Running with triplets: How slepton masses change with doubly-charged Higgs bosons

N. Setzer and S. Spinner

Department of Physics and Center for String and Particle Theory, University of Maryland, College Park, Maryland 20742, USA

(Received 16 January 2007; published 1 June 2007)

We examine the slepton masses of SUSYLR models and how they change due to the presence of light doubly-charged Higgs bosons. We discover that the measurement of the slepton masses could bound and even predict the value of the third generation Yukawa coupling of leptons to the $SU(2)_R$ triplets. We also consider the unification prospects for this model with the addition of left-handed, $B - L = 0$ triplets—a model we call the triplet extended supersymmetric standard model (TESSM).

DOI: 10.1103/PhysRevD.75.117701 PACS numbers: 12.60.Jv, 12.60.Cn, 14.80.Ly

I. INTRODUCTION

The minimal supersymmetric standard model (MSSM) [1] is an appealing solution to numerous standard model ills, but it still doesn't address neutrino masses. Though ν_R may be added as a singlet, a more compelling solution is to extend the gauge group to $SU(3)^c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ (G_{3221}). This provides niceties such as the seesaw mechanism [2], why $M_R \ll M_{\rm Pl}$ [3], and even allows M_R to be predicted [4]. Such models, which are well known [5–8], may contain $SU(2)_R$ triplets—permitting the additional features of R-parity conservation and potentially light doubly-charged particles [4,9–11]. Specifically, the right-handed triplets must couple to the leptons through the term $f_c L^{cT} \tau_2 \Delta^c L^c$ to give the large Majorana mass to the right-handed neutrinos. This coupling then forces the doubly-charged particles to couple to the sleptons, and because they survive to the TeV scale they alter the slepton renormalization group equations (RGEs) and hence their masses [12].

The slepton mass running is highly dependent on f_c, and we will demonstrate in Sec. II that one may bound f_c by limits on the stau mass. In fact, one can do better than bound f_c: a measurement of a right-handed selectron mass in excess of the MSSM's result, combined with a measurement of the $\tilde{\tau}_1$ mass, would yield a value for the third generation f_c (provided G_{3221} is the correct theory). We think that this is an important result to emphasize since probing the TeV scale slepton masses will then yield an indication of the physics roughly 8 orders of magnitude higher.

The remnant doubly-charged Higgs bosons also affect unification. This was considered in [12], but we adopt a different approach (Sec. III). Further complications arise in unification from requiring that the right-handed coupling remain perturbative. We find in Sec. III that α_R will only remain perturbative up until about 10^{12} GeV. Because of this, we focus our discussion of unification on gauge-mediated SUSY breaking scenarios.

II. SLEPTON MASSES WITH LIGHT DOUBLY-CHARGED HIGGSES

In this section we consider the MSSM + DC, the minimal supersymmetric standard model (MSSM) plus doubly-charged Higgs superfields (DC)—an effective theory of well-motivated high energy theories with the DC singlets of all but the hypercharge group ($Y = 4$). We investigate the low-energy consequences, adding to the work done in [12]. We also use this section as a springboard into Sec. III, which will have similar phenomenology. We begin by stating the additional interactions, denoting the DC's as Δ^{--} and $\bar{\Delta}^{++}$:

$$\Delta W_{\rm MSSM+DC} = e^c f_c e^c \Delta^{--} + \mu H_u H_d + \mu_\Delta \Delta^{--} \bar{\Delta}^{++}, \tag{2.1}$$

$$\Delta V_{\rm soft} = (\tilde{e}^c a_{f_c} \tilde{e}^c \Delta^{--} + b H_u H_d + b_\Delta \Delta^{++} \Delta^{--} + \text{c.c.}) + m_\Delta^2 |\Delta^{--}|^2 + m_{\bar{\Delta}}^2 |\bar{\Delta}^{--}|^2, \tag{2.2}$$

where, as usual, generational, color, and isospin indices have been suppressed. Furthermore the common practice of replacing Yukawa matrices in the MSSM by a scalar representing the third generation coupling can be applied to f_c based on results from muonium oscillations and flavor violating decays [13,14] (which constrain all but the $\tau\tau$ component [12]).

The gauge-mediated SUSY breaking (GMSB) scenario seems to be the most compatible with this model because of its low breaking scale (Sec. III), and so it is used to generate the soft terms (for a review see [1,15]; we follow this notation). Our RGE results are obtained at the one-loop level and yield MSSM values that are at most 3% off from ISAJET [16]. Since we are interested in percent difference between our MSSM and our MSSM + DC values, this should not be an issue.

The RGEs for a general SUSYLR can be found in [17], and we utilize those (with appropriate changes). The presence of the DC alter the hypercharge running to

$$\frac{d\alpha_1^{-1}}{dt} = -\frac{3}{5}\frac{19}{2\pi} \tag{2.3}$$

and so cause a larger hypercharge gauge coupling at the messenger scale; however, the increase is muted by the fact that the values also run down faster—causing a smaller difference at the observable scale.

Yet there will be a substantial difference for parameters whose only gauge coupling dependence at the messenger

TABLE I. Sparticle masses with standard parameters chosen as SPS8: $\Lambda = 100$ TeV, $\tan\beta = 15$, $N_5 = 1$, $M_{\text{mess}} = 200$ TeV, sgn $\mu = +1$, and $Q = 1$ TeV [18], with Q the scale at which the masses are quoted. We also use $b_\Delta = 0$ and $\mu_\Delta = 800$ GeV (a parameter that has no appreciable effects on these masses). To elucidate the f_c dependence, we show masses for $f_c(M_Z) = 0.1$ and 0.6. The masses are reported in GeV at M_{SUSY}.

Sparticle	MSSM	MSSM + DC $f_c = 0.1$	Δ%	MSSM + DC $f_c = 0.6$	Δ%
$\tilde{\tau}_1$	163	183	12%	118	28%
$\tilde{e}_R$	171	191	12%	191	12%
$\tilde{e}_L$	367	369	1%	369	1%
$\tilde{\chi}_1^0$	132	128	3%	128	3%

boundary is due to the hypercharge coupling, such as $\tilde{e}_R$ and $\tilde{\tau}_1$, whose boundary conditions are given by

$$m^2_{\tilde{e}^c,\tilde{\tau}^c} = 2\Lambda^2\left[\frac{3}{5}\left(\frac{\alpha_1}{4\pi}\right)^2\right]. \tag{2.4}$$

A more dramatic change will result from large f_c due to the running of $m^2_{\tilde{\tau}^c}$:

$$16\pi^2\frac{dm^2_{\tilde{\tau}^c}}{dt} = 4|y_\tau|^2(m^2_{\tilde{\tau}^c} + m^2_{H_d} + m^2_L) + 8|f_c|^2(2m^2_{\tilde{\tau}^c} + m^2_\Delta) + 4|a_\tau|^2 + 8|a_c|^2 - 4\pi\left(\frac{24}{5}\alpha_1|M_1|^2\right) \tag{2.5}$$

which will drive the $\tilde{\tau}^c$ mass parameter down causing $\tilde{\tau}_1$ mass to decrease. This can be seen in Table I.

To further illuminate the dependence of $m_{\tilde{\tau}_1}$ on f_c, we include Fig. 1, which shows that the mass can be driven to its limit depending on the value of f_c. Furthermore, it reveals that for low f_c, the stau mass is larger than the MSSM value and by the same percentage as $\tilde{e}_R$; for f_c bigger than about 0.4 the stau is distinctly lighter than its

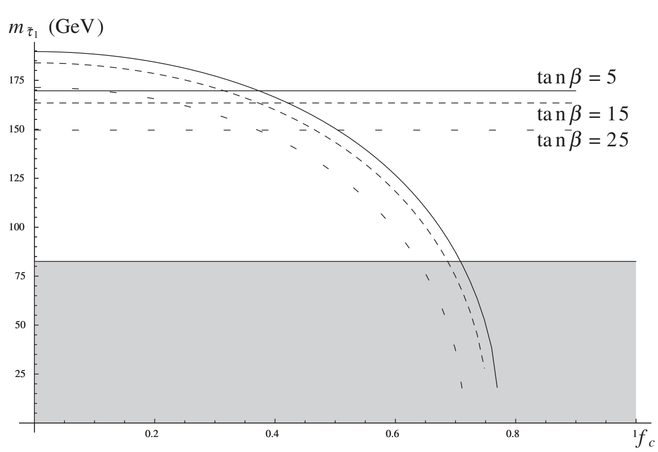

FIG. 1. The lightest stau mass as a function of f_c in the MSSM (straight lines) and MSSM + DC (curves). The shaded region is excluded by LEP II. The graph clearly demonstrates that for a given $\tan\beta$ and Λ there is an upper bound on f_c.

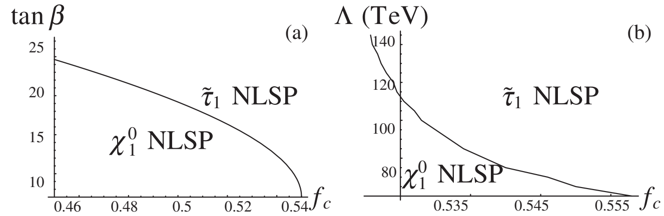

FIG. 2. (a) The dividing line between a neutralino LSP and a stau LSP as a function of $\tan\beta$ and f_c for $\Lambda = 100$ TeV. (b) The dividing line as a function of Λ and f_c for $\tan\beta = 15$.

MSSM value. This also means that the stau can become the next lightest supersymmetric particle (NLSP) which has important phenomenological consequences. Such a scenario is possible in the MSSM for larger values of $\tan\beta$ and low values of N_5 but the parameter space is greatly increased in the MSSM + DC. Figure 2 expands on this by showing which regions of the $\Lambda - f_c$ and the $\tan\beta - f_c$ planes produce a stau NLSP.

Detection of the DCs themselves would be a smoking gun for this model; however, assuming that these are not easily detectable, we can elucidate how to analyze a potential signal. We found that the only masses that change are $\tilde{e}_R$ and $\tilde{\tau}_1$ which will be larger (smaller) for small (large) f_c. Such a disproportional change in the right-handed sector would most naturally arise from additions in the hypercharge sector of the theory only. For larger f_c it would be clear that the new particle content must couple to the right-handed taus thereby suggesting the Yukawa term in Eq. (2.1) and DC. Specifically it would imply a hypercharge 4 singlet of both $SU(2)_L$ and $SU(3)^c$ since a higher representation of these could not couple at tree level only to the right-handed leptons. For smaller cases though, the nature of the new content would not be as clear and validation might have to wait for DC detection or some other signal. Regardless of the size of f_c, once the DC were discovered, measuring the mass of $\tilde{\tau}_1$ will yield a value for f_c, a parameter which has implications in the neutrino sector. For general GMSB phenomenology and limits, see (for example) [15,19,20].

III. UNIFICATION AND THE TRIPLET EXTENDED SUPERSYMMETRIC STANDARD MODEL

The gauge coupling unification of models with DC Higgses has been discussed [12], and it was pointed out that the couplings may be chosen to unify at around 10^{12} GeV; however, when [12] considered unification, the authors chose to have two additional Higgs doublets at 10 TeV. We will present an alternative solution that maintains the usual two Higgs doublets at low scales and requires the additional particle content to have masses at the TeV scale.

To motivate our solution, we first note that the DC Higgs bosons only affect hypercharge—causing a drastic increase in the running. Since we wish to unify to G_{3221}

this presents a major problem: if the left- and right-handed couplings run the same way, then both will run too slowly and force $\bar{\alpha}_{BL}^{-1}$ to be nonperturbtive or even less than zero at the right-handed scale. It follows that a $B - L = 0$ field would be a viable solution, so we choose to add the simplest higher representation: $SU(2)_L$ triplets. We name this the triplet extended supersymmetric standard model (TESSM). The couplings can then be made to unify in two ways.

The first example of unification is where all the G_{3221} couplings unify at a scale of $M_{\rm GUT} = 1.3 \times 10^{12}$ GeV. This scale is far too low for $SO(10)$ (due to proton decay constraints), but any group that conserves baryon number would suffice. To achieve this scenario it is only necessary to add one $Y = 0$ triplet, so in this sense it is the minimal model and the one on which we will focus our detailed analysis.

Before diving into TESSM let us briefly note that the couplings may alternatively unify to $SU(5)_L \times SU(5)_R$—this group being attractive because it requires the gauge couplings to be unequal at the unification scale [21]. If the $SU(5)^2$ unification is at M_U, it is quickly noticed that $v_R \simeq M_U$ so we take $v_R = M_U$. The couplings then unify at $M_U = v_R = 2.5 \times 10^{11}$ GeV. To realize this unification requires the inclusion of two $B - L = 0$ triplets to the model, thus making it in some sense "less minimal" than the previous scenario.

The Higgs sector of TESSM has been discussed previously in [22], though the addition of left-handed triplets was *ad hoc*. The authors of [22] do a thorough analysis of the vacuum structure and Higgs masses; however, since they do not assume any higher scale physics, their parameters are largely unconstrained. Our investigations show that the assumption of unification limits the parameter space to exclude the scenarios considered in [22]. We derive the RGEs for this model from [23].

To see the origins of these constraints, we start with the superpotential

$$W = W_{\rm MSSM+DC} + \mu_\delta \,{\rm Tr}\delta^2 + i y_\delta H_u^T \tau_2 \delta H_d \tag{3.1}$$

and the soft breaking terms

$$V_{\rm soft,T} = m_\delta^2 \,{\rm Tr}|\delta|^2 + [b_\delta \,{\rm Tr}\delta^2 + i a_\delta H_u^T \tau_2 \delta H_d + {\rm H.c.}]. \tag{3.2}$$

These new terms modify the MSSM minimization conditions, but, more interestingly, add the new constraint[1]

$$4\mu_\delta^2 + m_\delta^2 + 2b_\delta + \frac{1}{4} y_\delta^2 v^2 + \frac{1}{2}\frac{v^2}{v_\delta} y_\delta \mu - \frac{1}{2}\frac{v^2}{v_\delta}\left(y_\delta \mu_\delta + \frac{1}{2} a_\delta\right)\sin 2\beta = 0. \tag{3.3}$$

Additionally, they alter the stability requirements to include

$$4\mu_\delta^2 + m_\delta^2 > |2b_\delta|. \tag{3.4}$$

Electroweak precision measurements imply that $v_\delta \ll v$, so that the terms involving v/v_δ in Eq. (3.3) are much larger than the SUSY breaking scale. GMSB, meanwhile, predicts that the trilinear A terms are very small, and so approximately zero. Rewriting Eq. (3.3) keeping only the important terms (and assuming $\sin 2\beta \approx 1$, $y_\delta \approx 1$) gives

$$4\mu_\delta^2 + b_\delta - \frac{1}{2}\frac{v^2}{v_\delta}\mu_\delta + \frac{1}{2}\frac{v^2}{v_\delta}\mu = 0. \tag{3.5}$$

The last term is large and positive, and forces b_δ to be negative with a large magnitude (since μ_δ comes with terms of opposite sign, its contributions mostly cancel each other). With $|b_\delta|$ large, the stability condition of Eq. (3.4) requires that μ_δ also be large (given $m_\delta^2 \sim M_{\rm SUSY}^2$). It is therefore necessary for the new $Y = 0$ triplets to be "heavy," and our numerical analysis indicates they are around 5 TeV.

We consider now the slepton masses. The expressions for their masses remain the same as MSSM + DC except for the stau mass matrix, which is now given by

$$m_{\tilde{\tau}}^2 = \begin{pmatrix} m_{\tilde{L}3}^2 + y_\tau^2 v_d^2 + \frac{\pi}{2}\left(\frac{3}{5}\alpha_1 - \alpha_2\right)(v_d^2 - v_u^2) & \frac{1}{\sqrt{2}}\left(a_\tau v_d + \frac{1}{2} y_\delta y_\tau v_u v_\delta - v_u \mu y_\tau\right) \\ \frac{1}{\sqrt{2}}\left(a_\tau v_d + \frac{1}{2} y_\delta y_\tau v_u v_\delta - v_u \mu y_\tau\right) & m_{\tilde{\tau}^c}^2 + |y_\tau|^2 v_d^2 - \pi\alpha_2(v_d^2 - v_u^2) \end{pmatrix}. \tag{3.6}$$

The resulting right-handed slepton spectrum is very similar to MSSM + DC where as the left-handed sleptons get an increase of about 5% over their MSSM values. As for the MSSM Higgs sector, there will be no new radiative mass corrections [22,24–26]; however, the Higgs sector is obviously expanded and there is a new vev, $\langle\delta\rangle$. This vev is constrained by the ρ parameter to be less than about 1.7 GeV [14], so we take it to be around 1 GeV. The extended Higgs sector is composed of a neutral scalar H_δ^0, a neutral pseudoscalar B_δ^0, and two singly charged scalars $H_{\delta 1}^+$ and $H_{\delta 2}^+$. These fields will not mix very much with the MSSM fields because of the large μ_δ value. We take a quick peek at their typical tree-level masses (in TeV) for the SPS8 point and with $f_c = 0.1$, $\mu_\Delta = 800$ GeV, and $y_\delta = 0.1$:

$$M_{H_\delta^0, H_{\delta 1}^+} = 0.74 \text{ TeV}, \qquad M_{B_\delta^0, H_{\rm delta\ 2}^+} = 14.2 \text{ TeV}. \tag{3.7}$$

It is worth noting that the above sets of masses will remain equal even after radiative corrections because of the low coupling to quarks. Still, the lighter fields can be paired produced via W boson fusion and have electroweak-magnitude cross sections at the LHC. If produced, they will decay into MSSM Higgs fields or two electroweak bosons depending on the size of y_δ. Signatures in linear colliders for this model are discussed in [22,25,26].

IV. CONCLUSION

We have considered an extension of the MSSM with light doubly-charged Higgs bosons and showed that the right-handed slepton masses will be significantly different. Verifications of these mass deviations would be a good signal for this model—even if the doubly-charged Higgses are beyond the reach of future accelerators. Furthermore, measurements of the lightest stau mass will have implications in the neutrino sector and the parameter space where it is the NLSP is increased.

We also showed that unification requires left-handed triplets and results in phenomenology that includes all of the features of the MSSM + DC. The vev of the additional Higgses is suppressed by the ρ parameter, which leads to rigid constraints on the parameters. This effectively forces half of the new Higgses to be well outside the reach of future colliders, but potentially leaves the other half within the LHC's grasp (depending on the parameters).

ACKNOWLEDGMENTS

We would like to thank Rabindra Mohapatra for useful conversations regarding this model. This work is supported by the National Science Foundation Grant No. Phy-0354401.

[1] See, for instance, S. P. Martin, arXiv:hep-ph/9709356.

[2] P. Minkowski, Phys. Lett. B **67**, 421 (1977); M. Gell-Mann, P. Ramond, and R. Slansky, *Supergravity. Proceedings of a Workshop at Stony Brook*, edited by P. van Nieuwenhuizen *et al.*, (North-Holland, Amsterdam, 1980), p. 315; T. Yanagida, in Proceedings of the Workshop on the Baryon Number of the Universe and Unified Theories, Tsukuba, Japan, 1979 (unpublished); S.L. Glashow, in *Proceedings of the 1979 Cargese Summer Institute on Quarks and Leptons*, edited by M. Levy *et al.*, (Plenum, New York, 1980), p. 687; R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 912 (1980).

[3] R. N. Mohapatra and R. E. Marshak, Phys. Rev. Lett. **44**, 1316 (1980); **44**, 1644(E) (1980).

[4] R. N. Mohapatra, N. Setzer, and S. Spinner, Phys. Rev. D **73**, 075001 (2006).

[5] M. Cvetic and J. C. Pati, Phys. Lett. B **135**, 57 (1984); R. Kuchimanchi and R. N. Mohapatra, Phys. Rev. D **48**, 4352 (1993); C. S. Aulakh, A. Melfo, A. Rasin, and G. Senjanovic, Phys. Rev. D **58**, 115007 (1998).

[6] R. M. Francis, M. Frank, and C. S. Kalman, Phys. Rev. D **43**, 2369 (1991).

[7] K. S. Babu, B. Dutta, and R. N. Mohapatra, Phys. Rev. D **65**, 016005 (2001).

[8] K. S. Babu, B. Dutta, and R. N. Mohapatra, Phys. Rev. D **60**, 095004 (1999).

[9] C. S. Aulakh, A. Melfo, and G. Senjanovic, Phys. Rev. D **57**, 4174 (1998); Z. Chacko and R. N. Mohapatra, Phys. Rev. D **58**, 015003 (1998); S. Dar, Q. Shafi, and A. Sil, Phys. Lett. B **632**, 517 (2006).

[10] M. L. Swartz, Phys. Rev. D **40**, 1521 (1989); J. F. Gunion, Int. J. Mod. Phys. A **11**, 1551 (1996); M. Raidal, arXiv:hep-ph/9809370; K. Huitu, P. N. Pandita, and K. Puolamaki, arXiv:hep-ph/9904388; arXiv:hep-ph/9910504; S. Godfrey, P. Kalyniak, and N. Romanenko, Phys. Lett. B **545**, 361 (2002); J. Maalampi and N. Romanenko, Phys. Lett. B **532**, 202 (2002).

[11] B. Mukhopadhyaya and S. K. Rai, Phys. Lett. B **633**, 519 (2006).

[12] B. Dutta and R. N. Mohapatra, Phys. Rev. D **59**, 015018 (1998).

[13] T. M. Huber *et al.*, Phys. Rev. D **41**, 2709 (1990); B. E. Matthias *et al.*, Phys. Rev. Lett. **66**, 2716 (1991); L. Willmann *et al.*, Phys. Rev. Lett. **82**, 49 (1999).

[14] W. M. Yao *et al.* (Particle Data Group), J. Phys. G **33**, 1 (2006).

[15] G. F. Giudice and R. Rattazzi, Phys. Rep. **322**, 419 (1999).

[16] F. E. Paige, S. D. Protopopescu, H. Baer, and X. Tata, arXiv:hep-ph/0312045.

[17] N. Setzer and S. Spinner, Phys. Rev. D **71**, 115010 (2005).

[18] B. C. Allanach *et al.*, Eur. Phys. J. C **25**, 113 (2002); econf C010630, P125 (2001).

[19] J. Abdallah *et al.* (DELPHI Collaboration), Eur. Phys. J. C **27**, 153 (2003).

[20] R. Barate *et al.* (ALEPH Collaboration), Phys. Lett. B **433**, 176 (1998).

[21] R. N. Mohapatra, Phys. Rev. D **54**, 5728 (1996).

[22] J. R. Espinosa and M. Quiros, Nucl. Phys. **B384**, 113 (1992).

[23] S. P. Martin and M. T. Vaughn, Phys. Rev. D **50**, 2282 (1994).

[24] J. R. Espinosa and M. Quiros, Phys. Lett. B **279**, 92 (1992).

[25] O. Felix-Beltran, Int. J. Mod. Phys. A **17**, 465 (2002).

[26] E. Barradas-Guevara, O. Felix-Beltran, and A. Rosado, Phys. Rev. D **71**, 073004 (2005).