We propose a UV completion of the inverse see-saw scenario using fermion SU(2)$_L$ triplet representations. Within this framework, a variation of the standard thermal leptogenesis is achievable at the $\mathcal{O}(\text{TeV})$ scale, owing to the presence of a viable dark matter candidate. This baryogenesis scenario is ruled out if a triplet fermion is observed at the LHC. The dark matter is given by the lightest neutral component of a complex scalar SU(2)$_L$ triplet, with mass $m_{\text{DM}}$ in the TeV range. The scalar sector, which is enriched in order to account for the small neutrino masses, is treated in detail and shows potentially sizable Higgs boson $h \rightarrow \gamma\gamma$ rates together with large $h$ invisible branching ratios.

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I. INTRODUCTION

The recent discovery of a new scalar resonance by ATLAS [1] and CMS [2] experiments at the Large Hadron Collider (LHC) may finally allow us to disclose the mechanism of electroweak symmetry breaking (EWSB) and fully test the Standard Model (SM) scalar sector. On the other hand, new physics beyond the SM must be advocated if we want to explain neutrino masses and mixing, dark matter (DM) and the baryon asymmetry of the universe (BAU).

In this article, we propose a simple renormalizable extension of the SM where all these issues are addressed:

1. The small neutrino masses are explained by an inverse see-saw model, constructed by the addition of fermion and scalar fields to the SM particle content.

2. The DM corresponds to the neutral component of a complex triplet scalar field, explaining the observed relic abundance: \( \Omega_{\text{DM}}h^2 = 0.112 \pm 0.006 \) [3].

3. The BAU \( \Omega_Bh^2 = 0.0226 \pm 0.006 \) [3] is explained through a variation of the standard leptogenesis mechanism we proposed in [4].

The model Lagrangian is invariant under a global \( U(1)_X \) symmetry, which is spontaneously broken at the electroweak scale to a remnant \( Z_2 \) symmetry. The role of the additional \( U(1)_X \) is twofold: through its breaking, it ensures the stability of the triplet DM candidate and provides tiny light neutrino masses. With the new fields typically at the TeV scale, the model actually provides a UV completion of the inverse see-saw scenario [5]. We list in Table I the fermion and scalar content of our model and their charges under \( SU(2)_L \times U(1)_Y \times [U(1)_X] \). It is clear from this table that the \( U(1)_X \) charge effectively corresponds to a generalization of the baryon-lepton charge \( B-L \).

The heavy right-handed (RH) neutrinos \( N_{1,2} \) we introduce in Table I are charged under \( U(1)_X \), which is conserved above the EWSB. Therefore, thermal leptogenesis [6, 7] cannot be realized as in the standard type I [8] or type III [9] see-saw extensions of the SM. Similarly to scenarios where the \( B-L \) charge is conserved [10–15], a nonzero lepton number asymmetry can be generated by out of equilibrium decays of a heavy scalar/fermion particle, while a (model-dependent) mechanism prevents the washout of the total lepton charge \(^1\). The generated lepton asymmetry is converted into a nonzero baryon number by the rapid nonperturbative sphaleron processes [18, 19].

In the present case, we consider a variant of our work [4], using—instead of singlets—\( SU(2)_L \) triplet representations for the RH neutrinos \( N_{1,2} \) and the scalar \( S \) \(^2\). A singlet Majorana fermion \( N_3 \) is introduced, and as in type I leptogenesis its decays produce asymmetries in \( N \) and \( S \). Successful leptogenesis is possible due to the transfer \(^3\) of the \( U(1)_X \) charge asymmetry of \( N \) to the SM leptons, through (fast) neutrino Yukawa interactions.

The inverse see-saw completion we propose requires several additional scalar particles. Such an enlarged scalar sector induces deviations from SM expectations, and after the observation of a \( \sim 126 \) GeV boson [1, 2], a reanalysis of the scalar spectrum of [4] is in due order. As we show below, a large Higgs diphoton rate is possible, together with a large Higgs invisible branching ratio.

This paper is organized as follows: in Sec. II we describe the scalar sector of the model and compare its predictions to current ATLAS and CMS data. In Sec. III we outline the phenomenology of the triplet scalar dark matter scenario. The neutrino mass generation and the leptogenesis mechanism are reported in Sec. IV. Finally, we summarize the main results of this work in the concluding section.

II. SCALAR SECTOR

We extend the Standard Model particle content with new scalar representations of the SM gauge group, which are listed in Table I:

i) a \( SU(2)_L \) doublet \( H^T_2 \equiv (H^+_2, H^0_2) \), with hypercharge \( Y=1/2 \), in addition to the SM Higgs doublet \( H^T_1 \equiv (H^+_1, H^0_1) \),

ii) a complex singlet \( \phi \),

iii) a complex \( SU(2)_L \) triplet \( S \) with zero hypercharge.

---

1 In the case of Dirac leptogenesis, very small neutrino Yukawa couplings prevent equilibration between left and right-handed lepton asymmetries, until long after EWSB when sphalerons are no longer active. A similar role for the sphaleron decoupling epoch is used in [14], in interplay with lepton flavour effects [16, 17].

2 Triplet \( SU(2)_L \) fermion representations are used in the inverse see-saw mechanism, e.g. in [20, 21].

3 A similar \( B-L \) conserving Yukawa-driven transfer mechanism for leptogenesis is studied in [13], based on a radiative model of neutrino masses.
All these new fields have a nonzero charge under a global U(1)\textsubscript{X} symmetry. In this scenario, the presence of \(H_2\) and \(\phi\) is motivated by the requirement of generating light neutrino masses through the (inverse) see-saw mechanism and the possibility of realizing successful leptogenesis (see Sec. IV). We list in Table I the particle content of the model with their charge assignments under SU(2)\textsubscript{L}, U(1)\textsubscript{Y} and the global U(1)\textsubscript{X}.

The most general scalar potential which is invariant under the SU(2)\textsubscript{L}×U(1)\textsubscript{Y}×[U(1)\textsubscript{X}] is conveniently split into two parts: \(\mathcal{V}_{\text{SC}} \equiv \mathcal{V}_{\text{SB}} + \mathcal{V}_{\text{DM}}\). The scalar potential pertaining to DM, \(\mathcal{V}_{\text{DM}}\), is discussed in Sec. III. The symmetry breaking scalar potential \(\mathcal{V}_{\text{SB}}\) is [4]:

\[
\mathcal{V}_{\text{SB}} = -\mu_1^2 (H_1^\dagger H_1) + \lambda_1 (H_1^\dagger H_1)^2 - \mu_2^2 (H_2^\dagger H_2) + \lambda_2 (H_2^\dagger H_2)^2 - \mu_3^2 \phi^\dagger \phi + \lambda_3 (\phi^\dagger \phi)^2
+ \kappa_{12} (H_1^\dagger H_1) (H_2^\dagger H_2) + \kappa_{13} (H_1^\dagger H_1) (H_2^\dagger H_1) + \kappa_{23} (H_1^\dagger H_2) \phi^\dagger \phi
- \frac{\mu'}{\sqrt{2}} \left( (H_1^\dagger H_2) \phi + (H_2^\dagger H_1) \phi^\dagger \right). \tag{2.1}
\]

Only the scalar fields with even U(1)\textsubscript{X} charge acquire a nonzero (real) vacuum expectation value (vev): the two SU(2)\textsubscript{L} doublets \(H_{1,2}\) and the singlet \(\phi\) [4], their respective vev being \(v_{1,2}\) and \(v_\phi\). With such a charge assignment, the global U(1)\textsubscript{X} is broken down to a \(Z_2\) which stabilizes the DM, odd under U(1)\textsubscript{X}. Given the field content and the charge assignment reported in Table I, this model can be thought as a variant of the type I two-Higgs doublet model (see e.g. [22]), augmented with a further complex scalar \(\phi\) \(^4\); both doublets \(H_{1,2}\) couple to gauge bosons and while only \(H_1\) couples to SM fermions via Yukawa couplings, \(H_2\) and \(\phi\) couple to the new particles \(S\), \(N_D\) and \(N_3\) (see the following sections).

We first discuss the scalar spectrum arising from the spontaneous breaking of SU(2)\textsubscript{L}×U(1)\textsubscript{Y}×[U(1)\textsubscript{X}] to U(1)\textsubscript{em}×[Z_2] and we postpone the analysis of the triplet scalar \(S\) to the next section. The spectrum consists of [4]:

- 1 charged scalar \(H^\pm\),
- 3 CP-even neutral scalars \(h^0\), \(H^0\) and \(h_A\),
- 2 pseudoscalars \(A^0\) and \(J\).

The latter is the Goldstone boson associated with the breaking of the global U(1)\textsubscript{X} symmetry and is usually dubbed Majoron in theories with spontaneously broken lepton charge [26]. Since it is a massless particle at the renormalizable level, strong constraints apply on its couplings to SM fermions: a hierarchical pattern for the vevs of the scalar fields, \(v_2 \ll v_{1,1}\), \(v_\phi\), is required \(^5\). In the limit of negligible \(v_2\), the longitudinal gauge boson components are \(W^\pm_L \sim H^\pm_1\) and \(Z_L \sim \sqrt{2} \text{Im} (H^0_1)\), while the scalar mass eigenstates are to a good approximation:

\[
H^\pm \sim H^\pm_2, \quad h_A \sim \sqrt{2} \text{Re} (H^0_2), \quad A_0 \sim \sqrt{2} \text{Im} (H^0_2), \quad \text{and} \quad J \sim \sqrt{2} \text{Im} (\phi),
\]

while the two neutral scalars \(h^0\) and \(H^0\) are related to the interaction fields \(H_1\) and \(\phi\) by a rotation

\[
\begin{pmatrix}
\theta^0 \\
H^0
\end{pmatrix}
= R(-\theta)
\begin{pmatrix}
\sqrt{2} \text{Re} (H^0_1) \\
\sqrt{2} \text{Re} (\phi)
\end{pmatrix}, \tag{2.2}
\]

where \(\theta\) is a function of the vevs \(v_{1,\phi}\) and the quartic couplings of \(H_1\) and \(\phi\) in \(\mathcal{V}_{\text{SB}}\) [4]. Typically, we have \(v_2 \lesssim 10\) MeV, \(v_1 \simeq 246.2\) GeV, while \(v_\phi\) is free. Recalling that only \(H_1\) has Yukawa couplings to SM fermions (cf. Table I), \(h_A, A^0\) and \(H^\pm\) couple to the SM sector only through gauge interactions and via the scalar quartic couplings, while \(h^0\) and \(H^0\) can have \textit{a priori} sizable Yukawa couplings to SM fermions.

\(^4\) Such scalar spectrum is considered in a different context in [23].

\(^5\) A suppressed value of \(v_2\) is naturally realized from the minimization of the potential (2.1), due to the symmetries of the model [4].
A. LHC constraints

In [4] we performed a detailed analysis of the scalar spectrum and we constrained the parameter space given the experimental results available at that time. However, after the observation of a new scalar particle with mass $m_h \sim 126$ GeV by both ATLAS [1] and CMS [2] Collaborations, it is worth studying more carefully the scalar mass spectrum of the model.

We assume the new discovered particle corresponds to the lightest of the two CP-even scalars with possibly large couplings to SM fermions, that is $h^0$. We ought to verify whether the model can explain ATLAS and CMS data, while encompassing present phenomenological constraints. In order to proceed with the analysis, we consider the ratios between the production of a boson $H$ decaying into a generic final SM state $i$ to the corresponding SM prediction,

$$\mu_i(H) \equiv \frac{\sigma(pp \rightarrow H)_{i}}{\sigma(pp \rightarrow H)_{i}^{SM}} \times \frac{\text{Br}(H \rightarrow i)}{\text{Br}(h \rightarrow i)^{SM}}. \quad (2.3)$$

Depending on the decay products, Higgs searches target specific production channels, hence the labeling $\sigma(pp \rightarrow h)_{i}$. Notice that, in our model, the Higgs signal strengths $\mu_i$s may be affected in several aspects and may differ with respect to the SM predictions.

Regarding the Higgs boson production, as no extra colored particles are introduced, there are no new contributions to the loop-induced gluon-gluon fusion process. Furthermore, from the definition, eq. (2.2), the couplings of $h^0$ ($H^0$) to the SM particles are those of $H_1$ times $\cos(\theta)$ ($\sin(\theta)$): all production channels are thus equally rescaled and we have

$$\frac{\sigma(pp \rightarrow h^0)_{i}^{SM}}{\sigma(pp \rightarrow h)_{i}^{SM}} = \cos^2(\theta), \quad \frac{\sigma(pp \rightarrow H^0)_{i}^{SM}}{\sigma(pp \rightarrow h)_{i}^{SM}} = \sin^2(\theta).$$

The branching ratios $\text{Br}(H \rightarrow i)$ reported in (2.3) are affected in three ways. First, as $h^0$ ($H^0$) couplings to SM particles are rescaled by $\cos(\theta)$ ($\sin(\theta)$) compared to the SM ones, tree level Higgs decays to SM particles are rescaled by $\cos^2(\theta)$ ($\sin^2(\theta)$). Second and ceteris paribus, the branching ratios are reduced with respect to the corresponding SM predictions, $\text{Br}(h \rightarrow i)^{SM}$, because of the presence of new decay channels. For the sake of simplicity, we assume that $h^0$ is the lightest massive neutral scalar particle, thus closing these decay channels. Moreover, the invisible decays $h^0/H^0 \rightarrow J J$ can be sizable in some regions of the parameter space. Third, deviations from the SM occur in loop-induced processes. As already stated, no extra colored particles are introduced, so Higgs decays to gluons are not affected. On the other hand, the diphoton decay channel $h^0$ ($H^0$) $\rightarrow \gamma \gamma$ is affected by the presence of new charged particles. Several works have emphasized possible deviations from SM expectations of the diphoton decay rate, due to the presence of extra fermion/scalar charged particles (see e.g. [27–31]). In our model, we potentially have to consider the effect of 5 extra charged particles: the scalar $H^\pm$ originating from the doublet $H_{1,2}$, the two scalars $S_T^\pm$ and $S_M^\pm$ coming from our DM scalar triplet (see Sec. III), and the two charged fermions $\Sigma_{1,2}^\pm$ coming from the triplet $N_D$ introduced in our type III inverse seesaw variant (see Sec. IV). As discussed in the next sections, the triplet particles have large masses, typically $O(\text{TeV})$: their contribution to the diphoton decay rate is therefore negligible. Oppositely and as in the case of type I 2HDM, the charged scalar $H^\pm$ can be sufficiently light to produce observable effects in the $h^0 \rightarrow \gamma \gamma$ decay rate: assuming $\text{Br}(H^+ \rightarrow c \bar{s}) + \text{Br}(H^+ \rightarrow \tau^+ \nu) = 1$, LEP2 derived a conservative lower bound on the mass of $H^\pm$, $m_{H^\pm} \gtrsim 80$ GeV [34]. The Higgs diphoton rate is given in our model by [35, 36] $^8$

$$\Gamma(h^0/H^0 \rightarrow \gamma \gamma) = \frac{G_F \alpha_2 m_{h^0/H^0}^3}{128 \sqrt{2} \pi^3} \sum_f N_e Q_f^2 \lambda_{f f}^{h^0/H^0} A_{1/2} \left( \frac{m_{h^0/H^0}^2}{4 m_f^2} \right) + \lambda_{WW}^{h^0/H^0} A_1 \left( \frac{m_{h^0/H^0}^2}{4 m_W^2} \right)$$

$$- \frac{\alpha_2^3}{2 m_{H^+}^2} \lambda_{H^+ H^-}^{h^0/H^0} A_0 \left( \frac{m_{h^0/H^0}^2}{4 m_{H^+}^2} \right)^2. \quad (2.4)$$

The first line is the contribution of SM fermions and $W$ boson running in the loop, while the second line contains the contribution from $H^\pm$. As we said, the couplings of $h^0$ to SM fermions and gauge bosons are equally rescaled, $\lambda_{f f}^{h^0} = \lambda_{V V}^{h^0} = \cos(\theta)$, while we have $\lambda_{H^+ H^-}^{h^0} \simeq - (\kappa_{12} \cos(\theta) + \kappa_{23} \sin(\theta)) v_\phi / v_1$. Depending on the sign of the

---

$^6$ Flavour physics observables in our case do not put constraints on $H^\pm$ mass, given the large vev hierarchy: $v_2 \ll v_1$.

$^7$ LHC constraints on $m_{H^\pm}$ are yet based on the decays $t \rightarrow H^+ b$: the values we obtain are well below the present bound: $\text{Br}(t \rightarrow H^+ b) \lesssim 10^{-2}$ [32, 33].

$^8$ We use the same conventions and notations as [36].
We only dispose of data up to 600 GeV, that we use as an upper bound on the scalar masses.

Table II. Best fit value $\hat{\mu}_i$ and symmetric errors given at 1 $\sigma_i$ used in our fit.

| Channel: | $\tau\tau$ | $bb$ | $WW$ | $ZZ$ | $\gamma\gamma$ |
|----------|-------------|------|------|------|--------------|
| $\hat{\mu}_i$ | 0.15 | 0.49 | 0.9 | 0.88 | 1.67 |
| $\sigma_i$ | 0.7 | 0.73 | 0.3 | 0.34 | 0.34 |

coupling $\lambda^{h^0}_{H^+ H^-}$, the diphoton rate $h^0 \rightarrow \gamma\gamma$ can be enhanced ($\lambda^{h^0}_{H^+ H^-} > 0$) or suppressed ($\lambda^{h^0}_{H^+ H^-} < 0$) \(^9\), the effect being more pronounced for light $m_{H^\pm}$.

1. Analysis

We assume $m_{h^0} = 126$ GeV and $m_{A^0} < m_{H^0} \lesssim 600$ GeV. \(^{10}\) As previously said, the other neutral scalar fields $h_A$ and $A^0$ are very weakly coupled to the SM particles: their contributions can be neglected.

To perform the analysis, we construct the Higgs signal strengths $\mu_i$, corresponding to the five channels $h^0(H^0) \rightarrow bb$, $\tau\tau$, $\gamma\gamma$, WW and ZZ. For the $b$ and $\tau$ channel, we use the combined results of CMS at 7 and 8 TeV \([37, 38]\) and the 7 TeV results of ATLAS \([39, 40]\), while for the gauge boson channels we use both ATLAS and CMS results combining 7 and 8 TeV observations \([41–46]\). We do not distinguish the decay products of the gauge bosons produced in $h^0/H^0$ decays, $h^0/H^0 \rightarrow VV$, ($V = W^{(*)}, Z^{(*)}$): for both $h^0$ and $H^0$, the resulting branching factors cancel out with the SM ones. For $h^0$ we use combined results, while for $H^0$ only the most constraining signal at a given mass is considered. Further, for $h^0$ we combined ATLAS and CMS results assuming a Gaussian distribution for the observed signal strengths $\hat{\mu}_i$ and symmetric errors: we summarize in Table II the central value and the symmetric errors for the different channels we use. More refined analyses have been done in e.g. \([23, 27–29, 47]\). We nevertheless do not expect that a more rigorous treatment of LHC data would yield significant deviations from the results we obtain.

Finally, we take into account the electroweak precision data. We construct the $S$, $T$ and $U$ parameters following the results of \([48]\), and used the values of the electroweak fit quoted in \([29]\):

$$S = 0.0 \pm 0.1, \quad T = 0.02 \pm 0.11, \quad U = 0.03 \pm 0.09.$$  

We then compute a $\chi^2$ defined by

$$\chi^2(\mu_i(h^0)) = \sum_{i=\gamma, Z, W, S, T, U} \frac{(\hat{\mu}_i(h^0) - \hat{\mu}_i)^2}{\sigma_i^2},$$

that we minimize over the $\gamma$, $Z$ and $W^\pm$ Higgs signal strengths and the three oblique parameters. As no significant excess of events has been seen so far in $b$ and $\tau$ channels, we do not include them in the definition of the $\chi^2$: we instead require that $\hat{\mu}_b$ and $\hat{\mu}_\tau$ are below their respective 95\% C.L. upper bound. We further ask that the $H^0$ signal strength $\mu_i(H^0)$ ($i = b$, $\tau$, $\gamma$, $Z$, $W$) to be smaller than the observed ones over the full $H^0$ mass range.

2. Results

Before addressing the results, a few comments are in order. The observables we consider are built upon a rich scalar sector: 9 free parameters are scanned over. We do not aim to constrain these parameters, but to highlight the main features our model exhibits.

Constraints on the invisible decay width can already be set: in the SM, the branching ratio $\text{Br}(h \rightarrow \gamma\gamma)$ of a 126 GeV Higgs is about $2.3 \times 10^{-3}$ \([49]\) and even if the diphoton rate is increased in our case, we can neglect it in a first approximation. Then, the total decay width of $h^0$ is approximately the sum of the rescaled SM channels plus the invisible one:

$$\Gamma(h^0)_{\text{tot}} \simeq \cos(\theta)^2 \Gamma(h^0)_{\text{SM}} + \Gamma(h^0 \rightarrow \text{inv}).$$  \hspace{1cm} (2.5)

For the 4 tree level channels $h^0 \rightarrow bb$, $\tau\tau$, WW and ZZ, we can write:

\(^9\) Enhancement is also possible for large negative values of $\lambda^{h^0}_{H^+ H^-}$, such that $H^\pm$ contribution largely dominates over SM ones. We do not consider this possibility.

\(^{10}\) We only dispose of data up to 600 GeV, that we use as an upper bound on the scalar masses.
The minimal upper bounded. At 95% C.L., we obtain $\text{Br}(h^0 \rightarrow i)$ the 1-2-3 $\sigma$ the 1-2-3 $\sigma$ the $\text{h}_\text{SM}$ case, $\Gamma(h^0)$ increases. Results of our scan are depicted by red dots, and we shade in dark-red the region within the 95% C.L. allowed range of the global fit. Similarly, the green area represents the 95% C.L. allowed range of each observable below their respective 95% C.L. bounds (upper bound).

Asking for example that the $h^0 \rightarrow W W$ signal strength is within its 2 $\sigma$ range, we have approximately an upper bound on the invisible width $\text{Br}(h^0 \rightarrow \text{inv}) \lesssim 0.69$. The global fit to the 6 observables will however provide a slightly different bound.

We display in the left panel of Fig. 1 the influence of the invisible branching ratio $h^0 \rightarrow J J$ on the goodness of the fit. The thin red dots are the values of $\Delta \chi^2$ per degree of freedom 11, and the blue lines represent the 1-2-3 $\sigma$ deviations from the best fit, from bottom to top. The green shaded area represents the allowed invisible width. Another important result is the large possible enhancement of the diphoton signal strength is within its 2 $\sigma$ range, we have approximately an upper bound on $h^0 \rightarrow \text{inv}$ of 0.77. This value is large, and only possible from the increase of the total $h^0 \chi_\text{SM}^2 = \chi^2 - \chi_{\text{min}}^2$ per degree of freedom 11, and the blue lines represent the 1-2-3 $\sigma$ deviations from the best fit, from bottom to top. The green shaded area represents the allowed invisible branching ratio asking each $\mu_i$ to be within its 95% C.L. allowed range (but for $\mu_b$ and $\mu_\tau$ channels which are only upper bounded). At 95% C.L., we obtain $\text{Br}(h^0 \rightarrow \text{inv}) \lesssim 0.77$. This value is large, and only possible from the increase of the total $h^0 \chi_\text{SM}$ mass. A large enhancement compared to the SM case is possible, although still compatible with current observations [41, 44].

As a last comment, it is clear that our model can mimic the SM Higgs sector, or alternatively significantly deviate from it. Further data will hopefully shed light on the nature of the observed 126 GeV excitation.

### III. SCALAR TRIPLET DARK MATTER

In our model the dark matter is the complex scalar field $S$ triplet of SU(2)$_L$, for which the most general scalar potential reads

$$
\mathcal{V}_\text{DM} = \mu_S^2 S^\dagger S + \lambda_S (S^\dagger S)^2 + \lambda'_S (S^\dagger T_0^a S) (S^\dagger T_0^a S) + \mathcal{F}_1 H_1^\dagger H_1 S^\dagger S + \mathcal{F}_2 H_2^\dagger H_2 S^\dagger S + \mathcal{F}_3 \phi^* \phi S^\dagger S \\
+ \mathcal{F}_4 \left( H_1^\dagger T_2^a H_1 \right) \left( S^\dagger T_0^a S \right) + \mathcal{F}_5 \left( H_2^\dagger T_2^a H_2 \right) \left( S^\dagger T_0^a S \right) + \mathcal{H} S^2 H_1^\dagger H_2 + \mathcal{H}^* S^\dagger S^3 H_1^\dagger H_1 \\
- \frac{\mu''}{\sqrt{2}} (S^2 \phi^* + S^\dagger \phi). 
$$

11 The minimal $\chi^2$ we obtain is $\sim 0.3$ for 6 degrees of freedom.
The unitary transformation from the Cartesian to the spherical basis is:

\[ S = (\cos(\theta_s) S^+_L + \sin(\theta_s) S^+_H, \ (S^+_L + iS^+_H) / \sqrt{2}, \ \cos(\theta_s) S^-_H - \sin(\theta_s) S^-_L) \].

Defining

\[ m^0_{S_{L(H)}} = (m^2_0 + \frac{\delta^2_0}{2})^{1/2}, \]
\[ m^\pm_{S_{L(H)}} = (m^2_0 + \frac{1}{2} \sqrt{\delta^2_0 + \delta^2_0})^{1/2}. \]

These expressions are valid as long as \( \sqrt{\delta^2_0 + \delta^2_0} \leq 2 m^2_0 \).

\[ U = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 & 0 & 1 \\ i & 0 & -i \\ 0 & \sqrt{2} & 0 \end{pmatrix}. \]

In the spherical basis \( S = (S^+, \ S^0, \ S^-) \) the three generators of SU(2)_L in the adjoint representation are:

\[ T^3_G = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & -1 & 0 \\ -1 & 0 & 1 \\ 0 & 1 & 0 \end{pmatrix}, \]
\[ T^2_G = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & i & 0 \\ -i & 0 & -i \\ 0 & i & 0 \end{pmatrix}, \]
\[ T^3_G = \begin{pmatrix} 1 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & -1 \end{pmatrix}. \]

Notice that in the spherical basis the potential given in (3.1) is still invariant under the symmetries of the Lagrangian, provided one replaces the operators \( S^2 \) with \( ST_3 S \), where the transformation matrix \( T_3 \) is given by:

\[ T_3 = \begin{pmatrix} 0 & 0 & 1 \\ 0 & 1 & 0 \\ 1 & 0 & 0 \end{pmatrix}. \]
FIG. 3. Maximal value of the coupling $F_1'$ in the $\delta_0/m_0$ vs $m_0$ plane, such that $m_{S_L}^0 < m_{\mp_L}^\pm$.

At tree level, the charged fields receive a mass contribution from $\delta_{\pm}$. As a consequence, the lightest charged scalar, $S_L^\pm$, is lighter at tree level than the lightest neutral component, $S_L^0$. However, at one loop, the charged-neutral mass splitting receives a positive contribution from gauge boson loops [24]:

$$\left(m_{S_L}^\pm(H) - m_{S_L}^0(H)\right)_{\text{1loop}} = \left(m_{S_L}^\pm(H) - m_{S_L}^0(H)\right)_{\text{tree}} + \delta_m,$$

with $\delta_m \simeq (166 \pm 1)$ MeV in the triplet case. Thus, $S_L^0$ constitutes a viable DM candidate if $\delta_{\pm}$ contribution is smaller than $\delta_m$. In the case $\delta_0 = 0$ we have:

$$F_1' \lesssim 2.2 \times 10^{-2} \left(\frac{m_0}{1 \text{ TeV}}\right),$$

(3.4)
a similar bound derived in [25]. However, if $\delta_0 \neq 0$, this upper limit can be relaxed, as we show in Fig. 3. When $\delta_0$ is comparable to $m_0$ (still smaller than $\sqrt{2} m_0$), $F_1'$ can take sizable values. Under these conditions, $S_L^0$ provides a viable DM candidate.

### A. Relic abundance

The study of the scalar triplet as a possible DM candidate has been investigated in detail in e.g. [24, 25]. Our results agree with theirs. The dominant annihilation channel depends on the DM mass: the lowest mass is reached in the pure gauge boson limit, when Higgs portals are negligible. $S$ annihilations into gauge bosons receive in this case contributions either from the quartic coupling $\propto g S S \dagger A \mu A^\mu$ or through the trilinear coupling $\propto g S \dagger S A^\mu$. For heavier DM, additional annihilation channels are mandatory in order to get sufficiently large annihilation cross sections. In our case $S$ can annihilate into the scalars $H_{1,2}$ and $\phi$ through the Higgs portal couplings $F_i$ and $F'_j$. Owing to the scalar potential we considered, our model encompasses the cases discussed in [24, 25]. The case of a complex triplet can be recovered if we enforce $\delta_0 = 0$, while the case $\delta_{\pm} = 0$ resembles the real case (as the term $F_1'$ would identically vanish in such case), although we have twice more degrees of freedom.

When the mass parameter $\delta_{\pm}$ in eq. (3.3) is suppressed compared to both $\delta_0$ and $m_0$, the mass splitting between the neutral and the charged component of each pair $m_{S_L}/m_{S_H}$ is negligible. Coannihilations between neutral and charged components become, then, important. In the regime $\delta_{\pm} \ll \delta_0, m_0$, if $\delta_0$ is comparable to $m_0$, $m_{S_L}^0 \sim m_{\mp_L}^\pm$ is

\[\text{[13] The case discussed in [25] actually corresponds to } \delta_0^2 = 0 \text{ and } v_2 = v_\phi = 0.\]
FIG. 4. **Left**: dark matter mass, $m_{DM}$, as function of the mass parameter $\delta_0/m_0$ (see the text). **Right**: dark matter annihilation branching ratio versus $m_{DM}$: in blue are shown the annihilations into gauge bosons while the red area corresponds to annihilations into scalars.

much lighter than $m^{0}_{SH} \sim m^{\pm}_{SH}$ and a lower bound $m^{0}_{SL} \gtrsim 1.8$ TeV is then obtained [24, 25]. If, oppositely, $\delta_0 \ll m_0$, the number of annihilating particles effectively double and the lower bound $m^{0}_{SL} \gtrsim 1.8/\sqrt{2}$ TeV is reached.

To verify the validity of $S^0_L$ as a DM candidate, we implement our complete model in FeynRules [50] to generate the CalcHep [51] files used for micrOMEGAs [52]. We then scan over 19 parameters $^{14}$, and demand the DM relic density to lie within the $1\sigma$ range [3]: $\Omega_{DM} h^2 = 0.112 \pm 0.006$. We found, numerically, the following lower bound at $1\sigma$:

$$m_{DM} \gtrsim 1290 \text{ GeV}, \quad (3.5)$$

in agreement with [25]. As remarked above, this lower bound is obtained when the 4 $Z_2$–odd scalars contribute to the relic density, i.e. in the regime of low splitting $\delta_0 \ll m_0$. We illustrate this in the left panel of Fig. 4, where the values of $m_{DM}$ are plotted against the ratio $\delta_0/m_0$, recalling eq. (3.3). When $\delta_\pm \lesssim \delta_0 \ll m_0$, the DM mass reaches its minimum, while when $\delta_0$ becomes comparable to $m_0$, the lowest DM mass is about $m_{DM} \sim 1860$ GeV, the lowest values reached for small $\delta_\pm$ (corresponding to the case $\lambda_3 = 0$ of [25]). We also show, in the right panel of Fig. 4, the range the annihilating branching ratios span as function of $m_{DM}$. In blue are depicted the annihilations $S S \to V V$ ($V = W, Z$), while in red we show the branching fraction for $S S \to \phi_i \phi_j$, with $\phi_i$ any of the $Z_2$–even scalars present in our model. Over the mass range depicted, annihilations to gauge bosons mostly fix the DM relic abundance, as explained in [24, 25]. However for $m_{DM} \gtrsim 2$ TeV, the scalar contribution is necessary to compensate the $m_{DM}$ suppression of the annihilation cross section.

**B. Probing a scalar triplet dark matter**

Probing the triplet nature of the DM could be achievable, in principle, in direct and indirect detection experiments, and in collider searches. The indirect detection would consist in the observation of charged cosmic rays; however, the predicted fluxes are too suppressed compared to the background to be measured [24, 25]. Direct detection is more promising, and forthcoming experiment XENON 1T [53] can probe a part of the parameter space, corresponding to the larger values of $F_1^{(i)}$. As for colliders, the case for triplet scalars have been studied in [24, 54]. In our case with complex fields, the heavy particles decay to the light one plus a Majoron, $S^{0(\pm)}_H \to S^{0(\pm)}_L + J$, while the lightest charged scalar $S^{\pm}_L$ decays almost exclusively to $S^{0}_L \to S^{0}_L + \pi^{\pm}$. However, these channels provide no hope for probing the triplet nature of DM [24, 54]. Therefore, only if the quartic couplings $F_1^{(i)}$ are large enough one can hope to observe a characteristic signal of the triplet scalar $S$.

$^{14}$ 9 free parameters for the symmetry breaking potential and 10 for the DM sector. The fermion sector is fixed.
IV. NEUTRINO MASSES AND LEPTOGENESIS

Along the lines of [4], we explain the light neutrino masses within an inverse see-saw realization [5], using a minimal field content. In this scenario, a vectorlike fermion, singlet under the SM gauge group is introduced. This field has a charge -1 under $U(1)_X$, a symmetry constructed to generalize the $B-L$ quantum number. By the addition of two scalar fields, the doublet $H_2$ and the singlet $\phi$ introduced in the previous sections, a consistent UV completion of the inverse see-saw is possible.

In the present scenario we introduce a vectorlike Dirac fermion, $N_D$, that is a triplet of $SU(2)_L$ and is constructed from the RH neutrinos $N_{1,2}$, whose quantum numbers are reported in Table I: $N_D \equiv P_R N_1 + P_L N_2^C$. The interaction field $N_D$ couples to SM leptons via $H_{1,2}$ and $\phi$, and to the DM triplet scalar $S$. The couplings to leptons allow for an inverse see-saw generation of neutrino masses similar to the singlet case, and the coupling between $N_D$ and $S$ is necessary for the production of a baryon asymmetry in this model.

Indeed, as outlined in [4] and further discussed below, a nonzero $N_D S$ number density is produced in the early universe if one postulates the existence of a SM singlet Majorana fermion, $N_3$, which decays out of equilibrium in $N_D$ and $S$. This asymmetry is then transferred to SM lepton doublets through fast neutrino Yukawa interactions and partially converted into nonzero baryon number by in-equilibrium sphaleron nonperturbative processes.

Below we discuss the resulting neutrino mass spectrum and the main features of the leptogenesis mechanism.

A. Type III inverse see-saw realization

The most general interaction Lagrangian involving the Dirac field $N_D = (N_D^0, N_D^2) \, (\text{in the Cartesian basis})$ and the Majorana singlet $N_3$, is

$$ \mathcal{L} \supset -m_N \bar{N}_D^T N_D^0 - \left( \frac{Y_N^{\alpha}}{\sqrt{2}} \bar{N}_D^T \tilde{H}_1^* (T_2^D)_{jk} \ell_k + Y_{\nu}^2 \bar{N}_D^T \tilde{H}_2^* (T_2^D)_{jk} \ell_k + \frac{\delta_N}{\sqrt{2}} \phi \bar{N}_D^T N_D^0 C + \text{h.c.} \right) $$

$$ - \frac{1}{2} M_3 \bar{N}_3 N_3^C - \left( h S^a \bar{N}_D^T N_3 + \text{h.c.} \right), $$

where $\ell_\beta = (\ell_e, \ell_\mu, \ell_\tau)^T$ ($\beta = e, \mu, \tau$), $N_D^C \equiv C \bar{N}_D^T (a = 1, 2, 3)$ and $\tilde{H}_k = -i \sigma_2 H_k^* (k = 1, 2)$. The parameter $\delta_N$ is made real through a phase transformation. The first line of (4.1) contains terms providing masses to neutrinos after EWSB, while the second line contains additional terms required for leptogenesis. In the spherical basis of the adjoint representation, the components of $N_D$ are

$$ N_D = (N_D^0, N_D^2, N_D^2) = (P_R N_1^0 + P_L (N_2^0)^C, \, P_R N_1^0 + P_L (N_2^0)^C, \, P_R N_1^0 + P_L (N_2^0)^C), $$

and in particular $(N_D^0)^C \neq N_D^2$, similarly to the complex scalar triplet $S$.

Notice that the Yukawa interactions $\propto Y_{\nu}^2$ couple $N_1, N_2$ to the SM leptons. Therefore, after the Higgs doublets acquire a nonzero vev, the SM lepton number (i.e. the generalized $X$ charge) is explicitly violated by $Y_{\nu}^2$ mediated interactions. Furthermore, while the Dirac type mass $m_N$ conserves the lepton number, the term $\propto \delta_N$ provides, after $\phi$ takes a nonzero vev, a Majorana mass term for both $N_1$ and $N_2$. In the case $m_N \gg \delta_N v_\phi$ we have, in general, the inverse see-saw mechanism [5]. Finally, as in usual type III see-saw scenarios [9], the charged components of the triplet and the SM leptons mix after EWSB, implying typically larger contribution to lepton flavour violation processes than the singlet RH neutrino case (see e.g. [55–58]).

In the chiral basis $(\nu_L, (N_1^0)^C_L, (N_2^0)^C_L)$, the $5 \times 5$ symmetric neutrino mass matrix reads:

$$ \mathcal{M}_\nu = \begin{pmatrix} 0_{3 \times 3} & y_1 v_1 & y_2 v_2 \\ y_1^T v_1 & \delta_N v_\phi & m_N \\ y_2^T v_2 & m_N & \delta_N v_\phi \end{pmatrix}. $$

In eq. (4.3) $0_{3 \times 3}$ is the null matrix of dimension 3 and we introduce the shorthand notation: $y_k \equiv (Y_{\nu e}^k, Y_{\nu \mu}^k, Y_{\nu \tau}^k)^T / 2 \sqrt{2}$ with $k = 1, 2$. The neutral spectrum, therefore, consists of one massless neutrino, two massive light Majorana neutrinos and two heavy Majorana neutrinos $N_{1,2}$. These two particles are quasidegenerate, with mass $M_{N_{1,2}} = m_N \mp \delta_N v_\phi$, and form a pseudo-Dirac pair in the limit $m_N \gg \delta_N v_\phi$ [59–61].

The effective light neutrino mass matrix has elements

$$ (M_\nu)^{\alpha \beta} \simeq -\frac{v_2 v_1}{m_N} \left( y_1^\alpha y_2^\beta + y_2^\alpha y_1^\beta \right) + \frac{\nu_1}{m_N} \left( y_1^\alpha y_1^\beta + y_2^\alpha y_2^\beta \frac{v_2^2}{v_1^2} \right). $$

(4.4)
This expression clearly highlights the different contribution to the light neutrino masses. The first term acts as a linear seesaw, and its suppression originates from the small vev \( v_2 \). The second term is typical of inverse seesaw scenarios, where the small ratio \( v_2 \delta_N/m_N \) suppresses the neutrino mass scale. We remark that, since only two RH neutrinos \( N_{i,2} \) are introduced, the linear seesaw contribution alone (i.e., neglecting \( v_2 \phi \) in (4.4)) allows us to fit all current neutrino oscillation data, while if \( v_2 = 0 \) the inverse seesaw scenario can only account for one massive light neutrino. Therefore, the complex scalar field \( \phi \), with vev \( v_\phi \neq 0 \), is not necessary in order to obtain two massive light neutrinos through the (linear) see-saw mechanism. On the other hand, \( v_\phi \neq 0 \) is a necessary condition to set a hierarchy between the Higgs doublet vevs, \( v_2 \ll v_1 \), without fine-tuning of the parameters [4]. Finally, we remark that the scalar field \( \phi \) with a coupling \( \delta_N \neq 0 \) allows us to implement successful leptogenesis through the “two-step” scenario discussed below.

Using as shorthand notation \( A = M_\nu M_\nu^\dagger \), the two nonzero neutrino masses are given by

\[
m_\nu^+ = \frac{1}{\sqrt{2}} \sqrt{\text{Tr}(A) \pm \sqrt{2 \text{Tr}(A^2) - \text{Tr}(A)^2}}.
\]

In terms of the see-saw parameters in the Lagrangian, we have

\[
\text{Tr}(A) = \frac{1}{m_N^2 (1 - \mu_N^2)^2} \left[ 2 \left( y_1^2 y_2^2 + |y_{12}|^2 \right) - 4 \mu_N (y_1^2 + y_2^2) \text{Re}(y_{12}) + \mu_N^2 \left( y_1^2 + y_2^2 + 2 \text{Re}(y_{12})^2 \right) \right],
\]

\[
2 \text{Tr}(A^2) - \text{Tr}(A)^2 = \frac{1}{(m_N^2 (1 - \mu_N^2)^2)^2} \times \left[ 4 \left| y_{12} \right|^2 - 4 \mu_N (y_1^2 + y_2^2) \text{Re}(y_{12}) + \mu_N^2 \left( y_1^2 + y_2^2 + 2 \text{Re}(y_{12})^2 \right) \right] \times \left[ 4 y_1^2 y_2^2 - 4 \mu_N (y_1^2 + y_2^2) \text{Re}(y_{12}) + \mu_N^2 \left( y_1^2 + y_2^2 + 2 \text{Re}(y_{12})^2 \right) \right] + 4 \mu_N^2 (2 - \mu_N^2) |y_{12}|^2,
\]

where we introduce for convenience \( y_i = \sqrt{y_{11}^i \cdot y_i} v_i, y_{12} = y_1^i \cdot y_2 v_1 v_2, \eta_{12} = y_1 \times y_2 v_1 v_2 \) and \( \mu_N = (\delta_N v_\phi)/m_N \).

In the inverse seesaw limit, \( \mu_N \ll 1 \), the neutrino masses have a simple expression

\[
m_\nu^+ \approx \frac{1}{m_N} \left( \frac{y_1^2 y_2^2 - \mu_N (y_1^2 + y_2^2) \text{Re}(y_{12}) \mp \sqrt{|y_{12}|^2 - \mu_N (y_1^2 + y_2^2) \text{Re}(y_{12})}}{y_1 y_2 |y_{12}|} \right),
\]

\[
\approx \frac{1}{m_N} \left( y_1 y_2 \pm |y_{12}| \right) \left( 1 \mp \frac{\mu_N (y_1^2 + y_2^2) \text{Re}(y_{12})}{2 y_1 y_2 |y_{12}|} \right).
\]

Notice that if the neutrino Yukawa vectors \( y_1 \) and \( y_2 \) are proportional, \( m_\nu^+ \) is exactly zero, in contrast with neutrino oscillation data. For a normal hierarchical spectrum, \( m_\nu^+ = \sqrt{|\Delta m^2_\text{ATM}|} \) and \( m_\nu^- = \sqrt{|\Delta m^2_\text{SOL}|} \), while in the case of inverted hierarchy we have \( m_\nu^+ = \sqrt{|\Delta m^2_\text{ATM}|} \) and \( m_\nu^- = \sqrt{|\Delta m^2_\text{SOL}| - |\Delta m^2_\text{SOL}|} \) and \( |\Delta m^2_\text{SOL}| \) being the atmospheric and solar neutrino mass square differences, respectively. It is easy to show that at leading order in \( \mu_N \), the neutrino mass parameters satisfy the following relation [4]: \( |y_1 \times y_2| v_1 v_2/m_N \approx (\Delta m^2_\text{ATM}/|\Delta m^2_\text{SOL}|)^{1/4} \). Hence, barring accidental cancellations, the size of the neutrino Yukawa couplings is

\[
|y_1| |y_2| \approx 10^{-8} (m_N/1 \text{ TeV})(10 \text{ MeV}/v_2).
\]

The model also accommodates two charged Dirac fermions, \( \Sigma_{1,2}^\pm \), which at tree level are degenerate with the heavy neutral fermions \( N_{1,2} \). Similarly to the triplet scalar case, the triplet gauge couplings induce a mass splitting \( \simeq 166 \text{ MeV} \) [24] between the charged and neutral components of the triplet.

Production cross section at LHC is dominated by the triplet gauge interactions and scales with the overall mass of the triplet, \( m_N \). A discovery of these charged (\( \Sigma_{1,2}^\pm \)) and neutral (\( N_{1,2} \)) fermions may be possible at LHC if \( m_N \lesssim 1 \text{ TeV} \) [62–66]. Current searches constrain the triplet mass to be above \( m_N \gtrsim (180-210) \text{ GeV} \) range [67].

### B. Leptogenesis mechanism

As anticipated at the beginning of this section, for a Majorana fermion \( N_3 \) heavier than the SU(2)\(_L\) triplet fields \( N_D \) and \( S \), successful leptogenesis can be realized in this model within the “two-step” scenario introduced in [4]:

1. When the temperature of the universe decreases below \( M_3 \), \( N_3 \) out of equilibrium decays generate asymmetries in \( S \) and \( N_D \) abundances. A nonzero CP asymmetry, \( \epsilon_{CP} \), originates from the interference between the tree level and one-loop correction of \( N_3 \) decay amplitude [4] (see Fig. 5). We emphasize that the presence of \( \phi \) with a coupling \( \delta_N \neq 0 \) in (4.1) is mandatory in order to generate a nonzero \( \epsilon_{CP} \). The expression of \( \epsilon_{CP} \) is in general lengthy, but shows a very simple dependence on the key parameters \( \delta_N \), \( h \) and \( \mu'' \) in the hierarchical limit \( M_3^2 \gg m_N^2, m_\nu^2 \).
Indeed, the triplets of sphaleron processes, hence heavy enough triplets are required. In our case, the asymmetry is generated initially by scatterings maintain fermion triplets in thermal equilibrium down to low temperatures, possibly after the decoupling is set \[69, 70\]. A similar bound is found here, although for very different reasons. In the standard type III case, gauge percentage. In the singlet case we found no actual lower bound on the \(\Delta N = 1\) washout scatterings involving \(S\) and the dark matter fields are controlled by the leptogenesis scale \(T \sim M_3\), the see-saw scale \(m_N\) and the dark matter dimensional parameter \(\mu_S\).

We compute the final baryon asymmetry \(Y_B\) within the present leptogenesis scenario for several values of the relevant mass scales and couplings of the model and we compare it with the observed value \[68\]: \(Y_B^\text{exp} = (8.77 \pm 0.21) \times 10^{-11}\). We solve for this a system of coupled Boltzmann equations for \(N_3, N_D, S\) and lepton asymmetries. Typical results of the numerical analysis are given in Fig. 6, where we show the variations of \(Y_B\) against several key parameters. In both plots we fix for illustration \(|h| = 10^{-6}\) and \(\mu_S = 2\) TeV. The value of \(h\) is about the maximal possible, while for \(\mu_S = 2\) TeV the DM mass is \(m_{\text{DM}} = (1930 \pm 70)\) GeV. In the left panel of Fig. 6, we highlight the dependence on the see-saw scale \(m_N\) and on \(M_3\), through a degeneracy parameter equal to \(1 - (m_N + \mu_S)/M_3\) (in percentage). In the singlet case we found no actual lower bound on \(m_N\) from BAU requirement: clearly in the triplet case the picture changes. As is well known, in type III leptogenesis, a lower bound on the see-saw scale \(m_N \gtrsim 1.6\) TeV is set \[69, 70\]. A similar bound is found here, although for very different reasons. In the standard type III case, gauge scatterings maintain fermion triplets in thermal equilibrium down to low temperatures, possibly after the decoupling of sphaleron processes, hence heavy enough triplets are required. In our case, the asymmetry is generated initially by the decay of a singlet \(N_3\), therefore not affected by such effects. Instead, the lower bound on \(N_D\) mass is related to the strength \(N_D\) interacts with the singlet and gauge bosons: for a fixed \(M_3\) in the TeV range, the lower the \(m_N\), the stronger the \(\Delta N = 1\) washout processes, which tend to maintain \(N_3\) in equilibrium. The dependence of \(Y_B\) on the \(N_3\) mass scale is manifest in Fig. 6: a degenerate spectrum, with \(M_3\) close to the threshold \(m_N + \mu_S\) is favored. \(^{15}\)

\(^{15}\) Indeed, the triplets \(S\) and \(N_D\) receive thermal mass corrections \(\propto g T\), and \(N_3\) decays are kinematically open at a temperature \(T_D\) actually smaller than \(M_3\). The closest \(M_3\) is to the \(T = 0\) threshold \(m_N + \mu_S\), the lowest is \(T_D\), resulting in a Boltzmann suppression of the washout scatterings involving \(N_3\).
An absolute lower bound on $m_N$ could be formally derived, however subject to a certain tuning of the parameters. Typically $m_N \gtrsim 1.5$ TeV allows our mechanism to work. Such heavy masses cannot be probed at LHC, or oppositely, if a triplet fermion of mass $\lesssim 1$ TeV is observed, the scenario discussed here is not responsible for the observed BAU.

The dependence of the baryon asymmetry on $\mu''$ and $\delta_N$ is presented in the right panel of Fig. 6. For illustration, we fix this time $m_N = \mu_S = 2$ TeV, while $M_3 = 4.4$ TeV (10% degeneracy). We see the strong influence of $\delta_N$ on $Y_B$ from the rise of the washout effects, implying $\delta_N \lesssim 10^{-6}$. In this case, $\epsilon_{CP}$ increases with $\mu''$ (eq. (4.10)), as well as the baryon asymmetry. Not illustrated in this plot is the closing of the successful region in green for values of $\mu''$ larger than 3 TeV, due to an increase of the washout processes. However, note that for a given $\mu_S$, $\mu''$ cannot be arbitrarily large, as it enters in DM mass expression (3.3).

As a concluding remark about this section, we emphasize here the possibility of a low-energy realization of the leptogenesis scenario, although with some tuning of the parameters. Going to higher mass scale for $N_3$ relaxes such tuning.

V. DISCUSSION AND CONCLUSIONS

In this work we consider the model of [4], implemented with triplet fermions instead of singlets. The Standard Model particle content is extended with additional fermion and scalar representations:

i) an extra complex Higgs doublet $H_2$;
ii) a singlet complex scalar $\phi$;
iii) one SU(2)$_L$ triplet Dirac fermion $N_D$;
iv) a complex scalar triplet of SU(2)$_L$ $S$;
v) a Majorana singlet $N_3$.

The particles listed in i) – iii) naturally realize a low scale inverse see-saw mechanism of neutrino mass generation. The scalar iv) provides a viable dark matter candidate, and with the two representations in iii) – v) allows us to explain the observed baryon asymmetry of the universe, through a variant of the standard thermal leptogenesis mechanism. The overall model is symmetric under a global U(1)$_X$ transformation, which is spontaneously broken by $H_2$ and $\phi$ vevs at the electroweak scale. The corresponding vevs break the generalized $B - L$ number (see Table I), implying nonzero neutrino masses. Furthermore, this symmetry is broken down to a remnant $Z_2$, which stabilizes the DM. A Goldstone boson emerges from the spontaneous breaking of the symmetry, the Majoron $J$.

This model is characterized by 6 $Z_2$-even scalar particles: 1 charged $H^\pm$, 3 CP-even neutral scalars $h^0$, $H^0$ and $h_A$, and 2 CP-odd scalar particles $A^0$ and $J$. The lightest CP-even scalar $h^0$ corresponds to ATLAS [1] and CMS [2] observations of a $m_{h^0} \simeq 126$ GeV scalar boson. The features of this model may allow us to discriminate it against the SM at LHC. In particular, the presence of a charged scalar $H^\pm$ can reduce or largely enhance the $h^0 \to \gamma \gamma$ decay width compared to the SM expectation. Similarly, the invisible decay channel $h^0 \to JJ$ can be sizable and affect the branching ratios of the Higgs boson to the SM particles. Observations of large deviations from the SM case, as well as of other neutral scalar excitations, would constitute the most distinctive signatures of the scalar spectrum presented
here. Additional experimental results could help distinguish this model from 2HDM, or other scalar extensions of the SM.

The remnant $Z_2$ symmetry guarantees the stability of the DM. We have a total of 4 $Z_2$-odd scalars, which arise from the complex triplet $S$: 2 charged $S_{L,H}^\pm$ and 2 neutral particles $S_{L,H}^0$. The lightest neutral scalar $S_L^0$ provides a viable DM candidate, as long as the quartic coupling $F_1$ contribution to its mass, which splits $m_{S_L}^0$ and $m_{S_L}^\pm$, is smaller than the one-loop gauge corrections to $m_{S_L}^\pm$. We obtain numerically that the lightest viable DM candidate has $m_{DM} \gtrsim 1290$ GeV, consistently with previous works on scalar triplet DM [24, 25].

Light neutrino masses originate from the spontaneous breaking of the generalized $B-L$ charge, instead of relying on explicit lepton number breaking. Two heavy Majorana particles, $N_{1,2}$ quasidegenerate, are obtained, in addition to two massive and one massless light neutrino. Heavy fermions with mass $m_N \lesssim O$(TeV) can in principle be probed at the LHC, thanks to their couplings to SU(2) gauge bosons.

The baryon asymmetry of the universe can be explained within a two-step leptogenesis mechanism [4], provided one Majorana fermion, $N_3$, is included with a direct coupling to both triplet representations $S$ and $N_D$, see eq. (4.1). We obtain that an approximate lower bound on the see-saw scale $m_N \gtrsim 1500$ GeV is required to comply with the observed baryon asymmetry.

The model we present can easily explain and overcome several issues of the standard theory: dark matter abundance, baryon asymmetry of the universe and neutrino oscillation data. The observation at LHC of a $O$(TeV) fermion triplet, even if consistent with $N_D$ properties, would exclude the baryogenesis mechanism proposed here, although not the inverse see-saw model itself. When confronted with additional experimental data, a systematic study of the phenomenology of the model may allow us to significantly constraint the parameter space and eventually rule out this entire scenario.

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[1] G. Aad et al. [ATLAS Collaboration], Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC, Phys. Lett. B [arXiv:1207.7214].
[2] S. Chatrchyan et al. [CMS Collaboration], Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC, Phys. Lett. B [arXiv:1207.2355].
[3] J. Beringer et al. [Particle Data Group Collaboration], Review of Particle Physics (RPP), Phys. Rev. D 86 (2012) 010001.
[4] F.X. Josse-Michaux and E. Molinaro, A Common Framework for Dark Matter, Leptogenesis and Neutrino Masses, Phys. Rev. D 84 (2011) 125021 [arXiv:1108.0482].
[5] R. N. Mohapatra and J. W. F. Valle, Neutrino Mass and Baryon Number Nonconservation in Superstring Models, Phys. Rev. D 34 (1986) 1642.
[6] M. Fukugita and T. Yanagida, Baryogenesis Without Grand Unification, Phys. Lett. B 174 (1986) 45.
[7] S. Davidson, E. Nardi and Y. Nir, Leptogenesis, Phys. Rept. 466 (2008) 105 [arXiv:0802.2962].
[8] P. Minkowski, Phys. Lett. B 67 (1977) 421; M. Gell-Mann, P. Ramond and R. Slansky, Proceedings of the Supergravity Stony Brook Workshop, New York 1979, eds. P. Van Nieuwenhuizen and D. Freedman; T. Yanagida, Proceedings of the Workshop on Unified Theories and Baryon Number in the Universe, Tsukuba, Japan 1979, eds A. Sawada and A. Sugamoto; R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. 44 (1980) 912.
[9] R. Foot, H. Lew, X. G. He and G. C. Joshi, Seesaw Neutrino Masses Induced By A Triplet Of Leptons, Z. Phys. C 44 (1989) 441; E. Ma, Pathways to naturally small neutrino masses, Phys. Rev. Lett. 81 (1998) 1171 [hep-ph/9805219].
[10] K. Dick, M. Lindner, M. Ratz and D. Wright, Leptogenesis with Dirac neutrinos, Phys. Rev. Lett. 84 (2000) 4039 [hep-ph/9907562].
[11] H. Murayama and A. Pierce, Realistic Dirac leptogenesis, Phys. Rev. Lett. 89 (2002) 271601 [hep-ph/0206177].
[12] P. H. Gu and H. -J. He, Neutrino Mass and Baryon Asymmetry from Dirac Seesaw, JCAP 0612 (2006) 010 [hep-ph/0610275];
[13] N. Sahu and U. Sarkar, Extended Zee model for Neutrino Mass, Leptogenesis and Sterile Neutrino like Dark Matter, Phys. Rev. D 78 (2008) 115013 [arXiv:0804.2072].
[14] M. C. Gonzalez-Garcia, J. Racker and N. Rius, Leptogenesis without violation of B-L, JHEP 0911 (2009) 079 [arXiv:0909.3518].
[15] K. Kohri, A. Mazumdar, N. Sahu and P. Stephens, Probing Unified Origin of Dark Matter and Baryon Asymmetry at PAMELA/Fermi, Phys. Rev. D 80 (2009) 061302 [arXiv:0907.0622].
[44] http://cdsweb.cern.ch/record/1460419
[45] http://cdsweb.cern.ch/record/1460664
[46] http://cdsweb.cern.ch/record/1460424
[47] D. Carmi, A. Falkowski, E. Kuflik and T. Volansky, Interpreting the Higgs, arXiv:1206.4201.
[48] W. Grimus, L. Lavoura, O. M. Ogreid and P. Osland, The Oblique parameters in multi-Higgs-doublet models, Nucl. Phys. B 801 (2008) 81 [arXiv:0802.4353].
[49] http://twiki.cern.ch/twiki/bin/view/LHCPhysics/CERNYellowReportPageBR2
[50] N. D. Christensen and C. Duhr, FeynRules - Feynman rules made easy, Comput. Phys. Commun. 180 (2009) 1614 [arXiv:0806.4194].
[51] A. Pukhov, CalcHEP 2.3: MSSM, structure functions, event generation, batches, and generation of matrix elements for other packages, hep-ph/0412191.
[52] G. Belanger, F. Boudjema, A. Pukhov and A. Semenov, micrOMEGAs : a tool for dark matter studies, arXiv:1005.4133.
[53] https://twiki.cern.ch/twiki/bin/view/LHCPhysics/CERNYellowReportPageBR2
[54] N. D. Christensen and C. Duhr, FeynRules - Feynman rules made easy, Comput. Phys. Commun. 180 (2009) 1614 [arXiv:0806.4194].
[55] A. Abada, C. Biggio, F. Bonnet, M. B. Gavela and T. Hambye, Low energy effects of neutrino masses, JHEP 0712 (2007) 061 [arXiv:0707.4058].
[56] A. Abada, C. Biggio, F. Bonnet, M. B. Gavela and T. Hambye, $\mu \rightarrow e\gamma$ and $\tau \rightarrow \ell\gamma$ decays in the fermion triplet seesaw model, Phys. Rev. D 78 (2008) 033007 [arXiv:0803.0481].
[57] D. N. Dinh, A. Ibarra, E. Molinaro and S. T. Petcov, Low Energy Signatures of the TeV Scale See-Saw Mechanism, Phys. Rev. D 84 (2011) 013005 [arXiv:1005.4133].
[58] D. N. Dinh, A. Ibarra, E. Molinaro and S. T. Petcov, The $\mu - e$ Conversion in Nuclei, $\mu \rightarrow e\gamma$, $\mu \rightarrow 3e$ Decays and TeV Scale See-Saw Scenarios of Neutrino Mass Generation, JHEP 1208 (2012) 125 [arXiv:1205.4671].
[59] L. Wolfenstein, Different Varieties of Massive Dirac Neutrinos, Nucl. Phys. B 186 (1981) 147.
[60] S. T. Petcov, On Pseudodirac Neutrinos, Neutrino Oscillations and Neutrinoless Double beta Decay, Phys. Lett. B 110 (1982) 245.
[61] G. C. Branco, W. Grimus and L. Lavoura, The Seesaw Mechanism in the Presence of a Conserved Lepton Number, Nucl. Phys. B 312, 492 (1989).
[62] R. Franceschini, T. Hambye and A. Strumia, Type-III see-saw at LHC, Phys. Rev. D 78 (2008) 033002 [arXiv:0805.1613].
[63] F. del Aguila and J. A. Aguilar-Saavedra, Distinguishing seesaw models at LHC with multi-lepton signals, Nucl. Phys. B 813 (2009) 22 [arXiv:0809.2489].
[64] F. del Aguila and J. A. Aguilar-Saavedra, Electroweak scale see-saw and heavy Dirac neutrino signals at LHC, Phys. Lett. B 672 (2009) 158 [arXiv:0809.2096].
[65] O. J. P. Eboli, J. Gonzalez-Fraile and M. C. Gonzalez-Garcia, Neutrino Masses at LHC: Minimal Lepton Flavour Violation in Type-III See-saw, JHEP 1112 (2011) 009 [arXiv:1108.0661].
[66] A. Arhrib, B. Bajc, D. K. Ghosh, T. Han, G.-Y. Huang, I. Puljak and G. Senjanovic, Collider Signatures for Heavy Lepton Triplet in Type I+III Seesaw, Phys. Rev. D 82 (2010) 053004 [arXiv:0904.2390].
[67] S. Chatrchyan et al. [CMS Collaboration], Search for heavy lepton partners of neutrinos in proton-proton collisions in the context of the type III seesaw mechanism, arXiv:1210.1797.
[68] E. Komatsu et al. [WMAP Collaboration], Seven-Year Wilkinson Microwave Anisotropy Probe (WMAP) Observations: Cosmological Interpretation, Astrophys. J. Suppl. 192 (2011) 18 [arXiv:1001.4538].
[69] A. Strumia, Sommerfeld corrections to type-II and III leptogenesis, Nucl. Phys. B 809 (2009) 308 [arXiv:0806.1630].
[70] D. Aristizabal Sierra, J. F. Kamenik and M. Nemevsek, Implications of Flavor Dynamics for Fermion Triplet Leptogenesis, JHEP 1010 (2010) 036 [arXiv:1007.1907].