Upper bound of the lightest Higgs boson in a supersymmetric $SU(3)_L \otimes U(1)_X$ gauge model

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A supersymmetric version of the $SU(3)_L \otimes U(1)_X$ gauge model is presented. The model is embedded in a $SU(6)_L \otimes U(1)_X$ gauge group and it has been built up completely chiral anomaly free and without any exotic charges in the fermionic content. The minimal supersymmetric standard model can be seen as an effective theory of this larger symmetry. We analyze how the spectrum of the new model can shift the upper bound on the CP-even Higgs boson and we find that it can be moved up to 143 GeV.

I. INTRODUCTION

The standard model (SM) has been established with great precision as the model that describes interactions of the gauge bosons with the fermions [1]. However, the Higgs sector, responsible for the symmetry breaking has not been tested yet. But in any case the SM still has many unanswered questions and there are some experimental results on solar and atmospheric neutrinos which suggest no standard interpretation. The observation of neutrino oscillations can be explained by introducing mass terms for these particles, in which case the mass eigenstates are different from the interaction eigenstates; it should be noticed that the oscillation phenomena imply that lepton family number is violated and such a fact leads us to consider the existence of physics beyond the SM [2].

Among the various candidates to physics beyond the SM, supersymmetric theories play a special role. Supersymmetric models are the most developed and most widely accepted. Although there is not yet direct experimental evidence of supersymmetry (SUSY), there are many theoretical arguments indicating that SUSY might be of relevance for physics beyond the SM. The main reason behind the phenomenological interest is that SUSY provides a solution for the naturalness problem which allows a fundamental scalar of the order of the electroweak scale if the mass of the superpartners are of the order of $\sim O(1 \text{ TeV})$. The most popular version, of course, is the supersymmetric version of the SM, usually called the minimal supersymmetric standard model (MSSM), which is a computable theory and it is compatible with all precision electroweak tests [3].

Up to now, one impressive phenomenological virtue in favor of SUSY is obtained from the unification of couplings in grand unified theories, but on the other hand the lack of SUSY signals at the CERN e^+e^- collider LEP and the lower limit on m_h pose problems for the MSSM [4,5]. About this point, unlike in the SM, in the MSSM the coefficient of the quartic potential of the Higgs sector is not an independent parameter, instead it is related with the ratio $\tan \beta$ and the $SU(2)_L \otimes U(1)_Y$ gauge couplings. As a consequence, the

value of the CP-even scalar mass is not controlled by the SUSY breaking parameters but by the electroweak breaking sector. This yields the tree level bound $m_h^2 \lesssim m_Z^2 \cos^2 2\beta \lesssim m_Z^2$, mass range already excluded by LEPII and could rule out the MSSM. But if the leading effects included in radiative corrections from the top-quark–top-squark sector are taken into account, the upper bound can be pushed up to 135 GeV [6]. Now, the limit coming from LEP on the SM Higgs boson is $m_h \gtrsim 114$ GeV [7], putting restrictions on the parameter space of the MSSM, strong enough to consider new ideas beyond the SM that can shift this bound. The physical Higgs boson mass can be increased by enhancing the quartic coupling through an extended gauge sector and/or new superpotential Higgs couplings, in other words through D terms or F terms [8].

Among the possibilities with the nice properties of SUSY theories and the possibility to relax the bound on the lightest Higgs boson is $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ as a local gauge theory. It has been studied previously by many authors in non-SUSY and SUSY versions, who have explored different spectra of fermions and Higgs bosons [9–11]. There are many considerations about this model but the most studied motivations of this large symmetry are the possibility to give mass to the neutrino sector [10], anomaly cancellations in a natural way in the 3-family version of the model, and an interpretation of the number of the fermionic families related with the anomaly cancellations [12]. Recently, Ref. [13] has presented a careful analysis of these kind of models without SUSY, taking into account the anomaly cancellation constraints. In fact, the version called model A in Ref. [13] was already supersymmetrized [11] and it was shown to be an anomaly free model, and a family independent theory. Reference [11] has also shown that the SUSY models based on the symmetric $SU(3)_L \otimes U(1)_X$ model as a E_6 subgroup can shift the upper bound on the lightest CP-even Higgs boson up to 140 GeV. In the present work, we supersymmetrize the version called model B in Ref. [13]. The models A and B are different, the first one being an E_6 subgroup in contrast to the second one which is a $SU(6)_L \otimes U(1)_X$ subgroup.

The model presented here is a supersymmetric version of the gauge symmetry $SU(3)_L \otimes U(1)_X$ but it is different from the versions considered previously [11,14]. The model considered does not introduce Higgs triplets in the spectrum to break the symmetry; instead they are included in the scalar

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sector of the chiral superfields. As we will show, the free parameters of the model are also reduced by using a basis where only one of the vacuum expectation values (VEV) of the neutral singlets of the spectrum generates the breaking of the larger symmetry to the SM symmetry. Moreover, the fermionic content presented here does not have any exotic charges. Finally, our aim is to study how the spectrum of the new model based on the symmetry $SU(3)_L \otimes U(1)_X$ can change the upper bound on the lightest CP-even Higgs boson.

This work is organized as follows. In Sec. II, we present the details of the non-SUSY version of the $SU(3)_L \otimes U(1)_X$ model. In Sec. III we discuss the SUSY version and the spontaneous symmetry breaking mechanism, as well as some phenomenological implications of the model. Section IV contains our conclusions.

II. NON-SUSY VERSION

We want to describe the supersymmetric version of the $SU(3)_C \otimes SU(3)_L \otimes U(1)_X$ gauge symmetry which is embedded into $SU(6)_L \otimes U(1)_X$ [13]. But in order to be clear, first of all we present the non-SUSY version of the model. We assume that the left-handed quarks (color triplets) and left-handed leptons (color singlet) transform as the $\bar{3}$ and 3 representations of $SU(3)_L$, respectively. In this model the anomalies cancel individually in each family as it is done in the SM.

The model thus ends up with the following anomaly free multiplet structure:

$$\chi_L = \begin{pmatrix} d \\ u \\ U \end{pmatrix}_L \qquad d_L^c \qquad u_L^c \qquad U_L^c$$

$$(3, \bar{3}, \tfrac{1}{3}) \qquad (\bar{3}, 1, \tfrac{1}{3}) \qquad (\bar{3}, 1, -\tfrac{2}{3}) \qquad (\bar{3}, 1, -\tfrac{2}{3}) \qquad (1)$$

$$\psi_{1L} = \begin{pmatrix} \nu_e \\ e^- \\ E_1^- \end{pmatrix}_L \qquad \psi_{2L} = \begin{pmatrix} N_2^0 \\ E_2^- \\ E_3^- \end{pmatrix}_L \qquad \psi_{3L} = \begin{pmatrix} E_2^+ \\ N_1^0 \\ \nu_e^c \end{pmatrix}_L$$

$$(1, 3, -\tfrac{2}{3}) \qquad (1, 3, -\tfrac{2}{3}) \qquad (1, 3, \tfrac{1}{3}) \qquad (2)$$

$$e_L^+ \qquad E_{1L}^+ \qquad E_{3L}^+$$
$$(1,1,1) \qquad (1,1,1) \qquad (1,1,1). \qquad (3)$$

The numbers in parenthesis refer to the $(SU(3)_C, SU(3)_L, U(1)_X)$ quantum numbers respectively. The symmetry is broken following the pattern

$$SU(3)_C \otimes SU(3)_L \otimes U(1)_X \rightarrow SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$$

$$\rightarrow SU(3)_C \otimes U(1)_Q$$

and give masses to the fermion fields in the model. With this

in mind, it is necessary to introduce the following two Higgs triplets: $\phi_3(1,3,1/3)$ with a vacuum expectation value (VEV) aligned in the direction $\langle \phi_3 \rangle = (0,0,V)^T$ and $\phi_4(1,3,1/3)$ with a VEV aligned as $\langle \phi_4 \rangle = (0, v/\sqrt{2}, 0)^T$, with the hierarchy $V > v \sim 250$ GeV (the electroweak breaking scale). Note that ϕ_3 and ϕ_4 have the same quantum numbers but they get VEV's at different mass scales generating a large mass to the exotic U quark and a realistic mass to the known up quark. One more Higgs scalar $\phi_2(1,3,-2/3)$ is needed in order to give a mass to the down quark field in the model. With the same scalar fields the neutral and charged fermions get masses through the seesaw mechanism [15]. Diagonalizing the mass matrices there are a light Dirac neutrino and a heavy Dirac neutrino. Also there are a light charged lepton which corresponds to electron and three exotic heavy charged leptons.

The model has 9 gauge bosons. One of them B^μ is associated with $U(1)_X$ symmetry, and other 8 fields are associated with the $SU(3)_L$ symmetry. The expression for the electric charge generator in $SU(3)_L \otimes U(1)_X$ is a linear combination of the three diagonal generators of the gauge group

$$Q = T_{3L} + \frac{1}{\sqrt{3}} T_{8L} + X I_3, \qquad (4)$$

where $T_{iL} = \lambda_i/2$ with λ_i the Gell-Mann matrices and I_3 the unit matrix.

After breaking the symmetry, we get mass terms for the charged and the neutral gauge bosons. By diagonalizing the matrix of the neutral gauge bosons we get the physical mass eigenstates which are defined through the mixing angle θ_W given by $\tan \theta_W = \sqrt{3} g_1 / \sqrt{3g^2 + g_1^2}$ with g_1 and g the coupling constants of the $U(1)_X$ and $SU(3)_L$ respectively. Also we can identify the Y^μ hypercharge field associated with the SM gauge boson as

$$Y^\mu = \frac{\tan \theta_W}{\sqrt{3}} A_8^\mu + (1 - \tan \theta_W^2/3)^{1/2} B^\mu. \qquad (5)$$

In the SM the coupling constant g' associated with the hypercharge $U(1)_Y$, can be given by $\tan \theta_W = g'/g$ where g is the coupling constant of $SU(2)_L$ which in turn can be taken equal to the $SU(3)_L$ coupling constant. Using the $\tan \theta_W$ given by the diagonalization of the neutral gauge boson matrix we obtain the matching condition

$$\frac{1}{g'^2} = \frac{1}{g_1^2} + \frac{1}{3g^2}. \qquad (6)$$

We shall use this relation to write g_1 as a function of g' in order to find the Higgs potential of the $SU(3) \otimes U(1)_X$

SUSY model at low energies and compare it with the MSSM one. In particular, we will show that it reduces to the MSSM in this limit.

III. SUSY VERSION

In the SUSY version the above content of fermions should be written in terms of chiral superfields, and the gauge fields will be in vector supermultiplets as it is customary in SUSY theories. In this case we should introduce two more fermion fields and one more scalar. They will be arranged in the corresponding superfields, in such way that they cancel chiral anomalies automatically, and additionally the Higgs scalar fields break the symmetry and generate the fermion and gauge boson masses. Therefore the contents of superfields are

$$
\hat{\chi}_L = \begin{pmatrix} \hat{d} \\ \hat{u} \\ \hat{U} \end{pmatrix}_L \qquad \hat{d}_L^c \qquad \hat{u}_L^c \qquad \hat{U}_L^c
$$

$$
(3,\bar{3},\tfrac{1}{3}) \qquad (\bar{3},1,\tfrac{1}{3}) \quad (\bar{3},1,-\tfrac{2}{3}) \quad (\bar{3},1,-\tfrac{2}{3}) \tag{7}
$$

$$
\hat{\psi}_{1L} = \begin{pmatrix} \hat{\nu}_e \\ \hat{e}^- \\ \hat{E}_1^- \end{pmatrix}_L \qquad \hat{\psi}_{2L} = \begin{pmatrix} \hat{N}_2^0 \\ \hat{E}_2^- \\ \hat{E}_3^- \end{pmatrix}_L \qquad \hat{\psi}_{3L} = \begin{pmatrix} \hat{E}_2^+ \\ \hat{N}_1^0 \\ \hat{\nu}_e^c \end{pmatrix}_L
$$

$$
(1,3,-\tfrac{2}{3}) \qquad\qquad (1,3,-\tfrac{2}{3}) \qquad\qquad (1,3,\tfrac{1}{3}) \tag{8}
$$

$$
\hat{\psi}_{4L} = \begin{pmatrix} \hat{E}_4^+ \\ \hat{N}_4^0 \\ \hat{N}_5^0 \end{pmatrix}_L \qquad \hat{\psi}_{5L} = \begin{pmatrix} \hat{E}_4^- \\ \hat{N}_7^0 \\ \hat{N}_6^0 \end{pmatrix}_L \qquad \hat{e}_L^+ \quad \hat{E}_{1L}^+ \quad \hat{E}_{3L}^+
$$

$$
(1,3,\tfrac{1}{3}) \qquad\quad (1,3,-\tfrac{1}{3}) \quad (1,1,1)\,(1,1,1)\,(1,1,1). \tag{9}
$$

Equations (7) and (8) plus singlets of Eq. (9) have the same fermionic content as the non-SUSY version which is anomaly free. In order to give masses in the non-SUSY model it was necessary to introduce a scalar Higgs boson ϕ_4. To supersymmetrize the model we have introduced a chiral superfield with the same quantum numbers $\hat{\psi}_{4L}$. But it introduces more fermions and in order to cancel the triangle anomalies a new chiral superfield is necessary $\hat{\psi}_{5L}$, in such a way that $\hat{\psi}_{4L}$ and $\hat{\psi}_{5L}$ are in a vector representation of the $SU(3)_L \otimes U(1)_X$ gauge symmetry and then the model is completely chiral anomaly free.

The scalar sector of the supermultiplets has the appropriate quantum numbers in order to be useful as a Higgs sector of the theory and is not necessary to introduce new scalar fields to get the spontaneous symmetry breaking to the SM symmetry. The superpotential can be built up by taking into account the non-SUSY Yukawa Lagrangian form and $\epsilon_{abc}\psi_i^a \psi_j^b \psi_k^c$ which is a singlet; we have

$$
\begin{aligned}
W =\ & h_1 \epsilon_{abc} \hat{\psi}_1^a \hat{\psi}_3^b \hat{\psi}_4^c + h_2 \epsilon_{abc} \hat{\psi}_3^a \hat{\psi}_2^b \hat{\psi}_4^c + \hat{\psi}_1(h_3 \hat{\psi}_5 \hat{e}^+ \\
& + h_4 \hat{\psi}_5 \hat{E}_1^+ + h_5 \hat{\psi}_5 \hat{E}_3^+) + \hat{\psi}_2(h_6 \hat{\psi}_5 \hat{e}^+ + h_7 \hat{\psi}_5 \hat{E}_1^+ \\
& + h_8 \hat{\psi}_5 \hat{E}_3^+) + h_u \hat{\chi}_L \hat{\psi}_4 \hat{u} + h_U \hat{\chi}_L \hat{\psi}_4 \hat{U} + h_u' \hat{\chi}_L \hat{\psi}_3 \hat{u} \\
& + h_U' \hat{\chi}_L \hat{\psi}_3 \hat{U} + h_d \hat{\chi}_L \hat{\psi}_2 \hat{d} + h_d' \hat{\chi}_L \hat{\psi}_1 \hat{d} + \mu_1 \hat{\psi}_4 \hat{\psi}_5 \\
& + \mu_2 \hat{\psi}_3 \hat{\psi}_5 .
\end{aligned} \tag{10}
$$

Once we have the superpotential W, the theory is defined and we can get the Yukawa interactions and the scalar potential. Here, we should mention that the lepton number is not conserved, in order to generate neutrino masses similar to R_p violating models [16].

Now, we will concentrate our attention on the scalar Higgs potential, which is given by

$$
V = \left| \frac{\partial W}{\partial A_i} \right|^2 + \frac{1}{2} D^a D^a + \frac{1}{2} D' D', \tag{11}
$$

where

$$
D^a = g A_i^\dagger \frac{\lambda_{ij}^a}{2} A_j ,
$$

$$
D' = g_1 A_i^\dagger X(A_i) A_i \tag{12}
$$

and A_i are the scalar components of the chiral supermultiplets. For our purpose, it is enough to consider only

$$
W' = h_1 \epsilon_{abc} \hat{\psi}_1^a \hat{\psi}_3^b \hat{\psi}_4^c + h_2 \epsilon_{abc} \hat{\psi}_3^a \hat{\psi}_2^b \hat{\psi}_4^c + \mu_1 \hat{\psi}_4 \hat{\psi}_5 + \mu_2 \hat{\psi}_3 \hat{\psi}_5 . \tag{13}
$$

As we already mentioned, the scalar Higgs bosons responsible of the symmetry breaking are the scalar sector of the fermionic chiral supermultiplets, thus we have $\phi_i = \tilde{\psi}_{iL}$. Now we focus on the scalar potential obtained from the superpotential W' plus the soft breaking terms which are

$$
V_{soft} = \sum_{i=1}^{5} m_i^2 \phi_i^\dagger \phi_i + \mu_1'(\phi_4^\dagger \phi_5 + \text{H.c.}) + \mu_2'(\phi_3^\dagger \phi_5 + \text{H.c.}). \tag{14}
$$

The D terms and F terms are

$$V = \frac{g^2}{8}\left\{\frac{4}{3}\sum_{a=1}^{5}(\phi_a^\dagger\phi_a)^2 + 4\sum_{a\neq b}^{5}(\phi_a^\dagger\phi_b)(\phi_b^\dagger\phi_a) - \frac{4}{3}\sum_{a\neq b}^{5}(\phi_a^\dagger\phi_a)(\phi_b^\dagger\phi_b)\right\} + \frac{g_1^2}{18}\{-2(\phi_1^\dagger\phi_1 + \phi_2^\dagger\phi_2) + (\phi_3^\dagger\phi_3 + \phi_4^\dagger\phi_4)$$

$$-\phi_5^\dagger\phi_5\}^2 + h_1^2\{(\phi_3^\dagger\phi_3)(\phi_4^\dagger\phi_4) + (\phi_1^\dagger\phi_1)(\phi_3^\dagger\phi_3) - (\phi_3^\dagger\phi_4)(\phi_4^\dagger\phi_3) - (\phi_1^\dagger\phi_3)(\phi_3^\dagger\phi_1)\} + h_2^2\{(\phi_3^\dagger\phi_3)(\phi_4^\dagger\phi_4)$$

$$+(\phi_2^\dagger\phi_2)(\phi_3^\dagger\phi_3) - (\phi_3^\dagger\phi_4)(\phi_4^\dagger\phi_3) - (\phi_2^\dagger\phi_3)(\phi_3^\dagger\phi_2)\} + h_1^2\{(\phi_1^\dagger\phi_1)(\phi_4^\dagger\phi_4) - (\phi_1^\dagger\phi_4)(\phi_4^\dagger\phi_1)\} + h_2^2\{(\phi_2^\dagger\phi_2)(\phi_4^\dagger\phi_4)$$

$$-(\phi_2^\dagger\phi_4)(\phi_4^\dagger\phi_2)\} + 2h_1h_2\{(\phi_1^\dagger\phi_2)(\phi_4^\dagger\phi_4) - (\phi_1^\dagger\phi_4)(\phi_2^\dagger\phi_4) + (\phi_1^\dagger\phi_2)(\phi_3^\dagger\phi_3) - (\phi_1^\dagger\phi_3)(\phi_2^\dagger\phi_3)\} + (\mu_1^2 + \mu_2^2)(\phi_5^\dagger\phi_5)$$

$$+\mu_1^2(\phi_4^\dagger\phi_4) + \mu_2^2(\phi_3^\dagger\phi_3) + \mu_1\mu_2(\phi_4^\dagger\phi_3 + \phi_3^\dagger\phi_4). \tag{15}$$

Now, we are ready to break down the symmetry $SU(3)_L \otimes U(1)_X$ to the SM symmetry $SU(2)_L \otimes U(1)_Y$. Writing the Higgs triplets of $SU(3)_L$ as branching rules of $SU(2)_L$, we have

$$\phi_1 = \begin{pmatrix}\tilde{l}\\\tilde{E}_1^-\end{pmatrix}, \quad \phi_2 = \begin{pmatrix}H_1\\\tilde{E}_3^-\end{pmatrix}, \quad \phi_3 = \begin{pmatrix}H_2\\\tilde{N}_1^0\end{pmatrix},$$

$$\phi_4 = \begin{pmatrix}H_4\\\tilde{N}_5^0\end{pmatrix}, \quad \phi_5 = \begin{pmatrix}H_5\\\tilde{N}_6^0\end{pmatrix}, \tag{16}$$

where $\tilde{l}$ and H_i are doublets of $SU(2)_L$ and $\tilde{E}_i^-$, $\tilde{N}_i^0$ are singlets. Thus the VEV's of $\langle\tilde{N}_1^0\rangle = u_1$, $\langle\tilde{N}_6^0\rangle = u_2$ and $\langle\tilde{N}_5^0\rangle = u$ make the job. But it is possible to choose only one of them different from zero, $u_1 = u_2 = 0$, $u \neq 0$ [17], and the would-be Goldstone bosons of the symmetry breaking $SU(3)_L \otimes U(1)_X / SU(2)_L \otimes U(1)_Y$ become degrees of freedom of the field ϕ_4. The terms involving H_1, H_2 doublets generate the fermion masses in the non-SUSY model after the electroweak symmetry breaking. Since we want to generate the MSSM at low energy as an effective theory we suppose that ϕ_1, ϕ_5 and $\tilde{N}_1^0$ are decoupled at low energies too. Therefore we only deal with a Higgs potential which involves H_1, H_2 and $\tilde{N}_5^0$. It is given by

$$V' = \frac{g^2}{8}\left\{\frac{4}{3}((H_1^\dagger H_1)^2 + (H_2^\dagger H_2)^2 + (\tilde{N}_5^2)^2) + 4(H_1^\dagger H_2)(H_2^\dagger H_1) - \frac{4}{3}[(H_1^\dagger H_1)(H_2^\dagger H_2) + (H_1^\dagger H_1)\tilde{N}_5^2 + (H_2^\dagger H_2)\tilde{N}_5^2]\right\}$$

$$+\frac{g_1^2}{18}\{\tilde{N}_5^2 + (H_2^\dagger H_2) - 2(H_1^\dagger H_1)\}^2 + h_1^2(H_2^\dagger H_2)\tilde{N}_5^2 + h_2^2\{(H_2^\dagger H_2)\tilde{N}_5^2 + (H_1^\dagger H_1)(H_2^\dagger H_2) - (H_1^\dagger H_2)(H_2^\dagger H_1) + (H_1^\dagger H_1)\tilde{N}_5^2\}$$

$$+\mu_1^2\tilde{N}_5^2 + \mu_2^2(H_2^\dagger H_2) + m_1^2(H_1^\dagger H_1) + m_2^2(H_2^\dagger H_2) + m_5^2\tilde{N}_5^2. \tag{17}$$

As it has been already emphasized [14], in the MSSM the quartic scalar couplings of the Higgs potential are completely determined in terms of the two gauge couplings, but it is not the case if the symmetry $SU(2)_L \otimes U(1)_Y$ is a remnant of a larger symmetry which is broken at a higher mass scale together with the SUSY. The scalar particle content needed to produce the spontaneous symmetry breaking determines the structure of the Higgs potential. In this way, the reduced Higgs potential would be a 2HDM-like potential, but its quartic couplings would not be those of the MSSM, instead they will be related to the gauge couplings of the larger theory and to the couplings appearing in its superpotential. Analyses of supersymmetric theories in this context have been given in the literature [11,14,18]. In particular, it has been studied widely for different versions of the left-right model, SUSY version of the $SU(3)_L \otimes U(1)_X$ model where exotic charged particles of electric charges $(-4/3,5/3)$ appear [17] and the version where $SU(3)_L \otimes U(1)_X$ is an E_6 subgroup [11].

Following this idea with the reduced Higgs potentials presented in the previous paragraph, we can obtain the effective quartic scalar couplings λ_i of the most general 2HDM potential. Since there are cubic interactions in V' involving $H_{1,2}$ and $\tilde{N}_5^0$, it generates two types of Feynman diagrams which contribute to the quartic couplings [11]. The Feynman rules from the potential for these couplings are

$$H_1 - H_1 - \tilde{N}_5^0 \rightarrow i\left(2h_2^2 - \frac{1}{9}(3g^2 + 4g_1^2)\right)u,$$

$$H_2 - H_2 - \tilde{N}_5^0 \rightarrow i\left(2(h_1^2 + h_2^2) - \frac{1}{9}(3g^2 - 2g_1^2)\right)u \tag{18}$$

and using them we obtain the effective couplings; thus they are given by

$$\frac{\lambda_1}{2} = \frac{2g_1^2}{9} + \frac{g^2}{6} - \frac{3}{8}\left(\left(\frac{4g_1^2}{9} + \frac{g^2}{3}\right)^2\right.$$

$$\left. - 4h_2^2\left(\frac{4g_1^2}{9} + \frac{g^2}{3}\right) + 4h_2^4\right)G^{-1},$$

$$\frac{\lambda_2}{2} = \frac{g_1^2}{18} + \frac{g^2}{6} - \frac{3}{8}\left(\left(\frac{2g_1^2}{9} - \frac{g^2}{3}\right)^2\right.$$

$$+ 4(h_1^2 + h_2^2)\left(\frac{2g_1^2}{9} - \frac{g^2}{3}\right)$$

$$\left. + 4(h_1^2 + h_2^2)^2\right)G^{-1},$$

$$\lambda_3 = -\frac{2g_1^2}{9} - \frac{g^2}{6} + \frac{3}{4}\left(\left(\frac{4g_1^2}{9} + \frac{g^2}{3}\right)\left(\frac{2g_1^2}{9} - \frac{g^2}{3}\right)\right.$$

$$- 4h_2^2(h_1^2 + h_2^2) + 2h_2^2\left(\frac{g^2}{3} - \frac{2g_1^2}{9}\right)$$

$$\left. + 2(h_1^2 + h_2^2)\left(\frac{g^2}{3} + \frac{4g_1^2}{9}\right)\right)G^{-1},$$

$$\lambda_4 = \frac{g^2}{2} - h_2^2,$$

$$\lambda_5 = 0 \tag{19}$$

where $G = (g_1^2/3 + g^2)$. Additionally, we realize that in the limit of $h_1 = 0$ and $h_2 = h_e$ we recover the $SU(3)_L \otimes U(1)_X$ model previously worked out [11].

We want to remark that this SUSY model has the MSSM as an effective theory when the new physics is no longer there, that means $h_i = 0$, $i = 1, 2$, and the coupling constants are running down to the electroweak scale. At this point we use the approach where the $SU(2)_L$ coupling behaves like g. In the limit $h_1 = h_2 = 0$, we obtain

$$\lambda_{1,2} = \frac{g^2(4g_1^2 + 3g^2)}{4(g_1^2 + 3g^2)}, \quad \lambda_3 = -\frac{g^2(4g_1^2 + 3g^2)}{4(g_1^2 + 3g^2)}, \quad \lambda_4 = \frac{g^2}{2} \tag{20}$$

and, if we assume the matching condition from Eq. (6), we reduce the effective couplings to those appearing in the MSSM, as expected,

$$\lambda_1 = \lambda_2 = \frac{1}{4}(g^2 + g'^2), \quad \lambda_3 = -\frac{1}{4}(g^2 + g'^2), \quad \lambda_4 = \frac{g^2}{2}. \tag{21}$$

But when new physics is present, the conditions coming from the analysis of the Higgs potential on the new parameters h_i are quite different. The requirement that the potential be bounded from below, implies that $\lambda_1 \geqslant 0$, $\lambda_2 \geqslant 0$ and either $\lambda_3 + \lambda_4 \geqslant 0$ or if $\lambda_3 + \lambda_4 < 0$ then $|\lambda_3 + \lambda_4| \leqslant \sqrt{\lambda_1 \lambda_2}$, this is not a simple region, and therefore in order to get a bound
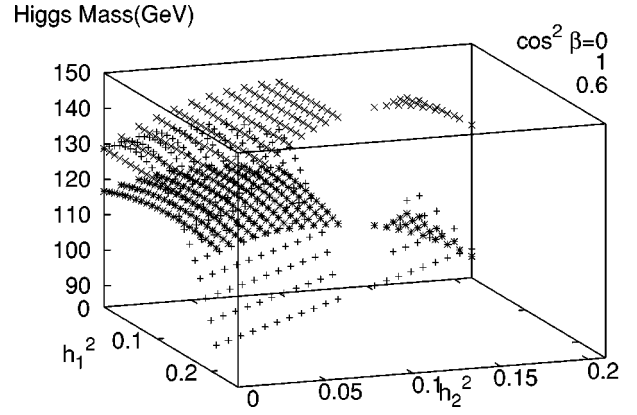

FIG. 1. The upper bound on the lightest *CP*-even Higgs boson as a function of the parameters h_i^2 from the superpotential. The different kind of points correspond to $\cos^2\beta = \{0, 0.6, 1\}$.

on the lightest Higgs boson, we evaluate numerically point by point the conditions, and the results are shown in Fig. 1. Furthermore, we have added the term coming from radiative corrections due to the top quark and its two supersymmetric partners [6,8,14], taking $m_t = 174.3$ GeV [19] and $\tilde{m} = 1$ TeV. Figure 1 is a 3D plot for three different values of $\cos^2\beta$ where $\tan\beta \equiv \langle H_2 \rangle_0 / \langle H_1 \rangle_0$, and the lowest bound of MSSM can be moved up to 143 GeV. The range of values for h_1^2 and h_2^2 are presented in such a way that they satisfy the inequalities coming from the requirement that the potential be bounded from below, and we note that for $h_1^2 = 0$ the largest value for h_2^2 is 0.21.

On the other hand, the upper bound of the Higgs boson mass has been calculated driving the VEV's of the SM singlets in such way that only one VEV has been included, but we can explore the behavior of this bound considering more singlets in the model which have couplings in the same form that the first one $\tilde{N}_5^0$. Doing that, we note that the upper bound of the lightest Higgs boson is not affected pretty much, but the constraints on the parameters h_i, are. In Fig. 2, we show the plane h_1^2-h_2^2 considering one, three and ten singlets in the model. The figure implies that the new physics
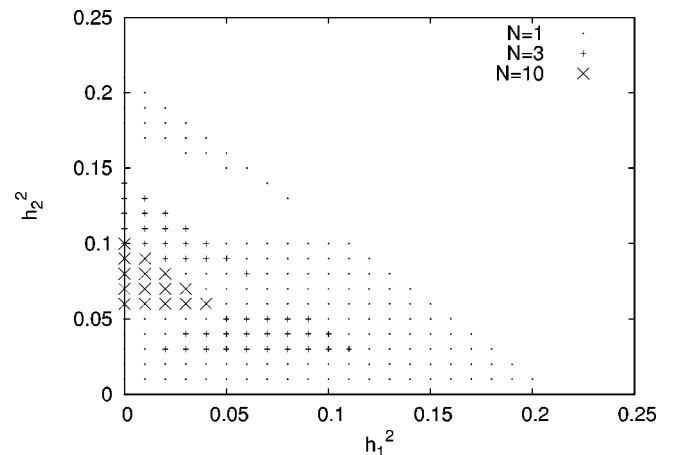

FIG. 2. The allowed region for the plane h_1^2-h_2^2 with additional SM singlets, for $N = 1$, $N = 3$ and $N = 10$.

related to the parameters h_i is more constrained when the number of singlets involved increases.

IV. CONCLUSIONS

We have presented a new supersymmetric version of the symmetry $SU(3)_L \otimes U(1)_X$ which is a $SU(6)_L \otimes U(1)_X$ subgroup, where the Higgs bosons are in the scalar sector of the superfields, and it is triangle anomaly free. We have also shown that using the limit when the parameter $h_i = 0$ and the matching condition [Eq. (6)], we obtain the SUSY constrains for the Higgs potential as in the MSSM. Therefore, if we analyze the upper bound for the mass of the lightest CP-even Higgs boson in this limit, we find the same bound of around 128 GeV for the MSSM. However, since in general $h_i = 0$, $i = 1,2$, such upper bound can be moved up to around 143 GeV. This fact can be an interesting alternative to take into account in the search for the lightest CP-even SUSY Higgs boson mass. This model is different from the one considered in [11] and only can be recovered when $h_1 = 0$ and $h_2 = h_e$. By considering the experimental value from LEP for the lightest Higgs boson mass $m_h \geq 114$ GeV it is possible to get bounds for the parameters of the model h_i^2; they are $0 \leq h_1^2 \leq 0.21$, $0 \leq h_2^2 \leq 0.21$. But if we include more singlets of the SM symmetry like $\tilde{N}_5^0$, the upper bound on the Higgs boson mass does not move considerably, only up to 145 GeV, but the allowed values in the plane h_1^2-h_2^2 do move. These allowed values are reduced when the number of singlets added are increased.

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