts with low energy neutrinos is a problem. Searches at LHC seem to be even more problematic, but we note that Kaplan and Rehermann have proposed searching for Higgs decays through neutralino LSP's into multi jet final states using the LHCb experiment [21]. LHCb catches events highly boosted in the forward direction, has maximal b acceptance, and has a p_T trigger which can be as low as 2 GeV. As in the B violating 6 jet decays of the Higgs [4], it may be possible to search for lepton number violating decays of the gaugino at LHCb, with missing energy and multiple b tags.

3.4 Direct searches for neutralinos

Direct searches for neutralinos with R-parity violation [22] put bounds on the neutralino cross-section for direct three body decays. Since our model leads to a new production mechanism for neutralinos, it is natural to investigate these bounds. The Higgs boson production cross-section and the neutralino pair production cross-section at an e^+e^- collider are around 300 fb each [23]. Since the direct searches for neutralinos with R-parity violation give an upper bound of about 2 pb for the neutralino pair production cross-section, our model escapes direct neutralino search bounds easily.

4. Neutrino masses

The lepton number violating operators, which we have invoked to hide the decays of a light Higgs boson, might also be the source of neutrino masses. There is a large literature (see [24] and references therein) on the use of renormalizable L violating terms in the MSSM to generate neutrino masses. Indeed, some of the strongest constraints on the L violating couplings we have used come from the requirement that the neutrino masses they generate not be too large. For instance, the numerical example of section 2.3 leads to a $\lambda'\lambda'$ loop neutrino mass of about $1.5 \cdot 10^{-3}$ eV for squark of about 250 GeV [24]. Obviously

one needs a more involved flavor structure to generate all neutrino masses since only the third generation neutrino mass is generated in our model.

A survey of the literature indicates that bilinear L violating terms of the form $L_i H_u$ are the dominant source of neutrino masses in a generic model. However, to be consistent one should require that all B and L violating terms which could lead to unobserved processes are forbidden by a symmetry. We do not know how to make a general analysis of such symmetries without committing ourselves to a specific model. Thus we will restrict our attention to the Pentagon model [8], though we expect that a similar analysis could be done for any specific model of gauge mediation. We will find, that within the context of the Pentagon model, the symmetries we utilize will forbid terms of the form $L_i H_u$ but allow $L_i H_u S$ (where S is the singlet of the Pentagon). If S has a vacuum expectation value (VEV), this will generate a tree level mass for one neutrino. The dominant contribution to the other two neutrino masses comes from loop corrections involving the $LQ\bar{D}$ couplings that hide the Higgs decay. There is thus a potential understanding of a $2-1$ hierarchy among the three neutrino masses, as seems to be indicated by experiment. However, we emphasize that both the magnitude of the $L_i H_u S$ term, and the $LQ\bar{D}$ couplings is determined by high energy physics beyond the range of the present analysis. Therefore, a proper understanding of the structure of the neutrino mass matrix really requires unification scale physics.

The original Pentagon model was designed to eliminate all baryon and lepton number violating operators of dimension ≤ 5, except for the neutrino seesaw term. This led to a $\mathbb{Z}_4$ R symmetry with two possible generation independent charge assignments. In order to admit renormalizable lepton number violating terms we must change the symmetry and the charge assignments. We will assume an R symmetry group $\mathbb{Z}_N$. Therefore in the following all equations are understood modulo N and the R-charge of a given field is denoted by the name of the field itself.

4.1 Independent R-charges

The aim of this subsection is to express the R-charges of all the fields of the Pentagon in terms of the R-charges of a minimal set of fields. The appropriate restricted set is somewhat arbitrary but a rather convenient one comes naturally from the model. First, the crucial $SP\tilde{P}$ and $SH_u H_d$ terms lead to

$$SP\tilde{P} \Rightarrow P + \tilde{P} = 2 - S \tag{4.1}$$

$$SH_u H_d \Rightarrow H_u = 2 - S - H_d. \tag{4.2}$$

The important Yukawa couplings give

$$LH_d\bar{E} \Rightarrow \bar{E} = 2 - L - H_d \tag{4.3}$$

$$QH_u\bar{U} \Rightarrow \bar{U} = 2 - Q - H_u \tag{4.4}$$

$$QH_d\bar{D} \Rightarrow \bar{D} = 2 - Q - H_d. \tag{4.5}$$

Thus one can rewrite everything as a function of the restricted set $\{S, L, Q, H_d\}$ as

$$P + \tilde{P} = 2 - S \tag{4.6}$$

$$H_u = 2 - S - H_d \tag{4.7}$$

$$\bar{E} = 2 - L - H_d \tag{4.8}$$

$$\bar{U} = S - Q + H_d \tag{4.9}$$

$$\bar{D} = 2 - Q - H_d. \tag{4.10}$$

This set is dubbed extended since anomaly conditions will generate relations between the four different R-charges of the set.

4.1.1 Anomaly conditions

The anomaly conditions of the Pentagon model are

$$\mathrm{SU}(5)_P \Rightarrow 5(P + \tilde{P}) = 0 \tag{4.11}$$

$$\mathrm{SU}(3)_C \Rightarrow 6Q + 3(\bar{U} + \bar{D}) + 5(P + \tilde{P}) = 0 \tag{4.12}$$

$$\mathrm{SU}(2)_L \Rightarrow H_u + H_d + 9Q + 3L + 5(P + \tilde{P}) = 0. \tag{4.13}$$

Using the relations obtained from the restricted set $\{S, L, Q, H_d\}$ of independent R-charges these can be rewritten as

$$\mathrm{SU}(5)_P \Rightarrow 5(S - 2) = 0 \tag{4.14}$$

$$\mathrm{SU}(3)_C \Rightarrow 3(S + 2) = 0 \tag{4.15}$$

$$\mathrm{SU}(2)_L \Rightarrow 2 - S + 9Q + 3L = 0. \tag{4.16}$$

The last anomaly condition leads to an unextended restricted set of independent R-charges by removing one R-charge in the extended restricted set. Due to the modulo N form of the equations the easiest one to remove is S but it is more convenient to keep everything written in function of the extended restricted set $\{S, L, Q, H_d\}$. Indeed one can easily solve the anomaly conditions as a function of S. Thus it is more practical to eliminate Q instead as shown later. The first two anomaly conditions can be combined as

$$0 = 5(S - 2) = 3(S + 2) + 2(S - 8) = 2(S - 8). \tag{4.17}$$

For $N = 2n + 1$ one has $S = 8$ and the first two anomaly conditions force $N|30$ thus $N = \{3, 5, 15\}$. For $N = 2n$ one has $S = 8$ or $S = 8 - n$. For the case $S = 8$ the first two anomaly conditions force $N|30$ thus $N = \{2, 6, 10, 30\}$. For the case $S = 8 - n$ the first two anomaly conditions lead to

$$N \mid (30 - 5n) \tag{4.18}$$

$$N \mid (30 - 3n) \tag{4.19}$$

thus $N = \{4, 12, 20, 60\}$.

4.2 Superpotential terms and RPV terms

Using the extended restricted set the possible S^3 superpotential term gives

$$S^3 \Rightarrow 3S - 2 \tag{4.20}$$

– 13 –

where the r.h.s. has to be zero modulo N if and only if the term is allowed. For the RPV terms, the trilinear lepton number violating (TLNV) terms (including the useful SLH_u) give

$$LL\bar{E} \Rightarrow 2L + (2 - L - H_d) - 2 = L - H_d \tag{4.21}$$

$$LQ\bar{D} \Rightarrow L + Q + (2 - Q - H_d) - 2 = L - H_d \tag{4.22}$$

$$SLH_u \Rightarrow S + L + (2 - S - H_d) - 2 = L - H_d. \tag{4.23}$$

and the bilinear lepton number violating (BLNV) term LH_u leads to

$$LH_u \Rightarrow L + (2 - S - H_d) - 2 = L - H_d - S. \tag{4.24}$$

Finally the trilinear baryon number violating (TBNV) term gives

$$\bar{U}\bar{D}\bar{D} \Rightarrow (S - Q + H_d) + 2(2 - Q - H_d) - 2 = S - 3Q - H_d + 2. \tag{4.25}$$

The no-go theorem here states that one cannot allow only specific TLNV terms since all TLNV terms are allowed when any one is allowed.

4.3 Dimension five baryon number violating operators

Dimension five baryon number violating (D5BNV) operators and D-terms lead to

$$QQQL \Rightarrow 3Q + L - 2 \tag{4.26}$$

$$QQQH_d \Rightarrow 3Q + H_d - 2 \tag{4.27}$$

$$\bar{U}\bar{U}\bar{D}\bar{E} \Rightarrow 2(S - Q + H_d) + (2 - Q - H_d) + (2 - L - H_d) - 2$$
$$= 2S - 3Q - L + 2 \tag{4.28}$$

$$\text{D-term} \Rightarrow Q + \bar{U} - L = Q + (S - Q + H_d) - L = S - L + H_d \tag{4.29}$$

$$\text{D-term} \Rightarrow \bar{U} + \bar{E} - \bar{D} = (S - Q + H_d) + (2 - L - H_d) - (2 - Q - H_d)$$
$$= S - L + H_d \tag{4.30}$$

where the last two equations come from D-terms.

4.4 Overall solutions

From the last two sections one can group together terms that lead to the same equation as a function of the R-charges of the extended restricted set. One has seven different sets labeled G_1 to G_7,

$$G_1 = \{S^3\} \Rightarrow 3S - 2 \tag{4.31}$$

$$G_2 = \{LL\bar{E}, LQ\bar{D}, SLH_u\} \Rightarrow L - H_d \tag{4.32}$$

$$G_3 = \{LH_u, \text{D-terms}\} \Rightarrow L - H_d - S \tag{4.33}$$

$$G_4 = \{\bar{U}\bar{D}\bar{D}\} \Rightarrow S - 3Q - H_d + 2 \tag{4.34}$$

$$G_5 = \{QQQL\} \Rightarrow 3Q + L - 2 \tag{4.35}$$

$$G_6 = \{QQQH_d\} \Rightarrow 3Q + H_d - 2 \tag{4.36}$$

N	S	$SU(2)_L$	G_1	G_3	G_4	G_5	G_6	G_7
2	0	$Q=L$	0	$L-H_d$	$L-H_d$	0	$L-H_d$	0
3	2	none	1	$L-H_d-2$	$1-H_d$	$L-2$	H_d-2	$-L$
4	2	$Q=L$	0	$L-H_d-2$	$L-H_d$	2	$-(L-H_d-2)$	2
5	3	$Q=3L-1$	2	$L-H_d-3$	$L-H_d-2$	0	$-(L-H_d)$	1
6	2	$3Q=3L$	4	$L-H_d-2$	$3L-H_d-2$	$-2(L+1)$	$3L+H_d-2$	$2L$
10	8	$Q=3L+4$	2	$L-H_d-8$	$L-H_d-2$	0	$-(L-H_d)$	6
12	2	$3Q=3L$	4	$L-H_d-2$	$4-3L-H_d$	$4L-2$	$3L+H_d-2$	$-4L+6$
15	8	$9Q=6-3L$	7	$L-H_d-8$	$10-3Q-H_d$	$3Q+L-2$	$3Q+H_d-2$	$-3Q-L+3$
	8	$3Q=2-L$	7	$L-H_d-8$	$L-H_d+8$	0	$-(L-H_d)$	1
	8	$3Q=7-L$	7	$L-H_d-8$	$L-H_d+3$	5	$-(L-H_d-5)$	-4
	8	$3Q=12-L$	7	$L-H_d-8$	$L-H_d-2$	10	$-(L-H_d-10)$	-9
20	18	$9Q=16-3L$	12	$L-H_d-18$	$-3Q-H_d$	$3Q+L-2$	$3Q+H_d-2$	$-3Q-L+18$
	18	$3Q=12-L$	12	$L-H_d-18$	$L-H_d-12$	10	$-(L-H_d-10)$	6
30	8	$9Q=6-3L$	22	$L-H_d-8$	$10-3Q-H_d$	$3Q+L-2$	$3Q+H_d-2$	$-3Q-L+18$
	8	$3Q=2-L$	22	$L-H_d-8$	$L-H_d+8$	0	$-(L-H_d)$	16
	8	$3Q=12-L$	22	$L-H_d-8$	$L-H_d-2$	10	$-(L-H_d-10)$	6
	8	$3Q=22-L$	22	$L-H_d-8$	$L-H_d-12$	20	$-(L-H_d-20)$	-4
60	38	$9Q=36-3L$	52	$L-H_d-38$	$40-3Q-H_d$	$3Q+L-2$	$3Q+H_d-2$	$-3Q-L+18$
	38	$3Q=12-L$	52	$L-H_d-38$	$L-H_d+28$	10	$-(L-H_d-10)$	6
	38	$3Q=32-L$	52	$L-H_d-38$	$L-H_d+8$	30	$-(L-H_d-30)$	-14
	38	$3Q=52-L$	52	$L-H_d-38$	$L-H_d-12$	50	$-(L-H_d-50)$	-34

Table 5: Solutions for all allowed S R-charges.

$$G_7 = \{\bar{U}\bar{U}\bar{D}\bar{E}\} \Rightarrow 2S - 3Q - L + 2 \qquad (4.37)$$

where group G_2 consists of all TLNV terms exclusively, group G_4 of the TBNV term and groups G_5 to G_7 of D5BNV terms. Table 5 shows the solutions using the previous sets and the extra $SU(2)_L$ anomaly relation $2 - S + 9Q + 3L$ to eliminate Q. Notice that the TLNV set $G_2 \Rightarrow L - H_d$ does not simplify and no extra relation comes from the $SU(2)_L$ anomaly condition for $N = 3$. The removal of Q from the extended restricted set is more subtle for the cases $N = \{15, 20, 30, 60\}$. For $N = 15$ one has $9Q + 3L = 15k + 6 \Rightarrow 3Q + L = 5k + 2$ ($k \in \mathbb{Z}$) thus $3Q = \{2 - L, 7 - L, 12 - L\}$. For $N = 20$ one has $9Q + 3L = 20k' + 16 \Rightarrow 3Q + L = \frac{1}{3}(20k' + 16) = 20k + 12$ with $k' = 3k + 1$ ($k \in \mathbb{Z}$) thus $3Q = 12 - L$. For $N = 30$ one has $9Q + 3L = 30k + 6 \Rightarrow 3Q + L = 10k + 2$ ($k \in \mathbb{Z}$) thus $3Q = \{2 - L, 12 - L, 22 - L\}$. Finally for $N = 60$ one has $9Q + 3L = 60k + 36 \Rightarrow 3Q + L = 20k + 12$ ($k \in \mathbb{Z}$) thus $3Q = \{12 - L, 32 - L, 52 - L\}$.

Looking at the previous table one sees that only the cases $N = \{2, 4\}$ allow for the S^3 term set G_1. In the case of interest to us, i.e. allowing TLNV terms set G_2 (thus $H_d = L$) while prohibiting sets G_3 to G_7, one can find a solution only for $N = \{12, 15, 20, 30, 60\}$ (notice that G_3 is not a problem unless $N = 2$). For example the case $N = 3$ is not a solution since prohibiting G_5 and G_7 forces $L = 1$ which allows the unwanted G_4 while the case $N = 12$ is a solution since the sets G_5, G_6 and G_7 do not constrain L but G_4 forces $L \neq \{1, 4, 7, 10\}$ which is possible. The specific R-charges for the five possible cases are then computable. In this framework it is therefore impossible to allow only TLNV terms set G_2

along with the S^3 term set G_1. It is however possible to allow only TLNV terms set G_2.

4.5 Constraints on $\langle S \rangle$

In light of the previous computations one can engineer the appropriate Pentagon superpotential

$$\mathcal{W} = (m_{\text{Iss}} + g_S SY)P\tilde{P} + g_\mu SH_uH_d + \lambda_L H_d L\bar{E} + \lambda_u H_u Q\bar{U} + \lambda_d H_d Q\bar{D}$$
$$+\frac{1}{2}\lambda LL\bar{E} + \lambda'LQ\bar{D} + g_\epsilon SLH_u. \tag{4.38}$$

If S gets a VEV then *one* neutrino mass is mostly due to the SLH_u term while the Higgs is hidden by the $LQ\bar{D}$ term. This comes from the specific form of the tree level neutrino mass matrix (rank = 1) and thus only one neutrino is massive which is good to generate a hierarchy. Loop diagrams from the $LL\bar{E}$ and $LQ\bar{D}$ terms give masses to the other neutrinos (see [24]). There is also an effective μ term and thus no light higgsinos.

On the other hand, it may be a challenge to give a VEV to S if there is no S^3 term. If there is no S VEV then we could get both neutrino masses and Higgs decay to jets plus missing energy from the $LQ\bar{D}$ term, but we are likely to have an unacceptable light higgsino. A model without a VEV for S could generate all neutrino masses through loops involving the $LQ\bar{D}$ couplings, and could hide the Higgs via these same couplings. However, it would probably have an unacceptable light higgsino.

5. Conclusions

We have seen that gauge mediated models with lepton number violation can in principle hide the Higgs and generate an acceptable neutrino mass spectrum simultaneously. Our attempt to find a model in which the appropriate couplings followed from a discrete symmetry of the low energy theory was not completely successful.

The problem we encountered was specific to embedding the lepton violating scenario into the framework of the Pentagon model, but we anticipate some general features. In particular, it seems hard to find models where low energy symmetries allow $LQ\bar{D}$ operators, but forbid $LL\bar{E}$ operators. One has to rely on a high energy Froggatt-Nielsen mechanism, combined with SUSY non-renormalization theorems, to explain the suppression of the latter, which are significantly more constrained.

There will also be an inevitable connection between the origin of neutrino masses in R-parity violating models and the μ term of the MSSM. Our analysis indicates that it may be hard to explain the value of μ in terms of a low energy singlet VEV in these lepton number violating models.

Nonetheless, we think that gauge mediated models with renormalizable lepton number violation could offer considerable insight into two puzzles of the standard model. We have barely scratched the surface of this general class of models, and they deserve further investigation.

Acknowledgments

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