Doubly charged Higgsino-mediated lepton flavor violating decays

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We analyze the supersymmetric contribution to the lepton flavor decays due to the presence of a doubly charged Higgsino in a general left-right model. We include contributions to $l \to l' \gamma$, μ-e conversion in nuclei, muonium-antimuonium decay, $l^- \to l_1^- l_2^- l_3^+$ and flavor changing $Z \to l_1 \overline{l_2} + l_2 \overline{l_1}$. We present a complete set of bounds on the couplings and masses of the doubly charged Higgsino and we discuss which of these processes will be most sensitive to the presence of such an exotic fermion. Lepton-flavor violating decays are shown to be a promising clue to an extended gauge structure in supersymmetry.

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I. INTRODUCTION

The conservation of lepton number and lepton flavor is among the most stringently tested laws of physics. In the standard model all three lepton flavors are exact global symmetries and are conserved separately. However, this is a consequence of vanishing neutrino masses, which has come into serious question lately [1]. Motivated by the data from SuperKamiokande, indicating oscillations of atmospheric and solar neutrinos, there has been a renewed interest in lepton flavor violation. This motivation is further enhanced by improved experimental upper limits on several interesting decays, such as $\mu \to e\gamma$, $\tau \to \mu\gamma$, $\mu^+ \to e^+e^+e^-$ as well as $\mu^-\text{Ti} \to e^-\text{Ti}$. Projects are currently underway to improve several of these upper limits by a few orders of magnitude. Coupled with neutrino masses and oscillations, looking for charged lepton number and flavor violation seems to be a promising signal for looking for physics beyond the standard model.

On the other hand, the tantalizing existence of neutrino masses is a definite reason to look at scenarios beyond the standard model in order to accommodate small neutrino masses. In that regard, the most elegant solution to neutrino masses is the generation of such masses through the seesaw mechanism, present in the left-right symmetric model. Choosing to break the extended symmetry $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ to the standard model through a triplet Higgs representation is the natural way to generate small neutrino masses. If, in addition, one would want to provide solutions to two seemingly unrelated problems of the standard model, the stability of the Higgs boson mass and the origin of the electroweak symmetry breaking, one requires a supersymmetric extension of that model, the left-right supersymmetric (LRSUSY) model. A careful analysis of the phenomenological consequences of this model reveals that one in fact gets solutions to more than the problems outlined above, such as exact *R*-parity conservation in the superpotential [2] and a solution to the strong and electroweak *CP* problem [3].

In the supersymmetric version of the left-right model, one encounters not just triplet Higgs bosons, but their supersymmetric partners as well, the Higgsinos. The conditions imposed on these Higgs bosons to give mass to the neutrinos through the seesaw mechanisms therefore introduce four exotic doubly charged Higgsinos which couple to leptons only and as such are an interesting source of lepton flavor violation. Before examining the problem closer, we want to point out that extensive analyses have been done regarding doubly charged scalar particles [4], either their production at various colliders or their lepton flavor violating interactions. Besides the complex $Y = \pm 2$ triplet Higgs representations of the left-right supersymmetric model, the doubly charged Higgs scalars appear in many models, either by augmenting the standard model with additional Higgs representations or as a natural consequence of extended gauge structures [5]. Any supersymmetric extensions of such models will contain exotic doubly charged Higgsinos. Of course, in assessing the attractiveness of a particular choice of Higgs represenations, one must consider the severity of constraints needed to be satisfied. For triplet and higher representations containg a neutral member, limits on the latter's vacuum expectation value needed to maintain $\rho = 1$ are generally very restrictive. Even imposing $\rho = 1$ at the tree level requires fine-tuning to maintain this relation at the one loop level. In order to avoid this, one must choose either representations that do not have a neutral member or in which the expectation value of the neutral member is zero, or very small, as is the case of the LRSUSY left-handed Higgs triplet. Another reason to favor the doubly charged Higgsinos of the LRSUSY model is the requirement of coupling constant unification. This is quite difficult to maintain with an arbitrary Higgs structure, but is possible for the LRSUSY model with intermediate scales comfortably adjusted such that coupling constant unification is achieved.

For these reasons, we will concentrate on the study of the doubly charged Higgsinos of the left-right supersymmetric theory. Analyses of some flavor-violating decays appeared earlier [6]. We propose to take a more general point of view here, also to combine all sources of lepton flavor violation and discuss their relative importance. The study has some general features, in that we attempt to put constraints on the masses and Yukawa couplings of these exotic Higgsinos by assuming the simple one parameter coupling charged-lepton–charged-lepton interaction. Given their exotic quantum numbers, the masses and couplings of the doubly charged Higgsinos are free parameters of the model and are

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unaffected by interactions in other sector. In that respect the study of their interactions is ‘‘clean’’ and will apply to any other such objects which couple similarly to two leptons only.

Our paper is organized as follows: we will first present a brief review of the LRSUSY model, with particular emphasis on the doubly charged Higgsino sector (Sec. II). We then proceed to analyze the lepton flavor violating decays $l \to l' \gamma$ (Sec. III), μ-e conversion (Sec. IV), $l \to l_1 l_2 l_3$ (Sec. V), muonium-antimuonium conversion (Sec. VI) and $Z \to \overline{l} l' + \overline{l'} l$ (Sec. VII). We analyze these results and comment on their relative suitability for investigating lepton flavor violation in Sec. VIII, and reach our conclusion in Sec. IX.

II. LEFT-RIGHT SUPERSYMMETRIC MODEL AND DOUBLY CHARGED HIGGSINOS

The LR SUSY model, based on $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, has matter doublets for both left- and right-handed fermions and the corresponding left- and right-handed scalar partners (sleptons and squarks) [7]. In the gauge sector, corresponding to $SU(2)_L$ and $SU(2)_R$, there are triplet gauge bosons $(W^{+,-},W^0)_L$, $(W^{+,-},W^0)_R$ and a singlet gauge boson V corresponding to $U(1)_{B-L}$, together with their superpartners. The Higgs sector of this model consists of two Higgs bi-doublets, $\Phi_u(\frac{1}{2},\frac{1}{2},0)$ and $\Phi_d(\frac{1}{2},\frac{1}{2},0)$, which are required to give masses to both the up and down quarks. The phenomenology of the doublet Higgs is similar to the non-supersymmetric left-right model [8], except that the second pair of Higgs doublet fields, which provide new contributions to the flavor-changing neutral currents, must be heavy, in the 5–10 TeV range, effectively decoupling from the low-energy spectrum [9]. The spontaneous symmetry breaking of the group $SU(2)_R \times U(1)_{B-L}$ to the hypercharge symmetry group $U(1)_Y$ is accomplished by the vacuum expectation values of a pair of Higgs triplet fields $\Delta_L(1,0,2)$ and $\Delta_R(0,1,2)$, which transform as the adjoint representation of $SU(2)_R$. The choice of the triplets (versus four doublets) is preferred because with this choice a large Majorana mass can be generated (through the seesaw mechanism) for the right-handed neutrino and a small one for the left-handed neutrino [8]. In addition to the triplets $\Delta_{L,R}$, the model must contain two additional triplets $\delta_L(1,0,-2)$ and $\delta_R(0,1,-2)$, with quantum number $B-L=-2$ to ensure cancellation of the anomalies that would otherwise occur in the fermionic sector. Given their strange quantum numbers, the δ_L and δ_R do not couple to any of the particles in the theory, so their contribution is negligible for any phenomenological studies.

The superpotential for the LRSUSY model is

$$W = \mathbf{h}_q^{(i)} Q^T \tau_2 \Phi_i \tau_2 Q^c + \mathbf{h}_l^{(i)} L^T \tau_2 \Phi_i \tau_2 L^c + i(\mathbf{h}_{ij} L^{iT} \tau_2 \Delta_L L^j + \mathbf{h}_{ij} L^{icT} \tau_2 \Delta_R L^{jc}) + M_\Delta [\mathrm{Tr}(\Delta_L \overline{\Delta}_L + \mathrm{Tr}(\Delta_R \overline{\Delta}_R)] + \mu_{ij} \mathrm{Tr}(\tau_2 \Phi_i^T \tau_2 \Phi_j) + W_{NR} \tag{1}$$

where W_{NR} denotes (possible) non-renormalizable terms arising from higher scale physics or Planck scale effects [10]. The presence of these terms ensures that, when the SUSY breaking scale is above M_{W_R}, the ground state is R-parity conserving [11]. In writing the superpotential we have assumed strict LR symmetry.

As in the standard model, in order to preserve $U(1)_{EM}$ gauge invariance, only the neutral Higgs fields acquire nonzero vacuum expectation values (VEV’s). These values are

$$\langle \Delta_L \rangle = \begin{pmatrix} 0 & 0 \\ v_L & 0 \end{pmatrix}, \quad \langle \Delta_R \rangle = \begin{pmatrix} 0 & 0 \\ v_R & 0 \end{pmatrix}, \quad \langle \Phi \rangle = \begin{pmatrix} \kappa & 0 \\ 0 & \kappa' e^{i\omega} \end{pmatrix}.$$

$\langle \Phi \rangle$ causes the mixing of W_L and W_R bosons with CP-violating phase ω. In order to simplify, we will take the VEV’s of the Higgs fields as $\langle \Delta_L \rangle = 0$ and

$$\langle \Delta_R \rangle = \begin{pmatrix} 0 & 0 \\ v_R & 0 \end{pmatrix}, \quad \langle \Phi_u \rangle = \begin{pmatrix} \kappa_u & 0 \\ 0 & 0 \end{pmatrix}, \quad \langle \Phi_d \rangle = \begin{pmatrix} 0 & 0 \\ 0 & \kappa_d \end{pmatrix}.$$

Choosing $v_L = \kappa' = 0$ satisfies the more loosely required hierarchy $v_R \gg \max(\kappa,\kappa') \gg v_L$ and also the required cancellation of flavor-changing neutral current. The Higgs fields acquire nonzero VEV’s to break both parity and $SU(2)_R$. In the first stage of breaking, the right-handed gauge bosons W_R and Z_R acquire masses proportional to v_R and become much heavier than the usual (left-handed) neutral gauge bosons W_L and Z_L, which pick up masses proportional to κ_u and κ_d at the second stage of breaking.

The supersymmetric sector of the model, while preserving left-right symmetry, has eight singly charged charginos, corresponding to $\tilde{\lambda}_L, \tilde{\lambda}_R, \tilde{\phi}_u, \tilde{\phi}_d, \tilde{\Delta}_L^-, \tilde{\Delta}_R^-, \tilde{\delta}_L^-$ and $\tilde{\delta}_R^-$. The model also has 11 neutralinos, corresponding to $\tilde{\lambda}_Z, \tilde{\lambda}_{Z'}, \tilde{\lambda}_V, \tilde{\phi}_{1u}^0, \tilde{\phi}_{2u}^0, \tilde{\phi}_{1d}^0, \tilde{\phi}_{2d}^0, \tilde{\Delta}_L^0, \tilde{\Delta}_R^0, \tilde{\delta}_L^0$, and $\tilde{\delta}_R^0$. Although Δ_L is not necessary for symmetry breaking [12] and is introduced only for preserving left-right symmetry, both Δ_L^{--} and its right-handed counterpart Δ_R^{--} play very important roles in phenomenological studies of the LRSUSY model. It has been shown that these bosons, and possibly their fermionic counterparts, could be light [10]. The production of doubly charged Higgs and their corresponding Higgsinos has been studied extensively. It has also been shown in the past that their presence enhances lepton-flavor violating decays, but a complete analysis of their interactions is still lacking. We propose to remedy this in the present paper.

The two-component mass terms for the doubly charged Higgsinos are derived from the superpotential,

$$\mathcal{L}_{mass} = -M_{LR} \tilde{\Delta}_L^{++} \tilde{\overline{\Delta}}_L^{--} - M_{LR} \tilde{\Delta}_R^{++} \tilde{\overline{\Delta}}_R^{--}, \tag{2}$$

and the Yukawa interaction of these fields is given by

$$\mathcal{L}_Y = -2 h_{ij} \overline{L_{iL}^c} \tilde{\Delta}_L^{++} \tilde{L}_{jL} - 2 h_{ij} \overline{L_{iR}^c} \tilde{\Delta}_R^{++} \tilde{L}_{jR}. \tag{3}$$

The advantages of studying the lepton-flavor violation with doubly charged Higgsinos is evident. Their masses and interactions do not depend on parameters in other gaugino-Higgsino sectors. In addition, $\tilde{\Delta}_L^{--}$ and $\tilde{\Delta}_R^{--}$ do not mix with each other, so the interactions do not give rise to graphs with mixed slepton states, making the interactions free of other parameters such as the trilinear coupling A_l. The only parameters are the masses ($M_{\tilde{\Delta}_L}$, $M_{\tilde{\Delta}_R}$) and the coupling constants h_{ij}.

III. $l\to l'\gamma$

The decays of the form $l\to l'\gamma$ have been long considered the best channel to search for lepton flavor violation (LFV). In particular, the decay $\mu\to e\gamma$, which proceeds through a simple dipole transition and which is very strongly constrained experimentally, has been the subject of numerous analyses [13]. The experimental bounds on such processes are [14,15]

$$BR(\mu\to e\gamma)<1.2\times 10^{-11} \tag{4}$$

$$BR(\tau\to e\gamma)<2.9\times 10^{-6} \tag{5}$$

$$BR(\tau\to \mu\gamma)<1.1\times 10^{-6}. \tag{6}$$

The amplitude of the $l\to l'\gamma$ transition can be written in the form of the usual dipole-type interaction:

$$\mathcal{M}_{l\to l'\gamma}=\frac{1}{2}\overline{\psi}_{l'}(d_LP_L+d_RP_R)\sigma^{\mu\nu}F_{\mu\nu}\psi_l. \tag{7}$$

It leads to the branching ratio

$$\Gamma_{l\to l'\gamma}=\frac{1}{16\pi}\tau_l(|d|_L^2+|d|_R^2)m_l^3. \tag{8}$$

Comparing it with the standard decay width, $\Gamma_l=(1/192\pi^3)G_F^2m_l^5$, and using the experimental constraint on the branching ratio, we get a limit on the dipole amplitude:

$$|d|=\sqrt{(|d_L|^2+|d_R|^2)/2}. \tag{9}$$

This process was discussed in [6] for $\mu\to e\gamma$, where the h_{ij} couplings were taken to be diagonal and the branching ratio was considered as a function of the (unknown) mixings between the scalar muons and scalar electrons. We will reanalyze this process here by not making any assumptions about slepton mixing (therefore eliminating several unknown parameters) but considering the Yukawa mixings as nondiagonal. The limits on the couplings will depend only on the values of the masses of doubly charged Higgsinos and scalar leptons. We include in this analysis the tighter experimental bound, as well as an analysis of the decays $\tau\to e\gamma$ and $\tau\to\mu\gamma$, which bound a different combination of couplings. The Feynman diagrams that contribute to this process are shown in Fig. 1. We obtain, for the dipole amplitude,

$$|d_{L,R}|=\frac{h_{li}h_{il'}}{(4\pi)^2}\frac{m_l}{M_{\tilde{l}}^2}[f_M(r_{L,R})-2g_M(r_{L,R})] \tag{10}$$

where $r_{L,R}=M^2_{\tilde{\Delta}_{L,R}}/M^2_{\tilde{l}}$ and $\tilde{l}$ is the scalar lepton in the loop, and where the $l\to l'+\gamma$ loop functions are

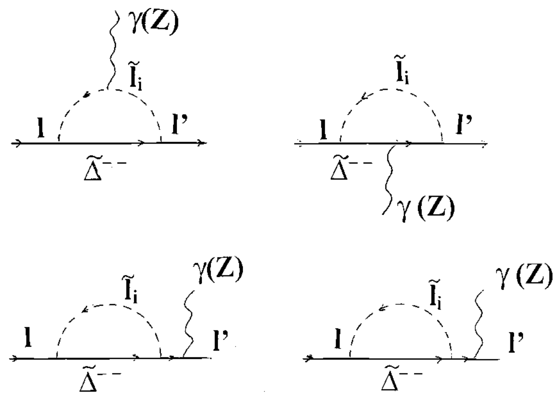

FIG. 1. One-loop Feynman graphs generating the effective vertices $ll'\gamma(Z)$, $l\neq l'$. Here $\tilde{l}_i^{\pm}$ represents a slepton state of flavor i, $i=e,\mu,\tau$ and $\tilde{\Delta}^{--}$ is the doubly charged Higgsino.

$$f_M(r)=\frac{1}{6(1-r)^4}(2r^3+3r^2-6r+1-6r^2\log r) \tag{11}$$

$$g_M(r)=\frac{1}{6(1-r)^4}(r^3-6r^2+3r+2+6r\log r). \tag{12}$$

IV. μ-e CONVERSION

The lepton flavor violating neutrinoless conversion of a bound 1 s muon to electron in the field of a nucleus,

$$(A,Z)+\mu^-\to e^-+(A,Z)^*, \tag{13}$$

is known as one of the best probes to search for muon and electron flavor violation. So far the experiments have put an upper bound on the branching ratio for muon-electron conversion relative to the total rate for muon capture:

$$R_{\mu e^-}=\Gamma(\mu^-\to e^-)/\Gamma(\mu^-\to\nu_\mu). \tag{14}$$

The upper limits extracted at PSI by the SINDRUM II experiments are [16,17]

$$R_{\mu e^-}<6.1\times 10^{-13} \quad \text{for } ^{48}\text{Ti target} \tag{15}$$

$$R_{\mu e^-}<4.6\times 10^{-11} \quad \text{for } ^{208}\text{Pb target.} \tag{16}$$

At present the planned MECO experiment at Brookhaven is being launched using ^{27}Al targets, and the expected sensitivity on $R_{\mu e^-}$ is [18]

$$R_{\mu e^-}<2\times 10^{-17} \quad \text{for } ^{27}\text{Al target,} \tag{17}$$

which implies an improvement over the existing limits of about four orders of magnitude. If this happens, the bounds from μ-e conversion would be stronger than the bounds found from $\mu\to e\gamma$ or $\mu\to eee$. Therefore, whatever the mechanisms considered for μ-e conversion, it can be expected that this process will become the principal test of muon number conservation.

On the theoretical side, many mechanisms of the μ-e conversion have been studied. These mechanisms fall into two categories: photonic and non-photonic. The relevant conversion effective Lagrangian can be expressed as

$$\mathcal{L}_{eff}=\frac{4\pi\alpha}{q^2}j^{\lambda}_{ph}J^{ph}_{\lambda}+\frac{G}{\sqrt{2}}\sum_{i,j}(j^{\lambda}_{i\,nph}J^{i\,nph}_{\lambda}+j_{m\,nph}J^{m\,nph}) \tag{18}$$

where the first term describes photon conversion and the second leptonic and hadronic non-photonic conversion. Here q^2 denotes the photon momentum transfer, G the effective coupling constant, j_i the vector currents and J_i the scalar currents. The photonic mecahnism is enhanced at small q^2 and can occur at the one loop level. The non-photonic mechanism is significant if it can occur at the tree level (which is not the case here), or can be enhanced by decoupling of heavy neutral fermions, such as massive neutrinos. If not, in general the non-photonic contribution is supressed. The leptonic current for the photonic mechanism can be parametrized as

$$j^{\lambda}_{ph}=\bar{e}\left[(f_{E0}+\gamma_5 f_{M0})\gamma_{\nu}\left(g^{\lambda\nu}-\frac{q^{\lambda}q^{\nu}}{q^2}\right)+(f_{M1}+\gamma_5 f_{E1})i\sigma^{\lambda\nu}\frac{q^{\nu}}{m_{\mu}}\right]. \tag{19}$$

In addition one must calculate coherent μ-e conversion nuclear form factor. We will follow previous approaches [19]. The relevant conversion branching ratio can be written as

$$R_{\mu e^-}=C\frac{8\pi\alpha^2}{q^4}p_eE_e\frac{|F(p_e)|^2}{\Gamma_{capt}}\xi_0^2 \tag{20}$$

where

$$\xi_0^2=|f_{E0}+f_{M1}|^2+|f_{E1}+f_{M0}|^2 \tag{21}$$

is the dependence of the matrix element for μ-e conversion on the form factors. $|F(p_e)|^2$ is the nuclear matrix element squared in the local approximation, and C is the correction factor to the approximation. We take the latest numerical values from [20]. The form factors which appear in ξ_0^2 are

$$f_{E0}=\frac{-q^2}{(4\pi)^2M^2_{\tilde{l}}}h_{\mu i}h_{je}[f_E(r_{L,R})-2g_E(r_{L,R})] \tag{22}$$

$$f_{M0}=f_{E1}=0 \tag{23}$$

$$f_{M1}=\frac{-q^2}{(4\pi)^2M^2_{\tilde{l}}}h_{\mu i}h_{je}[f_M(r_{L,R})-2g_M(r_{L,R})] \tag{24}$$

where the additional loop functions for the conversion are

$$f_E(r)=\frac{1}{(1-r)^4}(-18r^3+18r^2-9r+1+6r^3\log r) \tag{25}$$

$$g_E(r)=\frac{1}{(1-r)^4}[-7r^3+36r^2-45r+16+6(2-3r)\log r]. \tag{26}$$

Note that this conversion depends on a different combination of form factors than the decay $l\to l'\gamma$. In particular, only the form factor f_{M1} is responsible for the dipole decay and the form factor f_{E0} is logarithmically enhanced by a factor of $|\ln(m_l^2/M_{\tilde{l}}^2)|\sim\mathcal{O}(10)$ for typical fermion masses. This phenomenum has been discussed previously in the literature [21] and will play an important role here only if $M_{\tilde{l}}\gg M_{\tilde{\Delta}}$.

V. THREE-BODY LEPTON FLAVOR VIOLATING DECAYS

These decays are important sources of LFV, especially the decay $\mu\to eee$, which is severely constrained experimentally. The experimental constraints on these decays are [15]

$$BR(\mu^-\to e^-e^-e^+)<1\times10^{-12} \tag{27}$$

$$BR(\tau^-\to e^+\mu^-\mu^-)<2.9\times10^{-6} \tag{28}$$

$$BR(\tau^-\to e^-e^-e^+)<1.5\times10^{-6} \tag{29}$$

$$BR(\tau^-\to e^-\mu^-\mu^+)<1.8\times10^{-6} \tag{30}$$

$$BR(\tau^-\to e^-e^-\mu^+)<1.5\times10^{-6} \tag{31}$$

$$BR(\tau^-\to \mu^-\mu^-\mu^+)<1.9\times10^{-6} \tag{32}$$

$$BR(\tau^-\to \mu^-e^-e^+)<1.7\times10^{-6}. \tag{33}$$

The branching ratios for these three body decays can be expressed as

$$BR(l\to l_1l_2l_3)=\frac{(h_{li}h_{jl_1})^2}{(4\pi)^4}\frac{m_l}{\Gamma_l}\frac{m_W^4}{m^4_{\tilde{\Delta}}}|\mathcal{F}_{ll_1l_2l_3}|^2. \tag{34}$$

The transition amplitude of the decay $l(p)\to l_1(p_1)l_2(p_2)\overline{l_3}(p_3)$ receives contributions from γ- and Z-mediated graphs shown in Fig. 1 and box diagrams given in Fig. 2. We distinguish three different cases, depending on the flavors of the final states. These three different amplitudes are conveniently written down as follows:

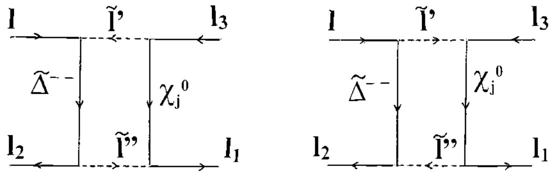

FIG. 2. One-loop Feynman box diagrams contributing to the decays $l\to l_1 l_2 l_3$ decays, as well as $\mu^+e^-\to\mu^-e^+$ conversion. Here $\tilde{l}'(\tilde{l}'')$ represents a slepton state of flavor e,μ,τ; χ_j^0, $j=1,\ldots,11$, represents a neutralino state; and $\tilde{\Delta}^{--}$ is the doubly charged Higgsino.

$$B(l^-\to l_1^- l_2^- l_3^+,l_1\neq l_3,l_2=l_3)=\frac{\left[\sum_{i,j,k=1,2,3} h_{ll_i}h_{l_jl_k}\right]^2}{(4\pi)^4}\frac{m_W^4}{M_{\tilde{\Delta}}^4}\frac{m_l}{\Gamma_l}\Bigg\{|F_{Box}^{ll_1l_2l_2}+F_Z^{ll_1}-2s_w^2(F_Z^{ll_1}-F_\gamma^{ll_1})|^2+4s_w^4|F_Z^{ll_1}-F_\gamma^{ll_1}|^2+8s_w^2\,\Re[(F_Z^{ll_1}+F_{Box}^{ll_1l_2l_2})G_\gamma^{ll_1\,*}]-32s_w^4\,\Re[(F_Z^{ll_1}-F_\gamma^{ll_1})G_\gamma^{ll_1\,*}]+32s_w^4|G_\gamma^{ll_1}|^2\!:\left[\ln\frac{m_l^2}{m_{l_2}^2}-3\right]\Bigg\},\quad(35)$$

where $s_w=\sin\theta_w$ and F_γ^{ll}, $G_\gamma^{ll_1}$ are form factors associated with the functions $2f$ and g associated with the phton vertex, $F_Z^{ll_1}$ is the form factor associated with the Z vetex, and $F_{Box}^{ll_1l_2l_3}$ is the form factor associated with the box diagram. All these composite form factors are defined explicitly in [22]. The decays in this category are $\tau^-\to e^-\mu^-\mu^+$ and $\tau^-\to e^-\mu^-e^+$. The second type of decay is

$$B(l^-\to l_1^- l_2^- l_3^+,l_3=l_1=l_2)=\frac{\left[\sum_{i,j,k=1,2,3} h_{ll_i}h_{l_jl_k}\right]^2}{(4\pi)^4}\frac{m_W^4}{M_{\tilde{\Delta}}^4}\frac{m_l}{\Gamma_l}\Bigg\{2|\tfrac{1}{2}F_{Box}^{ll_1l_1l_1}+F_Z^{ll_1}-2s_w^2(F_Z^{ll_1}-F_\gamma^{ll_1})|^2+4s_w^4|F_Z^{ll_1}-F_\gamma^{ll_1}|^2+16s_w^2\,\Re[(F_Z^{ll_1}+\tfrac{1}{2}F_{Box}^{ll_1l_1l_1})G_\gamma^{ll_1\,*}]-48s_w^4\,\Re[(F_Z^{ll_1}-F_\gamma^{ll_1})G_\gamma^{ll_1\,*}]+32s_w^4|G_\gamma^{ll_1}|^2\left[\ln\frac{m_l^2}{m_{l_1}^2}-\frac{11}{4}\right]\Bigg\}\quad(36)$$

Decays of this type are $\mu^-\to e^-e^-e^+$, $\tau^-\to\mu^-\mu^-\mu^+$ and $\tau^-\to e^-e^-e^+$. Finally for the third type of decay we have

$$B(l^-\to l_1^- l_2^- l_3^+,l_2\neq l_3,l_1\neq l_3)=\frac{\left[\sum_{i,j,k=1,2,3} h_{ll_i}h_{l_jl_k}\right]^2}{(4\pi)^4}\frac{m_W^4}{M_{\tilde{\Delta}}^4}\frac{m_l}{\Gamma_l}|F_{Box}^{ll_1l_2l_3}|^2.\quad(37)$$

The decays in this category are $\tau^-\to e^-e^-\mu^+$ and $\tau^-\to\mu^-\mu^-e^+$. We can approximate the branching ratios for the first two types of decays as

$$B(l^-\to l_1^- l_2^- l_3^+,l_1\neq l_3,l_2=l_3)=\frac{\left[\sum_{i,j,k=1,2,3} h_{ll_i}h_{l_jl_k}\right]^2}{(4\pi)^4}\frac{m_W^4}{M_{\tilde{l}}^4}\frac{m_l}{\Gamma_l}32s_w^4|G_\gamma^{ll_1}|^2\left[\ln\frac{m_l^2}{m_{l_2}^2}-3\right]\quad(38)$$

and

$$B(l^-\to l_1{}^- l_2^- l_3^+,l_3=l_1=l_2)=\frac{\sum_{i,j,k=1,2,3} h_{ll_i}h_{l_jl_k}}{(4\pi)^2}\frac{m_W^4}{M_{\tilde{\Delta}}^4}\frac{m_l}{\Gamma_l}32s_w^4|G_\gamma^{ll_1}|^2\left[\ln\frac{m_l^2}{m_{l_1}^2}-\frac{11}{4}\right].\quad(39)$$

Decays of the form $l^-\to l_1^- l_2^- l_3^+,l_2\neq l_3,l_1\neq l_3$ can occur through only the box diagram and their contribution is too small to place any constraints on masses and couplings of the $\tilde{\Delta}_{L,R}^{--}$. The most stringent bounds from the three body decays will come from $\mu^-\to e^-e^-e^+$, and also from $\tau^-\to e^-e^-e^+$ and $\tau^-\to\mu^-e^-e^+$.

VI. MUONIUM-ANTIMUONIUM CONVERSION

The effective Lagrangian for M-$\bar{M}$ conversion comes from the lepton box amplitude of Fig. 2. Therefore, the structure of the effective Hamiltonian density for M-$\bar{M}$ has the same $(V-A)\times(V-A)$ as in the original papers [23]:

$$\mathcal{H}=G_{M\bar{M}}\bar{\psi}_\mu\gamma_\lambda(1-\gamma_5)\psi_e\bar{\psi}_\mu\gamma_\lambda(1-\gamma_5)\psi_e.\quad(40)$$

The constant $G_{M\bar{M}}$ contains information on physics beyond the standard model. In our case, it contains the box amplitude for the process $\mu^+e^-\to\mu^-e^+$, forbidden in the standard model, but which proceeds through the box diagram of Fig. 2:

$$G_{M\bar{M}}=(|h_{e\mu}|^2+h_{\mu\mu}h_{ee}^*)\frac{\alpha_Z}{M_{\tilde{l}}^2}F_{box}^{\mu ee\mu}.\quad(41)$$

The quantity measured by experiments is the conversion probability $P(M\to\bar{M})$ which is constrained to be

$$P(M\to\bar{M})<1\times10^{-9}.\quad(42)$$

The conversion probability is related to the constant $G_{M\bar{M}}$ by

$$P(M\to\bar{M})=\frac{\delta^2}{2\Gamma_\mu^2}\quad(43)$$

where

$$\frac{\delta}{2}=\langle\overline{M}|H|M\rangle=\frac{16G_{M\overline{M}}}{\pi a^3} \tag{44}$$

is the transition matrix element between muonium and antimuonium, a is the radius of the muonium atom, and Γ_μ is the total decay width of the muon. As in many other extended gauge structures, the M-$\overline{M}$ conversion is not a good place to search for doubly charged Higgsino-mediated lepton-flavor violation. The amplitude for this process depends on the non-diagonal μ-e mixing, which is much more strongly constrained by processes $\mu\rightarrow e\gamma$, $\mu\rightarrow e$ conversion and $\mu\rightarrow eee$. The fact that the M-$\overline{M}$ conversion can only occur through the box diagram also decreases its chances of being the place to look for LVF. If the results of the previous three processes are improved by a factor of n, one needs a comparable improvement of factor of n^2 for the muonium conversion to be comparable in searching for LFV.

VII. $Z\rightarrow L_1\overline{l_2}+L_2\overline{l_1}$

In addition to the low-energy constraints discussed in the previous sections, the doubly charged Higgsino will give rise to sizable lepton flavor violating decays of the Z boson. In particular, it was found in [24] that the non-observation of such signals may impose constraints on the supersymmetric spectrum. The analysis in [24] took the point of view that the Yuakawa couplings h_{ij} were all diagonal and of order unity. This would allow one to put stringent bounds on the masses of the doubly charged Higgsinos. We will consider here, as in previous sections, that the masses and couplings are free parameters, the couplings nondiagonal, and attempt to bound them from the Z decays. The experimental limits on $Z\rightarrow l_1\overline{l}_2+l_2\overline{l}_1$ are [15]

$$BR(Z\rightarrow e^-\mu^+ + \mu^- e^+)<1.7\times10^{-6} \tag{45}$$

$$BR(Z\rightarrow e^-\tau^+ + \tau^- e^+)<9.8\times10^{-6} \tag{46}$$

$$BR(Z\rightarrow \tau^-\mu^+ + \mu^- \tau^+)<1.2\times10^{-5}. \tag{47}$$

The amplitude of the decay $Z\rightarrow ll'$ proceeding through a doubly charged Higgsino can be parametrized as

$$\mathcal{T}(Z\rightarrow l_1 l_2)=\frac{ig_w h_{l_1 i}h_{il_2}}{2(4\pi)^2 c_w}\mathcal{F}_Z^{l_1 l_2}\epsilon_Z^\mu \overline{u}_{l_2}\gamma_\mu(1-\gamma_5)v_{l_2} \tag{48}$$

where $c_w=\cos\theta_w$. The form factor is induced by the Feynman diagrams of Fig. 1. The branching ratio for this decay mode is

$$BR(Z\rightarrow\overline{l_1}l_2+\overline{l}_1 l_2)=\frac{\alpha_W}{3\cos\theta_W^2}\frac{(h_{l_1 i}h_{il_2})^2}{(4\pi)^3}\frac{M_Z}{\Gamma_Z}|f_Z^{l_1 l_2}|^2 \tag{49}$$

where $\Gamma_Z=2.490$ GeV is the experimental value of the total Z width. The form factor is

$$f_Z^{l_1 l_2}=r\left[\frac{f_1(r)}{2}+\frac{f_2(r)}{4}-\frac{r_Z f_3(r)}{4}+\frac{1}{8}\right] \tag{50}$$

where

$$f_1(r)=\int_0^1\int_0^1\frac{y\,dx\,dy}{(1-y)r+y[1-r_Z xy(1-x)]} \tag{51}$$

$$f_2(r)=\frac{r}{2(1-r)}+\frac{r^2\ln r}{2(1-r)^2} \tag{52}$$

$$f_3(r)=\int_0^1\int_0^1\frac{xy^3(1-x)\,dx\,dy}{(1-y)r+y[1-r_Z xy(1-x)]} \tag{53}$$

with $r=m_{\tilde{\Delta}}^2/m_{\tilde{l}}^2$ and $r_Z=m_Z^2/m_{\tilde{l}}^2$. The bounds obtained from the LFV decays of the Z boson are weaker than those from $\mu\rightarrow e\gamma$, but comparable to some of the three-body decays.

VIII. NUMERICAL ANALYSIS AND DISCUSSION

We proceed now to investigate quantitatively the predictions of doubly charged Higgsino-mediated lepton violating decays. We set as parameters the diagonal and non-diagonal couplings h_{ij} and the mass of the doubly charged Higgsino, $M_{\tilde{\Delta}_{L,R}}$. We assume that the strength of the coupling is the same for the left and right sectors; also we assume for simplicity that the masses of the doubly charged Higgsinos $\tilde{\Delta}_L^{--}$ and $\tilde{\Delta}_R^{--}$ are the same. The first assumption is consistent with left-right symmetry. With regards to the masses, various scenarios are possible. The authors of [12] neglect the left-handed $\tilde{\Delta}_L$ since it is not absolutely necessary for symmetry breaking to the standard model. A different point of view is taken by [25], who find the existence of $\tilde{\Delta}_L$ essential for the study of spontaneous parity and R parity breaking. We shall adopt this latter point of view and take, for simplicity, $M_{\tilde{\Delta}_L}\approx M_{\tilde{\Delta}_R}=M_{\tilde{\Delta}}$. The doubly charged components of the Higgs bosons do not acquire masses of order v_R, but their masses arise through the non-renormalizable operators and are of order v_R^2/M_{Planck}. It is not unreasonable to expect therefore that their masses could be of the same order of magnitude and light, and the same would be the case for their fermionic partners. (If one allows a coefficient of proportionality, rather than approximate equality between the masses, one introduces an extra parameter into the results which will unnecessarily, given the level of precision in the masses, complicate the bounds). In addition to the couplings and masses of the doubly charged Higgsinos, the results will depend on the mass (scale) of the scalar lepton mass $M_{\tilde{l}}$. In fact, we find that the LFVs are sensitive to the ratio $M_{\tilde{\Delta}}^2/M_{\tilde{l}}^2$. The explicit constraints obtained from the LVF processes discussed in previous sections are

$$h_{\mu i}h_{ie}^*<\frac{2.29\times10^{-10}M_{\tilde{\Delta}}^2}{rf_\Delta(r)}\quad \text{from }\ \mu\rightarrow e\gamma \tag{54}$$

$$h_{\tau i}h_{ie}^{*}<\frac{1.09\times10^{-8}M_{\tilde{\Delta}}^{2}}{rf_{\Delta}(r)}\quad \text{from}\quad \tau\rightarrow e\gamma \tag{55}$$

$$h_{\tau i}h_{i\mu}^{*}<\frac{1.15\times10^{-8}M_{\tilde{\Delta}}^{2}}{rf_{\Delta}(r)}\quad \text{from}\quad \tau\rightarrow \mu\gamma \tag{56}$$

where we put

$$f_{\Delta}(r)\equiv f_{M}(r)-2g_{M}(r)=\frac{1}{6(1-r)^{4}}[15r^{2}-12r-3 -6r(r+2)\ln r] \tag{57}$$

and $r=M_{\tilde{\Delta}}^{2}/M_{\tilde{l}}^{2}$. The constraints obtained from μ-e conversion are

$$h_{\mu i}h_{ie}^{*}<\frac{2.85\times10^{-11}M_{\tilde{\Delta}}^{2}}{rF_{\Delta}(r)}\quad \text{for}\quad {}^{27}\text{Al target} \tag{58}$$

$$h_{\mu i}h_{ie}^{*}<\frac{3.68\times10^{-9}M_{\tilde{\Delta}}^{2}}{rF_{\Delta}(r)}\quad \text{for}\quad {}^{48}\text{Ti target}, \tag{59}$$

$$h_{\mu i}h_{ie}^{*}<\frac{2.23\times10^{-8}M_{\tilde{\Delta}}^{2}}{rF_{\Delta}(r)}\quad \text{for}\quad {}^{208}\text{Pb target} \tag{60}$$

where

$$F_{\Delta}(x,r)\equiv f_{E}(r)-2g_{E}(r)+f_{\Delta}(r) =\frac{1}{2(1-r)^{4}}[-8r^{3}-103r^{2}+158r-61 +2(12r^{3}-r^{2}+36r-26)\ln r] \tag{61}$$

with $r=M_{\tilde{\Delta}}^{2}/M_{\tilde{l}}^{2}$. Similarly we get the following bounds from $l\rightarrow l_{1}l_{2}l_{3}$:

$$h_{\mu e}h_{ee}^{*}<\frac{4.7\times10^{-9}M_{\tilde{\Delta}}^{2}}{rf_{\Delta}(r)}\quad \text{from}\quad \mu\rightarrow e^{+}e^{-}e^{-} \tag{62}$$

$$h_{\tau e}h_{ee}^{*}<\frac{4.4\times10^{-6}M_{\tilde{\Delta}}^{2}}{rf_{\Delta}(r)}\quad \text{from}\quad \tau\rightarrow e^{+}e^{-}e^{-} \tag{63}$$

$$h_{\tau\mu}h_{\mu\mu}^{*}<\frac{1.07\times10^{-5}M_{\tilde{\Delta}}^{2}}{rf_{\Delta}(r)}\quad \text{from}\quad \tau\rightarrow \mu^{+}\mu^{-}\mu^{-} \tag{64}$$

$$\Re[h_{\tau e}h_{e\mu}^{*}+h_{\tau\mu}h_{ee}^{*}] <\frac{4.73\times10^{-6}M_{\tilde{\Delta}}^{2}}{rf_{\Delta}(r)}\quad \text{from}\quad \tau\rightarrow e^{+}e^{-}\mu^{-} \tag{65}$$

$$\Re[h_{\tau\mu}h_{\mu e}^{*}+h_{\tau e}h_{\mu\mu}^{*}] <\frac{1.09\times10^{-5}M_{\tilde{\Delta}}^{2}}{rf_{\Delta}(r)}\quad \text{from}\quad \tau^{+}\rightarrow \mu^{+}e^{-}\mu^{-}\quad . \tag{66}$$

Finally the bounds from the Z lepton flavor violating decays are

$$\Re[h_{\mu i}h_{ie}^{*}]<\frac{8.9\times10^{-6}M_{\tilde{\Delta}}^{2}}{f_{Z}^{ll'}}\quad \text{from}\quad Z\rightarrow e^{-}\mu^{+}+e^{+}\mu^{-} \tag{67}$$

$$\Re[h_{\tau i}h_{ie}^{*}]<\frac{2.16\times10^{-5}M_{\tilde{\Delta}}^{2}}{f_{Z}^{ll'}}\quad \text{from}\quad Z\rightarrow e^{-}\tau^{+}+e^{+}\tau^{-} \tag{68}$$

$$\Re[h_{\tau i}h_{i\mu}^{*}]<\frac{2.39\times10^{-5}M_{\tilde{\Delta}}^{2}}{f_{Z}^{ll'}}\quad \text{from}\quad Z\rightarrow \tau^{-}\mu^{+}+\tau^{+}\mu^{-}. \tag{69}$$

Exactly which process dominates will be determined by the ratio $M_{\tilde{\Delta}}^{2}/M_{\tilde{l}}^{2}$. In Table I below we present bounds on coupling constants obtained for some typical mass ratios.

A general feature of these bounds is that the most important ones experimentally come from μ-e conversion, $\mu\rightarrow e\gamma$, $\tau\rightarrow e\gamma$ and $\tau\rightarrow\mu\gamma$. The rest are only of somewhat academic importance, in particular the three body decays, although they bound different specific products. Compared to the $\mu\rightarrow e\gamma$, the branching ratio of $\mu^{+}\rightarrow e^{+}e^{+}e^{-}$ is

$$\frac{\Gamma(\mu^{+}\rightarrow e^{+}e^{+}e^{-})}{\Gamma(\mu\rightarrow e\gamma)}\approx 6\times10^{-3} \tag{70}$$

compared to an improvement of only 10^{-1} in the measured branching ratios. That does not necessarily mean that $\mu^{+}\rightarrow e^{+}e^{+}e^{-}$ is uninteresting experimentally, because the experimental detection and background are very different for the two decays. Also, the reaction $\mu^{+}\rightarrow e^{+}e^{+}e^{-}$ has a much richer structure than the $\mu\rightarrow e\gamma$ decay and can take place in cases in which $\mu\rightarrow e\gamma$ is forbidden. That is, the penguin diagrams with an intermediate off-shell photon, which are dominant in this case, might be forbidden.

For the μ-e conversion the rough estimate ratio

$$\frac{R(\mu^{+}\text{Ti}\rightarrow e^{+}\text{Ti})}{BR(\mu\rightarrow e\gamma)}\approx\frac{\alpha}{3\pi}\frac{E_{e}p_{e}}{m_{\mu}^{2}}\frac{ZF^{2}}{\Gamma_{capt}}\approx 5.6\times10^{-3}, \tag{71}$$

which shows a relative suppression of about two orders of magnitude [26], is not universal and in fact the μ-e conversion, when taking the loop functions correctly into account, dominates over $\mu\rightarrow e\gamma$ throughout the parameter space un-

TABLE I. Constraints on the couplings of the doubly charged Higgsinos from lepton flavor violating decays

Coupling	$M_{\tilde{\Delta}}=100$ GeV $M_{\tilde{l}}=1$ TeV	$M_{\tilde{\Delta}}=100$ GeV $M_{\tilde{l}}=200$ GeV	$M_{\tilde{\Delta}}=100$ GeV $M_{\tilde{l}}=100$ GeV	$M_{\tilde{\Delta}}=200$ GeV $M_{\tilde{l}}=100$ GeV	Process
$h_{\mu i}h^*_{ie}$	$<5.15\times 10^{-4}$	$<4.53\times 10^{-5}$	$<6.18\times 10^{-7}$	$<1.05\times 10^{-4}$	$\mu\rightarrow e\gamma$
$h_{\tau i}h^*_{ie}$	$<2.45\times 10^{-2}$	$<2.15\times 10^{-3}$	$<2.9\times 10^{-5}$	$<4.89\times 10^{-3}$	$\tau\rightarrow e\gamma$
$h_{\tau i}h^*_{i\mu}$	$<2.59\times 10^{-2}$	$<2.3\times 10^{-3}$	$<3.1\times 10^{-5}$	$<5.26\times 10^{-3}$	$\tau\rightarrow \mu\gamma$
$h_{\mu i}h^*_{ie}$	$<3.1\times 10^{-7}$	$<3.8\times 10^{-8}$	$<4\times 10^{-20}$	$<1.92\times 10^{-7}$	$\mu-e$ in Al
$h_{\tau i}h^*_{ie}$	$<4\times 10^{-5}$	$<4.9\times 10^{-6}$	$<5.25\times 10^{-18}$	$<2.47\times 10^{-5}$	$\mu-e$ in Ti
$h_{\mu e}h^*_{ee}$	$<1.06\times 10^{-2}$	$<9.3\times 10^{-4}$	$<1.27\times 10^{-4}$	$<2.15\times 10^{-3}$	$\mu\rightarrow 3e$
$h_{\tau e}h^*_{ee}$	<9.9	$<8.7\times 10^{-1}$	$<1.19\times 10^{-2}$	<2.01	$\tau\rightarrow 3e$
$h_{\tau\mu}h^*_{\mu\mu}$	<24.06	<2.1	$<2.89\times 10^{-1}$	<4.9	$\tau\rightarrow 3\mu$
$\Re[h_{\tau e}h^*_{e\mu}+h_{\tau\mu}h^*_{ee}]$	<10.63	$<9.34\times 10^{-1}$	$<1.28\times 10^{-2}$	<2.16	$\tau\rightarrow \mu ee$
$\Re[h_{\tau\mu}h^*_{\mu e}+h_{\tau e}h^*_{\mu\mu}]$	<24.51	<2.15	$<2.94\times 10^{-2}$	$<4.99\times 10^{-2}$	$\tau\rightarrow e\mu\mu$
$h_{\mu i}h^*_{ie}$	$<1.59\times 10^{-1}$	$<1.93\times 10^{-1}$	$<2.23\times 10^{-1}$	<9.33	$Z\rightarrow \mu^{\pm}e^{\mp}$
$h_{\tau i}h^*_{ie}$	$<3.85\times 10^{-1}$	$<4.68\times 10^{-1}$	$<5.4\times 10^{-1}$	<22.7	$Z\rightarrow \tau^{\pm}e^{\mp}$
$h_{\tau i}h^*_{i\mu}$	$<4.26\times 10^{-1}$	$<5.17\times 10^{-1}$	$<6\times 10^{-1}$	<25.1	$Z\rightarrow \tau^{\pm}\mu^{\mp}$

der consideration. Note that the strength of the bounds on the Yukawa couplings of the doubly charged Higgsinos at $M_{\tilde{\Delta}}=M_{\tilde{l}}$ in μ-e conversion is due to the fact that $F_{\Delta}(x)$ is singular at $x=1$. In the case of $Z\rightarrow l^{\pm}l'^{\mp}$, this graph is determined entirely by the penguin graph with the photon replaced by the Z boson. The bounds are much weaker than those coming from $l\rightarrow l'\gamma$, but the loop function is different. In fact the Z branching ratios are more predictive (less parameter dependent) because of the slow variation of the function $f_Z^{ll'}$ with $M^2_{\tilde{\Delta}}/M^2_{\tilde{l}}$.

Lepton-flavor violating processes have been studied in the context of a (non-supersymmetric) left-right theory. We present in Table II, for comparison, the bounds on the Yukawa couplings h_{ij} of the triplet $\Delta_{L,R}$ bosons.

Some of these bounds are based on old experimental data; however, some processes such as $\mu\rightarrow eee$ and the corresponding three body τ decays with Higgs bosons can occur at the tree level, and they benefit from having light fermions in the loop. In view of this, the bounds obtained from the doubly charged Higgsinos are very good.

If we assume that the off-diagonal couplings h_{ij}, $i\neq j$, are much smaller than the diagonal couplings h_{ii}, we can obtain better bounds on specific products of couplings rather than on the sum of products. This assumption is supported by bounds obtained from lepton flavor conserving processes, such as Δa_{μ} which restricts $h_{\mu\mu}\leqslant 5.9\times 10^{-3}/(f_{\Delta})^{1/2}M_{\tilde{\Delta}}$. If in addition we suppose that $h_{\mu e}h_{ee}\approx h_{\mu\mu}h_{\mu e}\gg h_{\mu\tau}h_{\tau e}$, we can obtain bounds on $h_{ee}h_{e\mu}$ as a function of scalar lepton and doubly charged Higgsino masses.

Figure 3 shows the variation of the couplings $\ln(h_{\mu e}h_{ee})$ as a function of the doubly charged Higgsino mass $M_{\tilde{\Delta}}$ for the heavy squark scenario, $M_{\tilde{l}}=1$ TeV. In this parameter space there is a local minimum of the function $f_{\Delta}(x)$ at $M_{\tilde{l}}=725$

TABLE II. Previous constraints on the couplings of the doubly charged Higgs bosons from lepton flavor violating decays for $M_{\Delta}=100$ GeV [4].

Coupling	Bound	Process
$h_{\mu e}h_{\mu\mu}$	$<2\times 10^{-6}$	$\mu\rightarrow e\gamma$
$h_{\tau e}h_{\tau\mu}$	$<5\times 10^{-6}$	$\mu\rightarrow e\gamma$
$h_{\mu e}h_{ee}$	$<3.2\times 10^{-7}$	$\mu\rightarrow 3e$
$h_{\tau e}h_{\mu e}$	$<5.5\times 10^{-3}$	$\tau\rightarrow 2e\mu$
$h_{\tau e}h_{ee}$	$<4.3\times 10^{-3}$	$\tau\rightarrow 3e$
$h_{\tau e}h_{ee}$	$<5.5\times 10^{-3}$	$\tau\rightarrow e2\mu$

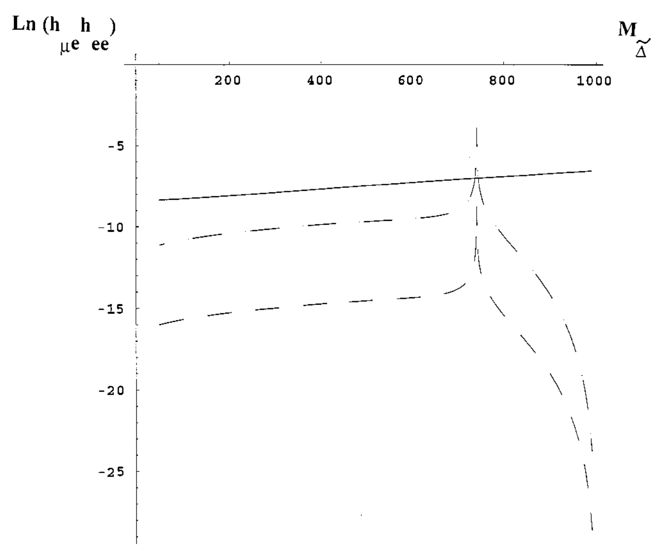

FIG. 3. $\ln(h_{\mu e}h_{ee})$ as a function of $M_{\tilde{\Delta}}$ for $M_{\tilde{l}}=1$ TeV in the case in which the off-diagonal couplings are smaller than the diagonal couplings. The solid curve represents the restriction coming from $\mu\rightarrow e\gamma$, the dashed the restrictions from the expected sensitivity of μ-e conversion in ^{27}Al, and the dot-dashed curve from the present sensitivity of $\mu-e$ conversion in ^{48}Ti.

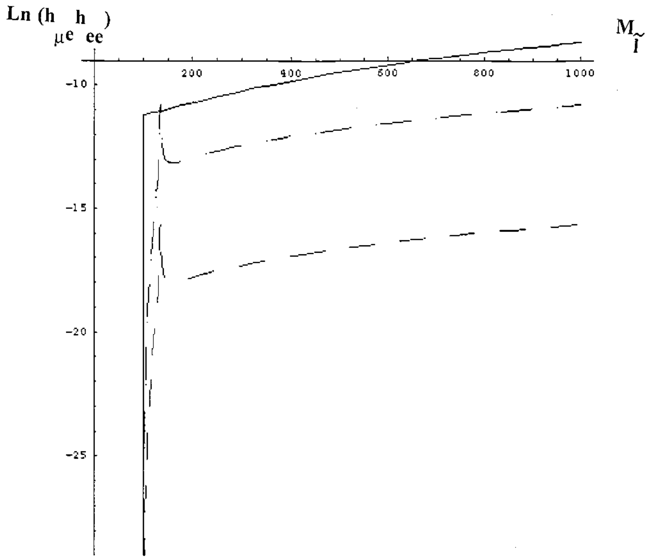

FIG. 4. $\ln(h_{\mu e}h_{ee})$ as a function of $M_{\tilde{l}}$ for $M_{\tilde{\Delta}}=100$ GeV in the case in which the off-diagonal couplings are smaller than the diagonal couplings. The solid curve represents the restriction coming from $\mu\rightarrow e\gamma$, the dashed the restrictions from the expected sensitivity of μ-e conversion in ^{27}Al, and the dot-dashed curve from the present sensitivity of μ-e conversion in ^{48}Ti.

GeV. One could see that the bound from μ-e conversion is dominant over the bound from $\mu\rightarrow e\gamma$ not just for the expected sensitivity reached for ^{27}Al, but even for the present sensitivity in ^{48}Ti experiments. In both these cases the bound from $\mu\rightarrow e^-e^-e^+$ is tighter by at least one order of magnitude.

In Fig. 4 we present the variation of the couplings $\ln(h_{\mu e}h_{ee})$ for a light doubly charged Higgsino, $M_{\tilde{\Delta}}=100$ GeV, as a function of the slepton mass $M_{\tilde{l}}$. The bounds from μ-e conversion dominate the whole parameter space. The threshold effects are seen here as well for $M_{\tilde{\Delta}}\approx M_{\tilde{l}}$.

IX. CONCLUSION

Doubly charged Higgsinos, which are present in supersymmetric theories with exotic Higgs representations, and occur naturally in left-right supersymmetric models, can have lepton flavor violating couplings h_{ij}. Since the supersymmetric left-right theory accommodates naturally neutrino masses and mixings, it is natural to look at the consequences of such mixing phenomena in the charged lepton sector. It is possible, and indeed expected in most versions of the theory, that the doubly charged Higgsinos will be light. Charged lepton flavor violation induced by these Higgsinos will in that case be important and possibly provide a clear signal of exotic particles and physics beyond the standard model. We have studied the bounds on the couplings imposed by a variety of lepton flavor violating decays and found that the most stringent bounds come from μ-e conversion (for all explored values of charged slepton and doubly charged Higgsino mases). Bounds on products of the form $h_{e\mu}h_{\mu e}$ can be restricted to as low as $10^{-7}-10^{-8}$ for $M_{\tilde{\Delta}}\approx 100$ GeV. The bounds obtained are as good and often better than for the lepton flavor violating decays of the corresponding bosons in left-right theories. We might conclude that either these off-diagonal couplings are extremely small, or the doubly charged Higgsinos are heavier than presently believed. Either way, the lepton flavor violating decays are an interesting and very restrictive window into an extended gauge structure.

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