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Double Type-II Seesaw, Baryon Asymmetry and Dark Matter for Cosmic e^{\pm} Excesses

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We construct a new realization of type-II seesaw for neutrino masses and baryon asymmetry by extending the standard model with one light and two heavy singlet scalars besides one Higgs triplet. The heavy singlets pick up small vacuum expectation values to give a suppressed trilinear coupling between the triplet and doublet Higgs bosons after the light singlet drives the spontaneous breaking of lepton number. The Higgs triplet can thus remain light and be accessible at the LHC. The lepton number conserving decays of the heavy singlets can generate a lepton asymmetry stored in the Higgs triplet to account for the matter-antimatter asymmetry in the Universe. We further introduce stable gauge bosons from a hidden sector, which obtain masses and annihilate into the Higgs triplet after spontaneous breaking of the associated non-Abelian gauge symmetry. With Breit-Wigner enhancement, the stable gauge bosons can simultaneously explain the relic density of dark matter and the cosmic positron/electron excesses.

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I. INTRODUCTION

Observations of solar, atmospheric, reactor and accelerator neutrino oscillations have strongly pointed to tiny but nonzero neutrino masses [1]. The smallness of neutrino masses is naturally explained by the seesaw [2] extension of the standard model (SM). In the seesaw scenario, the observed baryon asymmetry in the Universe can be generated through the leptogenesis [3, 4]. The type-II seesaw models [5] generically contain a trilinear coupling between the triplet and doublet Higgs bosons, which is a unique source for the lepton number violation. This interaction is the key ingredient to realize leptogenesis in the type-II seesaw [6]. If this lepton number violation is very weak, a lepton asymmetry stored in the Higgs triplet can survive and thus be transferred to the lepton doublets. Subsequently the sphaleron process [7] can partially convert this lepton asymmetry to the observed baryon asymmetry in the Universe. When the type-II seesaw is implemented in models of large extra dimensions with lepton number breaking in a distant brane, it can lead to interesting collider phenomenology by predicting a light Higgs triplet with significant couplings to leptons [8].

In this paper, we construct a new double type-II seesaw scenario to simultaneously generate a suppressed coupling of the Higgs triplet to the Higgs doublet and a lepton asymmetry stored in the Higgs triplet. This is realized by introducing two heavy and one light singlet scalars to the type-II seesaw model, where the Higgs triplet has a TeV-scale mass and its Yukawa couplings to the lepton doublets are naturally $\mathcal{O}(1)$, so it is within the discovery reach of the LHC. When the light singlet develops a vacuum expectation value (VEV), the lepton number will be spontaneously broken at the TeV scale. The heavy singlets will then pick up small VEVs to give the suppressed coupling between the Higgs triplet and doublet. These heavy singlets are also responsible for generating the lepton asymmetry stored in the Higgs triplet through their out-of-equilibrium decays which violate CP but conserve lepton number.

On the other hand, strong evidences for the nonbaryonic dark matter relic abundance [9] require supplementing additional new ingredients to the existing theory. The dark matter may also be responsible for the positron/electron excesses in the cosmic rays as observed by ATIC [10], PPB-BETS [11], PAMELA [12], HESS [13] and Fermi/LAT [14] collaborations. This indicates that the dark matter should mostly annihilate or decay into leptons with large cross section or long lifetime. Such type of dark matter may have special relation with the neutrino mass-generation [15]. In this study, we extend the double type-II seesaw model for the neutrino masses to explain the relic abundance of dark matter [16] as well as the cosmic positron/electron excesses. In a bosonic hidden sector, the vector bosons associated with a non-Abelian gauge group can be stable since they are forbidden to mix with the gauge bosons of the SM due to the non-Abelian character. Such type of hidden vector bosons can explain the relic density of dark matter in the Universe [17]. In the present construction, these hidden vector bosons will annihilate into the Higgs triplet in our double type-II seesaw scenario and the Higgs triplet dominantly decays into lepton pairs. The annihilation process of the hidden vector bosons into the Higgs triplets invokes an s-channel exchange of a Higgs boson (which is also responsible for the mass generation of dark matter). With the Breit-Wigner resonant enhancement [18, 19, 20, 21],

the cross section could be large enough to account for the cosmic-ray positron/electron excess without introducing any other additional boost factor. In this way the hidden vector bosons can naturally serve as a new candidate of leptonic dark matter [15, 22, 23].

II. DOUBLE TYPE-II SEESAW

The type-II seesaw mechanism is realized by extending the SM with a Higgs triplet. Under the SM gauge group, this Higgs triplet is allowed to have a Yukawa coupling with the lepton doublets,

$$\mathcal{L}_{\rm Y} \supset -\frac{1}{2} f \,\overline{\psi_L^c} i \tau_2 \xi \psi_L + \text{h.c.} \,, \tag{1}$$

and a trilinear interaction with the Higgs doublet in the scalar potential,

$$V \supset -\mu_0 \phi^T i \tau_2 \xi \phi + \text{h.c.}, \qquad (2)$$

where ψ_L and ϕ denote the lepton and Higgs doublets, respectively, while ξ is the Higgs triplet. In the presence of trilinear interaction (2), the Higgs triplet ξ can develop a small VEV once the Higgs doublet ϕ acquires its VEV to break the electroweak symmetry. Such a small VEV of ξ will thus naturally generate tiny neutrino masses through the Yukawa coupling (1). Conventionally, we assign the Higgs triplet ξ a lepton number L = -2 as the lepton doublet ψ_L has L = +1. In consequence, the Yukawa interaction (1) is lepton-number conserving, while the trilinear interaction (2) softly and explicitly breaks the lepton number because of L = 0 for the Higgs doublet ϕ .

To explain the origin of the lepton number violation in the type-II seesaw model, it is desirable to naturally start with a lepton number conserving Lagrangian and then break it spontaneously. This can be achieved by introducing a singlet scalar with a global lepton number L = +2, similar to the singlet majoron model [24]. Then the cubic coupling in the lepton number violating interaction (2) should be proportional to the breaking scale of lepton number.

In the present study, we propose a more attractive scenario by extending the type-II seesaw model with heavy and light singlet scalars. For simplicity, we will not give the complete scalar potential, instead we write the relevant part for our analysis as follows,

$$V \supset -m_1^2 (\sigma^{\dagger} \sigma) + \lambda_1 (\sigma^{\dagger} \sigma)^2 - m_2^2 (\phi^{\dagger} \phi) + \lambda_2 (\phi^{\dagger} \phi)^2 + M_{\chi}^2 (\chi^{\dagger} \chi) - (\mu \chi \sigma^2 + \text{h.c.}) + m_{\xi}^2 \text{Tr} (\xi^{\dagger} \xi) - (\kappa \chi \phi^T i \tau_2 \xi \phi + \text{h.c.}), \qquad (3)$$

where σ and χ are the light and heavy singlets, respectively. For the lepton number conservation in the

TABLE I: Summary of relevant quantum number assignments in the present model, where the new gauge boson X^{μ}_{μ} and Higgs doublet η under the hidden gauge group $SU(2)_h$ will be explicitly defined in Sec. IV.

| | $SU(2)_L$ | $U(1)_Y$ | $U(1)_{\text{lepton}}$ | $SU(2)_h$ |
|-----------|---------------|----------------|------------------------|---------------|
| Fields | (I_{L3}) | (Y) | (L^{π}) | (I_{h3}) |
| ψ_L | $\frac{1}{2}$ | $-\frac{1}{2}$ | 1 | 0 |
| ϕ | $\frac{1}{2}$ | $-\frac{1}{2}$ | 0 | 0 |
| ξ | 1 | 1 | -2 | 0 |
| χ | 0 | 0 | 2 | 0 |
| σ | 0 | 0 | -1 | 0 |
| η | 0 | 0 | 0 | $\frac{1}{2}$ |
| X^a_μ | 0 | 0 | 0 | 1 |

scalar potential (3), we have assigned L = +2 for χ and L = -1 for σ . A summary of the relevant quantum number arrangement of our model is given in Table. 1.

With proper choice of parameters, the singlet scalar σ is expected to develop a vacuum expectation value of the order of TeV¹ to drive the spontaneous symmetry breaking of the global lepton number. The vacuum expectation value $\langle \sigma \rangle$ will induce a small VEV for the singlet χ in the presence of the large mass term and the trilinear coupling, as shown in the second line of Eq. (3),

$$\langle \chi \rangle \simeq \frac{\mu \langle \sigma \rangle^2}{M_\chi^2} \,.$$
 (4)

It is evident that the vacuum expectation value $\langle \chi \rangle$ is highly suppressed for $M_{\chi} \gtrsim \mu \gg \langle \sigma \rangle$. This clearly shares the essential feature with the traditional seesaw mechanism. The small vacuum expectation value $\langle \chi \rangle$ will then induce a suppressed coupling of the Higgs triplet ξ to the Higgs doublet ϕ from the last terms of Eq. (3). According to Eq. (2), we thus have,

$$\mu_0 = \kappa \langle \chi \rangle = \kappa \frac{\mu \langle \sigma \rangle^2}{M_\chi^2} , \qquad (5)$$

which is naturally small for $\kappa \leq \mathcal{O}(1)$.

Subsequently, the Higgs doublet ϕ develops a VEV, $\langle \phi \rangle \simeq 174 \text{ GeV}$, to break the electroweak gauge symmetry $SU(2)_L \otimes U(1)_Y$. So the Higgs triplet ξ can pick up a small VEV through the type-II seesaw mechanism,

$$\langle \xi \rangle \simeq \frac{\mu_0 \langle \phi \rangle^2}{m_{\xi}^2}.$$
 (6)

¹ It is clear that, given $\langle \sigma \rangle = \mathcal{O}(\text{TeV})$ and without fine-tuning the parameters, the mass of the physical Higgs boson naturally lies at the TeV scale.



FIG. 1: The new double type-II seesaw for the generation of Majorana neutrino masses.

We note that the VEV $\langle \xi \rangle$ is naturally small even for a weak-scale triplet mass $m_{\xi} = \mathcal{O}(\text{TeV}) \gtrsim \mathcal{O}(\langle \phi \rangle)$, thanks to the highly suppressed trilinear coupling μ_0 in Eq. (5) through the tiny VEV $\langle \chi \rangle$ in Eq. (4). At this stage, the neutrinos eventually obtain tiny Majorana masses through their Yukawa couplings with the Higgs triplet as shown in Eq. (1),

$$m_{\nu} = f\langle \xi \rangle \,. \tag{7}$$

Therefore, the small neutrino masses $m_{\nu} \ll \langle \phi \rangle$ can be naturally realized for the Yukawa couplings $f = \mathcal{O}(1)$. Remarkably, our new construction includes an additional seesaw step to realize the type-II seesaw mechanism as shown in Fig. 1, and may thus be called "double type-II seesaw" to reflect this new ingredient².

For illustration, we give a numerical sample here. Let us take the typical inputs, $M_{\chi} = 10^{14} \,\text{GeV}, \ \mu = 3 \times 10^{12} \,\text{GeV}, \ \langle \sigma \rangle = 1.2 \,\text{TeV}, \ m_{\xi} = 0.5 \,\text{TeV}, \ \text{and} \ \kappa = 1$. We thus deduce, $\mu_0 = \langle \chi \rangle \simeq 0.43 \,\text{eV}$, and

$$\langle \xi \rangle \simeq 0.05 \,\mathrm{eV} \,. \tag{9}$$

With the Yukawa coupling $f = \mathcal{O}(1)$, we see from (7) that this naturally generates the small neutrino masses, $m_{\nu} = f\langle \xi \rangle = \mathcal{O}(0.1) \,\mathrm{eV}$, consistent with the oscillation data [1].

In the present construction, the cubic coupling (2) between the triplet and doublet Higgs scalars is highly suppressed so that the Higgs triplet ξ can lie at the weak

$$V \supset -m_1^2(\sigma^{\dagger}\sigma) + \lambda_1(\sigma^{\dagger}\sigma)^2 - m_2^2(\phi^{\dagger}\phi) + \lambda_2(\phi^{\dagger}\phi)^2 + M_{\chi}^2 \text{Tr}(\chi^{\dagger}\chi) - \left[\mu\sigma\text{Tr}(\xi^{\dagger}\chi) + \text{h.c.}\right] + m_{\xi}^2 \text{Tr}(\xi^{\dagger}\xi) - (\kappa\sigma^{\dagger}\phi^T i\tau_2\chi\phi + \text{h.c.}).$$
(8)

scale, with a mass around $\mathcal{O}(0.5-1)$ TeV, although its Yukawa couplings to the leptons are of $\mathcal{O}(1)$. Such a weak scale Higgs triplet will thus lead to interesting phenomenology at the LHC and future lepton colliders (such as the ILC [26] or CLIC [27]). This Higgs triplet contains neutral and charged Higgs bosons $(\xi^0, \xi^{\pm}, \xi^{\pm\pm})$. At the LHC, they can be produced via their gauge couplings with photon or W^{\pm}/Z^0 bosons and then decay *dominantly into di-leptons*. For instance, these triplet Higgs bosons (with masses around 0.5-1 TeV) can be detected at the LHC via *s*-channel pair productions,

$$pp \to (\gamma^*, Z^*) \to \xi^{++} \xi^{--} \to \ell^+ \ell^+ \ell'^- \ell'^-$$
, (10a)

$$pp \to (Z^*) \to \xi^0 \xi^0 \to \ell^- \ell^+ \ell'^- \ell'^+ ,$$
 (10b)

$$pp \to (W^{+*}) \to \xi^{++}\xi^- \to \ell^+ \ell^+ \ell'^- \bar{\nu}',$$
 (10c)

$$pp \to (W^{+*}) \to \xi^+ \xi^0 \to \nu \ell^+ \ell'^- \ell'^+$$
, (10d)

giving rise to distinct 4-lepton signatures or 3-lepton plus missing energy signals. In particular, from (10a) we see that the like-sign 4-leptons of the type $e^{\pm}e^{\pm}\mu^{\mp}\mu^{\mp}$ with large invariant masses (M_{ee} and $M_{\mu\mu}$) and large transverse momenta are striking signals for the LHC discovery. Another interesting process is the pair production,

$$pp \to jj\xi^{++}\xi^0 \to jj\ell^+\ell^+\ell'^+\ell'^-, \qquad (11)$$

via the quartic vertex $W^{\pm}W^{\pm}\xi^{\mp\mp}\xi^{0}$ or the *t*-channel ξ exchange, with two energetic forward-jets and 4-leptons (such as $e^{+}e^{+}\mu^{+}\mu^{-}$ and $\mu^{+}\mu^{+}e^{+}e^{-}$) in the final state. Furthermore, at a future TeV lepton collider (such as ILC or CLIC) operating in the $e^{-}e^{-}$ mode, the $\xi^{\pm\pm}$ can be directly produced via $e^{-}e^{-} \rightarrow \xi^{--} \rightarrow \ell^{-}\ell^{-}$, with $\ell^{-}\ell^{-} = \mu^{-}\mu^{-}$ for instance, which has very clean background.³

We also note that since lepton number is a global symmetry, its spontaneous breaking induced by the nonzero $\langle \sigma \rangle$ will lead to a massless Nambu-Goldstone boson. Although this Nambu-Goldstone boson has a component from the Higgs triplet, there is no problem with the low-energy phenomenology including the LEP constraints because this triplet fraction is highly suppressed by $\langle \xi \rangle / \langle \sigma \rangle \sim \mathcal{O}(10^{-13})$. Also, the imaginary parts of the Higgs triplet ξ and the doublet ϕ have mixing, and one of their combinations results in a physical pseudoscalar which has a mass controlled by the heavier mass m_{ξ} and thus escapes from the low-energy constraints.

 $^{^2}$ We can replace the heavy singlets by heavy triplets to give another double type-II seesaw model,

Here χ denotes the heavy triplets with the lepton number L = -1, other notations coincide with the present model. This model also has some similarity with the lepton number soft-breaking models [25].

³ If such double-charged Higgs bosons $\xi^{\pm\pm}$ fall in the mass-range of 300 - 500 GeV as could be seen at the LHC [28], a 500 GeV ILC operated in the e^-e^- mode is expected to confirm the ξ^{--} discovery via the *s*-channel production, and further measure its properties precisely.

III. BARYOGENESIS VIA LEPTOGENESIS

In the present model, the heavy singlets χ have two decay channels as shown in Fig. ,

$$\chi \to \xi^* \phi^* \phi^*$$
 and $\chi \to \sigma^* \sigma^*$. (12)

If CP is not conserved, the $\chi \to \xi^* \phi^* \phi^*$ process and its CP-conjugate can generate a lepton asymmetry stored in the Higgs triplet after they go out of equilibrium. At the same time, there will emerge an equal but opposite lepton asymmetry stored in the light singlet since the lepton number is conserved in the χ decays and the sum of lepton asymmetries from the two decay channels in (12)vanishes. Note that in the leptogenesis, the sphaleron action has no effect on the light singlet. So we can focus on the lepton asymmetry stored in the Higgs triplet, which has been decoupled from the lepton asymmetry stored in the light singlet and will be rapidly transferred to the lepton doublets. After the lepton number is spontaneously broken at the TeV scale by the VEV of the singlet scalar σ , the Higgs triplet ξ will develop a trilinear coupling μ_0 to the Higgs doublet ϕ , as shown in Eqs. (2) and (5). Due to the smallness of this trilinear coupling μ_0 , the induced lepton number violating processes take place so slowly that they will not reach equilibrium until the temperature falls well below the electroweak scale, where the sphaleron process has become very weak. In consequence, the lepton asymmetry stored in the Higgs triplet can be partially converted to a baryon asymmetry.

It is clear that at least two such heavy singlets are needed to have CP violation and realize an interference between the tree-level diagram and the loop-order selfenergy. Here we minimally introduce two heavy singlets $\chi_{1,2}$. For convenience, we choose a basis of the heavy singlets by a proper rotation, which gives real, diagonal scalar mass-matrix $M_{\chi}^2 = \text{diag}(M_{\chi_1}^2, M_{\chi_2}^2)$ and two real cubic scalar-couplings $\mu = (\mu_1, \mu_2)$. Consequently, we only need to keep $\kappa = (\kappa_1, \kappa_2)$ complex in the Higgs potential. For illustration, consider the case where the two heavy singlets $\chi_{1,2}$ have hierarchical mass-spectrum. In this case, the final lepton asymmetry stored in the Higgs triplet will mainly come from the decay of the lighter one. Without the loss of generality, we choose χ_1 to be the lighter singlet and χ_2 the heavier one, i.e., $M_{\chi_1} \ll M_{\chi_2}$. We then compute the CP-asymmetry for the Higgs triplet,

$$\varepsilon_{1} = 2 \frac{\Gamma(\chi_{1} \to \xi^{*} \phi^{*} \phi^{*}) - \Gamma(\chi_{1}^{*} \to \xi \phi \phi)}{\Gamma_{1}}$$
$$\simeq \frac{1}{2\pi} \frac{\mathrm{Im}(\kappa_{1}^{*} \kappa_{2})}{|\kappa_{1}|^{2}} \frac{\mu_{1} \mu_{2}}{M_{\chi_{2}}^{2} - M_{\chi_{1}}^{2}} \frac{\frac{3}{32\pi^{2}} |\kappa_{1}|^{2}}{\frac{\mu_{1}^{2}}{M_{\chi_{1}}^{2}} + \frac{3}{32\pi^{2}} |\kappa_{1}|^{2}}$$

$$= \frac{\sin \delta}{2\pi} \left| \frac{\kappa_2}{\kappa_1} \right| \frac{\mu_1 \mu_2}{M_{\chi_2}^2 - M_{\chi_1}^2} \frac{\frac{3}{32\pi^2} |\kappa_1|^2}{\frac{\mu_1^2}{M_{\chi_1}^2} + \frac{3}{32\pi^2} |\kappa_1|^2}, \quad (13)$$

where we have defined $\kappa_i = |\kappa_i|e^{i\delta_i}$ and $\delta \equiv \delta_2 - \delta_1$ is the difference of the two phase angles which controls the above CP-asymmetry parameter ε_1 . In (13), Γ_i denotes the total decay width of χ_i or χ_i^* ,

$$\Gamma_{i} = \Gamma(\chi_{i} \to \xi^{*} \phi^{*} \phi^{*}) + \Gamma(\chi_{i} \to \sigma^{*} \sigma^{*})$$

$$= \Gamma(\chi_{i}^{*} \to \xi \phi \phi) + \Gamma(\chi_{i}^{*} \to \sigma \sigma)$$

$$= \frac{1}{8\pi} \left(\frac{\mu_{i}^{2}}{M_{\chi_{i}}^{2}} + \frac{3}{32\pi^{2}} |\kappa_{i}|^{2} \right) M_{\chi_{i}}. \quad (14)$$

Here the second equality is guaranteed by the unitarity and the CPT conservation.

When the lepton number is spontaneously broken by the VEV $\langle \sigma \rangle$, there will emerge a trilinear coupling of the Higgs triplet to the Higgs doublet as shown in Eqs. (2) and (5). For $\langle \sigma \rangle = 1.2 \text{ TeV}$, the phase transition may occur at the temperature $T_c \lesssim m_{\xi} = 0.5 \text{ TeV}$, where the lepton number violating processes have been already decoupled because

$$\Gamma(\xi \to \phi^* \phi^*) = \frac{1}{16\pi} \frac{|\mu_0|^2}{m_{\xi}} \ll H(T) \Big|_{T \simeq m_{\xi}}.$$
 (15)

Here,

I

$$H(T) = \left(\frac{8\pi^3 g_*}{90}\right)^{\frac{1}{2}} \frac{T^2}{M_{\rm Pl}}$$
(16)

is the Hubble constant, with the relativistic degrees of freedom $g_*\simeq 100$ and the Planck mass $M_{\rm Pl}\simeq 10^{19}\,{\rm GeV}.$ We then derive the final baryon asymmetry by computing the ratio of the baryon number density n_B over the entropy density s,

$$\frac{n_B}{s} = \frac{28}{79} \frac{n_B - n_{L_{SM}}}{s} = -\frac{28}{79} \frac{n_{L_{SM}}}{s}$$
$$\simeq -\frac{28}{79} \varepsilon_1 \frac{n_{\chi_1}^{eq}}{s} \Big|_{T \simeq M_{\chi_1}} \simeq -\frac{1}{15} \frac{\varepsilon_1}{g_*}, \quad (17)$$

where $n_{\chi_1}^{eq}$ is the thermal equilibrium density of χ_1 . Note the above solution is only valid for the weak washout regime with

$$\Gamma_{\chi_1} \lesssim H(T) \Big|_{T \simeq M_{\chi_1}}$$
 (18)

Let us consider the sample inputs,

$$\begin{split} M_{\chi_1} &= 0.1 \, M_{\chi_2} = 10^{14} \, \text{GeV} \,, \\ \mu_1 &= \mu_2 = 3 \times 10^{12} \, \text{GeV} \,, \\ |\kappa_1| &= |\kappa_2| = 1 \,, \quad \sin \delta = -0.11 \,, \end{split} \tag{19}$$



FIG. 2: The lepton number conserving decays of the heavy singlets χ at the tree level and the loop orders for generating a lepton asymmetry stored in the Higgs triplet ξ , and an equal but opposite amount stored in the light singlet σ . Here the CP-conjugate diagrams are not shown for simplicity.

with which we can estimate the CP-asymmetry,

$$\varepsilon_1 \simeq -1.3 \times 10^{-7} \,. \tag{20}$$

For convenience, we express the baryon asymmetry in terms of the ratio of n_B over the photon density n_{γ} ,

$$\frac{n_B}{n_{\gamma}} = 7.04 \frac{n_B}{s} \simeq 6.3 \times 10^{-10} \,, \tag{21}$$

which is consistent with the five-year observations of the WMAP collaboration [29], $n_B/n_\gamma = (6.225\pm0.170)\times10^{-10}$.

In the present model, the leptogenesis mechanism is different from that in the conventional seesaw models. Now the amount of lepton asymmetry does not depend on the parameters in the neutrino mass matrix. As a result, there is no DI bound [30] on the decaying particles. So we are flexible to lower the leptogensis scale and avoid the gravitino problem in a supersymmetric extension of this model, although we do not elaborate on this point in the present paper.

IV. LEPTONIC DARK MATTER

The recent PAMELA experiment [12] found an anomalous rise of the $e^+/(e^+ + e^-)$ fraction in cosmic rays, while the HESS [13] and Fermi/LAT [14] observations further exhibit an excess over the conventional backgroundpredictions for the cosmic ray fluxes. Systematical data fitting shows [31] that for the dark matter particles of mass around 3 TeV, their annihilations into the leptonic final states $(\mu^+\mu^-, \tau^+\tau^-, 4\mu, 4\tau)$, depending on the dark matter profile) can give a good fit to the excess; in general, a dark matter having mass lighter than about a TeV is excluded. In this section, we will extend our model to include a dark matter sector, in which the vector dark-matter particles annihilate dominantly into the μ 's and τ 's. With the aid of Breit-Wigner enhancement and for a hidden vector boson mass around 3 TeV, we find that our model can naturally explain the cosmic e^{\pm} excesses based on the general analysis in [31].

To construct the dark matter sector, we consider a hidden $SU(2)_h$ gauge theory with a complex Higgs doublet field η , but with no extra fermions (cf. Table 1). So, the hidden sector Lagrangian for η doublet can be written as,

$$\mathcal{L}_{\eta} = \left(D_{\mu} \eta \right)^{\dagger} \left(D^{\mu} \eta \right) - V(\eta) , \qquad (22)$$

where

$$D_{\mu}\eta = \left(\partial_{\mu} - ig_X \frac{\vec{\tau}}{2} \cdot \vec{X}_{\mu}\right)\eta, \qquad (23)$$

$$V(\eta) = -m_3^2 \left(\eta^{\dagger} \eta\right) + \lambda_3 \left(\eta^{\dagger} \eta\right)^2 .$$
 (24)

After the hidden $SU(2)_h$ is spontaneously broken by the vacuum expectation value $\langle \eta \rangle$, we are left with three degenerate massive vector bosons \vec{X}_{μ} as well as one neutral physical Higgs boson ζ ,

$$\mathcal{L}_{\eta} \supset \frac{1}{2} m_X^2 |\vec{X}_{\mu}|^2 + \frac{1}{2} m_{\zeta}^2 \zeta^2 + \frac{1}{8} g_X^2 \zeta^2 |\vec{X}_{\mu}|^2 + \frac{1}{2\sqrt{2}} g_X^2 \langle \eta \rangle \zeta |\vec{X}_{\mu}|^2 , \quad (25)$$

with

$$m_X = \frac{1}{\sqrt{2}} g_X \langle \eta \rangle, \quad m_\zeta = \sqrt{2} m_3.$$
 (26)



FIG. 3: The hidden vector bosons annihilate into Higgs triplets through the *s*-channel exchange of the hidden physical Higgs boson.

The hidden $SU(2)_h$ forbids its associated vector bosons to mix with the SM gauge bosons because of its non-Abelian character. Consequently the hidden vector bosons are stable, with no decay channel.

The hidden sector can communicate with the visible sector through the following quartic scalar interactions,

$$\mathcal{L} \supset -\alpha \left(\eta^{\dagger} \eta\right) \operatorname{Tr}\left(\xi^{\dagger} \xi\right) - \beta \left(\eta^{\dagger} \eta\right) \left(\phi^{\dagger} \phi\right), \qquad (27)$$

where the α -term involves the interaction between the hidden Higgs doublet η and the Higgs triplet ξ , while the β -term links η to the SM Higgs doublet ϕ . Once η develops a vacuum expectation value $\langle \eta \rangle$, the α -term and β -term will induce the cubic scalar couplings for $\eta - \xi - \xi$ and $\eta - \phi - \phi$ vertices, respectively. In the unitary gauge, the Higgs doublet η (or ϕ) reduces to the neutral physical Higgs boson ζ^0 (or h^0). Because of the $\zeta - X^a_\mu - X^{a\mu}$ cubic interaction in (25), we have the dark matter an-nihilation processes⁴, $X^a X^a \to \xi \xi^*$ and $X^a X^a \to hh$, with the s-channel ζ -exchange, where $\xi \in (\xi^0, \xi^{\pm}, \xi^{\pm\pm})$ contains 6 real degrees of freedom. For the scattering energy much larger than the masses of ξ and h, we find that the ratio of the inclusive cross sections for the final states $\xi\xi^*$ and hh is about $6(\alpha/\beta)^2$. Let us consider the natural parameter space of $\beta \leq (0.3 - 0.5)\alpha$ without any fine-tuning. We can thus estimate the ratio of the two inclusive cross sections,

$$\frac{\sigma[X^a X^a \to hh]}{\sigma[X^a X^a \to \xi\xi^*]} \simeq \frac{1}{6} \left(\frac{\beta}{\alpha}\right)^2 \leqslant 0.5\% - 4\%, \quad (28)$$

indicating that the process $X^a X^a \to \xi \xi^*$ fully dominates the dark matter annihilation cross section. Due to the large triplet Yukawa couplings $f = \mathcal{O}(1)$ in (1) and tiny trilinear scalar coupling $\mu_0 = \mathcal{O}(0.1 - 1) \text{eV}$ in (2), our dark matter annihilation $X^a X^a \to \xi \xi^*$ will naturally lead to leptonic decay products⁵ and thus nicely agree with the exciting signals from the recent cosmic ray experiments [10, 11, 12, 13, 14], as mentioned at the beginning of this section. On the other hand, for the process $X^a X^a \to hh$, the SM Higgs boson h will decay preferably into $b\bar{b}$ or WW/ZZ final states, but will not cause visible signals in the present cosmic ray data due to its too small cross section in (28).

So, we consider the hidden vector boson annihilation into the Higgs triplets,⁶ $X^a X^a \rightarrow \xi \xi^*$ (a = 1, 2, 3), through the s-channel exchange of the hidden Higgs boson ζ , as shown in Fig.3. The form of the cubic scalar vertex $\zeta - \xi - \xi$ can be derived from the α -term in (27),

$$\mathcal{L} \supset -\alpha \left(\eta^{\dagger} \eta \right) \operatorname{Tr} \left(\xi^{\dagger} \xi \right)$$
$$\supset -\sqrt{2} \alpha \langle \eta \rangle \zeta \operatorname{Tr} \left(\xi^{\dagger} \xi \right) , \qquad (29)$$

To compute the cross section of this process, we will average over the initial state polarizations and gauge indices. Thus, we can derive the unpolarized cross section for $X^a X^a \to \xi \xi^*$,

$$\sigma v = \frac{g_X^4 \alpha^2}{72\pi s} \left[2 + \frac{\left(s - 2m_X^2\right)^2}{4m_X^4} \right] \frac{\langle \eta \rangle^4}{(s - m_\zeta^2)^2 + m_\zeta^2 \Gamma_\zeta^2},$$
(30)

where v is the relative velocity of the initial state vector bosons,

$$v = 2\sqrt{1 - \frac{4m_X^2}{s}},$$
 (31)

while the decay width of the Higgs boson ζ is given by the channel $\zeta \to \xi \xi^*$ (for $m_{\zeta} < 2m_X$),

$$\Gamma_{\zeta} = \frac{3\alpha^2}{8\pi} \frac{\langle \eta \rangle^2}{m_{\zeta}}.$$
(32)

We now verify that the hidden vector bosons can indeed serve as the dark matter. For this purpose, we need to thermally average the cross section (30) and determine their relic density. Following Refs. [18, 19, 20, 21], we consider the Breit-Wigner resonant case,

$$\delta = 1 - \frac{m_{\zeta}^2}{4m_X^2} \ll 1.$$
 (33)

⁴ We may also have decay processes for non-thermal production of dark matter [32]. For instance, the heavy singlets χ may decay into the hidden vector bosons X due to the quartic interaction between χ and the hidden Higgs doublet η , $(\chi^{\dagger}\chi) (\eta^{\dagger}\eta)$. However, this decay is highly suppressed because of the tiny VEV $\langle \chi \rangle$ in (4) and the heavy mass M_{χ} in (3). So, the induced relic density of dark matter is too small to be relevant.

⁵ As variations of our current construction, we may also consider the dark matter annihilation into other leptonic scalars like those in Zee model [33] or Zee-Babu model [34].

⁶ As the present work is being prepared, a new preprint [23] considered scalar dark matter with their annihilations into Higgs triplets.



FIG. 4: Predictions of the dark matter density $\Omega_{\rm DM}h^2$ as a function of hidden vector boson mass m_X in our model. The curves from right to left correspond to $g_X/\alpha = 10, 20, 30, 35, 40, 50, 60, 70$, respectively, where we have set $\delta = 0.4 \times 10^{-4}$ and $g_X = 0.34$. The shaded region (yellow) is allowed by the cosmological bound, $0.1029 < \Omega_{\rm DM}h^2 < 0.1233$ (3 σ level), as extracted from the WMAP data [29].

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The relic density of the dark matter in our model can be solved from the Boltzmann equation [35] with the resonant effect [20],

$$\Omega_{\rm DM} h^2 \simeq \frac{1.07 \times 10^9 x_f \,{\rm GeV}^{-1}}{\sqrt{g_*} \, M_{\rm Pl} \langle \sigma v \rangle_0} \times {\rm BF}$$
$$\simeq \frac{0.1 \times 10^{-9} {\rm GeV}^{-2}}{\langle \sigma v \rangle_0} \frac{x_f}{\sqrt{g_*}} \times {\rm BF}, \quad (34)$$

where BF is the effective boost factor due to Breit-Wigner resonant effect [20], as will be discussed shortly. The parameter $x_f = m_{\rm DM}/T$ is the freeze-out temperature [35],

$$x_f = \ln X - \frac{1}{2} \ln \ln X$$
, (35a)

$$X = 0.038 (g/\sqrt{g_*}) M_{\rm Pl} \langle \sigma v \rangle_0 , \qquad (35b)$$

with $g_*(g)$ the degrees of freedom for massless particles (dark matter). Here we have $g_* \simeq 100$ and $g = 3 \times 3 = 9$ in the present model. The x_f is typically of $\mathcal{O}(10-30)$. In the above we have defined $\langle \sigma v \rangle_0 = \langle \sigma v \rangle |_{T=0}$. For the

s-wave dark matter annihilation in the non-relativistic limit, the thermally averaged cross section equals the non-averaged one, $\langle \sigma v \rangle \simeq \sigma v$, so we can deduce from (30),

$$\langle \sigma v \rangle_0 \simeq \frac{g_X^2}{72m_X^2} \frac{\gamma}{\delta^2 + \gamma^2},$$
 (36a)

$$\gamma = \frac{\Gamma_{\zeta}}{m_{\zeta}} = \frac{3\alpha^2}{16\pi g_X^2} \frac{1}{1-\delta} \,. \tag{36b}$$

where, besides $\delta \ll 1$, we also note $\gamma \ll 1$ for $\alpha/g_X < 0.1$. Finally, the effective boost factor BF due to Breit-Wigner resonant effect is given by [20],

$$BF \simeq \frac{\max[\delta, \gamma]^{-1}}{x_f}, \qquad (37)$$

which is expected to be around $\mathcal{O}(10^2 - 10^3)$. The effective boost factor for the annihilation cross section may alternatively arise from Sommerfeld enhancement [36], which will not be explored in the present study.

For illustration, let us consider a sample input,

 $\delta = 0.4 \times 10^{-4}, \quad g_X = 40 \alpha = 0.34 \,, \quad m_X = 3 \, {\rm TeV}. \ (38)$

With this we can thus derive,

$$x_f \simeq 27, \quad \gamma \simeq 0.37 \times 10^{-4}, \quad (39a)$$

$$BF \simeq 0.93 \times 10^3 \,, \tag{39b}$$

$$\langle \sigma v \rangle_0 \simeq 2.2 \times 10^{-6} \,\mathrm{GeV}^{-2} \,, \tag{39c}$$

and the relic density for dark matter,

$$\Omega_{\rm DM} h^2 \simeq 0.11 \,, \tag{40}$$

which is consistent with the measured value given by the WMAP observations (combined with the distance measurements from Type-Ia Supernovae and the Baryon Acoustic Oscillations in the distributions of galaxies)[29], $\Omega_{\rm DM}h^2 = 0.1131 \pm 0.0034$.

To further explore the parameter space, we have plotted Fig. 4, where the relic dark matter density $\Omega_{\rm DM}h^2$ is shown as a function of the hidden vector boson mass m_X for possible values of the coupling-ratio g_X/α . The present WAMP constraint, $0.1029 < \Omega_{\rm DM}h^2 < 0.1233$ (3σ level), is imposed as the shaded region in the parameter space.

In the galactic halo, the annihilating dark matter has a relative velocity about $v \sim 10^{-3}$, and its thermally averaged cross section is well described by that at the zero temperature for $v^2 \ll \delta, \gamma$. From the parameter choice (38), we see that $v^2 \ll \gamma < \delta$ holds. It is known that a very small δ does not appear technically so natural [18]-[21], while the Breit-Wigner enhancement can give the desired large annihilation cross section of $\mathcal{O}(10^{-6})\,\mathrm{GeV}^{-2}$ as in (39c), and is consistent with the recent cosmic-ray signals [12, 13, 14]. In our model, the anomalies from the PAMELA[12], HESS[13] and Fermi/LAT [14] observations can be understood in two steps: (i) first, the three degenerate vector bosons X^a with mass around 3 TeV, annihilate into the Higgs triplet ξ with the enhanced cross section (39c); (ii) subsequently, the Higgs triplet ξ rapidly and mostly decays into the leptons (the μ 's and τ 's).

V. SUMMARY

We have constructed a new double type-II seesaw scenario where the coupling of the Higgs triplet to the Higgs doublet is highly suppressed after spontaneous breaking of the global lepton number. In our model, the small neutrino masses can be naturally realized, while at the same time the Yukawa couplings of the TeV-scale Higgs triplet to the lepton doublets are large enough, leading exciting phenomenology at the LHC and future TeV lepton colliders (such as the ILC or CLIC). Furthermore, a lepton asymmetry stored in the Higgs triplet can be generated in the lepton number conserving decays and then be rapidly transferred to the lepton doublets. Hence, the matterantimatter asymmetry in the Universe can be naturally explained via leptogenesis in our model. We have further extended the model to realize the leptonic dark matter by including stable vector bosons associated with a hidden non-Abelian gauge symmetry. The stable vector bosons can dominantly annihilate into the Higgs triplets, which decay mostly into leptons in our double type-II seesaw

scenario. In this annihilation, there is an s-channel process mediated by a hidden Higgs boson, which is responsible for the mass generation of dark matter. Furthermore, the Breit-Wigner resonant enhancement makes it possible to have a large annihilation rate of the dark matter particles into the μ 's and τ 's and thus explain the cosmic e^{\pm} excesses from our model.

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