Higgs-Mediated $\tau \to 3\mu$ in the Supersymmetric Seesaw Model

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(Received 10 July 2002; published 22 November 2002)

Recent observations of neutrino oscillations imply nonzero neutrino masses and lepton flavor violation (LFV), most economically explained by the seesaw mechanism. Within the context of supersymmetry, LFV among the neutrinos can be communicated to the sleptons and from there to the charged leptons. We show that LFV can appear in the couplings of the neutral Higgs bosons, an effect that is strongly enhanced at large $\tan\beta$. We calculate the branching fraction for $\tau \to 3\mu$ and $\mu \to 3e$ mediated by Higgs and find they can be as large as 10^{-7} and 5×10^{-14}, respectively. These modes, along with $\tau \to \mu\gamma$ and $\mu \to e\gamma$, can provide key insights into the neutrino mass matrix.

DOI: 10.1103/PhysRevLett.89.241802 PACS numbers: 13.35.Dx, 11.30.Hv, 12.60.Jv, 14.60.Pq

Over the past several years, evidence from a number of experiments, notably Super-Kamiokande and SNO, has pointed conclusively to the existence of neutrino oscillations in atmospheric [1] and solar neutrinos [2] and, by implication, to nonzero neutrino masses. Within the context of the standard model (SM), the most attractive explanation for the observed neutrino masses is the "seesaw" mechanism [3]: right-handed neutrinos are introduced in order to couple with left-handed neutrinos through $SU(2) \times U(1)$-violating Dirac mass terms, m_D, while also receiving large, $SU(2) \times U(1)$-invariant Majorana masses, M_R. The resulting spectrum consists of heavy neutrinos with masses $\sim M_R$ which are primarily right-handed, and neutrinos with extremely small masses $m_\nu \sim m_D^2/M_R$ which are primarily left-handed.

Within grand unified theories (GUTs) such as SO(10), the Dirac neutrino masses are predicted to be of order the corresponding up-quark masses; for example, $(m_D)_{\nu_\tau}$ would be roughly 100 to 200 GeV. Atmospheric neutrino data favors a ν_τ mass of about 0.04 eV [4]. Thus, one finds a right-handed Majorana mass M_R of order 10^{14} GeV several orders of magnitude below the Planck scale ($M_{\rm Pl}$).

Majorana neutrino masses and neutrino oscillations imply lepton flavor violation (LFV). But within the SM, flavor violation in charged lepton processes is necessarily generated by irrelevant operators and is, therefore, suppressed by powers of $1/M_R$. Within supersymmetric extensions of the SM, however, this is no longer true. In the minimal supersymmetric standard model (MSSM) (which we henceforth take to be augmented with heavy right-handed neutrinos, ν_R), LFV can be communicated directly from ν_R to the sleptons by relevant operators and from there to the charged leptons. LFV is then suppressed by powers of $1/M_{\rm SUSY}$ instead of $1/M_R$, with $M_{\rm SUSY} \ll M_R$. The initial communication is done most economically through renormalization group flow of the slepton mass matrices at energies between $M_{\rm Pl}$ and M_R. Though the scale M_R is far above the weak scale, the presence of the ν_R at scales above M_R leaves an imprint on the mass matrices of the sleptons which is preserved down to the weak scale. This effect has been used to predict large branching fractions for $\tau \to \mu\gamma$ and $\mu \to e\gamma$ within the MSSM [5–7].

In this Letter, we demonstrate a new way in which the imprint of LFV on the slepton mass matrices can be communicated to charged leptons through the exchange of Higgs bosons, allowing for observable flavor violation in the leptonic sector. We demonstrate that the decays $\tau \to 3\mu$ and $\mu \to 3e$ are particularly sensitive probes of LFV at large $\tan\beta$ (the ratio $\langle H_u\rangle/\langle H_d\rangle$), with branching fractions scaling as $\tan^6\beta$.

Flavor violation among the sleptons.—In the leptonic sector, we begin with a Lagrangian:

$$-\mathcal{L} = \overline{E}_R Y_E L_L H_d + \overline{\nu}_R Y_\nu L_L + \tfrac{1}{2}\nu_R^{\top} M_R \nu_R + \text{H.c.}, \tag{1}$$

where E_R, L_L, and ν_R represent 3×1 matrices in flavor space of right-handed charged leptons, left-handed lepton doublets, and right-handed neutrinos, and Y_E, Y_ν, and M_R are 3×3 matrices in flavor space; for example, $E_R = (e_R, \mu_R, \tau_R)^{\top}$. This Lagrangian clearly violates both family and total lepton number due to the presence of the ν_R Majorana mass term. We can choose to work in a basis in which both Y_E and M_R have been diagonalized, but Y_ν remains an arbitrary, complex matrix.

Within the SM, $O(1)$ flavor violation in the neutrinos does not translate into appreciable flavor violation in the charged lepton sector due to $1/M_R$ suppressions. But this is not true in the slepton sector of the MSSM. The supersymmetry (SUSY)-breaking slepton masses are unprotected by chiral symmetries and are, therefore, sensitive to physics at all mass scales between $m_{\tilde{L}}$ and the scale, M, at which SUSY breaking is communicated to the visible sector. This can be seen by examining the renormalization group equation for $m_{\tilde{L}}^2$ at scales above M_R:

$$\frac{d}{d\log Q}(m_{\tilde{L}}^2)_{ij} = \left(\frac{d}{d\log Q}(m_{\tilde{L}}^2)_{ij}\right)_{\rm MSSM} + \frac{1}{16\pi^2}[m_{\tilde{L}}^2 Y_\nu^\dagger Y_\nu + Y_\nu^\dagger Y_\nu m_{\tilde{L}}^2 + 2(Y_\nu^\dagger m_{\tilde{\nu}R}^2 Y_\nu + m_{H_u}^2 Y_\nu^\dagger Y_\nu + A_\nu^\dagger A_\nu)]_{ij}, \quad (2)$$

where the first term represents the (L-conserving) terms present in the usual MSSM at scales below M_R. Because Y_ν is off diagonal, it will generate flavor mixing in the slepton mass matrix. We can solve this equation approximately for the flavor-mixing piece:

$$(\Delta m_{\tilde{L}}^2)_{ij} \simeq \xi (Y_\nu^\dagger Y_\nu)_{ij}, \quad (3)$$

where

$$\xi = -\frac{\log(M/M_R)}{16\pi^2}(6 + 2a^2)m_0^2, \quad (4)$$

m_0 is a common scalar mass evaluated at the scale $Q = M$, a is $O(1)$, and $i \neq j$. In the simplest SUSY-breaking scenarios, where gravity is the messenger, $M = M_{\rm Pl}$, and the logarithm in Eq. (4) is roughly 10.

What does experiment tell us about the values of these matrices? Global fits to neutrino data favor large mixing between the ν_μ and ν_τ, and also between ν_e and ν_μ [4]. The following approximate form for the light neutrino mass matrix, m_ν, provides an excellent fit to existing neutrino data and can be motivated by theory [8]:

$$m_\nu \propto \begin{pmatrix} \epsilon & \epsilon & \epsilon \\ \epsilon & 1 & 1 \\ \epsilon & 1 & 1 \end{pmatrix}, \quad (5)$$

where ϵ is a small parameter ~ 0.1. If we further assume that M_R is an identity matrix, then $Y_\nu^\dagger Y_\nu$ will also have the form of Eq. (5). Another interesting possibility is provided by GUT models with lopsided mass matrices for charged leptons [9]; such models have $(Y_E)_{32} \simeq (Y_E)_{33}$ and lead to a light neutrino mass matrix as in Eq. (5) with $(Y_\nu)_{32} \simeq (Y_\nu)_{33} \simeq y_t$, where y_t is the top Yukawa coupling. In either case, the $Y_\nu^\dagger Y_\nu$ has $O(1)$ flavor violation in the ν_τ-ν_μ sector if $M_R \simeq 10^{14}$ GeV.

Higgs-mediated flavor violation.—Unlike the SM, the MSSM is not protected against the possibility of flavor-changing neutral currents (FCNCs) mediated by neutral Higgs bosons. Though the MSSM is a type-II two-Higgs doublet model at tree level, this structure is not protected by any symmetry. In particular, the presence of a nonzero μ term, coupled with SUSY breaking, is enough to induce nonholomorphic Yukawa interactions for the quarks and leptons. In the quark sector this was discovered by Hall, Rattazzi, and Sarid [10]; terms of the form $\overline{Q}_L u_R H_d^\dagger$ and $\overline{Q}_L d_R H_u^\dagger$ were found, the latter providing a significant correction to the b-quark mass at large $\tan\beta$. We have shown in a previous Letter [11] that these terms allow the neutral Higgs bosons to mediate FCNCs, in particular $B \to \mu\mu$. There we argued that branching fractions predicted at large $\tan\beta$ can be probed at Run II of the Tevatron. (For recent analyses of Higgs-mediated $B \to \mu\mu$, see Ref. [12].)

The two leading diagrams considered in Refs. [10,11] as a source for nonholomorphic quark couplings are not present in the leptonic sector since they involve gluinos and top squarks inside the loops. However, there are additional diagrams which are present in the leptonic sector [13] involving loops of sleptons and charginos/neutralinos; a subset of these is shown in Fig. 1. Given a source of nonholomorphic couplings and LFV among the sleptons, Higgs-mediated LFV is unavoidable.

The effective Lagrangian for the couplings of the charged leptons to the neutral Higgs fields can be written as:

$$-\mathcal{L} = \overline{E}_R Y_E E_L H_d^0 + \overline{E}_R Y_E (\epsilon_1 \mathbf{1} + \epsilon_2 Y_\nu^\dagger Y_\nu) E_L H_u^{0*} + \text{H.c.} \quad (6)$$

The first term is the usual Yukawa coupling, while the second term arises from the nonholomorphic loop corrections. LFV results from our inability to simultaneously diagonalize Y_E and the $\epsilon_2 Y_E Y_\nu^\dagger Y_\nu$ term; as $\epsilon_2 \to 0$, LFV in the Higgs sector will disappear.

The diagrams which contribute to ϵ_2 are shown in Fig. 1. Each diagram contains a single insertion of $\Delta m_{\tilde{L}}^2$ which introduces LFV into the process. Without this insertion, these diagrams would have a trivial flavor structure and would not contribute to ϵ_2 or to LFV. But the $\Delta m_{\tilde{L}}^2$ insertion introduces a $Y_\nu^\dagger Y_\nu$ into the diagram, yielding a contribution to ϵ_2. We can approximate the contributions of the diagrams in Fig. 1 by inserting a single $\Delta m_{\tilde{L}}^2$ mass insertion onto each of the internal $\tilde{E}_L$ lines. We also treat the higgsinos and gauginos as approximate mass eigenstates. The four diagrams in Fig. 1 contribute to ϵ_2 as:

$$\epsilon_2 \simeq \frac{\alpha'}{4\pi}\xi\mu M_1 f_2(M_1^2, m_{\tilde{\ell}L}^2, m_{\tilde{\tau}L}^2, m_{\tilde{\ell}R}^2) - \frac{\alpha'}{8\pi}\xi\mu M_1 f_2(\mu^2, m_{\tilde{\ell}L}^2, m_{\tilde{\tau}L}^2, M_1^2) + \frac{\alpha_2}{4\pi}\xi\mu M_2 f_2(\mu^2, m_{\tilde{\nu}\ell}^2, m_{\tilde{\nu}\tau}^2, M_2^2) + \frac{\alpha_2}{8\pi}\xi\mu M_2 f_2(\mu^2, m_{\tilde{\ell}L}^2, m_{\tilde{\tau}L}^2, M_2^2). \quad (7)$$

In these equations, $M_{1,2}$ are the U(1) and SU(2) gaugino masses, ξ is defined in Eq. (4), and the function f_2 is defined such that

$$-f_2(a, b, c, d) \equiv \frac{a\log(a)}{(a-b)(a-c)(a-d)} + \text{cyclic}. \quad (8)$$

The function f_2 is positive definite and we note several interesting limits in its behavior. When $a = b = c = d$, $f_2(a, a, a, a) = 1/(6a^2)$; and when $a \gg b = c = d$, then $f_2(a, b, b, b) = 1/(2ab)$. Whether $\tilde{\ell} = \tilde{\mu}$ or $\tilde{e}$ depends on the decay process we are considering.

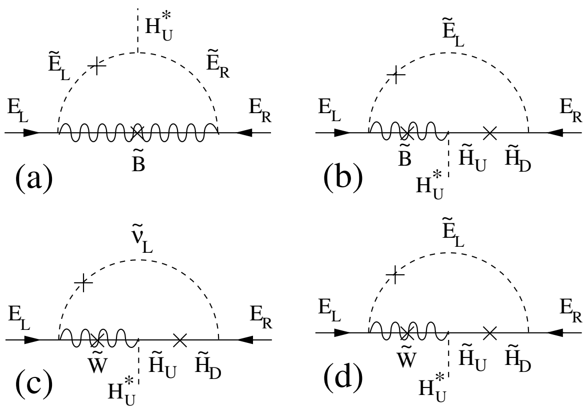

FIG. 1. Diagrams that contribute to ϵ_2. The crosses on the internal slepton lines represent LFV mass insertions due to loops of ν_R.

It is clear from the Lagrangian in Eq. (6) that the charged lepton masses cannot be diagonalized in the same basis as their Higgs couplings. This will allow neutral Higgs bosons to mediate LFV processes with rates proportional to ϵ_2^2. But in order to proceed, we choose a specific process: $\tau \rightarrow 3\mu$. Our discussions generalize to related processes (such as $\tau \rightarrow \mu ee$) very easily.

We only mention in passing the contributions to ϵ_1, since they do not induce LFV. The diagrams which contribute to ϵ_1 are (mostly) those of Fig. 1 without the slepton mass insertion. The contributions of these diagrams to ϵ_1 are found to be [13]:

$$\begin{aligned}\epsilon_1 = \frac{\alpha'}{8\pi}\mu M_1[&2f_1(M_1^2, m_{\tilde{\ell}L}^2, m_{\tilde{\ell}R}^2) - f_1(M_1^2, \mu^2, m_{\tilde{\ell}L}^2) \\ &+ 2f_1(M_1^2, \mu^2, m_{\tilde{\ell}R}^2)] \\ &+ \frac{\alpha_2}{8\pi}\mu M_2[f_1(\mu^2, m_{\tilde{\ell}L}^2, M_2^2) + 2f_1(\mu^2, m_{\tilde{\nu}}^2, M_2^2)],\end{aligned} \tag{9}$$

where

$$-f_1(a,b,c) \equiv \frac{ab\log(a/b) + bc\log(b/c) + ca\log(c/a)}{(a-b)(b-c)(c-a)}. \tag{10}$$

These terms generate a mass shift for the charged leptons that appears in our final formulas as a second-order effect.

Flavor-violating tau decays.—Extracting the $\overline{\tau}_R\,\mu_L$ terms in the effective Lagrangian of Eq. (6), the relevant LFV interaction has the form:

$$\begin{aligned}-\mathcal{L} \simeq (2G_F^2)^{1/4}\frac{m_\tau \kappa_{32}}{\cos^2\beta}(\overline{\tau}_R\,\mu_L) \\ \times[\cos(\beta-\alpha)h^0 - \sin(\beta-\alpha)H^0 - iA^0] \\ + \text{H.c.},\end{aligned} \tag{11}$$

where

$$\kappa_{ij} = -\frac{\epsilon_2}{\{1 + [\epsilon_1 + \epsilon_2(Y_\nu^\dagger Y_\nu)_{33}]\tan\beta\}^2}(Y_\nu^\dagger Y_\nu)_{ij}. \tag{12}$$

[The Lagrangian for $(\overline{\tau}_R\,e_L)$-Higgs can be derived from this by replacing κ_{32} with κ_{31}.]

Given an LFV $(\overline{\tau}_R\,e_L)$-Higgs interaction, the decay $\tau \rightarrow 3\mu$ can be generated via exchange of h^0, H^0, and A^0 as in Fig. 2. The diagram in Fig. 2 is straightforward to calculate. We derive a branching fraction for $\tau \rightarrow 3\mu$ of

$$B(\tau \rightarrow 3\mu) = \frac{G_F^2 m_\mu^2 m_\tau^7 \tau_\tau}{768\pi^3 m_A^4}\kappa_{32}^2\tan^6\beta, \tag{13}$$

where τ_τ is the τ lifetime. To derive this formula, we took the large m_A limit in which $\alpha \rightarrow \beta - \pi/2$.

How large can this branching fraction be? Consider the case in which $\mu = M_1 = M_2 = m_{\tilde{\ell}} = m_{\tilde{\nu}}$, $M_R = 10^{14}$ GeV, and $(Y_\nu^\dagger Y_\nu)_{32} = 1$. Then $\epsilon_2 \simeq 4 \times 10^{-4}$ and

$$B(\tau \rightarrow 3\mu) \simeq (1 \times 10^{-7}) \times \left(\frac{\tan\beta}{60}\right)^6 \times \left(\frac{100\text{ GeV}}{m_A}\right)^4, \tag{14}$$

which puts it into the regime that is experimentally accessible at B-factories over the next few years. At Large Hadron Collider and Super-KEKB, limits in the region of 10^{-9} should be achievable [14], allowing a deeper probe into the parameter space. We can also do better if $\mu \gg M_{1,2} \simeq m_{\tilde{\ell}}^2$. Then the bino contribution is enhanced by a factor μ/M_1; for $M_1 \simeq 100$ GeV and $\mu \simeq 1$ TeV, one can get $\epsilon_2 \simeq 8 \times 10^{-4}$, resulting in a branching fraction 4 times that stated above. However, we note that the value of $\epsilon_2 \simeq 4 \times 10^{-4}$ is remarkably stable to changes in the SUSY spectrum apart from this large-μ option.

Discussion.—We have demonstrated that LFV in the sleptons can generate large LFV in the couplings of leptons to neutral Higgs bosons. But it is already well known that sleptonic flavor violation can induce LFV in certain magnetic moment transitions such as $\tau \rightarrow \mu\gamma$. However, the two decays possess very different decoupling behavior so that either one could be large while the other is too small to observe. The effective operator for $\tau \rightarrow 3\mu$ is dimension-6: $(1/m_A^2)\overline{\tau}\mu\overline{\mu}\mu$. The $\tau \rightarrow \mu\gamma$ operator is formally dimension-5, but chiral symmetry requires an m_τ insertion, so that the operator is actually dimension-6: $(m_\tau/M_{\text{SUSY}}^2)\overline{\tau}\sigma^{\mu\nu}\mu F_{\mu\nu}$ where M_{SUSY} represents the heaviest mass scale to enter the slepton-gaugino loops. If sleptons and gauginos are light and A^0 is heavy, then $\tau \rightarrow \mu\gamma$ would tend to dominate; in the opposite limit and with $\tan\beta$ large, $\tau \rightarrow 3\mu$ would dominate. Because of this different decoupling behavior, it is impossible to correlate the two decays without choosing a specific model. Turning this around, observation of one or both of these decays can provide insight into the fundamental SUSY-breaking parameters.

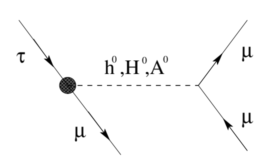

FIG. 2. The Feynman diagram contributing to $\tau \rightarrow 3\mu$. The shaded interaction vertex is the new vertex derived in Eq. (11).

The presence of the $\tau \to \mu\gamma$ operator can also lead to $\tau \to 3\mu$ if the photon goes off shell. However, for this operator the relation between the two branching fractions is roughly model independent [6]: $B(\tau \to 3\mu)/B(\tau \to \mu\gamma) \simeq 0.003$ without Higgs mediation. If a ratio much larger than 0.003 is discovered, then this would be clear evidence that some new process is generating the $\tau \to 3\mu$ decay, with Higgs mediation a leading contender.

Our calculation so far has also relied on Yukawa matrix *ansätze* which reproduce the mass matrix of Eq. (5). What about a different choice for $Y_\nu^\dagger Y_\nu$? Another popular option is the inverted mass hierarchy in which the $(1,2)$ and $(1,3)$ elements of $Y_\nu^\dagger Y_\nu$ would be $O(1)$ and the remainder $O(\epsilon)$ [8]. Such a matrix could lead to observable $\tau \to e\mu\mu$. The constraints coming from $\mu \to e\gamma$ are strong; for $M_{\rm SUSY} \simeq 100$ GeV one finds $(Y_\nu^\dagger Y_\nu)_{21}$ has to be $\lesssim 10^{-2}$ [15]. But in the large $M_{\rm SUSY}$, large $\tan\beta$ limit, this bound is weakened and $\tau \to e\mu\mu$ could dominate.

There is another interesting test of the inverted hierarchy models presented by Higgs mediation. The decay $\mu \to 3e$ can also proceed by neutral Higgs exchange. Though the electron Yukawa coupling is tiny, this is offset by the extreme precision of rare μ-decay searches. In particular, we find that for $(Y_\nu^\dagger Y_\nu)_{21} = 1$,

$$B(\mu \to 3e) \simeq (5 \times 10^{-14}) \times \left(\frac{\tan\beta}{60}\right)^6 \times \left(\frac{100\ \text{GeV}}{m_A}\right)^4. \tag{15}$$

Again, this process can be generated by the similar $\mu \to e\gamma$ operator taken off shell, but there the ratio is again model independent: $B(\mu \to 3e)/B(\mu \to e\gamma) \simeq 0.006$ without Higgs mediation. Any deviation from this fixed ratio would, as for the taus, be a strong indication of new physics such as that found here.

Finally, we mention here that one could also calculate the rate for processes involving LFV Higgs couplings at both vertices, though we leave this computation to a future work. For example, a second way to generate $\tau \to e\mu\mu$ would be to use the $\overline{\tau}_R\,\mu_L$-Higgs coupling that we have been considering in this Letter, along with a $\overline{\mu}_R e_L$-Higgs coupling. We can also generate $\tau \to ee\mu$ with $\overline{\tau}_R\,e_L$-Higgs and $\overline{\mu}_R e_L$-Higgs couplings. But as above, both processes are constrained by nonobservation of $\mu \to e\gamma$ since they require large $(\Delta m_{\tilde{L}}^2)_{21}$. It is notable that these processes have a remarkable $\tan^8\beta$ dependence; however, the additional powers are mitigated by additional loop suppressions.

We thank the scientific staffs of CERN and DESY for their hospitality while this work was completed. This research was supported by the National Science Foundation under Grant No. PHY00-98791, by the U.S. Department of Energy under Grants No. DE-FG03-98ER41076 and No. DE-FG02-01ER45684, and by a grant from the Research Corporation.

[1] Super-Kamiokande Collaboration, S. Fukuda *et al.*, Phys. Rev. Lett. **85**, 3999 (2000).

[2] Super-Kamiokande Collaboration, S. Fukuda *et al.*, Phys. Rev. Lett. **86**, 5656 (2001); SNO Collaboration, Q. R. Ahmad *et al.*, Phys. Rev. Lett. **89**, 011301 (2002).

[3] M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by P. van Nieuwenhuizen and D. Z. Freedman (North-Holland, Amsterdam, 1979); T. Yanagida, in *Proceedings of the Workshop on Unified Theory and Baryon Number in the Universe*, edited by O. Sawada and A. Sugamoto (KEK, Tsukuba, 1979); R. N. Mohapatra and G. Senjanović, Phys. Rev. Lett. **44**, 912 (1980).

[4] For a review, see M. C. Gonzalez-Garcia and Y. Nir, hep-ph/0202058.

[5] F. Borzumati and A. Masiero, Phys. Rev. Lett. **57**, 961 (1986); L. J. Hall, V. A. Kostelecky, and S. Raby, Nucl. Phys. **B267**, 415 (1986).

[6] J. Hisano, T. Moroi, K. Tobe, M. Yamaguchi, and T. Yanagida, Phys. Lett. B **357**, 579 (1995); J. Hisano, T. Moroi, K. Tobe, and M. Yamaguchi, Phys. Rev. D **53**, 2442 (1996).

[7] S. F. King and M. Oliveira, Phys. Rev. D **60**, 035003 (1999); J. Hisano and D. Nomura, Phys. Rev. D **59**, 116005 (1999); W. Buchmuller, D. Delepine, and F. Vissani, Phys. Lett. B **459**, 171 (1999); K. S. Babu, B. Dutta, and R. N. Mohapatra, Phys. Lett. B **458**, 93 (1999); J. R. Ellis *et al.*, Eur. Phys. J. C **14**, 319 (2000); J. Sato, K. Tobe, and T. Yanagida, Phys. Lett. B **498**, 189 (2001); S. Lavignac, I. Masina, and C. A. Savoy, Phys. Lett. B **520**, 269 (2001); J. A. Casas and A. Ibarra, Nucl. Phys. **B618**, 171 (2001).

[8] For reviews, see G. Altarelli and F. Feruglio, Phys. Rep. **320**, 295 (1999); S. M. Barr and I. Dorsner, Nucl. Phys. **B585**, 79 (2000).

[9] C. H. Albright, K. S. Babu, and S. M. Barr, Phys. Rev. Lett. **81**, 1167 (1998); J. Sato and T. Yanagida, Phys. Lett. B **430**, 127 (1998); N. Irges, S. Lavignac, and P. Ramond, Phys. Rev. D **58**, 035003 (1998).

[10] L. J. Hall, R. Rattazzi, and U. Sarid, Phys. Rev. D **50**, 7048 (1994).

[11] K. S. Babu and C. Kolda, Phys. Rev. Lett. **84**, 228 (2000).

[12] C. S. Huang, W. Liao, Q. S. Yan, and S. H. Zhu, Phys. Rev. D **63**, 114021 (2001); A. Dedes, H. K. Dreiner, and U. Nierste, Phys. Rev. Lett. **87**, 251804 (2001); G. Isidori and A. Retico, J. High Energy Phys. **0111**, 001 (2001); R. Arnowitt, B. Dutta, T. Kamon, and M. Tanaka, Phys. Lett. B **538**, 121 (2002); C. Bobeth *et al.*, hep-ph/0204225; H. Baer *et al.*, hep-ph/0205325.

[13] K. S. Babu, C. Kolda, J. March-Russell, and F. Wilczek, Phys. Rev. D **59**, 016004 (1999).

[14] F. Deppisch *et al.*, hep-ph/0206122.

[15] Y. G. Kim, P. Ko, J. S. Lee, and K. Y. Lee, Phys. Rev. D **59**, 055018 (1999); J. Ellis, J. Hisano, M. Raidal, and Y. Shimizu, hep-ph/0206110.