Is the 125 GeV Higgs the superpartner of a neutrino?

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Abstract

Recent LHC searches have provided strong evidence for the Higgs, a boson whose gauge quantum numbers coincide with those of a SM fermion, the neutrino. This raises the mandatory question of whether Higgs and neutrino can be related by supersymmetry. We study this possibility in a model in which an approximate $R$-symmetry acts as a lepton number. We show that Higgs physics resembles that of the SM-Higgs with the exception of a novel invisible decay into Goldstino and neutrino with a branching fraction that can be as large as $\sim 10\%$. Based on naturalness criteria, only stops and sbottoms are required to be lighter than the TeV with a phenomenology dictated by the $R$-symmetry. They have novel decays into quarks+leptons that could be seen at the LHC, allowing to distinguish these scenarios from the ordinary MSSM.
1 Introduction

The LHC has recently reported strong experimental evidence for the existence of the Higgs particle. This is the first discovered boson whose gauge quantum numbers are the same as those of an existing fermion, the neutrino. It is therefore tempting to speculate on the most minimal realization of supersymmetry, needed to protect the Higgs mass, corresponding to a situation where the Higgs and a neutrino (any of the three) belong to the same supermultiplet. In this article we study requirements and implications of this possibility.

If the Higgs is the neutrino superpartner, an approximate $R$-symmetry $U(1)_R$, that acts as a lepton symmetry 1, is necessary [2, 3]. This is needed to provide the neutrino with an approximate conserved lepton number that protects its mass, while leaving its supersymmetric partner, the Higgs, without lepton charge. In this way, the latter can acquire a nonzero vacuum expectation value (VEV) and break all symmetries under which it is charged, without breaking lepton number. As we will show, there are important implications of this approximate $R$-symmetry. Since the gravitino is $R$-charged, the Higgs can decay into a neutrino and a gravitino, with a branching ratio that can be as large as 10%. This gives an invisible decay width to the Higgs that could be indirectly detected by measuring a small reduction of all its visible branching ratios. Gauginos must get Dirac, rather than Majorana, masses and the wino mass must lie above the TeV in order to avoid large corrections to charged leptons couplings [3]. Therefore gauginos are not expected to be detectable during the first years of the LHC running. Another requirement of the model is that, if no extra Higgs superfields are present (and hence no Higgsinos), the up-quark Yukawa couplings must arise from a supersymmetry-breaking term. Interestingly, however, we will show that the soft-mass of the Higgs is insensitive (at the one-loop level) to this supersymmetry-breaking term that can have its origin in physics above the TeV, as we propose in the Appendix.

In a bottom-up approach to supersymmetry based on naturalness criteria, models with the Higgs as a neutrino superpartner have the most minimal low-energy supersymmetric spectrum, since no Higgsinos are present (hence avoiding the infamous $\mu$ problem). Below the TeV, only stops and sbottoms are required, but with a phenomenology very different from that of the Minimal Supersymmetric Standard Model (MSSM). In particular, stops and sbottoms exhibit leptoquark decays: $\tilde{t}_L \rightarrow b\ell^{-}, t\bar{\nu}$, $\tilde{t}_R \rightarrow t\nu$, while $\tilde{b}_L \rightarrow b\nu$. These decay channels can compete with decays into gravitino (a channel that is also present in the MSSM with low-scale supersymmetry breaking), thus allowing to differentiate this model from the MSSM. We will discuss the precise branching ratios, the present bounds and future searches to discriminate between these scenarios. If light enough to be produced at the LHC, we will show that first and second generation squarks could decay dominantly into 3-bodies including quarks, leptons and gauge/Higgs bosons, providing then distinctive novel signatures to be searched at the LHC.

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1The idea of an $R$-symmetry as a lepton symmetry was first proposed in [1]. In this original realization, however, the particle spectrum was not realistic.
\begin{center}
\begin{tabular}{|c|c|c|}
\hline
 & $SU(3)_c \times SU(2)_L \times U(1)_Y$ & $U(1)_R$ \\
\hline
$Q$ & $(3, 2)_{\frac{1}{6}}$ & $1 + B$ \\
$U$ & $(\bar{3}, 1)_{-\frac{2}{3}}$ & $1 - B$ \\
$D$ & $(\bar{3}, 1)_{\frac{1}{3}}$ & $1 - B$ \\
$L_{1,2}$ & $(1, 2)_{-\frac{1}{2}}$ & $1 - L$ \\
$E_{1,2}$ & $(1, 1)_{1}$ & $1 + L$ \\
$H \equiv L_3$ & $(1, 2)_{-\frac{1}{2}}$ & 0 \\
$E_3$ & $(1, 1)_{1}$ & 2 \\
$W_\alpha^a$ & $(8, 1)_0 + (1, 3)_0 + (1, 1)_0$ & 1 \\
$\Phi_a$ & $(8, 1)_0 + (1, 3)_0 + (1, 1)_0$ & 0 \\
$X \equiv \theta^2 F$ & $(1, 1)_0$ & 2 \\
\hline
\end{tabular}
\end{center}

Table 1: Superfield content and charge assignments under the SM gauge group and the $U(1)_R$ symmetry. The value of the $R$-charge ($q_R$) corresponds to the charge of the superfield and the scalar component, while the fermion component has charge $q_R - 1$ and the $F$-term has charge $q_R - 2$. $B$ and $L$ are arbitrary charge assignments.

2 The Higgs as a lepton superpartner

We consider a model that, differently from the MSSM, does not contain the two Higgs superfields $H_u$ and $H_d$. Instead, the SM scalar Higgs doublet is assumed to be one of the three lepton superpartners. The corresponding chiral superfield is denoted by

$$H \equiv L_3 = (H, l_L),$$

where we label by $l_L = (l_L^-, \nu_L)$ one of the three left-handed leptons, either the electron, muon or tau doublet. The other two are embedded in the chiral superfields $L_{1,2} \equiv (\tilde{L}_{1,2}, l_{1,2})$. The full spectrum of the theory is given in table 1. Notice that this theory does not have Higgsinos and is of course anomaly free, since the only extra fermions beyond the SM are all in adjoint representations.

Any theory beyond the SM must preserve an approximate lepton number in order to avoid large neutrino masses. In our model this lepton symmetry cannot commute with supersymmetry, otherwise the Higgs $H$, being in the same supermultiplet as the leptons, would carry lepton number and this would be broken when the Higgs gets a VEV. For this reason lepton number can only be defined as an $R$-symmetry $U(1)_R$ under which $H$ is neutral but $l_L$ is charged. The $R$-charges for this model are given in table 1. Few comments are in order. Since gauginos must carry nonzero $R$-charges, they cannot get Majorana masses. Nevertheless, they can get Dirac-type masses by marrying with additional fermions coming from adjoint chiral superfields $\Phi_a$. Notice also that there is a certain freedom in the symmetry properties of quarks and $l_{1,2}$ leptons, depending on whether or not they transform under the $U(1)_R$ ($B, L \neq 0$). A non-vanishing charge $B \neq 0$ corresponds to a non-vanishing $U(1)_R$ charge for protons and neutrons that can be used to protect proton decay. Indeed, for $B \neq |L|$
the proton decays to neutrinos or positrons are forbidden by the R-symmetry, as well as the decay into (anti)gravitinos (of R-charge \( \mp 1 \)) if \( |B| \neq 1/3 \). Also for \( L \neq 0 \) the R-symmetry protect the masses of all the three neutrinos, and for \( L \neq 1 \) the superpotential terms \( L_i L_j E_k \) and \( Q_i L_j D_k \), which are strongly constrained by lepton-flavor violating processes [4,5], are not allowed.

Working with \( B \neq 1/3 \) and \( L \neq 1 \), the only superpotential terms that can be written in this model at the renormalizable level are, including only matter fields,

\[
W = Y_d H Q D + Y_{e ij} H L_i E_j ,
\]

where indexes \( i, j = 1, 2 \) are summed over and \( Y_d \) is a matrix in flavor space. As it stands, the superpotential Eq. (2) does not generate up-type quark masses, gaugino masses, nor a mass for the \( l_L^- \) lepton (the latter is forbidden since \( SU(2)_L \) indices in Eq. (2) are summed antisymmetrically, meaning that the term \( H H E_3 \) vanishes). These must originate as supersymmetry-breaking effects. We can write these in a supersymmetry preserving notation by means of a spurion field \( X \), whose \( F \)-component is nonzero \( X = \theta^2 F \). To preserve the R-symmetry, \( X \) must have R-charge 2. The masses of the up-type quarks can be written as

\[
\int d^4 \theta \frac{y_u X^\dagger H Q U}{M} = \int d^2 \theta \frac{Y_u H Q U}{\Lambda} ,
\]

where \( y_u \) are dimensionless couplings and \( Y_u = y_u F/(M \Lambda) \) are the Yukawa couplings of the up-type quarks. Notice that we have defined two scales, \( M \) and \( \Lambda \), that could have different origin: \( M \) is the scale at which the supersymmetry-breaking effects are mediated to the SM superpartners, while \( \Lambda \) is the scale at which the higher-dimensional operator Eq. (3) is generated. Explicit examples for the origin of this operator are given in the Appendix. Since, as we will see, the soft masses of SM superpartners are of order \( F/M \), naturalness requires \( F/M \lesssim \text{TeV} \). On the other hand, since the Yukawa coupling of the top is of order one, \( Y_t \sim 1 \), we need \( \Lambda \sim y_u F/M \lesssim 4\pi \text{TeV} \). The mass for the lepton \( l_L^- \) can originate from supersymmetry-breaking terms as well. Indeed, we can have [6]

\[
\int d^4 \theta \frac{y_3 X^\dagger X H D^a H D_a E_3}{M^2} \frac{1}{\Lambda^2} ,
\]

where \( D_a \) is the superspace derivative. This term generates a Yukawa for \( l_L \) equal to \( Y_t = y_3 F^2/(M^2 \Lambda^2) \).

Gauge boson superpartners must also get masses from supersymmetry-breaking terms. Dirac-type gaugino masses can arise from

\[
\int d^2 \theta \frac{D^a X^\dagger W^a}{M} \phi_a ,
\]

that induces gaugino masses of order \( F/M \). There are important constraints on these masses since, after electroweak symmetry breaking (EWSB), charged winos mix with \( l_L^- \) [3] as they have equal R-charges. This mixing affects the coupling of \( Z \) to \( l_L^- \) as

\[
\delta g_{V,A}^l = - \frac{m_W^2}{M_W^2 + 2m_W^2} ,
\]
where $M_{\tilde{W}}$ is the wino mass. Taking the bounds on $\delta g_{V,A}^l$ from [7], we obtain at 99% C.L. the following lower bounds $^2$

$$M_{\tilde{W}} \gtrsim \begin{cases} 
2.5 \text{ TeV} & l_L^c = e_L \\
2 \text{ TeV} & l_L^c = \mu_L \\
1.8 \text{ TeV} & l_L^c = \tau_L 
\end{cases} \quad (7)$$

which can be satisfied for $F/M \sim \text{ TeV}$. The term Eq. (5) does not give mass to the imaginary part of the scalar component in $\Phi_a$, but this can arise from other supersymmetry-breaking terms such as $\int d^4 \theta \, XX^\dagger \Phi_a^2 / M^2$.

Finally, the $R$-symmetry forbids the appearance of supersymmetry-breaking trilinear $A$-terms, implying that the stop one-loop corrections to the Higgs mass are not enough to give $m_h \sim 125$ GeV for stop masses below the TeV, as required by naturalness. New contributions to the $D$-term Higgs quartic are then needed. These can come from supersymmetry-breaking interactions of the type

$$\int d^4 \theta \, \lambda H X^\dagger X \frac{|H|^4}{M^2 \Lambda^2} = \delta \lambda_h h^4 + \ldots, \quad (8)$$

that could either be induced from integrating heavy vector fields (of mass $\Lambda$) that would give extra $D$-terms, or from coupling $H$ directly to the supersymmetry-breaking mediators [9]. In order to obtain $m_h = 125$ GeV, we need $\delta \lambda_h \sim 0.015$.

The $R$-symmetry cannot be an exact symmetry of the model. In order to adjust the cosmological constant to (almost) zero, a gravitino Majorana mass of order

$$m_{3/2} \sim \frac{F}{M_p} \approx 10^{-4} \text{ eV} \left( \frac{\sqrt{F}}{2 \text{ TeV}} \right)^2, \quad (9)$$

is needed. This breaks the $R$-symmetry explicitly and generates neutrino masses of order $m_{3/2}$, which can be in agreement with the experimental limits for $m_{3/2} \lesssim 10$ MeV (or, equivalently, $\sqrt{F} \lesssim 10^7$ GeV) [3, 5, 10]. These upper-bound however can be evaded in theories with emergent global supersymmetry [2, 11, 12] in which the supersymmetric SM (or part of it) arises from a strong sector at high-energies. The $R$-symmetry is an accidental symmetry of these models not broken at order $m_{3/2}$ but by much smaller effects. The gravitino mass can then be much heavier than TeV, and then irrelevant for the phenomenology of the model. Having this in mind, we will consider scenarios in which either a neutrino or the gravitino is the lightest $R$-charged particle.

Summarizing, the Higgs as a lepton superpartner requires, at least, the supersymmetry-breaking operators Eqs. (3), (4), (5) and (8). We will not elaborate here on how these supersymmetry-breaking terms could arise from a specific renormalizable theory, but just postulate that this is the case and study their implications. Nevertheless, we give in the Appendix possible ultraviolet (UV) completions of these Higgsinoless models.

$^2$Charged current universality is also affected [8] but this puts only a mild constraint on the bino mass $M_B \sim 500$ GeV.
2.1 The most natural supersymmetric spectrum

The presence of the operators Eqs. (3), (4), (5) and (8), generates at the loop level other operators. Therefore it is natural in a quantum field theory to include all of them. For example, from loop effects, as depicted in figs. 1-4, the following terms are expected:

\[ \int d^4 \theta \left\{ g_Q \frac{X^\dagger X}{M^2} Q^\dagger Q + g_U \frac{X^\dagger X}{M^2} U^\dagger U + g_H \frac{X^\dagger X}{M^2} H^\dagger H \right\} , \tag{10} \]

and similarly for the leptons $L_{1,2}$. These terms give supersymmetry-breaking (soft) masses for the Higgs $m^2_H = g_H F^2 / M^2$ and squarks $m^2_{Q,U} = g_{Q,U} F^2 / M^2$. It is then crucial to estimate their size, in order to identify the most natural superpartner mass-spectrum of the model. Let us start with the gauge contribution arising from the supersymmetry-breaking term Eq. (5). As it was first noticed in Ref. [13], the gauge loop of fig. 1 gives a finite contribution to the scalar soft masses, as can be seen by simple power counting of this diagram. One obtains [13]

\[ m^2_i = \sum_a C^a_i g^2 a M^2 a \frac{4 \pi^2}{\ln \frac{M^2 a}{M^2 a}} , \tag{11} \]

where for a scalar $i$ in the fundamental representation of $SU(3)_c \times SU(2)_L \times U(1)_Y$ we have $C^a_i = (4/3, 3/4, Y^2_i)$, while $g_a$ are the gauge couplings of group $a$, $M_a$ the gaugino masses and $M_{\Phi_a}$ the supersymmetry-breaking masses of the real part of the scalar component of $\Phi_a$.

On the other hand, squark masses arising from Eq. (3), as illustrated in fig. 2, are quadratically divergent. The contribution to stop soft masses is

\[ m^2_{\tilde{U}} = 2m^2_Q \simeq \frac{Y_t^2}{8 \pi^2} \Lambda^2 , \tag{12} \]
where we have identified the momentum cut-off with $\Lambda \sim \text{TeV}$, the scale at which the operator Eq. (3) is induced. Interestingly, the equivalent one-loop contribution for the Higgs soft-mass, the first diagram of fig. 3, vanishes. This can be understood as follows. If we are interested only in the scalar component of $H$, we can neglect the $\theta$-dependent part of $H$ and write the top Yukawa coupling as $\int d^2\theta \ Y_u H^t QU = Y_u H^t \int d^2\theta QU$ that is supersymmetric and then cannot generate soft-breaking terms. At the two-loop level, however, where the full Higgs superfield $H$ can propagate (see fig. 3), we do expect a nonzero Higgs soft-mass to be induced. Surprisingly, we find that the contribution arising from the second diagram of fig. 3 vanishes, and only the third diagram induces a nonzero $m_H^2$. The latter is proportional to the squark masses, and, as in the MSSM, diverges logarithmically:

$$m_H^2 \simeq -\frac{3Y_t^2}{16\pi^2} \left[ m_Q^2 \ln \frac{\Lambda^2}{m_Q^2} + m_U^2 \ln \frac{\Lambda^2}{m_U^2} \right].$$

(13)

There is also a contribution to the Higgs soft-mass arising from Eq. (8) (see fig. 4) that diverges quadratically:

$$m_H^2 \simeq \frac{3\delta\lambda}{2\pi^2} \Lambda^2.$$

(14)

We can then conclude that the natural values for the stop masses are

$$m_{Q,U}^2 \simeq (400 \text{ GeV})^2 \left[ \left( \frac{M_\tilde{g}}{2 \text{ TeV}} \right)^2 \ln \frac{M_{\tilde{g}}^2}{M_{3/2}^2} + (0.15, 0.3) \left( \frac{\Lambda}{2 \text{ TeV}} \right)^2 \right],$$

(15)

where we have used Eq. (11) and Eq. (12). For the Higgs soft-mass we expect

$$m_H^2 \simeq -(100 \text{ GeV})^2 \left[ 1.9 \left( \frac{m_Q}{300 \text{ GeV}} \right)^2 \ln \frac{\Lambda}{m_Q} - 3.2 \left( \frac{M_{\tilde{W}}}{2 \text{ TeV}} \right)^2 \ln \frac{M_{\tilde{W}}^2}{M_{3/2}^2} - (\delta\lambda_{0.015}) \left( \frac{\Lambda}{2 \text{ TeV}} \right)^2 \right],$$

(16)
where we have used Eq. (11), Eq. (12) and Eq. (14), and taken $m_U \sim m_Q$. This shows that EWSB can occur naturally at $\langle H \rangle = v \simeq 174$ GeV without a major tuning of parameters for $\Lambda$ and gaugino masses around 2 TeV, and stops and left-handed sbottoms around 400 GeV.

The rest of the scalars are expected also to get masses from at least the gaugino loops (fig. 1), although they could also have couplings of order one to $X/M$ such that their masses would then be of order TeV. As it is well known, this does not create naturalness problems [14]. This scenario would really correspond to the most minimal low-energy supersymmetric model with only the stops/sbottoms and (possibly) the gravitino below the TeV scale.

3 Phenomenological Implications

3.1 The 125 GeV Higgs

Differently from the MSSM, this supersymmetric model possesses only one Higgs scalar, identified with a neutrino superpartner, while the charged scalars in the same isospin multiplet are the Nambu-Goldstone bosons responsible for $W^\pm$ and $Z$ masses. At the renormalizable level, the Higgs couplings to the SM fermions and gauge bosons are the same as those of the SM Higgs, and deviations can only arise from loop effects or higher-dimensional operators. Potentially, the most important effects on the Higgs phenomenology come from i) loops mediated by light stops, ii) invisible decay into neutrino $\nu_L$ and gravitino, and iii) Higgs coupling modifications from higher-dimensional operators.

i) The only scalars that can give sizable modifications to the Higgs couplings are the stops. Other scalars, even if light, have a small impact on Higgs physics since their couplings to the Higgs are small. On the contrary, light stops can give sizable loop contributions to the effective Higgs couplings to gluons and photons. The Higgs decay width to photons is corrected as [15]

$$R_{\gamma\gamma} \equiv \frac{\Gamma_{h\gamma\gamma}}{\Gamma_{SM_{h\gamma\gamma}}} \simeq \left| 1 - 0.2 \sum_{L,R} \frac{D_{i_{L,R}} + m_i^2}{m_{i_{L,R}}^2} A_0(\tau_{i_{L,R}}) \right|^2 ,$$

(17)

where $D_{i_L} \equiv (1/2 - 2s_W^2/3)m_Z^2$, $D_{i_R} \equiv (2s_W^2/3)m_Z^2$, $\tau_{i_{L,R}} \equiv m_i^2/(4m_{i_{L,R}}^2)$ and, in the region of parameter space that we consider here, $A_0(\tau) \equiv \tau^{-2}(\arcsin^2 \sqrt{\tau} - \tau)$, which has the limit $A_0(\tau \to 0) = 1/3$. Similarly, the effective coupling to gluons, and hence the production cross-section, is modified as

$$R_{gg} \equiv \frac{\sigma_{hgg}}{\sigma_{SM_{hgg}}} \simeq \left| 1 + 0.7 \sum_{L,R} \frac{D_{i_{L,R}} + m_i^2}{m_{i_{L,R}}^2} A_0(\tau_{i_{L,R}}) \right|^2 ,$$

(18)

where the same formula holds for the decay width $\Gamma_{hgg}$. Notice that the effects of light scalars on $\Gamma_{h\gamma\gamma}$ are generally small as compared with the SM loop contribution (which includes $W^\pm$), while the effects on the production cross-section can be sizable. We show this in fig. 5 by
Figure 5: The ratio $R = \Gamma/\Gamma^{SM}$ for the partial width of $h \rightarrow gg$ and $h \rightarrow \gamma\gamma$ as a function of $m_{\tilde{t}_R}$ while keeping $m_{\tilde{t}_L} = 500\,\text{GeV}$. The dashed part corresponds to a region that is already excluded by direct searches [16, 17] (see later).

plotting the ratio of the width for $h \rightarrow gg, (\gamma\gamma)$ in our model as compared with the SM.

\textbf{ii)} A genuine property of models in which the Higgs and neutrino are superpartners is their interaction with the goldstino, that is fixed by supersymmetry to be

$$\mathcal{L} = \frac{1}{\sqrt{2}F} \bar{\nu}_L h \partial_\mu \partial_\rho \tilde{G} \sigma^\mu \bar{\sigma}^\rho \nu_L + \text{h.c.}$$

(19)

If the gravitino is light this coupling induces invisible Higgs decays into neutrino and gravitino:

$$\Gamma(h \rightarrow \tilde{G}\nu_L) \simeq \frac{1}{16\pi} \frac{m_h^5}{F^2},$$

(20)

where we have neglected the small masses of the final states. For a 125 GeV Higgs, this invisible width equals the decay width into $b\bar{b}$ for $F \simeq (700\,\text{GeV})^2$, while for $F \simeq (1\,\text{TeV})^2$ it induces an invisible fraction of about 10%. Therefore this invisible Higgs decay can be a striking feature of this supersymmetric scenario if supersymmetry is broken at the TeV. Invisible Higgs decays, however, are also present in other well-motivated scenarios such as, for instance, composite Higgs models in which the Higgs can decay to a composite dark matter [19].

\textbf{iii)} Modifications to Higgs couplings can also arise from higher-dimensional operators induced either from integrating out heavy superpartners or from the new physics at the scale $\Lambda$ responsible for Eq. (3) and Eq. (8). Of the first type, only those from integrating the wino can give a tree-level correction to the Higgs coupling to $l_L$, but this is quite small, of order $g^2 v^2 M_W^2 \lesssim 0.01$. Corrections from higher-dimensional operators suppressed by $\Lambda$

\footnote{This possibility has been already considered in ref. [18] in the context of non-linearly realized supersymmetry as arising from specific string constructions.}
can give effects of order $g_M^2 v^2/\Lambda^2$ where $g_M$ generically denotes the coupling of the Higgs to the new sector at $\Lambda$, and therefore can be larger if $g_M > 1$. The only higher-dimensional supersymmetry-preserving operator that can be written is \([20]\)

$$
  g_M^2 \int d^4\theta \frac{(H^\dagger e^V H)^2}{\Lambda^2},
$$

where $V$ denotes the SM vector superfields. This operator however contributes also to the $T$-parameter, which is strongly constrained by precision tests \([21]\), requiring $g_M v/\Lambda \lesssim 10^{-3}$ and therefore small corrections to the Higgs couplings \(^4\). It is important to notice that Eq. \((21)\) can only be generated at tree-level from integrating out heavy singlets, and then it is not generated in models in which only heavy doublets are present. In the Appendix we propose a simple UV completion of our model which involves extra heavy Higgs superfields. In this case, certain corrections to the Higgs couplings can be sizable (see Eq. \((43)\)) without conflicting with the $T$-parameter.

In summary, the most important effects that characterize a Higgs as a superpartner of the neutrino are its invisible decay to neutrino and gravitino and possible modifications of the effective $hgg(\gamma\gamma)$ couplings if the stops are light. It is interesting to compare these predictions with the experimental data recently extracted for the 125 GeV Higgs. This is done in fig. 6 where we show the 68%, 95% and 99% C.L. contours obtained after performing a $\chi^2$ analysis

\(^4\)In strongly-interacting Higgs models \([21]\) a custodial $SU(2)$ symmetry, under which $(H, H^c)$ transforms as a $2$, can be implemented to avoid large corrections to $T$. Nevertheless, this custodial symmetry does not commute with supersymmetry and thus cannot be used here to protect the $T$-parameter.
Goldstino interactions, as in the MSSM, arise from table 2. These decays can arise from the following interactions. From the superpotential term that Lorentz, electromagnetic and $U$ the possible decay modes of the squarks are dictated by symmetries. One can easily see consequence is that $\tilde{t}_L$ and right-handed squarks do not mix and are mass eigenstates. One important arguments suggest to be the lightest.

Noticably, the impact of a light stop is to worsen the Higgs coupling fit. Nevertheless, the presence of a nonzero $BR_{inv}$ tends to improve the fit and for $BR_{inv} \gtrsim 0.2$ the fit can be comparable with the SM even for light stops, as we show in the righthand panel where we plot the preferred regions in the parameter space of our model ($m_{\tilde{L}}, BR_{inv}$). In both plots we have kept $m_{\tilde{L}} = 530$ GeV since, as we will see later, this is the experimental lower-bound. Although present experimental data is not decisive, future data should be able to favor or disfavor this scenario.

### 3.2 Stops and sbottoms

Models in which the Higgs is the neutrino superpartner have a squark phenomenology different from the ordinary MSSM. We focus first on the third generation squarks which naturalness arguments suggest to be the lightest.

Since the $U(1)_R$ symmetry forbids supersymmetry-breaking trilinear $A$-terms, the left-handed and right-handed squarks do not mix and are mass eigenstates. One important consequence is that $\tilde{b}_L$ is always lighter than $\tilde{t}_L$, since their masses are related by

$$m_{\tilde{b}_L}^2 = m_{\tilde{t}_L}^2 - m_t^2 + m_b^2. \quad (22)$$

The possible decay modes of the squarks are dictated by symmetries. One can easily see that Lorentz, electromagnetic and $U(1)_R$ symmetry only allow the decay channels shown in table 2. These decays can arise from the following interactions. From the superpotential term $Y_b HQD$ in Eq. (2), we have contributions to

$$\tilde{t}_L \to b_R \tilde{l}_L, \quad \tilde{b}_L \to b_R \tilde{\nu}_L, \quad \text{and} \quad \tilde{b}_R \to b_L \nu_L, \tilde{t}_L \tilde{\nu}_L. \quad (23)$$

Goldstino interactions, as in the MSSM, arise from

$$\frac{1}{F} \partial_\mu \tilde{t}_L \partial_\nu \tilde{G} \sigma^\mu \sigma^\nu t_L \text{ on-shell} \quad (m_{\tilde{t}_L}^2 - m_{\tilde{b}_L}^2) \frac{1}{F} \tilde{t}_L \tilde{G} t_L, \quad (24)$$

Table 2: Decay modes for the (third family) squarks with the corresponding Lagrangian interaction.
and similarly for other squarks, that leads to

\[ \tilde{t}_R \rightarrow t_R \tilde{G}, \quad \tilde{t}_L \rightarrow t_L \tilde{G}, \quad \tilde{b}_R \rightarrow b_R \tilde{G}, \quad \tilde{b}_L \rightarrow b_L \tilde{G}. \] (25)

Exchanges of heavy winos and binos leads to effective interactions between (s)quarks and leptons, such as

\[ \frac{2g'^2 v}{3M_W^2} \tilde{t}_R \tilde{t}_R \tilde{\nu}_L, \quad \frac{g'^2 v}{2M_W^2} \tilde{t}_L \tilde{b}_L \tilde{\nu}_L. \] (26)

However, due to the Dirac nature of the gauginos, the structure of these interactions is such that the decay amplitudes are proportional to the final-state lepton mass \( \sim \sqrt{m_{\nu,L}/M_W^2} \) and are therefore very small. Such decays, however, could also arise from dimension-six operators that might be induced at the scale \( \Lambda \). For example

\[ \int d^4\theta \frac{1}{\Lambda^2} |H|^2 |Q|^2, \quad \int d^4\theta \frac{1}{\Lambda^2} |H|^2 |U|^2, \] (27)

induce

\[ \tilde{t}_R \rightarrow t_L \nu_L, \quad \tilde{t}_L \rightarrow t_R \tilde{\nu}_L, \] (28)

with an amplitude proportional to the top mass. Note that in these decays, in the limit \( m_t \ll m_{\tilde{t}_L,R} \), the top helicity is fixed: \( U(1)_R \) charge conservation requires, for \( \tilde{t}_R \), a top and a neutrino (rather than an anti-neutrino) in the final state, while spin conservation implies that, in the stop rest frame, the quark helicity be opposite to the neutrino helicity; and vice versa for \( \tilde{t}_L \). This offers an interesting way to differentiate between the squarks decays of Eq. (28) and those of Eq. (25) that are also present in the MSSM with low-scale supersymmetry breaking, since these latter produce final-state tops with opposite helicity.

Let us now discuss the size of the different branching ratios for stops and sbottoms. In fig. 7 we compare the branching ratios of \( \tilde{t}_L \) into different channels for \( \Lambda = \sqrt{F} = 2 \) TeV. We can see that the decays into gravitinos dominate, but the branching ratio into \( b \) and leptons is sizable enough to allow detection. For larger values of \( F \), or in models in which the gravitino...
is heavy (such as models with emergent supersymmetry), \( \tilde{t}_L \) can decay dominantly into \( b + \tilde{l}^- \). Indeed, for \( \sqrt{F} \gg \text{TeV} \), the ratio between the two dominant stop decay widths is given by

\[
\frac{\Gamma(\tilde{t}_L \to b \tilde{l}^-)}{\Gamma(\tilde{t}_L \to t \tilde{\nu})} \simeq \left( \frac{m_b^2}{m_{\tilde{t}_L}^2} \right)^2 \left( 1 - \frac{m_t^2}{m_{\tilde{t}_L}^2} \right)^{-2} \simeq 10 \left( \frac{\Lambda}{2 \text{ TeV}} \right)^4 ,
\]

showing that for \( \Lambda \gtrsim 1 \text{ TeV} \) the decay into \( b + \tilde{l}^- \) dominates.

For \( \tilde{b}_R \), on the other hand, the branching ratios into \( b_L \nu_L \) and \( t_L l^-\) are comparable, as both are controlled by the Yukawa \( Y_b \),

\[
\frac{\Gamma(\tilde{b}_R \to t_L l^-)}{\Gamma(\tilde{b}_R \to b_L \nu_L)} \simeq \left( 1 - \frac{m_t^2}{m_{\tilde{b}_R}^2} \right)^2 .
\]

Nevertheless, for small \( F \) and a light gravitino, the decays into gravitinos dominate:

\[
\frac{\Gamma(\tilde{b}_R \to b \tilde{G})}{\Gamma(\tilde{b}_R \to b_L \nu_L)} \simeq \left( \frac{m_{\tilde{b}_R}^4 \Lambda^4}{F^2 m_b^2} v^2 \right)^2 \simeq 7 \left( \frac{m_{\tilde{b}_R}}{500 \text{ GeV}} \right)^4 \left( \frac{2 \text{ TeV}}{\sqrt{F}} \right)^4 .
\]

This same expression Eq. (31) holds for \( \tilde{t}_R \), for which the decay into charged leptons is forbidden by symmetries. Finally, for \( \tilde{t}_R \) we find

\[
\frac{\Gamma(\tilde{t}_R \to t_R \tilde{G})}{\Gamma(\tilde{t}_R \to t_L \nu_L)} \simeq \left( \frac{m_{\tilde{t}_R}^4 \Lambda^4}{v^2 m_t^2 F^2} \right)^2 \simeq 70 \left( \frac{m_{\tilde{t}_R}}{500 \text{ GeV}} \right)^4 \left( \frac{\Lambda^2}{F} \right)^2 .
\]

**Searches**

As discussed above, many decay processes have neutrinos or gravitinos in the final state, resulting in signatures with missing energy, which resemble much those of the MSSM. For this reason we can adapt present LHC searches to our model. This is particularly true for \( \tilde{b}_L \) whose decay final state is always a bottom-quark plus missing energy. This has the same signature as the MSSM decay into bottom plus neutralino, in the limit where the neutralino is massless, and is presently searched for at the LHC [23,24]. Present exclusion bounds amount to

\[
m_{\tilde{b}_L} > 500 \text{ GeV} .
\]

From Eq. (22), bounds on the sbottom mass imply a bound on the \( \tilde{t}_L \) mass:

\[
m_{\tilde{t}_L} \gtrsim 530 \text{ GeV} .
\]

Similarly, searches for \( t \) and missing energy, motivated by the MSSM decay pattern \( \tilde{t} \to t \chi_0 \) with a massless neutralino, also cover \( \tilde{t}_R \) decays in our model. The mass range

\[
220 \text{ GeV} \lesssim m_{\tilde{t}_R} \lesssim 465 \text{ GeV} ,
\]
is already excluded by a combination of searches [17,25,26]. Searches for stops lighter than the top (or almost degenerate) are reputedly very hard [27–29] and, to our knowledge, the best bound that can be extrapolated gives [27,28]

\[ m_{\tilde{t}_R} \gtrsim 150 \text{ GeV}, \]  

in the low-mass range. As commented above, it can be possible to distinguish between \( \tilde{t}_R \) decays into gravitinos or into neutrinos by measuring the helicity of the final state tops. This is feasible if \( m_{\tilde{t}_R} \gg m_t \). Indeed, in this case the final-state tops are boosted (so boosted that helicity almost coincides with chirality), and they decay before hadronization so that the distribution of its decay products can be measured and the helicity extracted [30]. Another interesting feature that singles out this model is that, for \( m_{\tilde{t}_R} \lesssim m_t \) (a region not yet excluded by direct searches, cf. Eq. (36)), the distribution in momentum of the decay products of the top quark in the decay \( \tilde{t}_R \to t\tilde{G} \) is different from the distribution in the \( \tilde{t}_R \to t\nu_L \) decay, due to the derivatives in the gravitino interaction Eq. (24); this is illustrated in fig. 8.

![Invariant mass distribution of the W and b in the decay \( \tilde{t}_R \to W b\tilde{G} \) (dashed curves) and \( \tilde{t}_R \to W b\nu_L \) (solid curves) for \( m_{\tilde{t}_R} = 160 \) (172) GeV in black (red), taking \( m_t = 173 \) GeV.](image)

Figure 8: Invariant mass distribution of the W and b in the decay \( \tilde{t}_R \to W b\tilde{G} \) (dashed curves) and \( \tilde{t}_R \to W b\nu_L \) (solid curves) for \( m_{\tilde{t}_R} = 160 \) (172) GeV in black (red), taking \( m_t = 173 \) GeV.

The supersymmetry searches described above also cover \( \tilde{t}_L \) and \( \tilde{b}_R \) when their dominant decays are into quarks plus neutrinos or gravitinos. On the contrary, for \( \sqrt{F} \gtrsim \text{few TeV} \) or if the gravitino is heavy, \( \tilde{t}_L \) and \( \tilde{b}_R \) have sizable decay widths into bottom/top quarks and charged leptons, see Eqs. (29), (30) and fig. 7. In this case, also LHC searches for leptoquarks apply to our model and, depending on the flavor of \( l_L \), the present bounds for \( \tilde{t}_L \) are  

\[ m_{\tilde{t}_L} > \begin{cases} 
660 \text{ GeV} & l_L^{-} = e_L^{-} \quad [31] \\
685 \text{ GeV} & l_L^{-} = \mu_L^{-} \quad [32] \\
525 \text{ GeV} & l_L^{-} = \tau_L^{-} \quad [33] 
\end{cases}, \]  

\(^5\)Dedicated searches for leptoquarks \( \to \text{b-jets}+\mu/e \) with b-tagging could improve the sensitivity to our model for \( l_L^{-} = e_L^{-}, \mu_L^{-} \).
while for $\tilde{b}_R$, which decays with 50% probability into $b\nu_L$ and $u'_L$ (see Eq. (30)), the best bounds come from searches on the decay product $b\nu_L$ ($b$-quarks plus missing energy [23, 24]) that lead to $m_{\tilde{b}_R} > 500$ GeV.

Due to the lack of ordinary MSSM $R$-parity, $\tilde{t}_L$ and $\tilde{b}_R$ squarks can be singly produced in this model. Nevertheless, the production cross-sections for these processes are proportional to $Y_b^2$ and are then very small ($\ll$ fb at 14 TeV). Furthermore, their topology (with a final state including $t + l^- + b$-jet and missing energy) coincides with that of a double produced squarks when the two squarks decay differently. This leaves little hope to single out this feature in the early phases of LHC.

We conclude with a possible strategy to differentiate between third family squarks of ordinary supersymmetric models and of models where the Higgs is a neutrino superpartner. If a scalar resonance decaying into $b$-jets and missing energy is observed, it can be our $\tilde{b}_R$ only if also leptoquark decays are observed at the same mass. If no leptoquark decays are observed then it could still be our $\tilde{b}_L$, but from Eq. (22), this would imply that another scalar resonance, the $\tilde{t}_L$, must be observed at slightly heavier mass. On the other hand, the observation of a scalar decaying into $t+bZ_T$ could be attributed to our $\tilde{t}_L$ if also decays into $b + \tilde{t}_L$ are seen. If not, there is still the possibility to be our $\tilde{t}_R$. To know whether this is the case, we must discriminate between the decay $\tilde{t}_R \rightarrow t_L\nu_L$, typical of our model, and the decay $\tilde{t}_R \rightarrow t_RG$ common to many supersymmetric models. For $m_{\tilde{t}_R} \gg m_t$, this can be done by measuring the final-state top-quark helicity, while for $m_{\tilde{t}_R} \ll m_t$, we must look at the differences in the $Wb$ invariant mass distribution, fig. 8.

### 3.3 First and Second Generation Squarks and Sleptons

If the gravitino is light and $\sqrt{F} \sim$ TeV, then the first and second generation squarks, similarly to the third generation ones, decay mainly into gravitinos and light quarks. Searches for jets plus missing energy address these decays [34] and the present bound is

$$ m > 760 \text{ GeV} . $$

On the other hand, in models where the gravitino is heavy (or $\sqrt{F} \gg$ TeV), the situation is quite different from the third-generation squark phenomenology discussed above. The reason is that the 2-body decay into light quarks and leptons are proportional to the Yukawa couplings of the first and second generation squarks that are very small. In particular, we have

$$ \Gamma(\tilde{u}_L(\tilde{d}_L) \rightarrow d + \tilde{l}_L(\tilde{\nu}_L)) \simeq Y_u^2 \frac{m_{\tilde{u}_L(\tilde{d}_L)}}{16\pi} , \quad \Gamma(\tilde{d}_R \rightarrow u + l^- + d + \nu_L) \simeq Y_d^2 \frac{m_{\tilde{d}_R}}{16\pi} , \quad \Gamma(\tilde{u}_R \rightarrow u + \nu_L) \simeq Y_u^2 \frac{m_{\tilde{u}_R}v^4}{16\pi\Lambda^4} , $$

and similarly for the second-generation squarks. Therefore, 3-body decays can be important or even dominate since they are not chirality-suppressed. For example, the Dirac-gaugino-mediated decays into a 3-body final state made of a quark, a lepton and a gauge/Higgs boson,
have a partial width given by (neglecting final-state masses)

\[
\Gamma(q_{iL} \to q_j + \bar{l}_L/\bar{\nu}_L + h/Z/W) \simeq \frac{c_{h,Z,W}}{12288\pi^3} \left[ g^4 Y_{\tilde{q}_L}^2 \frac{m_{\tilde{q}_L}^5}{4 M_B^4} + g^4 c_{ij} \frac{m_{\tilde{q}_L}^5}{M_W^4} \right], \tag{41}
\]

\[
\Gamma(\tilde{q}_R \to q + l_L/\nu_L + h/Z/W) \simeq Y_{\tilde{q}_R}^2 g^4 \frac{c_{h,Z,W}}{49152\pi^3} \frac{m_{\tilde{q}_R}^5}{M_B^4}, \tag{42}
\]

with \( \tilde{q}_{iL} = \tilde{u}_L, \tilde{d}_L, \tilde{q}_R = \tilde{u}_R, \tilde{d}_R, q_j = u, d \) and \( c_{ij} = (2 - \delta_{ij})^2 \). We also have \( c_W = 2, c_{h,Z} = 1 \), while \( Y_{\tilde{q}} \) is the squark hypercharge. The same formula holds for the second generation. In table 3 we provide the dominant decay mode for each of the first and second generation squarks. We must notice, however, that decays into other squark/quark pairs, if kinematically allowed, could dominate over the decays of table 3 (beside being enhanced by a color factor, these channels receive contributions from gluino-exchanges, which are proportional to the strong coupling).

Finally, let us briefly discuss the phenomenology of the sleptons of \( L_{1,2} \) that, we recall, contain the other two non-Higgs-superpartner leptons, and those of \( E_{1,2,3} \). If the gravitino is light and \( \sqrt{F} \sim \text{TeV} \), the corresponding charged sleptons decay into charged leptons and gravitinos, giving missing energy (this topology is searched at the LHC in the context of MSSM decays of sleptons into leptons and (massless) neutralinos [39], excluding the region \( m \lesssim 200 \text{GeV} \)), while sneutrinos decay invisibly into neutrinos and gravitinos and can be searched for using similar strategies as for generic DM searches (monojets or dijets and missing energy). On the other hand, if the gravitino is heavy, the analogous of Eq. (41) applies and 3-body decays can dominate. In table 4 we show the dominant decay mode of the sleptons depending on their corresponding flavour.

We see that the phenomenology of squarks and sleptons is very rich in this model and requires a dedicated study which we plan to pursue in the forthcoming future.
4 Conclusions

An important question, stemming from the recent LHC discovery of a resonance at 125 GeV, is whether or not this could be the scalar superpartner of an existing fermion, hence providing the first evidence for supersymmetry. Since its quantum numbers coincide with those of a neutrino, we have therefore proposed a supersymmetric model in which the Higgs is identified with one of the neutrino superpartner. This can be realized if lepton number is also an $R$-symmetry such that this is not broken by the Higgs VEV.

We have shown that the phenomenology of this model is quite different from that of the MSSM. In the Higgs sector, a sizable ($\sim 10\%$) invisible branching ratio for Higgs decays into neutrinos and gravitinos is possible, together with small deviations in the Higgs couplings to gluons and photons, due to loop effects if the stop $\tilde{t}_R$ is light. These effects are not yet favored nor disfavored by the present LHC Higgs data, but could be seen in the near future by measuring a reduction of the visible Higgs BRs. Higgsinos are absent in this model, and gauginos must get Dirac masses above the TeV. Only third-generation squarks are required, by naturalness, to be below the TeV. We have shown that the $R$-symmetry implies that squarks decay mainly into quarks and either leptons or gravitinos. Therefore, evidence for models with the Higgs as a neutrino superpartner can be sought through the ongoing searches for events with third-generation quarks and missing energy (tailored for the MSSM with a massless neutralino) or through leptoquark searches for final states with heavy quarks and leptons. In the stop decays into tops and neutrinos, the determination of the top helicity will be crucial to unravel these scenarios.

Finally, if first and second generation squarks or sleptons are light enough, they can leave, via 3-body decays, interesting signatures at the LHC that deserve further study.

Note added: While this work was being finalized, Ref. [35] appeared where some of the squarks phenomenology of these models is also discussed.

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Appendix:
Possible UV completions of Higgsinoless models

Here we want to briefly discuss two possible UV completion of the model proposed in this article. The first possibility corresponds to the $R$-symmetric MSSM of ref. [3]. This model
contains two extra Higgs superfields w.r.t. our model, $H_u$ and $R_d$, with a supersymmetric mass given by $\int d^2\theta \mu H_u R_d$. As in the MSSM the superfield $H_u$ can have Yukawa terms with the up-quark sector, $\int d^2\theta y_u H_u QU$, and mix with $H$ (called $L_a$ in [3]) via a bilinear $(B_u)$ soft-term, that we write as $\int d^2\theta R_d H^\dagger X^\dagger /M$. For $\mu \gg v$, we can integrate out $H_u$ and $R_d$, generating the coupling Eq. (3), with the identification $\mu = \Lambda$, or equivalently $Y_u = y_u F/(\mu M)$. Unfortunately, in this limit also a soft-term for $H$ is generated at tree-level, $m_H \simeq F/M$ that, for $Y_u \sim 1$, implies $y_u \sim \mu/m_H$. Consequently, $\mu > m_H$ leads to $y_u > 1$, thus possibly leading to strong dynamics slightly above the TeV. To extrapolate to higher energies we could assume, along the lines of [36], that the Higgs or top are composite states of a strong group and use Seiberg dualities. The procedure of integrating out $H_u$ also generates corrections to the Higgs couplings that we can explicitly calculate. At $O(M^2_u/\mu^2)$, we find that only the Higgs coupling to the top is modified:

$$\frac{g_{h\alpha u}^{SM}}{g_{h\alpha t}} \simeq 1 + 2 \frac{m_h^2}{\mu^2}. \quad (43)$$

Another possible UV completion of our model corresponds to a situation in which either the left-handed or right-handed top is partly arising from a vector superfield. For example, we can have a massive vector superfield $V_{\pm}$ transforming under the SM as a $(3, 2)$ + $(\bar{3}, 2)$, and with the following couplings: $\int d^4\theta[M_V^2 V_+ V_- + g_V V_+ X^\dagger Q + g_V V_+ H^\dagger U]$. Integrating out $V_{\pm}$ gives Eq. (3) with $y_u \sim g_V^2$ and $\Lambda \sim M \sim M_V$. A soft-mass for $Q$ of order $m_Q \sim g_V F/M_V \sim Y_u M_V/g_V$ is also generated at tree-level and requires $g_V > 1$ if we want $m_Q < M_V \sim$ TeV. Theories of massive gauge bosons, however, need to be UV completed at energies $\sim 4\pi M_V/g_V$ either by incorporating them into a new strong sector or by a Higgs mechanism. In the second case, the vector $V_{\pm}$ must be promoted into gauge bosons. A possibility discussed in [37] is to have $V_{\pm}$ arising from an SU(5) gauge model. Notice that, as proposed in [37], we could take the limit in which the squarks are heavier than $M_V$ and have the vector component of $V_{\pm}$ to be the main superpartner of the $t_L$ that would be in this case mainly a gaugino. Other options are given in [38].

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