Triplet-singlet extension of the MSSM with a 125 GeV Higgs boson and dark matter

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We study the extension of the minimal supersymmetric standard model by adding one singlet and one hypercharge zero $SU(2)$ triplet chiral superfield. The triplet sector gives an additional contribution to the scalar masses, and we find that the lightest CP-even Higgs boson can have a mass of 119–120 GeV at tree level, and radiative correction raises the value to 125 GeV. In this model no significant contributions from stop loops is needed to get the required Higgs mass that alleviates the fine-tuning problem of fixing the stop mass to a high precision at the grand unified theory scale. In addition, this model gives a neutralino dark matter of mass around 100 GeV that is a mixture of Higgsino and triplino with a dark matter density consistent with WMAP observations. The spin-independent scattering cross section with nucleons is 10^{-43} cm^2, which makes it consistent with the bounds from direct detection experiments like XENON100 and others.

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I. INTRODUCTION

The ATLAS and CMS collaborations [1,2] have narrowed down the allowed range of a light Higgs mass to the region 115–131 GeV. In addition there are hints of the Higgs mass being near $m_h = 125$ GeV with standard model (SM)-like decay widths into 2γ and $4l$. A light Higgs is favored in supersymmetry although the minimal supersymmetric standard model (MSSM) predicts a tree-level upper bound on the lightest CP-even Higgs mass as $m_h < M_Z \cos 2\beta$. Within MSSM, loop corrections can give required large corrections to the Higgs mass provided the stop is heavier than 1 TeV or there is near maximal stop mixing. Implications of the 125 GeV Higgs for the MSSM and constrained-MSSM parameter space have been extensively studied [3]. Going beyond MSSM, in order to get a larger tree-level Higgs mass, the simplest extension is a singlet superfield in the next-to-minimal supersymmetric standard model (NMSSM) model [4]. The singlet interaction with the two Higgs doublet of MSSM is via the $\lambda S H_u \cdot H_d$ term. The Higgs mass is now given by the relation $m_h^2 = M_Z^2 \cos^2 2\beta + \lambda^2 v^2 \sin^2 2\beta + \delta m_h^2$, where δm_h^2 is due to radiative correction. Taking $\lambda = 0.7$ [larger values would make it flow to the nonperturbative regime much below the grand unified theory (GUT) scale] and $\tan\beta = 2$, the radiative correction needed to get a 125 GeV Higgs mass is $\delta m_h = 55$ GeV, which is an improvement over the $\delta m_h = 85$ GeV needed in the MSSM. However, fine-tuning of the stop mass is still required in NMSSM to get the required Higgs mass [5]. Also by extending the MSSM gauge group in a suitable way, the new Higgs sector dynamics can push the tree-level mass well above the tree-level MSSM limit if it couples to the new gauge sector [6]. In most of the cases the nondecoupling D-terms contribute nontrivially to increase the tree-level mass of the SM-like Higgs boson. Recent analysis of the supersymmetric (SUSY) model based on $SU(3)_C \times SU(2)_L \times U(1)_R \times U(1)_{B\text{-}L}$ gauge group [7] has shown that the tree-level physical Higgs mass can be at most 110 GeV and through the one-loop correction it can be raised considerably. Another recent work on MSSM extended by a $U(1)$ gauged Peccei-Quinn symmetry [8] where the new D-terms can raise the tree-level mass well enough to accommodate the 125 GeV Higgs boson without significant radiative correction and hence requires less fine-tuning.

An important aspect of the 125 GeV Higgs mass is that the parameter space of the thermal relic for dark matter is severely restricted. In the MSSM, the lightest supersymmetric particle (LSP) is a Higgsino at the TeV scale [9]. In NMSSM, the SUSY partner of the singlet scalar—the singlino—mixes with the neutralinos to provide a light dark matter [10,11]. Recent analysis [5] has shown that the benchmark parameters that give a 125 GeV Higgs mass also provide a neutralino dark matter candidate with mass in the range of 68–85 GeV. To our knowledge, the dark matter in triplet-extended MSSM has not been studied so far.

The extension of MSSM by extending it with $Y = 0$ and $Y = 0, \pm 1$ $SU(2)$ triplet superfields has been studied [12–14] where the tree-level contribution to the Higgs mass from the triplet Higgs sector has been calculated. It was shown in Ref. [14] that with the $Y = 0$ triplet superfield the tree-level Higgs mass can be raised to 113 GeV, which would still require substantial loop corrections from stops. Recently, the MSSM extended by two real triplets ($Y = \pm 1$) and one singlet [15] has been studied with a motivation to solve the μ-problem as well as to obtain a large correction to the lightest Higgs mass. The analysis of the dark matter sector of this model will be complicated as the LSP will be the lightest eigenstate of

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the 7×7 neutralino mass matrix that has not yet been done.

In this paper we explore the minimal extension of the MSSM that can give a tree-level Higgs mass of 119–120 GeV. We find that by extending the MSSM by adding a singlet superfield and a $Y = 0$ $SU(2)$ triplet superfield, this aim can be achieved. The upper bound on the tree-level mass of the lightest CP-even Higgs mass is given in Eq. (27). With this tree-level Higgs mass the stop mass need not be very heavy and this solves the fine-tuning problem of the Higgs mass in MSSM and NMSSM [5]. We also study the dark matter candidates in this model that are obtained by diagonalizing the 6×6 neutralino mass matrix. We find a viable dark matter with mass 100 GeV, which is a mixture of the Higgsino and triplino (the fermionic partner of the neutral component of the triplet Higgs mass). We fix two sets of benchmark parameters at the electroweak scale that would give the 125 GeV and dark matter relic density $\Omega h^2 = 0.1109 \pm 0.0056$ compatible with WMAP-7 measurements [16]. We find that the direct detection cross section of the dark matter is $\sigma_{\rm SI} \simeq 10^{-43}$ cm^2, which is compatible with the direct detection experiments like XENON100 [17].

In Sec. II we display the superpotential of our model, and we derive the scalar potential from the D-terms and F-terms and from the various soft-breaking terms. In Sec. III we give a detailed analysis of the Higgs sector, and we calculate the CP-even, CP-odd, and charged Higgs mass matrices. In Sec. IV, the neutralino and the chargino mass matrices are discussed. The numerical results based on this model are discussed in detail in Sec. V. We show the results for two sets of benchmark points that include the parameters like couplings, trilinear soft-breaking terms, soft masses, and the fermionic and scalar mass spectrum. We have also taken into account the one-loop corrections to the lightest physical Higgs mass and shown a quantitative improvement of the level of fine-tuning compared to other models. In Sec. VI we discuss the dark matter from the neutralino sector of this model and its phenomenology, which is one of the main results of this paper. In the concluding section we summarize the results and point out some directions for further study of the triplet-singlet model that will enable the model to be tested at the LHC.

II. MODEL

In this model, we have extended the superpotential of the minimal supersymmetric standard model by adding one singlet chiral superfield S and one $SU(2)$ triplet chiral superfields T_0 with hypercharge $Y = 0$. The most general form of the superpotential for this singlet-triplet-extended model can be written as

$$\mathcal{W} = (\mu + \lambda\hat{S})\hat{H}_d.\hat{H}_u + \frac{\lambda_1}{3}\hat{S}^3 + \lambda_2\hat{H}_d.\hat{T}_0\hat{H}_u + \lambda_3\hat{S}^2\,\mathrm{Tr}(\hat{T}_0) + \lambda_4\hat{S}\,\mathrm{Tr}(\hat{T}_0\hat{T}_0) + W_{\rm Yuk}, \qquad (1)$$

where $\hat{H}_{u,d}$ are the Higgs doublets of the MSSM and the Yukawa superpotential $W_{\rm Yuk}$ is given as

$$W_{\rm Yuk} = y_u\hat{Q}_L.\hat{H}_u\hat{U}_R + y_d\hat{Q}_L.\hat{H}_d\hat{D}_R + y_e\hat{L}_L.\hat{H}_d\hat{E}_R. \qquad (2)$$

In terms of the components, we have

$$\hat{H}_u = \begin{pmatrix} \hat{H}_u^+ \\ \hat{H}_u^0 \end{pmatrix}, \qquad \hat{H}_d = \begin{pmatrix} \hat{H}_d^0 \\ \hat{H}_d^- \end{pmatrix} \quad \text{and} \quad \hat{T}_0 = \begin{pmatrix} \frac{\hat{T}^0}{\sqrt{2}} & -\hat{T}_0^+ \\ \hat{T}_0^- & \frac{-\hat{T}^0}{\sqrt{2}} \end{pmatrix}.$$

Here, $(\hat{T}_0^-)^* \neq -\hat{T}_0^+$, which would not have been true for a real Higgs triplet in nonsupersymmetric models. We can solve the μ-problem by starting with a scale invariant superpotential, given as

$$W_{\rm sc\,inv} = \lambda\hat{S}\hat{H}_d.\hat{H}_u + \frac{\lambda_1}{3}\hat{S}^3 + \lambda_2\hat{H}_d.\hat{T}_0\hat{H}_u + \lambda_4\hat{S}\,\mathrm{Tr}(\hat{T}_0\hat{T}_0) + W_{\rm Yuk}, \qquad (3)$$

where the $SU(2)$ invariant dot product is defined as

$$\hat{H}_d.\hat{T}_0\hat{H}_u = \frac{1}{\sqrt{2}}(\hat{H}_d^0\hat{T}^0\hat{H}_u^0 + \hat{H}_d^-\hat{T}^0\hat{H}_u^+) - (\hat{H}_d^0\hat{T}_0^-\hat{H}_u^+ + \hat{H}_d^-\hat{T}_0^+\hat{H}_u^0). \qquad (4)$$

This superpotential (3) also has an accidental Z_3-symmetry, i.e., invariance of the superpotential on multiplication of the chiral superfields by the factor of $\frac{2\pi i}{3}$. By this choice we are eliminating the μ-parameter but an effective μ-term is generated when the neutral components of S and T_0 acquire vacuum expectation values (vev's) v_s and v_t, respectively,

$$\mu_{\rm eff} = \lambda v_s - \frac{\lambda_2}{\sqrt{2}}v_t. \qquad (5)$$

Therefore, in terms of the neutral components of the superfields, Eq. (3) sans $W_{\rm Yuk}$ can be rewritten as

$$W^{\rm neu} = -\lambda\hat{S}\hat{H}_u^0\hat{H}_d^0 + \frac{\lambda_1}{3}\hat{S}^3 + \frac{\lambda_2}{\sqrt{2}}\hat{H}_d^0\hat{T}^0\hat{H}_u^0 + \lambda_4\hat{S}\hat{T}^0\hat{T}^0. \qquad (6)$$

A. Scalar potential

The scalar potential involving only the Higgs field can be written as

$$V = V_{\rm SB} + V_F + V_D. \qquad (7)$$

In the above equation, $V_{\rm SB}$ consists of the soft-supersymmetry breaking term associated with the superpotential in Eq. (3), given by

$$V_{\rm SB} = m_{H_u}^2[|H_u^0|^2 + |H_u^+|^2] + m_{H_d}^2[|H_d^0|^2 + |H_d^-|^2] + m_S^2|S|^2 + m_T^2 {\rm Tr}(T_0^\dagger T_0) + \Big(-\lambda A_\lambda S H_u.H_d + \frac{\lambda_1}{3}A_{\lambda_1}S^3 + \lambda_2 A_{\lambda_2} H_d.T_0 H_u + \lambda_4 B_\lambda S {\rm Tr}(T_0^2) + {\rm H.c.}\Big). \quad (8)$$

In Eq. (7) V_F is the supersymmetric potential from F-terms, given by

$$\begin{aligned} V_F = &\left|-\lambda S H_d^0 + \frac{\lambda_2}{\sqrt{2}}H_d^0 T^0 - \lambda_2 H_d^- T_0^+\right|^2 + \left|-\lambda S H_u^0 + \frac{\lambda_2}{\sqrt{2}}H_u^0 T^0 - \lambda_2 H_u^+ T_0^-\right|^2 \\ &+ \left|\frac{\lambda_2}{\sqrt{2}}(H_u^0 H_d^0 + H_d^- H_u^+) + 2\lambda_4 S T^0\right|^2 + \left|\lambda(H_d^- H_u^+ H_u^0 H_d^0) + \lambda_1 S^2 + \lambda_4(T^{0^2} - 2T_0^+ T_0^-)\right|^2 \\ &+ \left|\lambda S H_d^- + \frac{\lambda_2}{\sqrt{2}}T^0 H_d^- - \lambda_2 H_d^0 T_0^-\right|^2 + \left|-\lambda_2 H_d^- H_u^0 - 2\lambda_4 S T_0^-\right|^2 \\ &+ \left|\lambda S H_u^+ + \frac{\lambda_2}{\sqrt{2}}T^0 H_u^+ - \lambda_2 H_u^0 T_0^+\right|^2 + \left|-\lambda_2 H_u^+ H_d^0 - 2\lambda_4 S T_0^+\right|^2, \end{aligned} \quad (9)$$

whereas the F-term for the neutral scalar potential can be derived from Eq. (6) as

$$V_{F_{\rm neu}} = \sum_i \left|\frac{\partial W_{\rm scalar}^{\rm neu}}{\partial \phi_i^0}\right|^2, \quad (10)$$

where ϕ_i^0 stands for H_u^0, H_d^0, S, T^0 and $W_{\rm scalar}^{\rm neu}$ is the scalar counterpart of the neutral superpotential $W^{\rm neu}$.

Finally, V_D is supersymmetric potential from D-terms in Eq. (7), given by

$$\begin{aligned} V_D = &\frac{g_1^2}{8}[|H_d^-|^2 + |H_d^0|^2 - |H_u^+|^2 - |H_u^0|^2]^2 + \frac{g_2^2}{8}[|H_d^-|^2 + |H_d^0|^2 - |H_u^+|^2 - |H_u^0|^2 + 2|T_0^+|^2 - 2|T_0^-|^2]^2 \\ &+ \frac{g_2^2}{8}[H_d^{0*}H_d^- + H_u^{+*}H_u^0 + \sqrt{2}(T_0^+ + T_0^-)T_0^* + {\rm H.c}]^2 - \frac{g_2^2}{8}[H_d^{-*}H_d^0 + H_u^{0*}H_u^+ + \sqrt{2}(T_0^+ - T_0^-)T_0^* + {\rm H.c}]^2. \end{aligned} \quad (11)$$

1. EWSB

After electroweak symmetry breaking, only the neutral components of the scalar fields acquire vev's, i.e.,

$$\langle H_u^0\rangle = v_u, \quad \langle H_d^0\rangle = v_d, \quad \langle S\rangle = v_s, \text{ and } \langle T^0\rangle = v_t.$$

The neutral scalar part of the chiral superfields can be decomposed into real and imaginary parts,

$$H_u^0 = (H_{u_R}^0 + v_u) + iH_{u_I}^0, \quad (12)$$

$$H_d^0 = (H_{d_R}^0 + v_d) + iH_{d_I}^0, \quad (13)$$

$$S = (S_R + v_s) + iS_I, \quad (14)$$

$$T^0 = (T_R^0 + v_t) + iT_I^0. \quad (15)$$

The minimization conditions are derived from the fact that

$$\frac{\partial V}{\partial v_u} = \frac{\partial V}{\partial v_d} = \frac{\partial V}{\partial v_s} = \frac{\partial V}{\partial v_t} = 0. \quad (16)$$

We can determine the soft-breaking mass parameters like $m_{H_u}^2$, $m_{H_d}^2$, m_T^2, and m_S^2 using the following minimization conditions:

$$\begin{aligned} m_{H_u}^2 = &\cot\beta\Big[A_{\rm eff} - \Big(\lambda^2 + \frac{\lambda_2^2}{2}\Big)\frac{v^2}{2}\sin 2\beta \\ &+ \lambda\lambda_4 v_t^2 - \sqrt{2}\lambda_2\lambda_4 v_t v_s - \frac{\lambda_2}{\sqrt{2}}A_{\lambda_2}v_t\Big] - \mu_{\rm eff}^2 \\ &+ \frac{1}{4}(g_1^2 + g_2^2)v^2\cos 2\beta, \end{aligned} \quad (17)$$

$$\begin{aligned} m_{H_d}^2 = &\tan\beta\Big[A_{\rm eff} - \Big(\lambda^2 + \frac{\lambda_2^2}{2}\Big)\frac{v^2}{2}\sin 2\beta \\ &+ \lambda\lambda_4 v_t^2 - \sqrt{2}\lambda_2\lambda_4 v_t v_s - \frac{\lambda_2}{\sqrt{2}}A_{\lambda_2}v_t\Big] \\ &- \mu_{\rm eff}^2 - \frac{1}{4}(g_1^2 + g_2^2)v^2\cos 2\beta, \end{aligned} \quad (18)$$

$$\begin{aligned} m_S^2 = &v^2\Big[\frac{v_t}{\sqrt{2}v_s}\lambda\lambda_2 + \lambda\lambda_1\sin 2\beta + \frac{1}{2v_s}\lambda A_\lambda \sin 2\beta - \lambda^2\Big] \\ &- [2\lambda_1^2 v_s + \lambda A_{\lambda_1}]v_s - \lambda_4 v_t^2[B_\lambda/v_s + 2\lambda_1 + 4\lambda_4] \\ &- \sqrt{2}\lambda_2\lambda_4 v_u v_d v_t/v_s, \end{aligned} \quad (19)$$

$$\begin{aligned} m_T^2 = &\Big[\frac{1}{\sqrt{2}}\lambda\lambda_2\frac{v_s}{v_t} - \frac{\lambda_2^2}{2} - \frac{\lambda_2}{2\sqrt{2}v_t}A_{\lambda_2}\sin 2\beta\Big]v^2 \\ &- 2\lambda_4^2 v_t^2 + 2\lambda\lambda_4 v_u v_d - \lambda_4 v_s^2[2B_\lambda/v_s + 2\lambda_1 + 4\lambda_4] \\ &- \sqrt{2}\lambda_2\lambda_4 v_u v_d v_s/v_t, \end{aligned} \quad (20)$$

where

$$A_{\rm eff} = \lambda v_s[A_\lambda + \lambda_1 v_s], \tag{21}$$

and $v_u^2 + v_d^2 = v^2 = (174)^2$ GeV^2, $\tan\beta = \frac{v_u}{v_d}$.

Because of the addition of the triplets, the gauge bosons receive additional contribution in their masses like

$$M_Z^2 = \frac{1}{2}(g_1^2 + g_2^2)v^2, \tag{22}$$

$$M_W^2 = \frac{1}{2}g_2^2(v^2 + 4v_t^2). \tag{23}$$

The ρ-parameter at the tree level is defined as

$$\rho = \frac{M_W^2}{M_Z^2\cos^2\theta_W} = 1 + 4\frac{v_t^2}{v^2}. \tag{24}$$

Clearly, the ρ-parameter deviates from unity by a factor of $4\frac{v_t^2}{v^2}$. Using the recent bound on the ρ-parameter at 95% C. L., we can determine the bound on the triplet Higgs vev v_t. ρ can be confined in the range 0.9799–1.0066 [13] and hence $v_t \lesssim 9$ GeV at 95% C. L.

III. HIGGS SECTOR

A. *CP*-even Higgs mass matrices

The symmetric *CP*-even Higgs mass matrix is written in the basis $(H_{u_R}^0, H_{d_R}^0, T_R^0, S_R)$ with 10 independent components. After electroweak symmetry breaking (EWSB) the entries of the squared mass matrix are

$$\begin{aligned}
M_{11}^2 &= \frac{1}{2}(g_1^2 + g_2^2)v^2\sin^2\beta + C_1\cot\beta + C_4,\\
M_{22}^2 &= \frac{1}{2}(g_1^2 + g_2^2)v^2\cos^2\beta + C_1\tan\beta + C_4,\\
M_{33}^2 &= 4\lambda_4^2 v_t^2 + \lambda_2 v^2[\lambda v_s - (A_{\lambda_2} + 2\lambda_4 v_s)\\
&\quad \times \sin\beta\cos\beta]/\sqrt{2}v_t,\\
M_{44}^2 &= \lambda_1 v_s[A_{\lambda_1} + 4\lambda_1 v_s] + \left[v_t\left(\lambda\lambda_2\frac{v^2}{\sqrt{2}} - \lambda_4 B_\lambda v_t\right)\right.\\
&\quad \left. + (\lambda A_\lambda - \sqrt{2}\lambda_2\lambda_4 v_t)v^2\sin\beta\cos\beta\right]\Big/ v_s,\\
M_{12}^2 &= -C_1 + \left[2\lambda^2 + \lambda_2^2 - \frac{(g_1^2 + g_2^2)}{2}\right]v^2\sin\beta\cos\beta,\\
M_{13}^2 &= v[C_2\cos\beta - \sqrt{2}\lambda_2\mu_{\rm eff}\sin\beta],\\
M_{14}^2 &= -v[C_3\cos\beta - 2\lambda\mu_{\rm eff}\sin\beta],\\
M_{23}^2 &= v[C_2\sin\beta - \sqrt{2}\lambda_2\mu_{\rm eff}\cos\beta],\\
M_{24}^2 &= -v[C_3\sin\beta - 2\lambda\mu_{\rm eff}\cos\beta],\\
M_{34}^2 &= 2\lambda_4 v_t[B_\lambda + 2v_s(\lambda_1 + 2\lambda_4)]\\
&\quad - \lambda_2 v^2(\lambda - 2\lambda_4\sin\beta\cos\beta)/\sqrt{2},
\end{aligned} \tag{25}$$

where C_i's are defined as

$$\begin{aligned}
C_1 &= A_{\rm eff} + \lambda\lambda_4 v_t^2 - \lambda_2 A_{\lambda_2}\frac{v_t}{\sqrt{2}} - \sqrt{2}\lambda_2\lambda_4 v_t v_s,\\
C_2 &= \frac{\lambda_2 A_{\lambda_2}}{\sqrt{2}} - 2\lambda\lambda_4 v_t + \sqrt{2}\lambda_4\lambda_2 v_s,\\
C_3 &= \lambda A_\lambda + 2\lambda\lambda_1 v_s - \sqrt{2}\lambda_2\lambda_4 v_t,\\
C_4 &= \lambda_2 v_t\left[\frac{\lambda_2 v_t}{2} - \sqrt{2}\lambda v_s\right],
\end{aligned} \tag{26}$$

and $A_{\rm eff}$ is defined in Eq. (21).

1. Bound on the lightest Higgs mass

The bound on the lightest Higgs mass is derived from the fact that the smallest eigenvalue of a real, symmetric $n \times n$ matrix is smaller than the smallest eigenvalue of the upper left 2×2 submatrix [12]. Using this we obtain an upper bound on the lightest *CP*-even Higgs mass,

$$m_h^2 \le M_Z^2\left[\cos^2 2\beta + \frac{2\lambda^2}{g_1^2 + g_2^2}\sin^2 2\beta + \frac{\lambda_2^2}{g_1^2 + g_2^2}\sin^2 2\beta\right]. \tag{27}$$

The bound on lightest Higgs mass has been considerably improved over the MSSM due to the additional contribution from the singlet and triplet gauge fields. Using Eq. (27) we can put constraints on the parameters like λ, λ_2 and $\tan\beta$ satisfying the recent bound on Higgs mass from ATLAS and CMS.

B. *CP*-odd Higgs mass matrices

The elements of the 4×4 *CP*-odd Higgs squared mass matrix, after EWSB, in the basis $(H_{d_I}^0, H_{u_I}^0, S_I, T_I^0)$ are

$$\begin{aligned}
M_{P_{11}}^2 &= C_1\tan\beta + C_4, \qquad M_{P_{22}}^2 = C_1\cot\beta + C_4,\\
M_{P_{33}}^2 &= -3\lambda_1 A_{\lambda_1}v_s - \lambda_4[B_\lambda + 4\lambda_1 v_s]\frac{v_t^2}{v_s} + D_1\left(\frac{v_t}{v_s}\right)\\
&\quad + [\lambda A_\lambda/v_s + 4\lambda\lambda_1]v^2\sin\beta\cos\beta,\\
M_{P_{44}}^2 &= -4\lambda_4 v_s[B_\lambda + \lambda_1 v_s] + D_1\left(\frac{v_s}{v_t}\right)\\
&\quad + \left[4\lambda\lambda_4 - \frac{1}{\sqrt{2}v_t}\lambda_2 A_{\lambda_2}\right]v^2\sin\beta\cos\beta,\\
M_{P_{12}}^2 &= A_{\rm eff} - \frac{v_t}{\sqrt{2}}\lambda_2 A_{\lambda_2} + \lambda_4 v_t[\lambda v_t - \sqrt{2}\lambda_2 v_s],\\
M_{P_{13}}^2 &= v\sin\beta[\lambda A_\lambda - 2\lambda\lambda_1 v_s + \sqrt{2}\lambda_2\lambda_4 v_t],\\
M_{P_{14}}^2 &= -v\sin\beta\left[2\lambda\lambda_4 v_t + \frac{1}{\sqrt{2}}\lambda_2(A_{\lambda_2} - 2\lambda_4 v_s)\right],\\
M_{P_{23}}^2 &= M_{P_{13}}^2/\tan\beta, \qquad M_{P_{24}}^2 = M_{P_{14}}^2/\tan\beta,\\
M_{P_{34}}^2 &= -2\lambda_4 v_t(B_\lambda - 2\lambda_1 v_s) - D_1,
\end{aligned} \tag{28}$$

where

$$D_1 = \frac{1}{\sqrt{2}}\lambda_2 v^2(\lambda + 2\lambda_4 \sin\beta\cos\beta). \tag{29}$$

This matrix always contains a Goldstone mode G^0 (gives mass to Z boson), which can be written as

$$G^0 = \cos\beta H^0_{d_I} - \sin\beta H^0_{u_I}, \tag{30}$$

and we rotate the mass matrix in the basis (G^0, A_1, A_2, A_3) where

$$\begin{pmatrix} A_1 \\ G^0 \\ A_2 \\ A_3 \end{pmatrix} = \begin{pmatrix} \cos\beta & \sin\beta & 0 & 0 \\ -\sin\beta & \cos\beta & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & 1 \end{pmatrix} \begin{pmatrix} H^0_{u_I} \\ H^0_{d_I} \\ S_I \\ T^0_I \end{pmatrix}. \tag{31}$$

After removing the Goldstone mode, we again rotate the remaining 3×3 mass matrix and finally obtain

$$\begin{aligned} P_1 &= \cos\alpha\sin\beta H_{d_I} + \cos\alpha\cos\beta H_{u_I} + \sin\alpha S_I, \\ P_1 &= -\sin\alpha\sin\beta H_{d_I} - \sin\alpha\cos\beta H_{u_I} + \cos\alpha S_I, \\ P_3 &= T_I, \end{aligned} \tag{32}$$

where P_1, P_2, P_3 are the massive modes.

C. Charged Higgs mass matrices

The charged Higgs sector comprises of a 4×4 symmetric matrix, written in the basis $(H_u^+, H_d^{-^*}, T_0^+, T_0^{-^*})$, which has 10 independent components (after EWSB) given by

$$\begin{aligned}
(M^2_\pm)_{11} &= E_1 v_d^2 + \left[\sqrt{2}\lambda\lambda_2 v_t v_s + \frac{\lambda_2^2}{2} v_t^2\right] + C_1 \cot\beta, \\
(M^2_\pm)_{12} &= A_{\rm eff} + E_1 v_u v_d + [\lambda_2 A_{\lambda_2} + \sqrt{2}\lambda v_t + 2\lambda_2 v_s]\frac{v_t}{\sqrt{2}}, \\
(M^2_\pm)_{13} &= E_2 v_d - 2\lambda_2 v_u\left[\lambda v_s + \frac{\lambda_2 v_t}{\sqrt{2}}\right], \\
(M^2_\pm)_{14} &= E_3 v_d + v_u \lambda_2 \mu_{\rm eff}, \\
(M^2_\pm)_{22} &= E_1 v_u^2 + \left[\sqrt{2}\lambda\lambda_2 v_t v_s + \frac{\lambda_2^2}{2} v_t^2\right] + C_1 \tan\beta, \\
(M^2_\pm)_{23} &= E_3 v_u + v_d \lambda_2 \mu_{\rm eff}, \\
(M^2_\pm)_{24} &= E_2 v_u - 2\lambda_2 v_d\left[\lambda v_s + \frac{\lambda_2 v_t}{\sqrt{2}}\right], \\
(M^2_\pm)_{33} &= \frac{g_2^2}{2}[v_u^2 - v_d^2] + \lambda_2^2 v_u^2 + E_4, \\
(M^2_\pm)_{34} &= [g_2^2 - 2\lambda_4^2]v_t^2 - 2\lambda_4 v_s[B_\lambda + \lambda_1 v_s] + 2\lambda\lambda_4 v_u v_d, \\
(M^2_\pm)_{44} &= \frac{g_2^2}{2}[v_d^2 - v_u^2] + \lambda_2^2 v_d^2 + E_4,
\end{aligned} \tag{33}$$

where E_i's are defined as

$$E_1 = \frac{g_2^2}{2} - \lambda^2 + \frac{\lambda_2^2}{2}, \qquad E_2 = \frac{g_2^2 v_t}{\sqrt{2}} + 2\lambda_2\lambda_4 v_s,$$
$$E_3 = \frac{g_2^2 v_t}{\sqrt{2}} - \lambda_2 A_{\lambda_2}, \qquad E_4 = g_2^2 v_t^2 + 4\lambda_4^2 v_s^2. \tag{34}$$

After diagonalization, we obtain one massless Goldstone state G^+ (gives mass to the $W^\pm$ boson, since $G^- \equiv G^{+*}$),

$$G^+ = \sin\beta H_u^+ - \cos\beta H_d^{-*} + \sqrt{2}\frac{v_t}{v}(T_0^+ - T_0^{-*}), \tag{35}$$

and three other massive modes like $H_1^\pm$, $H_2^\pm$, $H_3^\pm$.

IV. NEUTRALINOS AND CHARGINOS

The neutralino mass matrix extended by the singlet and triplet sector, in the basis $(\tilde{B}, \tilde{W}^0, \tilde{H}_d^0, \tilde{H}_u^0, \tilde{S}, \tilde{T}^0)$ is given by

$$\mathcal{M}_{\tilde{G}} = \begin{pmatrix} M_1 & 0 & -c_\beta s_w M_Z & s_\beta s_w M_Z & 0 & 0 \\ 0 & M_2 & c_\beta c_w M_Z & -s_\beta c_w M_Z & 0 & 0 \\ -c_\beta s_w M_Z & c_\beta c_w M_Z & 0 & -\mu_{\rm eff} & -\lambda v_u & \frac{\lambda_2}{\sqrt{2}} v_u \\ s_\beta s_w M_Z & -s_\beta c_w M_Z & -\mu_{\rm eff} & 0 & -\lambda v_d & \frac{\lambda_2}{\sqrt{2}} v_d \\ 0 & 0 & -\lambda v_u & -\lambda v_d & 2\lambda_1 v_s & 2\lambda_4 v_t \\ 0 & 0 & \frac{\lambda_2}{\sqrt{2}} v_u & \frac{\lambda_2}{\sqrt{2}} v_d & 2\lambda_4 v_t & 2\lambda_4 v_s \end{pmatrix}, \tag{36}$$

where M_1, M_2 are the soft-breaking mass parameters for Bino and Wino, respectively, and

$$c_\beta = \cos\beta, \quad s_\beta = \sin\beta, \quad c_w = \cos\theta_w, \quad \text{and} \quad s_w = \sin\theta_w.$$

The left-most 4×4 entries are exactly identical with that in the MSSM, except for the $\mu_{\rm eff}$-term that is defined in Eq. (5). As the triplet and the singlet fermion do not have any interaction with the neutral gauginos, the right-most 2×2 entries are zero.

The chargino mass terms in the Lagrangian can be written as

$$-\frac{1}{2}[\tilde{G}^{+T} M_c^T.\tilde{G}^- + \tilde{G}^{-T} M_c.\tilde{G}^+], \tag{37}$$

where the basis $\tilde{G}^+$ and $\tilde{G}^-$ are specified as

$$\tilde{G}^+ = \begin{pmatrix} \tilde{W}^+ \\ \tilde{H}_u^+ \\ \tilde{T}^+ \end{pmatrix}, \qquad \tilde{G}^- = \begin{pmatrix} \tilde{W}^- \\ \tilde{H}_d^- \\ \tilde{T}^- \end{pmatrix},$$

and the chargino matrix in the gauge basis is given by

$$M_c = \begin{pmatrix} M_2 & \frac{1}{\sqrt{2}} g_2 v_d & g_2 v_t \\ \frac{1}{\sqrt{2}} g_2 v_u & \lambda v_s + \frac{\lambda_2}{\sqrt{2}} v_t & -\lambda_2 v_d \\ -g_2 v_t & \lambda_2 v_u & 2\lambda_4 v_s \end{pmatrix}. \tag{38}$$

V. RESULTS AND DISCUSSIONS

The main results of this paper are shown in Tables I and II. We have specified the values of the parameters like couplings, soft-breaking parameters at the electroweak scale. The choice of $\tan\beta$, λ, and λ_2 are restricted from the bound on lightest Higgs mass (27). In Fig. 1 we show the relation between λ_2 and λ for different values of $\tan\beta$. As we increase $\tan\beta$, λ and λ_2 tend to shift towards the higher values. The plot in the right panel of Fig. 1 shows the dependence of m_h on $\tan\beta$ for some particular choices of $\lambda = 0.6$, 0.64 and $\lambda_2 = 0.75$, 1.02, which are consistent with $m_h = 125$ GeV (shown by the dotted line). In order to satisfy the bound on the Higgs mass, we can put a constraint on $\tan\beta$, i.e., $\tan\beta \leq 3.0$. The coupling λ_1 sets the mass for the singlino through the Yukawa term $2\lambda_1 S \chi_s \cdot \chi_s$. In order to have a light neutralino for satisfying the dark matter phenomenology, we choose small values of $\lambda_1 = 0.2$, 0.25 as our benchmark values. The choice of λ_4 is determined from the bounds on chargino masses. The other soft-breaking parameters A_λ, A_{λ_1}, A_{λ_2}, B_λ are chosen to fit the CP-even scalar masses specially to make the lightest Higgs mass close to 125 GeV. Finally, we have chosen $\mu_{\rm eff}$ to be $\mathcal{O}(200\ \text{GeV})$ and $v_t = 2$ GeV, which determines the choice of v_s from Eq. (5). The ratio of M_1 to M_2 at the electroweak scale is consistent with universal gaugino masses at the GUT scale and gravity mediated SUSY breaking.

TABLE I. Value of the parameters specified at the electroweak scale for two sets of benchmark points.

Parameters at EW scale	Point 1	Point 2
$\tan\beta$	2.0	3.0
λ	0.60	0.64
λ_1	0.20	0.25
λ_2	0.75	1.02
λ_4	0.17	0.20
$\mu_{\rm eff}$ [GeV]	200	200
A_λ [GeV]	400	500
A_{λ_1} [GeV]	−10	−10
A_{λ_2} [GeV]	600	700
B_λ [GeV]	500	600
v_t [GeV]	2	2
M_1 [GeV]	150	200
M_2 [GeV]	300	400

TABLE II. Mass spectrum and relic density for two sets of benchmark points.

Mass spectrum	Point 1	Point 2
Neutral Higgs spectrum		
$m_h^{\rm Tree}$ [GeV]	120.6	119.2
m_{H_1} [GeV]	145.5	156.8
m_{H_2} [GeV]	482.4	630.7
m_{H_3} [GeV]	825.2	707.9
m_{A_1} [GeV]	114.3	116.9
m_{A_2} [GeV]	487.8	629.9
m_{A_3} [GeV]	897.3	816.0
Charged Higgs spectrum		
$m_{H_1}^\pm$ [GeV]	208.4	239.9
$m_{H_2}^\pm$ [GeV]	280.5	320.6
$m_{H_3}^\pm$ [GeV]	496.3	647.1
Neutralino spectrum		
$m_{\tilde{\chi}_1^0}$ [GeV]	100.4	102.9
$m_{\tilde{\chi}_2^0}$ [GeV]	122.6	145.7
$m_{\tilde{\chi}_3^0}$ [GeV]	164.7	205.9
$m_{\tilde{\chi}_4^0}$ [GeV]	212.6	261.5
$m_{\tilde{\chi}_5^0}$ [GeV]	248.2	265.7
$m_{\tilde{\chi}_6^0}$ [GeV]	345.0	426.6
Chargino spectrum		
$m_{\tilde{\chi}_1^\pm}$ [GeV]	124.2	127.7
$m_{\tilde{\chi}_2^\pm}$ [GeV]	194.5	250.2
$m_{\tilde{\chi}_3^\pm}$ [GeV]	347.1	428.1
Relic density		
Ωh^2	0.117	0.08

The mass spectrum shown in Table II indicates all masses at tree level. The Higgs spectrum consists of four CP-even Higgs (h, H_1, H_2, H_3), three pseudoscalar Higgs (A_1, A_2, A_3), and three charged Higgs ($H_1^\pm$, $H_2^\pm$, $H_3^\pm$). We obtain significant contribution from the singlet and triplet sector at tree level which is highly appreciable, since this has raised the mass of the lightest CP-even Higgs boson to 125 GeV. Here we do not require a significant radiative contribution from the top-stop sector [5]. The components of the lightest physical Higgs for $\tan\beta = 2.0$ are given as

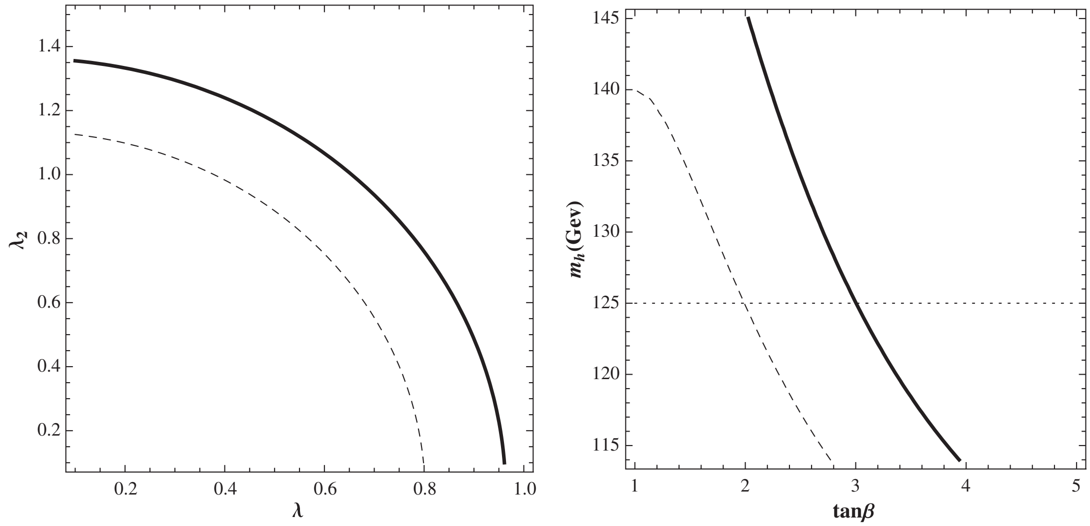

FIG. 1. Left: λ vs λ_2, for $\tan\beta = 2$ (dashed line), 3 (thick line) with $m_h = 125$ GeV. Right: m_h vs $\tan\beta$ for $\lambda = 0.6$, $\lambda_2 = 0.75$ (dashed line) $\lambda = 0.64$, $\lambda_2 = 1.02$ (thick line), and the dotted line shows the recent bound, i.e., $m_h = 125$ GeV.

$$h = 0.84205 H^0_{u_R} + 0.44422 H^0_{d_R} + 0.01977 T^0_R + 0.30533 S_R. \quad (39)$$

The lightest Higgs mass eigenstate has significant contribution from the singlet and some contribution from the triplet sectors. We obtain the lightest scalar Higgs mass for the two sets of benchmark points as 120.6 and 119.2 GeV, respectively. This will change the $h \to \gamma\gamma$ branching compared to the standard model and precise determination of the Higgs decay branchings at the LHC will be a good test of this model. In the pseudoscalar Higgs sector, we obtain one Goldstone boson exactly identified as Eq. (30), i.e., $G^0 = 0.4472 H^0_{d_I} - 0.8942 H^0_{u_I}$, for $\tan\beta = 2.0$ and $G^0 = 0.3163 H^0_{d_I} - 0.9487 H^0_{u_I}$, for $\tan\beta = 3.0$. All other Higgs masses are listed under Higgs spectrum in Table II.

The neutralino and the chargino sector consists of six and three mass eigenstates, respectively. The mass of the lightest neutralino being $\mathcal{O}(100\text{ GeV})$ is the LSP of this model. The prospects of the LSP to be identified as a dark matter candidate are discussed in detail in Sec. VI. The rest of the mass spectrum are shown in Table II.

A. One-loop correction to the lightest physical Higgs mass

The one-loop correction to m_h^2 is calculated by constructing the Coleman-Weinberg potential [18],

$$V_{\rm CW} = \frac{1}{64\pi^2} S\,{\rm Tr}\left[M^4\left(\ln\frac{M^2}{Q_r^2} - \frac{3}{2}\right)\right], \quad (40)$$

where M^2 are the field-dependent tree-level mass matrices and Q_r is the renormalization scale. STr is the supertrace that includes a factor of $(-1)^{2J}(2J+1)$ and summed over the spin degrees of freedom. The one-loop mass matrix can be derived from the above potential as follows:

$$(\Delta M_f^2)_{ij} = \frac{\partial^2 V_{\rm CW}(f)}{\partial f_i \partial f_j}\bigg|_{\rm vev} - \frac{\delta_{ij}}{\langle f_i\rangle}\frac{\partial V_{\rm CW}(f)}{\partial f_i}\bigg|_{\rm vev}, \quad (41)$$

where $f_{i,j}$ stands for all the real components of H_u^0, H_d^0, S, and T^0. Finally, the set of mass eigenvalues of the CP-even, CP-odd, charged Higgs, and neutralino-chargino mass matrices (all field dependent) enters the calculation. The dominant contribution in the one-loop correction comes from the top-stop sector and the triplet sector. We compute the corrections only numerically using the benchmark values assigned for the sets of parameters. The results we obtain are given below in Table III.

In both cases we do not require a large contribution from the radiative corrections to raise the lightest physical Higgs mass so as to satisfy the value of 125 GeV. This in turn implies that the contribution from the stop-top sector is not significant as in the case of the MSSM. In fact in absence of fine-tuning, the correction to lightest physical Higgs mass from the stop-top sector is given by

$$\delta m^2_{H_u}(Q) \simeq \frac{3m_t^2}{(4\pi)^2 v^2}\ln\frac{m_{\tilde t_1} m_{\tilde t_2}}{m_t^2}. \quad (42)$$

For $m_{\tilde t_1}$ and $m_{\tilde t_2}$ being $\mathcal{O}(200\text{ GeV})$, this amounts to a correction of only a few GeV.

TABLE III. Value of the lightest physical Higgs mass after one-loop correction for two sets of benchmark points.

Benchmark point	$m_h^{\rm Tree}$ [GeV]	$m_h^{\rm Tree+Loop}$ [GeV]
Point 1	120.6	124.9
Point 2	119.2	125.5

B. Fine-tuning in the electroweak sector

In this model, the lightest physical Higgs mass at tree level is boosted compared to the NMSSM and other triplet-extended models [14] as it gets contribution from both the singlet and triplet sectors (27). Therefore, we can obtain a Higgs boson close to 125 GeV even at tree level. After including the leading order radiative corrections from the stop-top and triplet sectors, we get

$$\delta m_{H_u}^2(Q) \simeq \frac{3y_t^2}{8\pi_2}(m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2 + A_t^2)\ln\left(\frac{Q}{M_z}\right) + \frac{3\lambda_2^2}{8\pi_2}(m_T^2 + A_{\lambda_2}^2)\ln\left(\frac{Q}{M_z}\right), \quad (43)$$

where $m_{\tilde{t}_1}$ and $m_{\tilde{t}_2}$ are the soft masses of the stops, A_t is the soft trilinear coupling, y_t is the Yukawa coupling, and Q is the fundamental scale of SUSY breaking.

The fine-tuning parameter can be quantified [19,20] as

$$\Delta_{\rm FT} \equiv \frac{m_{H_u}^2}{M_z^2}\frac{\partial M_z^2}{\partial m_{H_u}^2}. \quad (44)$$

In case of the MSSM [only the first term in Eq. (43) is present], we have

$$\Delta_{\rm FT}^{\rm Stop} \simeq \frac{3y_t^2}{8\pi_2}(m_{\tilde{t}_1}^2 + m_{\tilde{t}_2}^2 + A_t^2)\ln\left(\frac{Q}{M_z}\right), \quad (45)$$

but the tree-level bound on the Higgs mass is $m_h \leq M_z \cos 2\beta$. Therefore, one is forced to consider large values for $m_{\tilde{t}_1}$, $m_{\tilde{t}_2}$, and A_t, say 1 TeV, in order to raise the lightest physical Higgs boson mass up to 125 GeV. In this case, $\Delta_{\rm FT}^{\rm Stop} \simeq 80$ and thus it leads to maximal stop mixing.

In the NMSSM, the radiative correction needed to get a 125 GeV Higgs mass is $\delta m_h = 55$ GeV. There is no doubt an improvement over MSSM, but still fine-tuning is required in the stop-top sector [5]. In the model with one triplet [14], the lightest physical Higgs mass can be raised to 113 GeV. Here, the required value of radiative correction is $\delta m_h = 53$ GeV. Now, the fine-tuning due to the triplet sector is

$$\Delta_{\rm FT}^{\rm Trip} \simeq \frac{3\lambda_2^2}{8\pi_2}(m_T^2 + A_{\lambda_2}^2)\ln\left(\frac{Q}{M_z}\right), \quad (46)$$

where $\lambda_2 = 0.8, 0.9$. The value of $\Delta_{\rm FT}^{\rm Trip}$ can be as large as 40. Therefore, this model can no longer be considered as a zero fine-tuning model.

Now coming to our model, we require $\delta m_h \simeq 35$ GeV only—here we see a distinct improvement of 20–50 GeV compared to other models discussed so far. Also, $\lambda_2 = 0.75$, 1.02 being comparable to y_t, we do not need heavy stops or large stop-top mixing to get the required Higgs mass. For example, using $m_T = 200$ GeV, $A_{\lambda_2} = 700$ GeV, and $Q = 1$ TeV, we obtain $\Delta_{\rm FT}^{\rm Trip} \simeq 10$. Thus, we can achieve little fine-tuning compared to other models, since the lightest physical Higgs mass can be large at tree level and does not require a large contribution from the radiative corrections. Here we note that the Higgs-triplet-Higgs coupling λ_2 (0.75 and 1.02) becomes nonperturbative at the GUT scale. But, these choices of λ_2 actually help to raise the Higgs mass close to 125 GeV at tree level. Another alternative could be of course having small λ_2, but then we would require large radiative corrections. Therefore, we improve the level of fine-tuning at the cost of giving up perturbativity of λ_2 at the GUT scale.

VI. DARK MATTER

We have analyzed the neutralino sector where the lightest neutralino (LSP) is a mixture of the Higgsino-triplino and turns out to be a viable dark matter candidate. The components of $\tilde{\chi}_0$ (for $\tan\beta = 2.0$), i.e., the LSP, are

$$\tilde{\chi}_0 = -0.321\tilde{B} + 0.192\tilde{W}_3^0 - 0.323\tilde{H}_d^0 + 0.644\tilde{H}_u^0 - 0.213\tilde{S} + 0.544\tilde{T}^0. \quad (47)$$

Since the LSP has mass $\mathcal{O}(100\text{ GeV})$, there are two possibilities of final states into which it can annihilate, i.e. (i) fermion final states and (ii) gauge boson final states. For annihilation into fermions, except $t\bar{t}$ it can go to any other $f\bar{f}$ pairs via pseudoscalar Higgs, Z-boson exchange and sfermion exchange. But, if we consider the neutralino to be more like the triplino, then its coupling with the Z boson is forbidden. Generally, it can annihilate into gauge boson pairs via several processes like chargino exchange, scalar Higgs exchange, and Z-boson exchange. But the dominant contribution comes from annihilation into $W^\pm$ via chargino exchange, which finally leads to the relic density of 0.117, consistent with WMAP [16].

The scalar interaction between the dark matter (i.e., neutralino LSP) and the quark is given by

$$\mathcal{L}_{\rm scalar} = a_q \bar{\chi}\chi\bar{q}q, \quad (48)$$

where a_q is the coupling between the quark and the neutralino. The scalar cross section for the neutralino scattering off a target nucleus (one has to sum over the proton and neutrons in the target) is given by

$$\sigma_{\rm scalar} = \frac{4m_r^2}{\pi}(Zf_p + (A-Z)f_n)^2, \quad (49)$$

where m_r is the reduced mass of the nucleon and $f_{p,n}$ is the neutralino coupling to the proton or neutron [21,22], given by

$$f_{p,n} = \sum_{q=u,d,s} f_{Tq}^{(p,n)} a_q \frac{m_{p,n}}{m_q} + \frac{2}{27} f_{TG}^{(p,n)} \sum_{q=c,b,t} a_q \frac{m_{p,n}}{m_q}, \quad (50)$$

where $f_{Tu}^{(p)} = 0.020 \pm 0.004$, $f_{Td}^{(p)} = 0.026 \pm 0.005$, $f_{Ts}^{(p)} = 0.118 \pm 0.062$, $f_{Tu}^{(n)} = 0.014 \pm 0.003$, $f_{Td}^{(n)} = 0.036 \pm 0.008$, and $f_{Ts}^{(n)} = 0.118 \pm 0.062$ [23]. $f_{TG}^{(p,n)}$ is related to these values by

$$f_{TG}^{(p,n)} = 1 - \sum_{q=u,d,s} f_{Tq}^{(p,n)}. \quad (51)$$

The term in Eq. (50) that includes $f_{TG}^{(p,n)}$ results from the coupling of the weakly interacting massive particle (WIMP) to gluons in the target nuclei through a heavy quark loop.

We can approximate $\frac{a_q}{m_q} \simeq \alpha/(s - m_h^2)$ where α is the product of different couplings and mixings, m_q is the mass of the quark, and $s = 4m_\chi^2$ (m_χ being the dark matter mass). The parameter α plays a crucial role in determining the spin-independent cross section and is highly model dependent. Using this we estimate $\alpha \simeq 2 \times 10^{-4}$ GeV^{-1} and the value of the spin-independent cross section is 10^{-43} cm^2, which is below the exclusion limits of XENON100 [17] and other direct detection experiments.

VII. CONCLUSIONS

In this paper we have explored an extension of the MSSM where the Higgs sector is extended by a singlet and a $Y = 0$ triplet superfield. This is the minimal model that gives a tree-level Higgs mass of $\mathcal{O}$(119–120 GeV) and the one-loop correction can easily raise it to 125 GeV without significant contribution from the stop-top sector. However, $\lambda_2 = 0.75$, 1.02 (at the electroweak scale) becomes nonperturbative at the GUT scale, while all other couplings remain perturbative up to the GUT scale; on the other hand, this is the price we pay to retain small fine-tuning.

In addition, we see that the triplino and singlino contributions to the neutralino mass matrix gives a viable dark matter candidate with mass around 100 GeV that may be seen at the LHC from the missing transverse energy signals [24]. In the MSSM and NMSSM the problem for getting the correct relic density of dark matter is related to the necessity of choosing chargino and scalar masses to be in the multi-TeV scale to fit the Higgs mass from radiative corrections. The dark matter mass in the MSSM is around 700 GeV while in the NMSSM it is possible to obtain viable dark matter in the 100 GeV range. The main advantage of our model for the dark matter is that since the sparticle masses need not be very large compared to the electroweak scale, the "WIMP miracle" is restored, and we are able to get dark matter mass in the 100 GeV range over a large parameter space of our model.

The data from the LHC with integrated luminosity of 5 fb^{-1} has not only given an indication of the Higgs mass but there is also a measurement of the Higgs decay branching into different channels. Detailed analysis [25] of the 125 GeV Higgs branching fractions seen at the LHC indicates that the signal ratio for Higgs decay into two photons is larger than the SM prediction by a factor of 2.0 ± 0.5; decay into WW^* and ZZ^* channels is smaller than the SM by a factor of 0.5 ± 0.3; and into bb and $\tau\tau$ channels it is factor of 1.3 ± 0.5 consistent with the SM. The lightest CP-even Higgs (39) has a sizable fraction of the singlet, and the Higgs decay phenomenology will be distinguishable from the MSSM [26] and likely to be similar to the NMSSM scenario [27,28]. But, there will be some contribution from the triplet sector too. The phenomenological aspects of the real triplet-extended SM has been studied in Ref. [29]. More data from the LHC will pinpoint or rule out the extended Higgs sector models, and it would be useful to study the singlet-triplet-extended MSSM model in greater detail with emphasis on the LHC signal in the future.

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