Inverse seesaw model with global $U(1)_H$ symmetryUjjal Kumar Dey,^{1,*} Takaaki Nomura,^{2,†} and Hiroshi Okada^{3,4,‡}¹*Asia Pacific Center for Theoretical Physics (APCTP), Headquarters San 31, Hyoja-dong, Nam-gu, Pohang 790-784, Republic of Korea*²*School of Physics, KIAS, Seoul 02455, Republic of Korea*³*Asia Pacific Center for Theoretical Physics, Pohang 37673, Republic of Korea*⁴*Department of Physics, Pohang University of Science and Technology, Pohang 37673, Republic of Korea*

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We propose an inverse seesaw model based on hidden global symmetry $U(1)_H$ in which we realize tiny neutrino masses with rather natural manner taking into account relevant experimental bounds. The small Majorana mass for inverse seesaw mechanism is induced via the small vacuum expectation value of a triplet scalar field whose Yukawa interactions with standard model fermions are controlled by $U(1)_H$. We discuss the phenomenology of the exotic particles present in the model, including the Goldstone boson coming from breaking of the global symmetry, and explore testability at the Large Hadron Collider experiments.

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I. INTRODUCTION

Inverse seesaw mechanism [1,2] is one of the promising candidates to induce neutrino masses and their mixing, which are typically understood by an additional symmetry beyond the standard model (SM). The mechanism requires both chiralities of neutral fermions ($N_{L/R}$) that couple to the SM fermions in different manners, and their mass structures are also different. That is why an additional symmetry is needed [3–6]. Furthermore, it demands hierarchies among mass parameters associated with neutral fermions, including active neutrinos where, especially, the Majorana mass of N_L should be tiny. Indeed, realizing such hierarchies in a natural way is rather challenging from a model building perspective [4,5].

In this paper, we propose an inverse seesaw model under a hidden *global* $U(1)_H$ symmetry [7]¹ in which we try to realize natural hierarchies in a neutral fermion mass matrix, taking advantage of experimental constraints of electroweak precision to test our model. In our model, we introduce $SU(2)_L$ triplet scalar and exotic lepton doublets with nonzero $U(1)_H$ charges and scalar singlets.

The Yukawa interaction between the triplet and the exotic lepton doublet induces a small Majorana mass term required for inverse seesaw mechanism due to small vacuum expectation value (VEV) of the triplet, while the Yukawa interaction between the SM lepton doublet and the scalar triplet is forbidden by the $U(1)_H$ symmetry. In addition, as a result of the global symmetry that accommodates a physical Goldstone boson (GB) [13,14], our model can be well tested at future colliders since the GB couples to the SM fermions through their kinetic terms.²

This letter is organized as follows. In Sec. II, we introduce our model and formulate the scalar sector, charged-lepton sector, and neutral fermion sector, and we briefly discuss an electroweak precision test. In the scalar sector, we show hierarchies among VEVs that connect to our desired hierarchies among neutral fermions. Using electroweak precision data, we show the constraint on the VEV of an isospin triplet scalar at tree level, and this bound is directly related to the hierarchies of neutral fermion mass matrices. Then we formulate the charged-lepton mass matrix that encapsulates the mixing between SM charged leptons and newly introduced heavier ones. From this, we show that their mass hierarchies could be naturally realized while maintaining the experimental bound on the new, heavier leptons. Note that these hierarchies are also related to our neutral fermion mass matrix. In the last part of this section, we

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¹There are several studies utilizing hidden *gauged* symmetries to explain neutrino masses [8–12].

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²Even in the case of a gauged symmetry, a GB is induced when one introduces two or more two bosons that break the additional $U(1)$ symmetry; see, e.g., Ref. [9].

TABLE I. Charge assignments of the our lepton and scalar fields under local $SU(2)_L \times U(1)_Y$ and global $U(1)_H$ symmetries, where the upper index a is the number of family that runs over 1–3 and all of them are singlet under $SU(3)_C$.

	L_L^a	e_R^a	$L_L'^a$	$L_R'^a$	H	Δ	φ_1	φ_2
$SU(2)_L$	2	1	2	2	2	3	1	1
$U(1)_Y$	$-\frac{1}{2}$	-1	$-\frac{1}{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	1	0	0
$U(1)_H$	0	0	ℓ	2ℓ	0	2ℓ	ℓ	2ℓ

discuss the neutral fermion mass matrix and show how to realize these hierarchies. Also we discuss the bounds on the nonunitarity effect originating from the inverse seesaw model. In Sec. III, we discuss the phenomenology of our model. We mainly consider the decays of the exotic charged leptons, and they can leave their footprints in the multilepton signatures of the collider experiments, like LHC. Finally we summarize our results and conclude.

II. MODEL

In this section, we formulate our model, in which we introduce a global $U(1)_H$ symmetry. The fermionic sector is augmented by three families of vectorlike fermions $L'_{L/R} \equiv [N_{L/R}, E_{L/R}]^T$ with charge $(2, -1/2)$ under the $SU(2)_L \times U(1)_Y$ gauge symmetry, while under the global $U(1)_H$ right- and left-handed ones are assigned charges 2ℓ and ℓ , respectively. As for the scalar sector, we add an isospin triplet scalar Δ with hypercharge 1, and two singlet scalars $\varphi_{1,2}$ with zero hypercharge, while $(2\ell, \ell, 2\ell)$ are, respectively, assigned to $(\Delta, \varphi_1, \varphi_2)$ under the global $U(1)_H$ symmetry. The SM-like Higgs field is denoted by H . The VEVs of $(H, \varphi_1, \varphi_2, \Delta)$ are $(v/\sqrt{2}, v'_1/\sqrt{2}, v'_2/\sqrt{2}, v_\Delta/\sqrt{2})$, respectively. The field content and the respective quantum number assignments are shown in Table I. Note that the quark sector remains the same as that of SM. The renormalizable Yukawa Lagrangian under these symmetries is then given by

$$-\mathcal{L}_\ell = y_{\ell aa} \bar{L}_L^a H e_R^a + f_{ab} \bar{L}_L'^a L_R'^b \varphi_1^* + g_{Lab} (\bar{L}_L'^c)_L \tilde{\Delta} L_L'^b + y_{Dab} \bar{L}_L^a L_R'^b \varphi_2^* + \text{H.c.}, \quad (1)$$

where the indices $a, b (= 1-3)$ represent the number of families, $\tilde{\Delta} \equiv i\sigma_2 \Delta^\dagger$, and y_ℓ is assumed to be a diagonal matrix without loss of generality. After spontaneous symmetry breaking, one finds that the charged-lepton mass matrix $m_\ell = y_\ell v/\sqrt{2}$.

A. Scalar potential and VEVs

First, we define each scalar field as follows:

$$H \equiv \begin{bmatrix} h^+ \\ \frac{v+h+ia}{\sqrt{2}} \end{bmatrix}, \quad \Delta \equiv \begin{bmatrix} \frac{\delta^+}{\sqrt{2}} & \delta^{++} \\ \frac{v_\Delta + \delta_R + i\delta_L}{\sqrt{2}} & -\frac{\delta^+}{\sqrt{2}} \end{bmatrix},$$

$$\varphi_{1/2} \equiv \frac{v'_{1/2} + \varphi_{R1/2} + i\varphi_{I1/2}}{\sqrt{2}}. \quad (2)$$

The scalar potential in our model is given by

$$\begin{aligned} \mathcal{V} = & -\mu_h^2 H^\dagger H + M_\Delta^2 \text{Tr}[\Delta^\dagger \Delta] - \mu_1^2 |\varphi_1|^2 + M_2^2 |\varphi_2|^2 \\ & - (\mu_\varphi \varphi_2^* \varphi_1 \varphi_1 + \text{H.c.}) - [\lambda_0 H^T \tilde{\Delta} H \varphi_2 + \text{H.c.}] \\ & + \lambda_H (H^\dagger H)^2 + \lambda_{\varphi_1} |\varphi_1|^4 + \lambda_{\varphi_2} |\varphi_2|^4 + \lambda_{H\varphi_1} (H^\dagger H) |\varphi_1|^2 \\ & + \lambda_{H\varphi_2} (H^\dagger H) |\varphi_2|^2 + \lambda_{\varphi_1\varphi_2} |\varphi_1|^2 |\varphi_2|^2 + \mathcal{V}_{\text{trivial}}, \end{aligned} \quad (3)$$

where $\mathcal{V}_{\text{trivial}}$ indicates trivial quartic terms containing a scalar triplet. For simplicity, we assume all of the couplings to be real. The quartic coupling λ_0 plays a role in reducing the scale of VEV of Δ to $\mathcal{O}(1)$ GeV. Applying the minimization condition $\partial\mathcal{V}/\partial v = \partial\mathcal{V}/\partial v'_{1,2} = \partial\mathcal{V}/\partial v_\Delta = 0$, we obtain the VEVs approximately as

$$v \simeq \sqrt{\frac{\mu_h^2}{\lambda_H}}, \quad v'_1 \simeq \sqrt{\frac{\mu_1^2}{\lambda_{\varphi_1}}},$$

$$v'_2 \simeq \frac{\sqrt{2}\mu_\varphi v_1'^2}{2M_2^2 + \lambda_{\varphi_1\varphi_2} v_1'^2}, \quad v_\Delta \simeq \frac{\lambda_0 v^2 v'_2}{4M_\Delta^2}, \quad (4)$$

where we assume $v'_2 \sim v_\Delta \ll \{v'_1, v\}$ and $\lambda_{H\varphi_1}$ to be small. The above hierarchy of VEVs is motivated from the mass hierarchy in neutral fermions, as we discuss below. To obtain the VEV hierarchy and electroweak vacuum consistently, we need to choose appropriate parameters. When we impose $v'_1 \gtrsim 5 \times 10^6$ GeV as required by constraints from massless Goldstone boson interactions as discussed below, and taking $v'_2 \sim v_\Delta = \mathcal{O}(1)$ GeV and $M_2^2 \sim \lambda_{\varphi_1\varphi_2} v_1'^2 \sim M_\Delta^2 \sim v^2$, we should choose $\mu_\varphi \lesssim 10^{-8}$ GeV and $\lambda_{\varphi_1\varphi_2} \lesssim 10^{-8}$. Furthermore, $\lambda_{H\varphi_1}$ should be very small as $\mathcal{O}(10^{-8})$ or less when $v'_1 \gtrsim 5 \times 10^6$ GeV since $\sqrt{\lambda_{H\varphi_1}} v'_1 \lesssim \mathcal{O}(100)$ GeV in order to realize electroweak vacuum. Also the one loop contribution to $\lambda_{H\varphi_1}$ is roughly obtained as $\sim \lambda_{H\varphi_2} \lambda_{\varphi_1\varphi_2} / (4\pi)^2$, and the condition $\lambda_{H\varphi_1} \lesssim 10^{-8}$ requires that $\lambda_{H\varphi_2} \lambda_{\varphi_1\varphi_2} \lesssim 10^{-6}$.

Notice that smallness of μ_φ would be natural in the 't Hooft sense since we can restore a “lepton number” symmetry for the $\mu_\varphi \rightarrow 0$ limit if the quantum number is assigned as 1 for fields $\{L_L, e_R, L'_R, \varphi_1\}$, and 0 for $\{L'_L, \varphi_2, \Delta\}$.

Since we assume small $\lambda_{H\varphi_1}$ coupling, the mass matrix for CP -even scalars from the SM singlets $\varphi_{1,2}$ is derived approximately as

$$\frac{1}{2} \begin{pmatrix} \varphi_{R_1} \\ \varphi_{R_2} \end{pmatrix}^T \begin{pmatrix} 2\lambda_{\varphi_1} v_1'^2 & \lambda_{\varphi_1\varphi_2} v_1' v_2' - \sqrt{2}\mu_\varphi v_1' \\ \lambda_{\varphi_1\varphi_2} v_1' v_2' - \sqrt{2}\mu_\varphi v_1' & M_2^2 + \frac{1}{2}\lambda_{\varphi_1\varphi_2} v_1'^2 \end{pmatrix} \times \begin{pmatrix} \varphi_{R_1} \\ \varphi_{R_2} \end{pmatrix}. \quad (5)$$

Note that mixing between φ_{R_1} and φ_{R_2} would be small since off-diagonal elements are suppressed by v_2' and μ_φ unless values of diagonal elements are not too close. Thus we approximate φ_{R_1} and φ_{R_2} as mass eigenstates, and the corresponding mass eigenvalues are

$$m_{\varphi_{R_1}} \simeq \sqrt{2\lambda_{\varphi_1} v_1'^2}, \quad m_{\varphi_{R_2}} \simeq \sqrt{M_2^2 + \frac{1}{2}\lambda_{\varphi_1\varphi_2} v_1'^2}. \quad (6)$$

In the CP -odd scalar sector of the SM singlet $\varphi_{1,2}$, we have one massless Goldstone boson associated with breaking of the global $U(1)_H$ symmetry. Applying our assumption of $v_2' \ll v_1'$, we can approximately identify φ_{I_1} as a Goldstone boson, while φ_{I_2} is a massive scalar boson with mass value $m_{\varphi_{I_2}} \simeq \sqrt{M_2^2 + \lambda_{\varphi_1\varphi_2} v_1'^2/2}$. Thus φ_{R_2} and φ_{I_2} have approximately degenerate masses. In the following analysis, we consider that massive scalars from φ_1 and φ_2 have electroweak scale masses expecting collider phenomenology of these scalar bosons. We then impose $\lambda_{\varphi_1} \lesssim 10^{-8}$ since v_1' should be larger than $\sim 5 \times 10^6$ GeV from constraints of a Goldstone boson having an axionlike interaction, as discussed below. In addition, $\lambda_{\varphi_1\varphi_2} \lesssim 10^{-3}$ is required since φ_2 loop correction to λ_{φ_1} is roughly estimated to be $\sim \lambda_{\varphi_1\varphi_2}^2/(4\pi)^2$. Clearly with these choices of quartic couplings, one can have the massive scalars in the electroweak scale relevant to the collider phenomenology.

The SM Higgs and scalar triplet sectors are mostly the same as the usual scalar triplet model since we assume the mixing in the SM singlet scalar to be sufficiently small. The mass scale of scalar bosons in the triplet is given by M_Δ .

B. ρ parameter

The VEV of Δ is restricted by the ρ parameter at tree level that is given by [15]

$$\rho \approx \frac{v^2 + 2v_\Delta^2}{v^2 + 4v_\Delta^2}, \quad (7)$$

where the experimental value is given by $\rho = 1.0004_{-0.0004}^{+0.0003}$ at a 2σ confidence level [16]. On the other hand, we have $v_{\text{SM}} = \sqrt{v^2 + 2v_\Delta^2} \approx 246$ GeV. Therefore the upper bound on v_Δ is of the order $\mathcal{O}(1)$ GeV.³

³The theoretical origins are studied in Refs. [17–19].

C. Charged-lepton sector

The charged-lepton mass matrix consists of the component of the SM mass matrix and a heavier one, after the spontaneous symmetry breaking. We define the mass matrix M to be $f v_1'/\sqrt{2}$. Furthermore, we assume the mass matrices m_D and M to be diagonal for simplicity. Then, for each generation indicated by “ a ,” we can write the charged-lepton fermion mass matrix as

$$\begin{pmatrix} \bar{e}_L^a \\ \bar{E}_L^a \end{pmatrix}^T M_{E_a} \begin{pmatrix} e_R^a \\ E_R^a \end{pmatrix} = \begin{pmatrix} \bar{e}_L^a \\ \bar{E}_L^a \end{pmatrix}^T \begin{bmatrix} m_{\ell_a} & m_{D_a} \\ 0 & M_a \end{bmatrix} \begin{pmatrix} e_R^a \\ E_R^a \end{pmatrix}, \quad (8)$$

$$M_{E_a} M_{E_a}^\dagger = \begin{bmatrix} m_{\ell_a}^2 + m_{D_a}^2 & m_{D_a} M_a \\ m_{D_a} M_a & M_a^2 \end{bmatrix},$$

$$M_{E_a}^\dagger M_{E_a} = \begin{bmatrix} m_{\ell_a}^2 & m_{D_a} m_{\ell_a} \\ m_{D_a} m_{\ell_a} & M_a^2 + m_{D_a}^2 \end{bmatrix}. \quad (9)$$

The mass matrix is diagonalized by the transformation $(e_{L(R)}^a, E_{L(R)}^a) \rightarrow V_{L_a}^\dagger (e_{L(R)}^a, E_{L(R)}^a)$. Thus we can obtain diagonalization matrices V_{L_a} and V_{R_a} , which, respectively, diagonalize $M_{E_a} M_{E_a}^\dagger$ and $M_{E_a}^\dagger M_{E_a}$ as $V_{L_a} M_{E_a} M_{E_a}^\dagger V_{L_a}^\dagger = V_{R_a} M_{E_a}^\dagger M_{E_a} V_{R_a}^\dagger = \text{diag}(m_{\ell_a}^2, M_a^2)$, such that

$$V_{L_a} = \begin{pmatrix} \cos\theta_{L_a} & -\sin\theta_{L_a} \\ \sin\theta_{L_a} & \cos\theta_{L_a} \end{pmatrix}, \quad V_{R_a} = \begin{pmatrix} \cos\theta_{R_a} & -\sin\theta_{R_a} \\ \sin\theta_{R_a} & \cos\theta_{R_a} \end{pmatrix}, \quad (10)$$

$$\tan 2\theta_{L_a} = \frac{2m_{D_a} M_a}{M_a^2 - m_{\ell_a}^2 - m_{D_a}^2} \simeq \frac{2m_{D_a}}{M_a},$$

$$\tan 2\theta_{R_a} = \frac{2m_{D_a} m_{\ell_a}}{M_a^2 + m_{D_a}^2 - m_{\ell_a}^2} \simeq \frac{2m_{D_a} m_{\ell_a}}{M_a^2}, \quad (11)$$

where we have assumed that $m_{\ell_a}, m_{D_a} \ll M_a$. Then the mass eigenvalues for e^a and E^a are given simply by m_{ℓ_a} and M_a . Also their mixing angles θ_{R_a} and θ_{L_a} are very small and satisfy $\theta_{R_a} \ll \theta_{L_a}$ since the lower bound on the mass of the heavier leptons is about 100 GeV [16], which is suggested by the current experimental data at the LHC and the LEP. This is because hierarchies in the mass parameter $m_{\ell_a}, m_{D_a} \ll M_a$ are comparatively natural. Note also that the constraints from electroweak precision measurements (S, T, U parameters) can be avoided when the components of the exotic lepton doublet have a degenerate mass. In fact we consider such a case, where the masses of E and N are dominantly given by Dirac mass M .

D. Neutrino sector

After the spontaneous symmetry breaking, the neutral fermion mass matrix in the basis of $(\nu_L^c, N_R, N_L^c)^T$ is given by

$$M_N = \begin{bmatrix} 0 & m_D & 0 \\ m_D & 0 & M \\ 0 & M & \mu^* \end{bmatrix}, \quad (12)$$

where m_D , M are diagonal, while $\mu \equiv g_L v_\Delta / \sqrt{2}$ is a symmetric 3×3 mass matrix. Then the active neutrino mass matrix can be given as

$$m_\nu \approx \mu^* \left(\frac{m_D}{M} \right)^2. \quad (13)$$

Once we fix $m_D/M \sim \mathcal{O}(0.01)$, then $\mu \sim \mathcal{O}(10^{-7})$ GeV in order to satisfy the observed neutrino mass squared differences. It suggests that $g_L \sim \mathcal{O}(10^{-7})$ when we fix $v_\Delta \sim \mathcal{O}(1)$ GeV.⁴ The neutrino mass matrix is diagonalized by unitary matrix U_{MNS} ; $D_\nu = U_{MNS}^T m_\nu U_{MNS}$, where $D_\nu \equiv \text{diag}(m_1, m_2, m_3)$. The constraint from nonunitarity can simply be obtained by considering the Hermitian matrix $F \equiv M^{-1} m_D$. Combining several experimental results [24], the upper bounds are given by [25]

$$|FF^\dagger| \leq \begin{bmatrix} 2.5 \times 10^{-3} & 2.4 \times 10^{-5} & 2.7 \times 10^{-3} \\ 2.4 \times 10^{-5} & 4.0 \times 10^{-4} & 1.2 \times 10^{-3} \\ 2.7 \times 10^{-3} & 1.2 \times 10^{-3} & 5.6 \times 10^{-3} \end{bmatrix}. \quad (14)$$

Since F is assumed to be diagonal, the stringent constraint originates from the 2-2 component of $|FF^\dagger|$. It suggests that $|F| \sim m_D/M \lesssim 10^{-2}$ is always safe in our model thanks to the small VEV v'_2 .

III. PHENOMENOLOGY

In this section, we discuss the phenomenology of the model focusing on the decay and production of exotic particles. One specific property of our model is the existence of the physical GB as a result of global $U(1)_H$ symmetry breaking.⁵ Here we first derive interactions associated with the Goldstone boson, which is denoted

by α_G and identified as $\varphi_1 = e^{\frac{i\alpha_G}{v'_1}}(v'_1 + \varphi_{R1})/\sqrt{2}$ in the limit of $v'_2 \ll v'_1$. Then fields with a $U(1)_H$ charge are rewritten as $\varphi_2 \rightarrow e^{\frac{2i\alpha_G}{v'_1}} \varphi_2$, $\Delta \rightarrow e^{\frac{2i\alpha_G}{v'_1}} \Delta$, $L_L^a \rightarrow e^{\frac{i\alpha_G}{v'_1}} L_L^a$, and $L_R^a \rightarrow e^{\frac{2i\alpha_G}{v'_1}} L_R^a$. Then φ_1 and φ_2 interact with α_G as follows:

⁴These hierarchies could be explained by several mechanisms such as radiative models [20–22] and effective models with higher order terms [23].

⁵The physical GB would obtain a tiny mass due to breaking of the global $U(1)$ symmetry by gravitational effects [26]. In our analysis, we simply consider it a massless particle.

$$\begin{aligned} \mathcal{L} \supset & \frac{1}{v'_1} \varphi_{R1} \partial_\mu \alpha_G \partial^\mu \alpha_G + \frac{1}{2v'_1{}^2} \varphi_{R1} \varphi_{R1} \partial_\mu \alpha_G \partial^\mu \alpha_G \\ & + i \frac{2}{v'_1} \partial_\mu \alpha_G (\partial^\mu \varphi_{R2} \varphi_{I2} - \varphi_{R2} \partial^\mu \varphi_{I2}) \\ & + \frac{2v'_2}{v'_1{}^2} \varphi_{R2} \partial_\mu \alpha_G \partial^\mu \alpha_G + \frac{1}{v'_1{}^2} (\varphi_{R2} \varphi_{R2} + \varphi_{I2} \varphi_{I2}) \partial_\mu \alpha_G \partial^\mu \alpha_G. \end{aligned} \quad (15)$$

The covariant derivative of Δ is rewritten including the Goldstone boson as

$$\begin{aligned} D_\mu \Delta = & \partial_\mu \Delta + i \frac{2}{v'_1} \partial_\mu \alpha_G \Delta - i \frac{g}{\sqrt{2}} (W_\mu^+ [T^+, \Delta] + W_\mu^- [T^-, \Delta]) \\ & - i \frac{g}{\cos \theta_W} Z_\mu \left(\left[\frac{\sigma_3}{2}, \Delta \right] - \sin^2 \theta_W \hat{Q} \Delta \right) - ie A_\mu \hat{Q} \Delta, \end{aligned} \quad (16)$$

where g is the $SU(2)_L$ gauge coupling, $T^\pm = (\sigma_1 \pm i\sigma_2)/2$, θ_W is the Weinberg angle, and $\hat{Q}$ is the electric charge operator acting on each component of the multiplet. Then we can obtain interactions of α_G and Δ from the kinetic term $\text{Tr}[(D_\mu \Delta)^\dagger (D^\mu \Delta)]$. The triplet Higgs decay modes with α_G are suppressed by a v_Δ/v'_1 factor so that components in Δ decay via gauge or scalar potential interactions. Also exotic lepton L'^a interaction with α_G is derived from the kinetic term as

$$\begin{aligned} \mathcal{L} \supset & -\bar{L}'^a \gamma^\mu \left(\frac{1}{v'_1} \partial_\mu \alpha_G P_L + \frac{2}{v'_1} \partial_\mu \alpha_G P_R \right) L'^a \\ & + \bar{L}'^a \gamma^\mu \left(\frac{g}{\sqrt{2}} (W_\mu^+ T^+ + W_\mu^- T^-) \right. \\ & \left. + \frac{g}{\cos \theta_W} \left(\frac{\sigma_3}{2} - \sin^2 \theta_W \hat{Q} \right) Z_\mu + e \hat{Q} A_\mu \right) L'^a. \end{aligned} \quad (17)$$

Here we briefly discuss constraints on our physical GB. The GB can have axionlike photon coupling since assignment of the $U(1)_H$ charge for extra leptons L' depends on its chirality. Then we roughly obtain the effective interaction as [27]

$$\mathcal{L}_{\alpha_G \gamma \gamma} \sim \frac{e^2}{32\pi^2 v'_1} \alpha_G F_{\mu\nu} \tilde{F}^{\mu\nu}, \quad (18)$$

where $F_{\mu\nu}$ is the photon field strength. Then cooling of stars provides a constraint for this coupling such that [28,29]

$$\frac{e^2}{32\pi^2 v'_1} \lesssim 6.6 \times 10^{-11} \text{ GeV}^{-1} \rightarrow v'_1 \gtrsim 4.6 \times 10^6 \text{ GeV}. \quad (19)$$

We thus consider v'_1 to be of a very high scale such as 5×10^6 GeV. In addition, experiments which search for

new force mediated by axionlike particles provide another constraint [30,31], but it is less severe compared to the star cooling one. If we take a large v'_1 , the existence of our GB does not cause serious problem in cosmology since it does not couple to the SM particles via some direct coupling, except to the Higgs boson whose couplings are well suppressed by a $1/v'_1$ factor and controlled by the parameters in the potential, which we assume to be small. Thus the GB decouples from the thermal bath in a sufficiently early Universe since all hidden particles are heavier than the $\mathcal{O}(100)$ GeV scale, and it almost does not affect the number of relativistic degrees of freedom in the current Universe. Even if the GB interacts with a SM particle via the Higgs portal, the cosmological constraint can be avoided when the GB decouples from the thermal bath at a scale larger than the muon mass scale [13].

A. Decay of exotic particles

In this subsection, we discuss the decays of exotic particles in the model. At first, we write the Lagrangian relevant to decay of the exotic charged lepton E_a as follows,

$$\begin{aligned} \mathcal{L} \supset & \frac{1}{v'_1} \frac{m_{D_a}}{M_a} \bar{\ell}^a \gamma^\mu P_L E^a \partial_\mu \alpha_G + f_{aa} \frac{m_{D_a}}{M_a} \bar{\ell}^a P_R E^a \varphi_{R_1} \\ & + \frac{y_{D_{aa}}}{\sqrt{2}} \bar{\ell}^a P_R E^a (\varphi_{R_2} + i\varphi_{I_2}) \\ & + \frac{y_{D_{aa}}}{\sqrt{2}} \frac{m_{D_a} m_{\ell^a}}{M_a^2} \bar{\ell}^a P_R \ell^a (\varphi_{R_2} + i\varphi_{I_2}) + \text{H.c.}, \end{aligned} \quad (20)$$

where we have omitted subdominant terms. The partial decay widths are computed as

$$\Gamma_{E^{a-} \rightarrow \ell^{a-} \varphi_{R_1}} \simeq \frac{f_{aa}^2}{64\pi M_a} \frac{y_{D_{aa}}^2 v_2'^2}{M_a^2} (M_a^2 - m_{\varphi_2}^2), \quad (21)$$

$$\Gamma_{E^{a-} \rightarrow \ell^{a-} \varphi_{R_2}} = \Gamma_{E^a \rightarrow \ell^a \varphi_{I_2}} \simeq \frac{y_{D_{aa}}^2}{64\pi M_a} (M_a^2 - m_{\varphi_2}^2), \quad (22)$$

$$\Gamma_{E^{a-} \rightarrow \ell^{a-} \alpha_G} \simeq \frac{y_{D_{aa}}^2 v_2'^2}{32\pi v_1'^2} M_a, \quad (23)$$

where we have used $m_{D_a} = y_{D_{aa}} v'_2 / \sqrt{2}$. Thus $E^{a\pm}$ dominantly decay into $\ell^{a\pm} \varphi_{R_2(I_2)}$ since the other modes are suppressed by the small VEV v'_2 . Then we also estimate partial decay widths of $\varphi_{R_2}(\varphi_{I_2})$ as

$$\Gamma_{\varphi_{R_2} \rightarrow \alpha_G \alpha_G} \simeq \frac{m_{\varphi_2}^3 v_2'^2}{8\pi v_1'^4}, \quad (24)$$

$$\Gamma_{\varphi_{R_2} \rightarrow \ell^{a-} \ell^{a+}} = \Gamma_{\varphi_{I_2} \rightarrow \ell^a \bar{\ell}^a} \simeq \frac{y_{D_{aa}}^4 v_2'^2 m_{\ell^a}^2}{32\pi M_a^4} m_{\varphi_2}. \quad (25)$$

Thus we find that φ_{R_2} and φ_{I_2} dominantly decays into $\ell^{a+} \ell^{a-}$ when we take $v'_1 = 5 \times 10^6$ GeV. Therefore

dominant decay chains of E^{a-} are $E^{a-} \rightarrow \ell^{a-} \varphi_{I_2} \rightarrow \ell^{a-} \ell^{b-} \ell^{b+}$ and $E^{a-} \rightarrow \ell^{a-} \varphi_{R_2} \rightarrow \ell^{a-} \ell^{b-} \ell^{b+}$ with branching ratio (BR) of 0.5 for both modes. Notice also that here we took the coupling between φ_R and SM Higgs to be zero, but φ_R can decay SM particles through such a coupling.

B. Collider physics

In this subsection, we discuss collider signatures of our model. The exotic charged scalar bosons from Higgs triplet can be produced by electroweak production, and they dominantly decay into gauge bosons, as the triplet components have degenerate mass. This phenomenology is the same as the scalar triplet model with a relatively large triplet VEV case ($v_\Delta \sim 1$ GeV) where charged scalars in the triplet $\{\delta^\pm, \delta^{\pm\pm}\}$ decays into SM gauge bosons. The phenomenology of such a case can be found, for example, in Refs. [32–36].

Then we focus on exotic charged-lepton production at the LHC. The exotic charged-lepton pair can be produced by $pp \rightarrow Z/\gamma \rightarrow E^{a+} E^{a-}$ via the gauge interaction depicted in Eq. (16). Furthermore, the single production process can be induced through mixing between E^a and SM charged leptons. We obtain the relevant interaction by

$$\mathcal{L} \supset \frac{g}{\sqrt{2}} \frac{m_{D_a}}{M_a} \bar{E}^a \gamma^\mu P_L \nu^a W_\mu^- + \text{H.c.}, \quad (26)$$

where the mixing effect in the neutral current is canceled. Thus E^a can be singly produced as $pp \rightarrow W^{-(+)} \rightarrow E^{a-} \bar{\nu}^a (E^{a+} \nu^a)$ at the LHC. For these processes, cross sections are estimated by using CalcHEP [37] with the CTEQ6 parton distribution functions (PDFs) [38]. We show pair and single production cross sections as a function of the exotic lepton mass in the left and right panels in Fig. 1, where we fixed $m_{D_a}/M_a = 0.01$ in our calculation and where only the lightest exotic charged lepton $E^\pm$ is considered. Then we find that the single production cross section is much smaller than the pair production one due to the suppression by mixing factor m_{D_a}/M_a . The pair production cross section is larger than ~ 1 fb when the exotic lepton mass is $M_a \lesssim 700$ GeV for $\sqrt{s} = 14$ TeV. Since the exotic charged-lepton decays into a multilepton final state, the cross section can be constrained by the multilepton search at the LHC, where an inclusive search indicates a cross section producing more than three electrons/muon is required to be $\sigma \cdot BR \lesssim 1$ fb at LHC 8 TeV [39]. We thus require our exotic charged-lepton mass to be $M_a \gtrsim 500$ GeV.

One specific signal from the $E^+ E^-$ pair production is the multilepton process given by $E^+ E^- \rightarrow \ell^{a+} \ell^{b-} \varphi_{I_2, R_2} \varphi_{I_2, R_2} \rightarrow \ell^{a+} \ell^{b-} \ell^{c+} \ell^{c-} \ell^{d+} \ell^{d-}$. Notice that we require scalar bosons φ_{I_2, R_2} to have masses of $\mathcal{O}(100)$ GeV so that the decay process is dominant; if these scalar bosons are much heavier, the dominant decay mode of $E^\pm$ is $W^\pm \nu$ mode through mixing between the heavy extra lepton and the SM

FIG. 1. (Left panel) $E^+ E^-$ pair production cross section as a function of its mass. (Right panel) $E^\pm$ single production cross section as a function of its mass where we fix $m_D/M = 0.01$.

FIG. 2. (Left panel) Number of events for the invariant mass for $e^+ e^+ e^-$ with $m_\phi \equiv m_{\phi_{I_2, R_2}} = 100$ GeV and integrated luminosity 300 fb^{-1} at the LHC 14 TeV. (Right panel) Number of events for invariant mass for $e^+ e^-$.

leptons. Here signal events are generated by employing the event generator MADGRAPH/MADEVENT 5 [40], where the necessary Feynman rules and relevant parameters of the model are implemented by use of FeynRules 2.0 [41] and the NNPDF23LO1 PDF [42] is adopted. Then the PYTHIA 6 [43] is applied to deal with hadronization effects, the initial-state radiation and final-state radiation effects, and the generated events are also run through the PGS 4 for detector level simulation [44]. Then we select events with six charged leptons and impose cuts for lepton transverse momentum as $p_T(\ell) > 15$ GeV. In the left panel of Fig. 2, we show the number of events distribution for invariant mass for $e^+ e^+ e^-$, assuming that $E^\pm$ and ϕ_{I_2, R_2} dominantly decay into a mode including an electron and $m_{\phi_{I_2, R_2}} = 100$ GeV with an integrated luminosity 300 fb^{-1}

at the LHC 14 TeV.⁶ We see the clear peak at the exotic lepton mass and distributions around the peak as we cannot always choose the $e^+ e^+ e^-$ combination coming from the E^+ decay chain. In addition, we show the distribution for the invariant mass for $e^+ e^-$ in the right panel of Fig. 2, which shows the peak at the mass of ϕ_{I_2, R_2} . The signal is clean, and number background (BG) events can be suppressed, requiring a multiple charged-lepton final state. In fact the number of BG events is negligibly small, requiring six charged leptons in the final state, where the largest cross section from SM processes is $\sigma(pp \rightarrow ZZZ, Z \rightarrow \ell^+ \ell^-)$ and it is around 10^{-3} fb magnitude. Thus the signal can be

⁶The invariant mass for $e^- e^- e^+$ provides a similar distribution.

tested well at the future LHC experiments if the exotic lepton mass is less than 1 TeV. A more detailed simulation including sophisticated cut analysis is beyond the scope of this paper, and it is left for future work.

IV. SUMMARY AND CONCLUSIONS

We have constructed an inverse seesaw model based on $U(1)_H$ global symmetry in which we have introduced exotic lepton doublets L' and new scalar fields including scalar triplet with $U(1)_H$ charge. The exotic lepton doublets are vectorlike under gauge symmetry but chiral under $U(1)_H$ and Majorana mass of neutral components are zero before $U(1)_H$ symmetry breaking. Then $U(1)_H$ is spontaneously broken by VEVs of scalar fields, and only the neutral component of L'_L obtains Majorana mass term via triplet VEV due to appropriate $U(1)_H$ charge assignment. We thus realize neutral fermion mass matrix which has a structure similar to the inverse seesaw mechanism. Our advantage of this model is to realize the neutrino mass matrix with natural mass parameters in the framework of the inverse seesaw scenario. Therefore the lepton number is broken by exotic fermions, and its violation is restricted by the triplet VEV, which is of the order 1 GeV, originating from the constraint of the oblique parameters. Also m_D/M , which is also proportional to the neutrino mass matrix, is naturally suppressed when we appropriately assign the lepton number to the fields. This is because m_D is

proportional to $\mu_\phi \sim \mathcal{O}(1)$ GeV, which is the source of lepton number violation, which is expected to be small due to the 't Hooft sense, while M should be greater than 100 GeV, which is required by current experimental data such as that from the LHC and the LEP. We thus have naturally achieved a neutrino mass matrix less than the order of 10^{-4} GeV.

We have also discussed the phenomenology of the model focusing on decay and production processes of the exotic particles. The decay widths of exotic charged leptons $E^\pm$ are estimated by taking into account interactions with the Goldstone boson from $U(1)_H$ symmetry breaking. We have found that $E^\pm$ will provide multilepton final states from the decay chain. Then an $E^\pm$ production cross section has been estimated which is induced by electroweak interactions. Testable number of multilepton events could be obtained if mass of $E^\pm$ is less than the $\mathcal{O}(1)$ TeV scale.

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