The Higgs sector of the minimal SUSY $B-L$ model

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I review the Higgs sector of the $U(1)_{B-L}$ extension of the minimal supersymmetric standard model (MSSM). I will show that the gauge kinetic mixing plays a crucial role in the Higgs phenomenology. Two light bosons are present, a MSSM-like one and a $B-L$-like one, that mix at one loop solely due to the gauge mixing. After briefly looking at constraints from flavour observables, new decay channels involving right-handed (s)neutrinos are presented. Finally, even though recent data disfavour it, the enhancement of the Higgs-to-diphoton coupling will be described to show model features beyond the MSSM.

INTRODUCTION

The recently discovered Higgs boson is considered as the last missing piece of the standard model (SM) of particle physics. Nonetheless, several firm observations unanimously call for its extension. Mainly but not limited to, the presence of dark matter, the neutrino masses and mixing patterns, the stability of the SM vacuum, the hierarchy problem. Supersymmetry (SUSY) has long been considered as the most appealing framework to extend the SM. Its minimal realisations (MSSM and its constrained versions) start however to feel considerable pressure to accommodate the recent findings, especially the measured Higgs mass of 125 GeV. Despite not in open contrast with the MSSM, the degree of fine tuning required to achieve it is more and more felt as unnatural. In order to alleviate this tension, non-minimal SUSY realisations can be considered. One can either extend the MSSM by the inclusion of extra singlets (e.g. NMSSM) by extending its gauge group. Concerning the latter, one of the simplest possibilities is to add an additional Abelian gauge group. I will focus here on the presence of an $U(1)_{B-L}$ group which can be a result of an $E_8 \times E_8$ heterotic string theory (and hence M-theory) \cite{3,4}. This model, the minimal $R$-parity-conserving $B-L$ supersymmetric standard model (BLSSM) in short, was proposed in \cite{5,6} and neutrino masses are obtained via a type I seesaw mechanism. Furthermore, it could help to understand the origin of $R$-parity and its possible spontaneous violation in supersymmetric models \cite{7,8} as well as the mechanism of leptogenesis \cite{9,10}.

It was early pointed out that the presence of two Abelian gauge groups in this model gives rise to kinetic mixing terms of the form

$$-\chi_{ab} F^a_{\mu\nu} \hat{F}^b_{\mu\nu}, \quad a \neq b$$

that are allowed by gauge and Lorentz invariance \cite{11}, as $F^a_{\mu\nu}$ and $\hat{F}^b_{\mu\nu}$ are gauge-invariant quantities by themselves, see e.g. \cite{12}. Even if these terms are absent at tree level at a particular scale, they will in general be generated by RGE effects \cite{13,14}. These terms can have a sizable effect on the mass spectrum of this model, as studied in detail in Ref. \cite{15}, and on the dark matter, where several scenarios would not work if it is neglected, as thoroughly investigated in Ref. \cite{16}. In this work, I will review the properties of the Higgs sector of the model. Two light states exist, a MSSM-like boson and a $B-L$-like boson. After reviewing the model, I will show that a large portion of parameter space exists where the SM-like Higgs boson has a mass compatible with its measure, both in a “normal” ($M_{H_2} > M_{H_1} = 125$ GeV) and in an “inverted” hierarchy ($M_{H_2} < M_{H_1} = 125$ GeV), also in agreement with bounds from low energy observables and dark matter relic abundance. The phenomenological properties of the two lightest Higgs bosons will be systematically investigated, where once again the gauge mixing is shown to be fundamental. Finally, even though recent data disfavour it \cite{17}, the enhancement of the Higgs-to-diphoton coupling will be described to show model features beyond the MSSM.

THE MODEL

For a detailed discussion of the masses of all particles as well as of the corresponding one-loop corrections we refer to \cite{15}. Attention will be payed on the main aspects of the $U(1)_{B-L}$ kinetic mixing since it has important consequence for the scalar sector. For the numerical investigations that will be shown, we used the SPheno version \cite{18,19} created with SARAH \cite{20,21} for the BLSSM. For the standardised model definitions, see Ref. \cite{22}, while for a review of the model implementation in SARAH, see Ref. \cite{23}. This spectrum calculator performs a two-loop RGE evaluation and calculates the mass spectrum at one loop. In addition, it calculates the decay widths and branching ratios (BRs) of all SUSY and Higgs particles as well as low-energy observables like $(g-2)_\mu$. We will discuss the most constrainned scenario with a universal scalar mass $m_0$, a universal gaugino mass $M_{1/2}$ and trilinear
The superpotential is given by
\[ W = \sum Y_{ij} \bar{u}_i^c \hat{Q}_j H_u - Y_{ij} \bar{d}_i \hat{Q}_j \hat{H}_d - Y_{ij} \bar{c}_i \hat{L}_j \hat{H}_d + \mu \hat{H}_u \hat{H}_d + Y_{ij} \bar{\nu}_i^c \hat{L}_j \hat{H}_u - \mu' \bar{\eta} \bar{\eta} + Y_{ij} \bar{\nu}_i^c \bar{\eta} \bar{\nu}_j^c \] (2)

and we have the additional soft SUSY-breaking terms:
\[
\mathcal{L}_{SB} = \mathcal{L}_{MSSM} - \lambda_i \lambda_i^\dagger M B B^\dagger - \frac{1}{2} \lambda_i \lambda_i^\dagger M B' B'^\dagger - m_{ij}(\eta_i^2 - m_{ij}^2) \bar{\nu}_i^c \bar{\nu}_j^c - m_{ij}^2 \bar{\nu}_i^c \bar{\nu}_j^c - \eta \bar{\nu}_i^c \bar{\nu}_j^c + \bar{\nu}_i^c \bar{\nu}_j^c \bar{\nu}_k^c + T_{ij} \eta \bar{\nu}_i^c \bar{\nu}_j^c
\] (3)
i, j are generation indices. Without loss of generality one can take \( B_\mu \) and \( B_{\mu'} \) to be real. The extended gauge group breaks to \( SU(3)_C \otimes U(1)_em \) as the Higgs fields and bileptons receive vacuum expectation values (vevs):

\[
H'_d = \frac{1}{\sqrt{2}} (\sigma_d + v_d + i\phi_d), \quad H'^0_u = \frac{1}{\sqrt{2}} (\sigma_u + v_u + i\phi_u)
\]

\[
\eta = \frac{1}{\sqrt{2}} (\sigma_\eta + v_\eta + i\phi_\eta), \quad \bar{\eta} = \frac{1}{\sqrt{2}} (\sigma_\eta + v_\eta + i\phi_\eta)
\]

We define \( \tan \beta = v_\eta / v_u \) in analogy to the ratio of the MSSM vevs (\( \tan \beta = v_\eta / v_d \)).

### Gauge kinetic mixing

As already mentioned in the introduction, the presence of two Abelian gauge groups in combination with the given particle content gives rise to a new effect absent in the MSSM or other SUSY models with just one Abelian gauge group: gauge kinetic mixing. This can be seen most easily by inspecting the matrix of the anomalous dimension, which at one loop is given by

\[
\gamma_{ab} = \frac{1}{16\pi^2} \text{Tr} Q_a Q_b, \tag{4}
\]

where the indices \( a \) and \( b \) run over all \( U(1) \) groups and the trace runs over all fields charged under the corresponding \( U(1) \) group.

For our model we obtain

\[
\gamma = \frac{1}{16\pi^2} N \begin{pmatrix} 11 & 4 & 6 \\ 4 & 6 & 0 \\ 6 & 0 & 9 \end{pmatrix} N. \tag{5}
\]

and we see that there are sizable off-diagonal elements. \( N \) contains the GUT normalisation of the two Abelian gauge groups. We will take as in Ref. \[7\] \( \sqrt{3/5} \) for \( U(1)_Y \) and \( \sqrt{3/2} \) for \( U(1)_{B-L} \), i.e. \( N = \text{diag}(\sqrt{3/5}, \sqrt{3/2}) \).

Hence, we obtain finally

\[
\gamma = \frac{1}{16\pi^2} \left( \begin{smallmatrix} 4 & 3 & 6 \sqrt{3/5} \\ 3 & 6 & 6 \sqrt{3/2} \\ 6 \sqrt{3/5} & 6 \sqrt{3/2} & 9 \end{smallmatrix} \right). \tag{6}
\]

Therefore, even if at the GUT scale the \( U(1)_Y \) kinetic mixing terms are zero, they are induced via RGE evaluation at lower scales. In practice it turns out that it is easier to work with non-canonical covariant derivatives instead of off-diagonal field-strength tensors such as in eq. (11). However, both approaches are equivalent \[33\]. Hence in the following, we consider covariant derivatives of the form

\[
D_\mu = \partial_\mu - iQ_\phi^T G A \tag{7}
\]

where \( Q_\phi \) is a vector containing the charges of the field \( \phi \) with respect to the two Abelian gauge groups, \( G \) is the gauge coupling matrix

\[
G = \begin{pmatrix} g_{YY} & g_{YB} & g_{BB} \\ g_{YB} & g_{BB} \end{pmatrix} \tag{8}
\]
and $A$ contains the gauge bosons $A = (A^Y_\mu, A^B_\mu)$.

As long as the two Abelian gauge groups are unbroken, we have still the freedom to perform a change of basis: $A = (A^Y_\mu, A^B_\mu) \rightarrow A' = ((A^Y_\mu)' , (A^B_\mu)') = RA$ where $R$ is an orthogonal matrix. This freedom can be used to choose a basis such that electroweak precision data can be accommodated in an easy way. A convenient choice is the basis where $g_{BY} = 0$ as in this basis only the Higgs doublets contribute to the entries in the gauge boson mass matrix of the $U(1)_Y \otimes SU(2)_L$ sector and the impact of $\eta$ and $\bar{\eta}$ is only in the off-diagonal elements. Therefore we choose the following basis at the electroweak scale [34]:

$$
\begin{align*}
  g'_{Y Y} &= \frac{g_{Y Y} g_{B Y} - g_{Y B} g_{B Y}}{\sqrt{g_{B B} + g_{B Y}^2}} = g_1 \\
  g'_{B B} &= \sqrt{g_{B B}^2 + g_{B Y}^2} = g_{B L} \\
  g'_{Y B} &= \frac{g_{Y B} g_{B Y} - g_{Y Y} g_{B Y}}{\sqrt{g_{B B} + g_{B Y}^2}} = \tilde{g} \\
  g'_{B Y} &= 0
\end{align*}
$$

When unification at some large scale ($\sim 2 \cdot 10^{16}$ GeV) is imposed, i.e., $g_{Y B}^{GUT} = g_{Y Y}^{GUT} = g_{B L}^{GUT}$ and $g'_{B Y}^{GUT} = 0$, at SUSY scale we get

$$
\begin{align*}
  g_{BL} &= 0.548, \\
  \tilde{g} &\simeq -0.147.
\end{align*}
$$

**Tadpole equations**

The minimisation of the scalar potential is here described in the so-called tadpole method. We can solve the tree-level tadpole equations arising from the minimum conditions of the vacuum with respect to $\mu, B_\mu, \mu'$ and $B_{\mu'}$. Using $v_x^2 = v_\eta^2 + v_\bar{\eta}^2$ and $v^2 = v_d^2 + v_u^2$ we obtain

$$
\begin{align*}
  |\mu|^2 &= \frac{1}{8} \left( 2 g_{B L} v_\eta^2 \cos(2\beta') - 4 m_\eta^2 + 4 m_0^2 \right) \sec(2\beta) - 4 \left( m_\eta^2 + m_0^2 \right) - \left( g_1^2 + \tilde{g}^2 + g_2^2 \right) v^2 \\
  B_{\mu} &= -\frac{1}{8} \left( -2 g_{B L} v_\eta^2 \cos(2\beta') + 4 m_\eta^2 - 4 m_0^2 \right) + \left( g_1^2 + \tilde{g}^2 + g_2^2 \right) v^2 \cos(2\beta) \tan(2\beta) \\
  |\mu'|^2 &= \frac{1}{8} \left( -2 g_{B L} v_\eta^2 \cos(2\beta') + 2 m_0^2 - 2 m_\eta^2 + g_{B L} v^2 \cos(2\beta') \right) \sec(2\beta') \\
  B_{\mu'} &= \frac{1}{8} \left( -2 g_{B L} v_\eta^2 \cos(2\beta') + 2 m_0^2 - 2 m_\eta^2 + g_{B L} v^2 \cos(2\beta') \right) \tan(2\beta')
\end{align*}
$$

For the numerical results, the one-loop corrected equations are used, which lead to a shift of the solutions in $M_{Z'} \simeq g_{B L} x$ and, thus, we find an approximate relation between $M_{Z'}$ and $\mu'$

$$
M_{Z'}^2 \simeq -2|\mu'|^2 + \frac{4(m_0^2 - m_\eta^2 \tan^2 \beta') - v^2 g_{B L} \cos \beta (1 + \tan^2 \beta')}{2(\tan^2 \beta' - 1)}
$$

For the numerical results, the one-loop corrected equations are used, which lead to a shift of the solutions in $m_{h, T}^2$.

$$
m_{h, T}^2 = \begin{pmatrix}
  -m_0^2 s_\beta^2 + g_2^2 v_u^2 & -m_0^2 c_\beta s_\beta - g_2^2 v_d \eta \\
  -m_0^2 c_\beta s_\beta - g_2^2 v_u \eta & m_0^2 s_\beta^2 + g_2^2 v_d^2 \\
  -g_{B L}^2 v_u \eta & -g_{B L}^2 v_d \eta \\
  -g_{B L}^2 v_d \eta & g_{B L}^2 v_u \eta
\end{pmatrix}
$$

where we have defined $g_2^2 = \frac{1}{4} (g_1^2 + \tilde{g}^2 + g_2^2)$, $c_x = \cos(x)$, and $s_x = \sin(x)$ ($x = \beta, \beta'$), and used the masses of the physical pseudoscalars $A^0$ and $A_{\eta}$ given by

$$
\begin{align*}
  m_{A^0}^2 &= \frac{2B_{\mu}}{\sin 2\beta}, \\
  m_{A_{\eta}}^2 &= \frac{2B_{\mu'}}{\sin 2\beta'}.
\end{align*}
$$
For completeness we note that the mass of charged Higgs boson reads as in the MSSM as
\[ m_{H^+_R}^2 = B_\mu (\tan \beta + \cot \beta) + m_W^2. \] (22)

In this model, the CP-odd and charged Higgses are typically very heavy. In eq. (16) we see that compared to the MSSM, there is a non-negligible contribution from the gauge kinetic mixing. LHC searches limit tan \( \beta' < 1.5 \) and \( v_x \gtrsim 7 \text{ TeV} \), since \[ M_{Z'} \gtrsim 3.5 \text{ GeV}. \] (23)

This very large bound is in contrast with the non-SUSY version of the model, where the gauge couplings are free parameters and can be much smaller, hence yielding lower mass bounds. The latter need to be evaluated as a function of both gauge couplings (see, e.g., [33]). A consequence of this strong constraint in the BLSSM is that the first terms in eqs. (16)–(18) can be large, pushing for CP-odd and charged Higgs masses in the TeV range.

Concerning the CP-even scalars, we see in eq. (20) that the MSSM and bilepton sectors are almost decoupled, mixing exclusively due to the gauge kinetic mixing. In first approximation, the mass matrix is block-diagonal, and has mass eigenstates that mimic the MSSM case. In practice, it turns out that only two Higgs bosons are light (hereafter called \( H_1 \) and \( H_2 \), one per sector), while the other two are very heavy (above the TeV scale). The lightest scalars are well defined states, being either almost exclusively doublet-like or bilepton-like. It is worth stressing that their mixing is small (see Fig. 5) and solely due to the gauge kinetic mixing (see also Ref. [38]).

Next, the sneutrino sector is here described since it shows two distinct features compared to the MSSM. Firstly, it gets enlarged by the superpartners of the right-handed neutrinos. Secondly, even more drastically, a splitting between the real and imaginary parts of each sneutrino occurs resulting in twelve states: six scalar sneutrinos and six pseudoscalar ones \([33, 40]\). The origin of this splitting is the \( Y_{\nu}^j \nu_i \nu_j \) term in the superpotential, eq. (2), which is a \( \Delta L = 2 \) operator after the breaking of \( U(1)_{B-L} \). Therefore, we define
\[ \tilde{\nu}_L^j = \frac{1}{\sqrt{2}} (\sigma_L^j + i \phi_L^j) \quad \tilde{\nu}_R^j = \frac{1}{\sqrt{2}} (\sigma_R^j + i \phi_R^j) \] (24)

In the following we will denote the partners of the left-handed and right-handed neutrinos by L-sneutrinos and R-sneutrinos, respectively. The 6 \( \times 6 \) mass matrices of the CP-even \( (m_{\tilde{\nu}^e_\nu}^2) \) and CP-odd \( (m_{\tilde{\nu}^o_\nu}^2) \) sneutrinos can be written in the basis \( (\sigma_L, \sigma_R) \) respectively \( (\phi_L, \phi_R) \) as
\[ m_{\tilde{\nu}^e_\nu}^2 = \begin{pmatrix} m_{\nu_L}^R & m_{\nu_L}^{R,T} \\ m_{\nu_L}^{T,R} & m_{\nu_R}^L \end{pmatrix}, \quad m_{\tilde{\nu}^o_\nu}^2 = \begin{pmatrix} m_{\nu_L}^{R,L} & m_{\nu_L}^{R,T} \\ m_{\nu_R}^{T,L} & m_{\nu_R}^R \end{pmatrix}. \] (25)

While \( m_{\nu_L}^L = m_{\nu_R}^R = m_{\nu_L}^R \) holds \(^2\), the entries involving R-sneutrinos differ by a few signs. It is possible to express them in a compact form by
\[ m_{\tilde{\nu}^e_\nu}^L = \frac{1}{8} \left( 1 \left( g_1^2 + g_2^2 + g_4^2 \right) \left( -v_u^2 + v_d^2 \right) + \eta g_{BL} \left( -2v_d^2 + 2v_u^2 - v_u^2 + v_d^2 \right) ight. 
+ \left. 2g_2^{BL} \left( -v_u^2 + v_d^2 \right) + 8m_{\nu_\tau}^2 \right) + 4v_u^2 Y_{\nu}^T Y_{\nu}^* \right) \] (26)

The upper signs correspond to the scalar and the lower ones to the pseudoscalar matrices and we have assumed CP conservation. In the case of complex trilinear couplings or \( \mu \)-terms, a mixing between the scalar and pseudoscalar particles occurs, resulting in 12 mixed states and consequently in a 12 \( \times 12 \) mass matrix. It particular the term \( \sim v_\nu Y_{\nu} \mu^* \) is potentially large and induces a large mass mixing between the scalar and pseudoscalar states. Also the corresponding soft SUSY-breaking term \( \sim v_\nu T_x \) can lead to a sizable mass splitting in the case of large \( |T_x| \), e.g. for large \( |A_0| \) at the GUT-scale where \( T_x = A_0 Y_{\nu} \) holds.

To gain some feeling for the behaviour of the sneutrino masses we can consider a simplified setup: neglecting kinetic mixing as well as left-right mixing, the masses of the R-sneutrinos at the SUSY scale can be expressed as
\[ m_{\tilde{\nu}^o_\nu}^2 \simeq m_{\tilde{\nu}^e_\nu}^2 + M_{Z'}^2 \left( \frac{1}{4} \cos(2\beta') + \frac{2\sqrt{2}v_\nu^2}{g_{BL}} \sin \beta'^2 \right) + M_{Z'}^2 \left( \frac{2\sqrt{2}v_\nu^2}{g_{BL}} \left( A_x \sin \beta' - \mu' \cos \beta' \right) \right), \] (29)
\[ m_{\tilde{\nu}^e_\nu}^2 \simeq m_{\tilde{\nu}^o_\nu}^2 + M_{Z'}^2 \left( \frac{1}{4} \cos(2\beta') + \frac{2\sqrt{2}v_\nu^2}{g_{BL}} \sin \beta'^2 \right) - M_{Z'}^2 \left( \frac{2\sqrt{2}v_\nu^2}{g_{BL}} \left( A_x \sin \beta' - \mu' \cos \beta' \right) \right). \] (30)

In addition, we treat the parameters \( A_x, m_{\tilde{\nu}^o_\nu}^2, M_{Z'}, \mu', Y_{\nu} \) and tan \( \beta' \) as independent. The different effects on the sneutrino masses can easily be understood by inspecting eqs. \([29]\) and \([30]\). The first two terms give always a positive contribution whereas the third one gives either

\(^2\) We have neglected the splitting induced by the light-handed neutrinos as this is suppressed by powers of the light neutrino mass over the sneutrino mass.
a positive or a negative one depending on the sign of $A_x \sin \beta' - \mu' \cos \beta'$. For example choosing $Y_a$ and $\mu'$ positive, one find the CP-even (CP-odd) sneutrino is the lighter one for $A_x < 0$ ($A_x > 0$). This is pictorially shown in Fig. 1 as a function of the GUT-scale input parameter $A_0$, for a choice of the other parameters. One notices that the CP-even (CP-odd) sneutrino is the lightest one when the 125 GeV Higgs boson is predominantly $H_1$ ($H_2$). It is worth pointing out here that, as will be described in the following section, when $M_{H_1} = 125$ GeV, the next-to-lightest Higgs boson can decay into pairs of CP-even sneutrinos, but not into the similar channel with CP-odd sneutrinos. Being $H_2$ predominantly a bilepton field, when this decay is open it saturates its BRs, see Fig. 4. Regarding the decay into CP-odd sneutrinos, this channel is accessible (i.e. $\tilde{\nu}^P$ is light enough) only in the region where $H_2$ is the SM-like Higgs boson, i.e. mainly coming from the doublets. In this case however, this decay channel is mitigated by the small scalar mixing and is not overwhelming (unlike for $H_1$, now mainly from the bileptons).

Depending on the parameters, either type of sneutrinos can get very light. If the LSP, it can be a suitable dark matter candidate and yield extra fully invisible decay channels to the Higgs bosons, thereby increasing their invisible widths. In the case of the decay into the CP-odd sneutrino, since this can happen mainly for the SM-like Higgs boson, one should account for the constraints on the former. Eventually, the R-sneutrinos could also get tachyonic or develop dangerous $R$-parity-violating effects. The first possibility is taken into account in our numerical evaluation by SPheno. The second case has been studied in Ref. [41] and is not here considered.

The last important sector for considerations that will follow is the one of the charged sleptons. See Ref. [42] for further details. Their mass matrix reads, in the $(\tilde{e}_L, \tilde{e}_R)$ basis, as

$$m_{\tilde{e}}^2 = \begin{pmatrix} m_{LL}^e & \frac{1}{\sqrt{2}} (v_d T_e - v_u Y_e) \\ \frac{1}{\sqrt{2}} (v_d T_e - v_u Y_e) & m_{RR}^e \end{pmatrix},$$

with

$$m_{LL}^e = m_{LL}^2 + \frac{\eta_1^2}{2} Y_1 e^c + \frac{1}{8} (g_1^2 - g_2^2) (v_u^2 - v_d^2) + 2 g_2^2 (v_u^2 - v_d^2) \tilde{1},$$

$$m_{RR}^e = m_{E}^2 + \frac{\eta_1^2}{2} Y_1 e^c + \frac{1}{8} (g_1^2 - g_2^2) (v_u^2 - v_d^2) - 2 (g_2^2) (v_u^2 - v_d^2) \tilde{1}. \tag{32}$$

In these equations we have suppressed the terms coming from the kinetic mixing, since the latter plays a smaller role here than for the Higgs particles. Of course, in our numerical studies all terms are taken into account. In general, we can parametrise the D-term contributions as a function of the $Z'$ mass, $M_{Z'}$, and of $\tan \beta'$ as

$$\frac{Q_{B-L}^2}{2} \frac{M_{Z'} (\tan^2 \beta' - 1)}{1 + \tan^2 \beta'}. \tag{34}$$

Obviously, the D-term contributions from the $B-L$ sector are larger for the sleptons than for the squarks by a factor of 3 due to the different $B-L$ charges. Because of this, their impact is much more pronounced in the slepton sector and it is possible to get large effects for staus while keeping the impact on the squarks under control. This is depicted in Fig. 2 for an illustrative choice of parameters, where we show the mass difference of stau (top-panel) and stops (bottom-panel) in the BLSSM in comparison to the MSSM expectations for a fixed set of SUSY-scale parameters. The different sfermion masses in the BLSSM as compared to the MSSM has a net impact onto the Higgs phenomenology, in particular in enhancing the $h \gamma \gamma$ coupling while keeping unaltered the SM-like Higgs boson, i.e. mainly coming from the doublets. This mechanism has been recently reanalysed also in Ref. [43] in the very same model.

$p$...
BR(\(\mu \rightarrow e\gamma\)) and BR(\(\mu \rightarrow 3e\)). The present exclusions are BR(\(\mu \rightarrow e\gamma\)) < \(5.7 \cdot 10^{-13}\) \([14]\) and BR(\(\mu \rightarrow 3e\)) < \(1.10^{-12}\) \([17]\). In Fig. 3 we plot these branching ratios as a function of the lightest (in black) and next-to-lightest (in red) SM-like neutrinos, which display some pattern for evading the bounds. In particular, they are required to be rather light, below 0.5 eV, while the model, ought to the scans here performed, seems to prefer configurations with neutrinos heavier than 0.01 eV, hence the preferred region in between. Lighter mass values are nonetheless also allowed.

FIG. 3: (Upper plot) BR(\(\mu \rightarrow e\gamma\)) and (lower plot) BR(\(\mu \rightarrow 3e\)) as a function of the light neutrino masses in GeV (black: \(\nu_1\), red: \(\nu_2\)). The blue horizontal lines represent the actual experimental limits, from Refs. \([10]\) and \([17]\), respectively. The parameters have been chosen as \(m_0 \in [0.4, 2]\) TeV, \(M_{1/2} \in [1.0, 2.0]\) TeV, \(\tan \beta \in [5, 40]\), \(A_0 \in [-4.0, 4.0]\) TeV, \(\tan \beta' \in [1.05, 1.15]\), \(M_{Z'} \in [2.5, 3.5]\) GeV, \(y_t \in [0.002, 0.4]\), \(y_\ell \in [0.05, 5]\) \times 10^{-6}.

For convenience, the impact of satisfying the earlier bounds will be shown only in the inverted hierarchy case, due to the smaller density of configurations therein.

Regarding the long-lasting \((g-2)_{\mu}\) discrepancy, in the setup investigated here charginos and charged Higgses are too heavy, same for the \(Z'\) boson, while the neutralino and sneutrino are too weakly coupled, to give a significant enhancement over the SM prediction.

HIGGS PHENOMENOLOGY

We review here the phenomenology of the Higgs sector, showing a survey of its features. First, results when
normal hierarchy is imposed are presented. Then, we will show that the inverted hierarchy is also possible on a large portion of the parameter space. Finally, it is instructive to describe how the Higgs-to-diphoton branching ratio can be easily enhanced in this model, despite the experimental data now converging to a more SM-like behaviour than in the recent past.

Normal hierarchy

In this subsection we discuss the normal hierarchy case, with the lightest Higgs bosons being the SM-like one (i.e., predominantly from the doublets), and a heavier Higgs boson predominantly from the bilepton fields (those carrying $B - L$ number and responsible for its symmetry breaking). Their mixing is going to be small and solely due to the kinetic mixing.

In Fig. 4 we first inspect the heavy Higgs boson branching ratios. Besides the standard decay modes, the decay into a pair of SM Higgs boson exist, as well as two new characteristic channels of this model, comprising right-handed (s)neutrinos.

1. $H_2 \rightarrow H_1 H_1$. Its BR can be up to 40% before the top quark threshold, and around 30% afterwards;
2. $H_2 \rightarrow \nu_\alpha \nu_\beta$. A similar decay channel exists for the $Z'$ boson. The BR are $\mathcal{O}(10)\%$, up to 20% depending on the heavy Higgs and neutrino masses;
3. $H_2 \rightarrow \tilde{\nu}^S \tilde{\nu}^S$, where $\tilde{\nu}^S$ is the CP-even sneutrino and the LSP, hence providing fully invisible decays of the heavy Higgs. If kinematically open, it saturates the Higgs BRs. Notice that only points with very light CP-even sneutrinos are shown.

While the first two channels exist also in the non-SUSY version of the model (see, e.g., [48]), the last one, involving the CP-even sneutrino, is truly new and rather intriguing. This is because the sneutrino is light and it can be a viable LSP candidate if with mass lower than $H_2$, as in this case [16]. It however implies that the heavy Higgs is predominantly bidoublet, with a light Higgs very much SM-like. This can be seen in Fig. 5 where the points with large BR($H_2 \rightarrow \tilde{\nu}^S \tilde{\nu}^S$) (in red) have the lowest mixing between $H_2$ and the SM scalar doublet fields, of the order of 0.1%. It immediately follows that this channel will have very small cross section at the LHC, when considering SM-like Higgs production mechanisms. This is true for all heavy Higgs masses $M_{H_2} > 140$ GeV. The 125 GeV Higgs is well SM-like, with tiny reduction of its couplings to the SM particle content. On the other side, the heavy Higgs is feebly mixed with the doublets, suppressing its interactions with the SM particles, and hence its production cross section. This can be seen in Fig. 4 (top frame). Considering only the gluon fusion production mechanism, and multiply it by the relevant BR, we get the cross sections for a choice of channels displayed therein. The most constraining channels, $H \rightarrow WW \rightarrow t\bar{t}jj$ and $H \rightarrow WW \rightarrow 2\ell 2\nu$, are also compared to the exclusions at the LHC for $\sqrt{s} = 8$ TeV from Refs. [14] and [51], respectively. The $Z$-mediated channels are well below current exclusions, that are hence not shown.

We see that all the displayed configurations are al-

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**References:**

[4] However, in the non-SUSY $B - L$ model the Higgs mixing angle is a free parameter, directly impacting on these branching ratios.

[5] Starting from $M_{H_2} > 130$ GeV.
Typical cross sections range between $0 \ll \mathcal{A} \lesssim 10^2 \eta/\tilde{N}$ and it can yield $\pm M_{10} \lesssim 100$. Notice that for $H_2$ is SM-like one and a lighter Higgs boson exists.

![FIG. 6: Cross sections at $\sqrt{s} = 8\,\text{TeV}$ for (upper plot) the SM-like channels (lower plot) the new channels, as a function of the heavy Higgs mass. The solid lines above are the exclusion curves from [44, 54].](image)

followed by the current searches (the exclusions shown by solid curves of same color as the depicted channel). This is due to the suppression of the heavy Higgs cross sections coming from the small scalar mixing.

In the lower plot are displayed the cross sections for the new channels. Those pertaining to model configurations for which the heavy Higgs boson decays to the $CP$-even sneutrino (LSP), yielding a fully invisible decay mode, are displayed in red. Contrary to the all other cases, the production of the heavy Higgs for this channel is via vector boson fusion as searched for at the LHC [51]. Typical cross sections range between 0.1 fb and 1 fb. The $H_2 \rightarrow H_1 H_1$ channel is shown in blue and it can yield cross sections of $1 \div 10\,\text{fb}$ for $250 < M_{H_2} < 400$ GeV. Last is the $H_2 \rightarrow \nu_\mu \nu_\mu$ channel. It can be sizable only for very light $H_2$ masses: $\sim 10 \div 100\,\text{fb}$ for $140 < M_{H_2} < 160$ GeV, although the further decay chain of the heavy neutrinos have to be accounted for. The latter can give spectacular multi-leptonic final state of the heavy Higgs boson ($4\ell 2\nu$ and $3\ell 2\nu$) or high jet multiplicity ones ($2\ell 4j$), via $\nu_\mu \rightarrow e^+ W^+ \nu$ and $\nu_\mu \rightarrow \nu Z$ in a $2 : 1$ ratio (modulo threshold effects). Further, these decays are typically

seesaw suppressed and can therefore give rise to displaced vertices [52].

**Inverted hierarchy**

In this subsection we discuss the inverted hierarchy case, where $H_2$ is SM-like one and a lighter Higgs boson exists.

![FIG. 7: Branching ratios for the 125 GeV Higgs boson ($H_2$). The decay into heavy neutrinos is displayed with diamonds. All others with circles. Gray points are excluded by the low energy observables and by HiggsBounds. The decay into CP-odd sneutrinos is not shown.](image)

We start once again by presenting the BRs for the next-to-lightest Higgs boson in Fig. 4. This time however this is the SM-like boson, hence predominantly from the doublets. It has the same new channels as the heavy Higgs in the normal hierarchy, the only difference being the $CP$-odd R-sneutrino instead of the $CP$-even one. The configurations not allowed by the low energy observables or by HiggsBounds are displayed as gray points. We see that it may have sizable decays into pairs of the lighter Higgs boson. This decay is still allowed with rates up to a few percent. Further, rare decays into pairs of heavy neutrinos are also present, with BRs below the permil level. This channel can give rise to rare multi-lepton/jets decays for the SM-like Higgs boson, that are searched for at the LHC, even in combination with searches for displaced vertices [53]. The last available channel is the decay into pairs of $CP$-odd R-sneutrinos. Being the LSP, it will increase the invisible decay width and hence give larger-than-expected widths for the SM-like boson. Its rate is obviously constrained, and a precise evaluation of the allowed range is needed. It however goes beyond the scope of the present review and we postpone it to a future publication.

Regarding the lightest Higgs boson ($H_1$), this will obviously decay predominantly into $b\bar{b}$ pairs. Notice that
due to its large bilepton fraction it can also decay into pairs of very light RH neutrinos, at sizable rates depending on the neutrino masses. As in the in previous figure, the non-allowed configurations are displayed as gray points. We see that the pattern of decays is not affected by the inclusion of the constraints, in the sense that this channel stays viable. Once again, the latter will yield multi-lepton/jet final state, which will be very soft, and hence very challenging for the LHC. However, also in this case displaced vertices my appear.

![Graph](image1.png)

**Fig. 8:** Same as in Fig. 7 for the lightest Higgs boson ($H_1$).

As in the previous section, we show in Fig. 9 the mixing between the Higgs mass eigenstates and the doublet fields as a function of the light Higgs mass, to show that $H_2$ is here rather SM-like. Once more, the gray points displayed here are excluded by the low energy observables and by HiggsBounds.

Finally, the production cross sections for the lightest Higgs boson can be evaluated. In Fig. 10 we compare the direct production (for the main SM production mechanisms, gluon fusion and vector boson fusion) with the pair production via $H_2$ decays only via gluon fusion, $gg \to H_2 \to H_1 H_1$. When the latter channel is kinematically open, *i.e.* $2M_{H_1} < 125$ GeV, the lightest Higgs pair production has cross section up to 1 pb at the LHC at $\sqrt{s} = 8$ TeV. A thorough analysis of the phenomenology of the Higgs sector in the BLSSM for the upcoming LHC run 2, based on the first investigations shown here, will be performed soon.

**Enhancement of the diphoton rate**

Despite disfavoured by most recent data [12], it is still instructive to show how in the BLSSM model (and in general in gauge-extended MSSM models) the Higgs-to-diphoton decay can easily be enhanced. See Ref. [12] for further details.

![Graph](image2.png)

**Fig. 9:** Mixing between scalar mass eigenstates and Higgs doublets. (black: $H_1$, red: $H_2$) and scalar doublet fields, as a function of $M_{H_1}$. $ZH[i,J]$ is the scalar mixing matrix. Gray points are excluded by the low energy observables and by HiggsBounds.

![Graph](image3.png)

**Fig. 10:** Cross sections at $\sqrt{s} = 8$ TeV for different production mechanisms. Gluon-fusion (in red) and vector-boson-fusion (in green) mechanisms are displayed only for $M_{H_1} > 50$ GeV for simplicity. Gray points are excluded by the low energy observables and by HiggsBounds.

To start our discussion let us briefly review the partial decay width of the Higgs boson $h$ into two photons within the MSSM and its singlet extensions. This can be written as (see, e.g., [54])

$$\Gamma_{h \to \gamma \gamma} = \frac{\alpha^2 m_h^3}{128 \sqrt{2} \pi^3} \left| \sum_f N_f Q_f^2 g_{h f f} A_f^h (\tau_f) + g_{h WW} A_1^h (\tau_W) \right|^2,$$

$$+ m_h^2 \frac{g_{h H^+ H^-}}{2 m_{H^+} m_{H^-}} A_0^h (\tau_{H^\pm}) + \sum_{\chi_i^\pm} \frac{2 m_{\chi_i^\pm}}{m_{H^\pm}} \frac{g_{h \chi_i^\pm \chi_i^\mp}}{m_{H^\pm}} A_1^h (\tau_{\chi_i^\pm})$$

$$+ \sum_{\tilde{e}_i} \frac{g_{h \tilde{e}_i \tilde{e}_i}}{m_{\tilde{e}_i}^2} A_0^h (\tau_{\tilde{e}_i}) + \sum_{\tilde{q}_i} \frac{g_{h \tilde{q}_i \tilde{q}_i}}{m_{\tilde{q}_i}^2} 3 Q_i^2 A_0^h (\tau_{\tilde{q}_i}) \right|^2, \quad (35)$$
corresponding to the contributions from charged SM fermions, $W$ bosons, charged Higgs, charginos, charged sleptons and squarks, respectively. The amplitudes $A_i$ at lowest order for the spin–1, spin–$\frac{1}{2}$ and spin–0 particle contributions, can be found for instance in Ref. [54]. $g_{hXX}$ denotes the coupling between the Higgs boson and the particle in the loop and $Q_X$ is its electric charge. In the SM, the largest contribution is given by the $W$-loop, while the top-loop leads to a small reduction of the decay rate. In the MSSM, it is possible to get large contributions due to sleptons and squarks, although it is difficult to realise such a scenario in a constrained model with universal sfermion masses [55–57]. In singlet or triplet extension of the MSSM also the chargino and charged Higgs can enhance the loop significantly [58, 59]. However, this is only possible for large singlet couplings which lead to a cut-off well below the GUT scale. In contrast, it is possible to enhance the diphoton ratio in the BLSSM due to light staus even in the case of universal boundary conditions at the GUT scale. We show this by calculating explicitly the contributions of the stau:

$$A(\tau) = \frac{1}{3} \frac{\partial \det m^2}{\partial \log v} \quad (36)$$

$$\simeq -\frac{2}{3} \left( m_E^2 + D_L (m_L^2 + D_L) + m_Z^2 \mu \tan \beta (2 A_\tau - \mu \tan \beta) \right) \quad (37)$$

Here, $D_L$ and $D_R$ represent the D-term contributions of the left- and right-handed stau and we have neglected sub-leading contributions. Given that $2 A_\tau < \mu \tan \beta$, for fixed values of the other parameters, $D_R$ and $D_L$ can be used to enhance the $\gamma\gamma$ rate by suppressing the denominator. This is shown in Fig. 11 depending on $\tan \beta'$, it is possible to obtain enhanced branching ratios (BR) of the SM-like Higgs into photons even for much heavier soft masses than possible in the MSSM.

We turn now to a fully numerical analysis to demonstrate the mechanism to reduce the stau mass in comparison to the stop mass (to enhance the Higgs to diphoton rate) discussed above. In Table 11 we have collected two possible scenarios that provide a SM-like Higgs particle in the mass range preferred by LHC results with an enhanced diphoton rate. In the first point, the lightest $CP$-even scalar eigenstate is the SM-like Higgs boson while the light bilepton is roughly twice as heavy. In Fig. 12 we show that all the features arise from the extended gauge sector: it is sufficient to change only $\tan \beta'$ to obtain an enhanced diphoton signal $R_{\gamma\gamma} \equiv [\sigma(gg \to h_1) BR(h_1 \to \gamma\gamma)]_M / [\sigma(gg \to h_1) BR(h_1 \to \gamma\gamma)]_L$, and the correct dark matter relic density while keeping the mass of the SM-like Higgs nearly unchanged. The dark matter candidate in this scenario is the lightest neutralino, that is mostly
a bileptino (the superpartner of the bileptons). The correct abundance for \( \tan \beta' \simeq 1.156 \) is obtained due to a co-annihilation with the light stau. In the second point, the SM-like Higgs is accompanied by a light scalar around 98 GeV which couples weakly to the SM gauge bosons, compatibly with the LEP excess \([60–62] \). In this case, the LSP is a \( CP \)-odd sneutrino which annihilates very efficiently due to the large \( Y \). This usually results in a small relic density. To get an abundance which is large enough to explain the dark matter relic, the mass of the sneutrino has to be tuned below \( m_W \) \([16] \). This can be achieved by slightly increasing \( \tan \beta' \) and by tuning the Majorana Yukawa couplings \( Y \), that tends to increase the SM-like Higgs mass for the given point. It is worth mentioning that a neutralino LSP with the correct relic density in the stau co-annihilation region can also be found in this scenario. Notice that both points yield rates consistent with observations in the \( WW^*/ZZ^* \) channels (measured at the LHC) (being \( c_{hhZZ} \sim 1 \)).

\[
\begin{array}{|c|c|c|}
\hline
 & \text{Point I} & \text{Point II} \\
\hline
m_{h_1} [\text{GeV}] & 125.2 & 98.2 \\
m_{h_2} [\text{GeV}] & 186.9 & 123.0 \\
m_3 [\text{GeV}] & 267.0 & 237.3 \\
\hline
doublet fr. [%] & 99.5 & 8.7 \\
bilepton fr. [%] & 0.5 & 91.3 \\
\hline
c_{h_1,gg} & 0.992 & 0.087 \\
c_{h_1,ZZ} & 1.001 & 0.085 \\
c_{h_2,gg} & 0.005 & 0.911 \\
c_{h_2,ZZ} & 0.005 & 0.921 \\
\hline
\Gamma(h_1) [\text{MeV}] & 4.13 & 0.22 \\
R_{zz}^1 & 1.57 & 0.085 \\
R_{zz}^2 & 1.03 & 0.089 \\
R_{WW^*}^H & 0.98 & 0.05 \\
\hline
\Gamma(h_2) [\text{MeV}] & 4.8 & 3.58 \\
R_{zz}^2 & 0.005 & 1.79 \\
R_{ZZ}^1 & 0.006 & 0.95 \\
R_{WW^*}^Z & 0.01 & 0.88 \\
\hline
LSP mass [GeV] & 253.9 & 82.9 \\
\Omega h^2 & 0.1 & 10^{-2} \\
\hline
\end{array}
\]

\textbf{TABLE II}: The input parameter used: Point I: \( m_0 = 673 \text{ GeV} \), \( M_{1/2} = 2220 \text{ GeV} \), \( A_0 = -1842 \text{ GeV} \), \( \tan \beta = 42.2 \), \( \tan \beta' = 1.1556 \), \( M_{\tilde{\tau}} = 2550 \text{ GeV} \), \( Y_x = 1 - 0.42 \) (neutralino LSP). Point II: \( m_0 = 742 \text{ GeV} \), \( M_{1/2} = 1572 \text{ GeV} \), \( A_0 = 3277 \text{ GeV} \), \( \tan \beta = 37.8 \), \( \tan \beta' = 1.140 \), \( M_{\tilde{\tau}} = 2365 \text{ GeV} \), \( Y_x = \text{diag}(0.40, 0.40, 0.13) \) (\( CP \)-odd sneutrino LSP). \( c_{SVV} \) denotes the coupling squared of the Higgs fields to vector bosons normalised to the SM values.

\textbf{CONCLUSIONS}

In this review I described the \( U(1)_{B-L} \) extension of the MSSM, focusing in particular on the scalar sector, described in details. The fundamental role that the gauge kinetic mixing plays in this sector has been repetitively underlined.

The comparison to the most constraining low energy observables showed that a preferred region for the light neutrino masses exist to evade these bounds. Then, I presented a first systematic investigation of the phenomenology of the Higgs sector of this model, showing that both the normal hierarchy and the inverted hierarchy of the two lightest Higgs bosons are naturally possible in a large portion of the parameter space. Particular attention has been devoted to analyse the new decay channels comprising both the \( CP \)-even and \( CP \)-odd R-sneutrinos, which are a peculiarity pertaining the BLSSM. Based on these first findings, a thorough analysis of the Higgs sector in the BLSSM at the upcoming LHC run 2 will be soon prepared. The fit of the SM-like Higgs boson to the LHC data will also be performed with \textit{HiggsSignals} \([63] \).

Finally, I described how in the BLSSM model (and in general in gauge-extended MSSM models) the Higgs-to-diphoton decay can easily be enhanced. Despite disfavoured by most recent data, it is still instructive to show this mechanism to understand the model properties.

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