Natural and dynamical neutrino mass mechanism at the LHC

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We generalize the scalar triplet neutrino mass model, the type II seesaw. Our model displays compelling theoretical and experimental features. First, lepton number violation is dynamically generated, without introducing a naturalness problem, nor relying on arbitrarily small parameters. Second, we identify a smoking gun signature at the LHC that both probes the triplet structure of our seesaw mechanism and allows to disentangle it from the usual type II scenario. Additionally, we discuss other interesting phenomenological aspects of the model such as the presence of a massless Goldstone boson and deviations of standard model Higgs couplings.

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I. INTRODUCTION

The presence of nonzero neutrino masses, as inferred by neutrino oscillation experiments, is the only laboratory-based evidence of physics beyond the standard model [1,2]. Strictly speaking, neutrinos have no mass in the standard model (SM). There is no unique prescription of how neutrinos could become massive. Perhaps the simplest way of generating neutrino masses is via the seesaw framework. In its naïve realizations, seesaw types I, II and III [3–10], a large suppression of the electroweak breaking scale provides an explanation for the smallness of neutrino masses. Without a full underlying framework, like grand unified theories or supersymmetry, these mechanisms typically introduce a hierarchy problem due to the large mass gap [11] or rely on very small (but technically natural [12]) parameters.

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In general, the seesaw mechanism generates a small parameter from the ratio of two disparate physics scales, e.g., electroweak versus grand unification scales. Therefore, when we set the new heavy states to the weak scale (such as done in studies of type II seesaw at colliders [13,14]), the "seesaw" mechanism is exchanged by a small parameter. This can be appreciated in a model independent way by writing down schematically the Weinberg effective operator that generates neutrino masses [15], namely

$$\mathcal{L}_5 = \frac{c}{\Lambda} LLHH \tag{1}$$

(H and L are the Higgs and lepton doublets) and observing that if $\Lambda \sim \langle H \rangle$ then the Wilson coefficient c needs to be tiny in order to obtain sub-electronvolt neutrino masses. We will show in this Letter that a simple generalization of the type II seesaw, replacing the seesaw by a chain of seesaws, can completely avoid the postulation of a new scale. The model is entirely at the weak scale, and a small lepton number breaking is dynamically generated.

More concretely, in type II seesaw a scalar triplet

$$\Delta = \begin{pmatrix} \delta^+/\sqrt{2} & \delta^{++} \\ (v_\Delta + \delta + i a_\delta)/\sqrt{2} & -\delta^+/\sqrt{2} \end{pmatrix} \tag{2}$$

obtains its vacuum expectation value (vev) after electroweak symmetry breaking

$$v_\Delta \simeq \frac{\mu}{\sqrt{2}} \frac{v^2}{M_\Delta^2}, \tag{3}$$

where μ is a dimensionful lepton number breaking parameter of the scalar potential, $v = 246$ GeV is the Higgs doublet vev, and M_Δ is approximately the physical mass of Δ. Neutrino masses are given by $m_\nu = \sqrt{2} Y v_\Delta$, with Y being a matrix of Yukawa couplings.

We can immediately see that the smallness of neutrino masses can only be obtained by having small Yukawas, large M_Δ, and/or small *ad hoc* lepton number breaking parameter μ. For instance, if M_Δ is accessible at the LHC, say at the TeV scale, and the Yukawas are taken to be of order 1, we obtain

$$\mu \simeq 1.6 \text{ eV} \left(\frac{m_\nu}{0.1 \text{ eV}} \right). \tag{4}$$

Since $\mu = 0$ restores a symmetry of the Lagrangian, it is not generated by other couplings due to quantum corrections, thus being technically natural in the t'Hooft sense [12]. Nevertheless, it is unappealing to have this enormous hierarchy of scales $\mu/v \lesssim \mathcal{O}(10^{-11})$ put in arbitrarily. As suggested by the considerations made before regarding the Weinberg operator, this is not exclusive to type II seesaw.

In this paper we present a generalization of the type II seesaw scenario which dynamically generates a very low lepton number breaking scale from a small hierarchy. The model is naturally found at the weak scale, introducing no new fine-tuning neither arbitrarily small couplings. Our mechanism engenders a rich phenomenology, including deviations of SM Higgs couplings, presence of a massless Majoron, lepton flavor violation and a smoking gun signature at the LHC that distinguishes this model from the usual type II seesaw.

II. THE MECHANISM

The idea simply amounts to replicate the induced vev suppression mechanism with additional scalar singlets, as shown in Fig. 1. In our concrete setup, all mass parameters are near the electroweak scale and all dimensionless couplings are of similar order, thus yielding a natural model of neutrino masses accessible at the LHC. We will focus on a scenario with two extra scalar singlets, as this is the most minimal realization that successfully implements the mechanism and also exhibits all important phenomenological features of our framework.

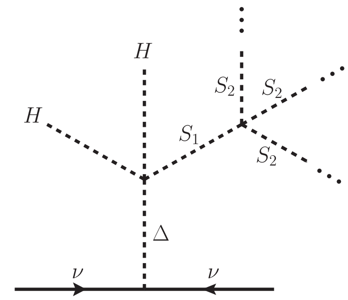

FIG. 1. Illustration of the generalized type II seesaw mechanism for neutrino mass generation.

First, we require dynamical lepton number breaking. To that end, we promote $U(1)_\ell$ lepton number to a global symmetry in which leptons have charge $\ell_{\text{leptons}} = +1/2$ (the normalization has been chosen for convenience) and quarks have no charge. The neutrino Yukawa coupling

$$\mathcal{L}_{\text{Yuk}}^\nu = -Y \overline{L^c} i\sigma_2 \Delta L + \text{H.c.} \tag{5}$$

(σ_2 is the second Pauli matrix, Y is a matrix of Yukawa couplings in flavor space, and ^{c} denotes charge conjugation) requires $\ell_\Delta = -1$, forbidding the triple coupling $\mu H^T i\sigma_2 \Delta^\dagger H$. We introduce the first complex SM singlet scalar S_1 with lepton number $\ell_1 = +1$ so its vev may play the role of lepton number violating parameter μ. Then, we generalize the type II seesaw model by invoking another extra scalar singlet with charge $\ell_2 = 1/3$, allowing for a term $S_1^* S_2^3$ in the scalar potential. All scalars but the Higgs and S_2 have positive bare mass terms. The crucial point is that when S_2 develops a vev spontaneously, it induces a suppressed vev for S_1, which then induces an even smaller vev for Δ. The model can easily be generalized for any number N of scalar singlets, see Appendix A. We identify the usual type II seesaw with a $N = 1$-step version of the generalized model in which S_1 is integrated out. Our model bears similarities with multiple seesaw and clockwork models (see, for instance, Refs. [16–24]).

As we will see later, a simple 2-step realization can lead to small neutrino masses given that some quartic couplings and neutrino Yukawas are of order 10^{-2}–10^{-3} (larger couplings can be obtained in realizations with extra steps). Without further ado, we write down the scalar potential

$$\begin{aligned} V = &-\frac{m_H^2}{2} H^\dagger H + m_\Delta^2 \langle \Delta^\dagger \Delta \rangle + m_1^2 S_1^* S_1 - \frac{m_2^2}{2} S_2^* S_2 + \frac{\lambda_H}{4} (H^\dagger H)^2 + \frac{\lambda_2}{4} (S_2^* S_2)^2 \\ &+ \lambda_{1H} (S_1^* S_1)(H^\dagger H) + \lambda_{2H} (S_2^* S_2)(H^\dagger H) + \left[\lambda_A H^T i\sigma_2 \Delta^\dagger H S_1^* - \frac{2}{3} \lambda'_{12} S_1^* S_2^3 + \text{H.c.} \right] \\ &\left. \begin{aligned} &+ \tfrac{\lambda_\Delta}{4} \langle \Delta^\dagger \Delta \rangle^2 + \tfrac{\lambda'_\Delta}{4} \langle \Delta^\dagger \Delta \Delta^\dagger \Delta \rangle + \tfrac{\lambda_1}{4} (S_1^* S_1)^2 + \lambda_{12} (S_1^* S_1)(S_2^* S_2) \\ &+ \lambda_{H\Delta} (H^\dagger H) \langle \Delta^\dagger \Delta \rangle + \lambda_{H\Delta}{}' \langle H^\dagger \Delta \Delta^\dagger H \rangle + \lambda_{1\Delta} \langle \Delta^\dagger \Delta \rangle (S_1^* S_1) + \lambda_{2\Delta} \langle \Delta^\dagger \Delta \rangle (S_2^* S_2), \end{aligned} \right\} \text{“incidental” terms} \end{aligned} \tag{6}$$

where the parameters more relevant for the mechanism and phenomenology are in the first two lines. Although the quartic couplings on third and fourth lines are important for the potential stability, they play almost no role otherwise (thus called "incidental"). The potential stability is not a primary concern of this manuscript, but it is important to note that the quartic couplings λ_A and λ'_{12} tend to destabilize the potential, and hence are expected to be small. For more considerations regarding stability see Appendix B. We define the neutral components of the fields as $H^0 = (v + h + ia)/\sqrt{2}$, $\Delta^0 = (v_\Delta + \delta + ia_\delta)/\sqrt{2}$ and $S_j = (v_j + s_j + ia_j)/\sqrt{2}$, for $j = 1, 2$.

The positive mass terms for Δ and S_1 ensure that if $\lambda_A = \lambda'_{12} = 0$ then the vevs for these fields are zero. These two quartic couplings are protected from loop corrections by accidental global $U(1)$ symmetries. Moreover, λ_A and λ'_{12} can be made real by rephasing the scalar singlet fields. As long as v_Δ and v_1 are much smaller than v and v_2, we can obtain the former vevs by treating H and S_2 as background fields. First, we obtain the approximate vevs of H and S_2 by setting other scalar fields to zero,

$$m_H^2 = \frac{1}{2}\lambda_H v^2 + \lambda_{2H} v_2^2, \qquad m_2^2 = \frac{1}{2}\lambda_2 v_2^2 + \lambda_{2H} v^2. \quad (7)$$

Then, by replacing H and S_2 by their vevs, we can easily calculate the vevs and spectrum of other scalars:

$$v_1 = \frac{\lambda'_{12} v_2^3}{3M_1^2}, \qquad v_\Delta = \frac{\lambda_A v^2 v_1}{2M_\Delta^2}, \quad (8)$$

and

$$M_h^2 = \frac{1}{2}\lambda_H v^2, \quad (9a)$$

$$M_1^2 = m_1^2 + \frac{1}{2}(\lambda_{1H} v^2 + \lambda_{12} v_2^2), \quad (9b)$$

$$M_2^2 = \frac{1}{2}\lambda_2 v_2^2, \quad (9c)$$

$$M_\Delta^2 = m_\Delta^2 + \frac{1}{2}[\lambda_{2\Delta} v_2^2 + (\lambda_{H\Delta} + \lambda'_{H\Delta})v^2]. \quad (9d)$$

The physical masses of the scalars are approximately given by the M's in Eqs. (9). Here we see the mechanism at work: λ'_{12} induces a suppression from v_2 to v_1, and λ_A induces a further suppression from v_1 to v_Δ. It is useful to write these quartics in terms of the scalar masses and vevs,

$$\lambda_A = 0.008\left(\frac{M_\Delta}{500\ \text{GeV}}\right)^2\left(\frac{v_\Delta/\text{keV}}{v_1/\text{MeV}}\right), \quad (10a)$$

$$\lambda'_{12} = 0.03\frac{(M_1/100\ \text{GeV})^2(v_1/\text{MeV})}{(v_2/10\ \text{GeV})^3}. \quad (10b)$$

These relations do not depend on the number of steps, as long as perturbation theory holds.

III. SPECTRUM AND MIXING PHENOMENOLOGY

The scalar spectrum of this 2-step scenario consists of the 4 aforementioned neutral scalars (h, δ, s_1, s_2), singly and doubly charged scalars δ^+ and δ^{++}, with masses approximately given by M_Δ, two massive pseudoscalar degrees of freedom (a_δ, a_1) with masses approximately given by M_Δ and M_1, and two massless Goldstone bosons. One of the Goldstones is the longitudinal polarization of the Z boson while the other is a massless Majoron, J [25–27]. We will analyze the Majoron phenomenology in the following section.

The mixings among CP even scalars will have important phenomenological impacts (see Table I for a summary). The mixings between $h - s_2$, $\delta - s_1$ and $h - s_1$ are given by

$$\theta_{h2} \simeq \frac{\lambda_{2H} v_2 v}{M_h^2 - M_2^2} \simeq 0.16\lambda_{2H}\left(\frac{v_2}{10\ \text{GeV}}\right)\beta_{h2}, \quad (11a)$$

$$\theta_{\delta 1} \simeq \frac{\lambda_A}{2}\frac{v^2}{M_1^2 - M_\Delta^2} \simeq 10^{-3}\left(\frac{v_\Delta/\text{keV}}{v_1/\text{MeV}}\right)\beta_{1\delta}, \quad (11b)$$

$$\theta_{h1} \simeq \frac{\lambda_{1H} v_1 v}{M_h^2 - M_1^2} \simeq 1.5\times 10^{-5}\lambda_{1H}\left(\frac{v_1}{\text{Mev}}\right)\beta_{h1}, \quad (11c)$$

where $\beta_{ab} \equiv (1 - M_b^2/M_a^2)^{-1}$. First, the Higgs mixing with s_2 could in principle be sizable. Observations of Higgs production and decay modes together with precision electroweak measurements constrain the mixing angle α with a scalar singlet to be about $\sin\theta_{h2} \lesssim 0.2$–$0.3$ for a 200–800 GeV singlet mass [28]. If the scalar is much lighter than the Higgs, for instance in the region $1 < M_2 < 10$ GeV, the constraints on the mixing range from $\sin\theta_{h2} \lesssim 10^{-3}$–$10^{-1}$ [29]. This Higgs-singlet mixing can lead to very interesting phenomenology, but it is not an exclusive signature of our model. For small values of v_2, the invisible Higgs decay to a pair of Majorons strongly constrains this mixing, as we will see later.

The mixing between δ and s_1 is quite special, as it leads to drastic deviations from usual type II seesaw phenomenology. For δ^{++}, a new decay channel may open up,

TABLE I. Sizable scalar mixings and their phenomenological impact.

Mixing	Phenomenology
$h - s_2$	Higgs observables, direct s_2 production
$\delta - s_1$	New LHC signatures, s_1 decay modes
$h - s_1$	s_1 decay modes
$s_1 - s_2$	Irrelevant

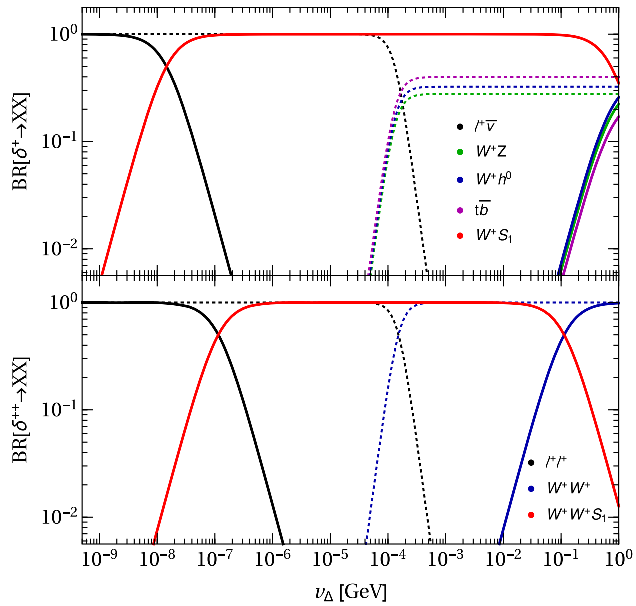

FIG. 2. Branching ratios of δ^+ (upper panel) and δ^{++} (lower panel) as a function of the triplet vev v_Δ for the usual type II seesaw model (dotted) and our generalized version (solid). We considered $m_0 = 0.1$ eV, as the lightest neutrino mass, $M_{\delta^+} = M_{\delta^{++}} = 500$ GeV, $M_1 = 100$ GeV, and $\theta_{\delta 1} = 0.005$.

$\delta^{++} \to W^+W^+s_1$ with s_1 typically decaying to neutrinos (via mixing with δ), quarks or gauge bosons (both via mixing with the Higgs) depending on its mass. Similarly, one can have $\delta^+ \to W^+s_1$ and $\delta \to hs_1$. Another distinctive feature is the possibility of having sizable visible pseudoscalar decays, $a_\delta \to Zs_1$. Differently from type II seesaw, these decays are controlled uniquely by gauge coupling and mixing angle $\theta_{\delta 1}$, see Table II. As can be seen in Fig. 2, the new decays can dominate a large region of parameter space in the generalized type II seesaw (solid lines) compared to the usual case (dotted lines). We stress that the mixing angle $\theta_{\delta 1}$ essentially drives the large, novel branching ratio, making the model more predictive. As we will see later, these new decays provide a smoking gun signature at the LHC, not only opening the possibility for discovering new particles, but also distinguishing the model from the type II seesaw.

TABLE II. Typical decay modes in type II seesaw and new modes in the generalized type II framework. In the last column it is indicated the most relevant parameters governing the partial widths.

Scalar	Type II	Generalized type II	Parameters
δ^{++}	$\ell^+\ell^+$, W^+W^+	$W^+W^+s_1$	v_Δ, $\theta_{\delta 1}$
δ^+	$\ell^+\nu$, W^+Z, W^+h, $t\bar{b}$	W^+s_1	v_Δ, $\theta_{\delta 1}$
δ	$\nu\nu$, W^+W^-, ZZ, hh	hs_1	v_Δ, $\theta_{\delta 1}$
a_δ	$\nu\nu$, $t\bar{t}$, Zh	Zs_1	v_Δ, $\theta_{\delta 1}$
s_1	not present	$\nu\nu$, $q\bar{q}$, W^+W^-, ZZ	v_Δ, $\theta_{\delta 1}$, θ_{h1}

Finally, the mixing between the Higgs and s_1 given in Eq. (11), although small, plays a significant role in the scalar phenomenology. The s_1 decay to charged fermions, driven by θ_{h1}, will compete with the invisible decay to neutrinos, sourced by $\theta_{\delta 1}$. By analyzing the ratio of these partial widths (see Appendix C for more details),

$$\frac{\Gamma_{s_1\to\nu\nu}}{\Gamma_{s_1\to ff}} \simeq \frac{3.1}{N_c}\left(\frac{\theta_{\delta 1}/10^{-3}}{\theta_{h1}/10^{-5}}\right)^2\left(\frac{m_\nu/0.1\ \text{eV}}{m_f/\text{GeV}}\right)^2\left(\frac{\text{keV}}{v_\Delta}\right)^2,$$

we can see that either visible or invisible s_1 decays can dominate in large natural regions of the parameter space. In this manuscript we will focus on the latter. Besides, there is some region of parameter space in which s_1 decays to b quarks and gives rise to displaced vertices at the LHC. We will nevertheless refrain from analyzing that possibility here.

IV. MAJORON PHENOMENOLOGY

Although a massless particle in the spectrum may at first seem problematic, its couplings to standard model fermions are extremely suppressed due to hierarchy of vevs. The Majoron field is the linear combination

$$J \simeq \frac{1}{\ell_2 v_2}\left(\ell_1 v_1 a_1 + \ell_2 v_2 a_2 + \frac{1}{2}v_\Delta a_\delta - \frac{v_\Delta^2}{v}a\right), \quad (12)$$

where $\ell_1 = 1$ and $\ell_2 = 1/3$ are lepton numbers of the corresponding scalars. It is straightforward to see that the Majoron has very small couplings to charged fermions given by

$$G_{Jff} = \frac{y_f}{\sqrt{2}}\frac{v_\Delta^2}{\ell_2 v_2 v} = \frac{1.6\times 10^{-18}}{\ell_2}\frac{(m_f/\text{GeV})(v_\Delta/\text{keV})^2}{(v_2/10\ \text{GeV})},$$
$$G_{J\nu\nu} = \sqrt{2}y_\nu\frac{v_\Delta}{\ell_2 v_2} = \frac{5\times 10^{-12}}{\ell_2}\frac{(m_\nu/0.1\ \text{eV})}{(v_2/10\ \text{GeV})}, \quad (13)$$

easily avoiding constraints from neutrinoless double beta decay with Majoron emission $G_{J\nu\nu} < (0.8-1.6)\times 10^{-5}$ [30], as well as astrophysical bounds $G_{Jee} < 4.3\times 10^{-13}$ [31]. Although a thermalized Majoron would contribute to increase the effective number of relativistic degrees of freedom by 4/7, the tiny coupling in this scenario leads to very little Majoron production in the early universe.

A stringent bound on Higgs-s_2 mixing comes from Higgs decaying invisibly to a pair of Majorons [32]. It is straightforward to obtain the approximate constraint [33]

$$\theta_{h2} < 1.5\times 10^{-3}\left[\frac{v_2}{10\ \text{GeV}}\right]\left[\frac{\Gamma_h}{4.2\ \text{MeV}}\frac{\text{BR}_{h\to\text{inv}}}{0.22}\right]^{1/2} \quad (14)$$

where Γ_h is the Higgs total width and $\text{BR}_{h\to\text{inv}}$ is its invisible branching ratio. The Higgs total width has only been measured indirectly, via comparison between on-shell and off-shell Higgs production, yielding the model-dependent bound $\Gamma_h^{\text{exp}} < 13$ MeV at 95% C.L. [34]. The

Higgs invisible branching ratio has been bounded to be below 0.22 [35,36]. This strong bound on θ_{h2} could be alleviated by raising v_2 to the TeV. Complementary constraints (though still not competitive) come from monojets at the LHC [37].

V. COLLIDER PHENOMENOLOGY

In this section, we study the collider phenomenology for the generalised type II seesaw model. The leading production channels for this framework remain the same as in the usual type II, i.e., the charged Higgs states will be dominantly produced in pairs via s-channel electroweak boson exchange, leading primarily to associated production of double and single charged Higgs bosons $\delta^{\pm\pm}\delta^{\mp}$, followed by double charged Higgs pair production $\delta^{++}\delta^{--}$.[1] Although these two production channels do not present differences in rate between the standard type II seesaw and our new model construction, their corresponding decays display new relevant phenomenological signatures. The δ—s_1 mixing engenders new interaction terms from the triplet kinetic term

$$\mathcal{L} \supset \mathrm{Tr}[(D_\mu\Delta)^\dagger D_\mu\Delta], \tag{15}$$

making the decays $\delta^{\pm\pm} \to W^\pm W^\pm s_1$ and $\delta^\pm \to W^\pm s_1$ available. Note that these partial widths do not present any v_Δ suppression, instead they depend only on gauge couplings, being equally large in a wide range of parameter space $v_\Delta \sim 10^{-7}$–10^{-1} GeV, distinctly from the usual type II, see Fig. 2.

Therefore, the $pp \to \delta^{\pm\pm}\delta^{\mp}$ production channel not only reveals the triplet structure nature of $\delta^{\pm\pm}$ and $\delta^\pm$ [13,14], but can also differentiate our construction from the usual type II model. To explore this phenomenology, we analyse the $pp \to \delta^{\pm\pm}\delta^{\mp}$ production at the $\sqrt{s} = 13$ TeV LHC, focusing on the trilepton plus missing energy signature, with two same flavor and same sign leptons, $e^\pm e^\pm \mu^\mp + \not{E}_T$ and $\mu^\pm\mu^\pm e^\mp + \not{E}_T$. The leptons arise from the W-boson decays and relevant extra sources of missing energy follow from the dominant s_1 decay, $s_1 \to \nu\bar{\nu}$.

Our model is implemented in FeynRules [38] and the signal sample is generated with MadGraph5 [39]. A Next-to-leading order QCD K-factor of 1.25 has been applied [40]. To obtain a robust simulation of the background components, that display large fake rates, our simulation follows the recent 13 TeV CMS study [41]. Although CMS targets a heavy neutral Majorana lepton N, it presents a set of search regions for the high mass regime $m_N > m_W$, leading to a more sizable $\not{E}_T$, that also applies to our model.

[1] We have checked that producing one triplet scalar in association with s_1 is typically subleading, as it is suppressed by the small mixing $\theta_{\delta 1}$. Thus, these production modes will be disregarded here.

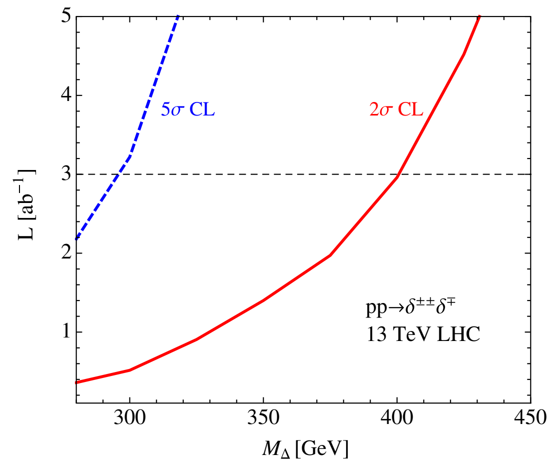

FIG. 3. Luminosity required to observe $pp \to \delta^{\pm\pm}\delta^{\mp}$ as a function of M_Δ at 2σ (red full) and 5σ (blue dashed) confidence level. We assume $M_1 = 100$ GeV and $v_\Delta = 10^{-6}$ GeV.

In this analysis, jets are defined with the anti-k_T clustering algorithm with $R = 0.4$, $p_{Tj} > 25$ GeV and $|\eta_j| < 2.4$ via FastJet [42]. Events with one or more b-jets are vetoed with 70% b-tagging efficiency and 1% mistag rate. Electrons and muons are defined with $|\eta_\ell| < 2.4$ and the three leptons must satisfy $p_{T\ell} > 55, 15, 10$ GeV. Finally, the events are divided in bins associated to three observables: (i) the trilepton mass system $m_{3\ell}$; (ii) minimum invariant mass of all opposite sign leptons $m_{2\ell OS'}^{\rm min}$; and (iii) transverse mass $m_T = \sqrt{2p_{T\ell}\not{E}_T(1-\cos\phi)}$, where $p_{T\ell}$ corresponds to the lepton which is not used in the $m_{2\ell OS'}^{\rm min}$ calculation and ϕ is the azimuthal angle between $\vec{p}_{T\ell}$ and $\vec{p}_T^{\rm miss}$.

Using the CMS background estimate, we perform a binned log-likelihood analysis based on the CL$_s$ method [43], exploring all search regions with $e^\pm e^\pm \mu^\mp + \not{E}_T$ and $\mu^\pm\mu^\pm e^\mp + \not{E}_T$ displayed by Ref. [41]. In Fig. 3, we present the luminosity required to observe $pp \to \delta^{\pm\pm}\delta^{\mp}$ as a function of M_Δ at 2σ and 5σ confidence level. At the high-luminosity LHC, $\mathcal{L} = 3\ \mathrm{ab}^{-1}$, we can discover charged Higgses at 5σ level up to $M_\Delta = 300$ GeV and exclude them at 2σ level up to $M_\Delta = 400$ GeV.

A final comment is in order regarding two phenomenological aspects beyond the ones discussed so far. First, our model may also induce lepton flavor violation processes, very similar to the usual type II seesaw scenario [44]. Second, although the model does not have enough CP violation, adding a second $SU(2)$ triplet scalar [45] may lead to successful leptogenesis. The study of such possibilities is beyond the scope of this manuscript.

VI. CONCLUSIONS

In this paper we have proposed a generalization of type II seesaw in which lepton number is broken dynamically and no hierarchy problem neither arbitrarily small parameters are present. The rich phenomenology of the model includes deviations of standard Higgs couplings, the presence of a massless neutral pseudoscalar and more importantly a novel smoking gun signature at the LHC. This distinctive new signature may reveal the triplet nature of the charged scalars and at the same time disentangle the framework from the usual type II seesaw model.

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APPENDIX A: *n*-STEP GENERALIZED TYPE II SEESAW

Here we present the generalization of our framework for an arbitrary number of scalar singlets n. We define the following scalar bilinears,

$$B_i \equiv S_i^* S_i, \qquad B_\Delta \equiv \mathrm{Tr}(\Delta^\dagger \Delta), \qquad B_H \equiv H^\dagger H, \tag{A1}$$

which allow to write the scalar potential in a compact form

$$\begin{aligned} V = & -\frac{m_H^2}{2} B_H + \sum_{\varphi}^{\Delta,1\ldots n-1} m_\varphi^2 B_\varphi - \frac{m_n^2}{2} B_n + \sum_{\varphi}^{\text{all}} \frac{\lambda_\varphi}{4} B_\varphi^2 \\ & + \sum_{\varphi,\varphi'>\varphi}^{\text{all}} \lambda_{\varphi\varphi'} B_\varphi B_{\varphi'} + \frac{\lambda'_\Delta}{4} \mathrm{Tr}(\Delta^\dagger \Delta \Delta^\dagger \Delta) + \lambda'_{H\Delta} H^\dagger \Delta \Delta^\dagger H \\ & + \left[\lambda_A H^T i\sigma_2 \Delta^\dagger H S_1^* - \frac{2}{3} \sum_{i=1}^{n-1} \lambda'_{i,i+1} S_i^* S_{i+1}^3 + \text{H.c.} \right]. \end{aligned} \tag{A2}$$

The notation in the sum of the first term of the second line indicates that permutations of $\lambda_{\varphi\varphi'}$ should not be taken (to avoid double counting). Without loss of generality, all $\lambda'_{i,i+1}$ and λ_A can be made real by rephasing the scalar singlet fields. The masses and vevs in the n-step realization are approximately given by

$$m_H^2 = \frac{1}{2} \lambda_H v^2 + \lambda_{nH} v_n^2, \tag{A3a}$$

$$m_n^2 = \frac{1}{2} \lambda_n v_n^2 + \lambda_{nH} v^2, \tag{A3b}$$

$$v_i = \frac{\lambda'_{i,i+1} v_{i+1}^3}{3 M_i^2}, \quad \text{for } i = 1, \ldots, n-1, \tag{A3c}$$

$$v_\Delta = \frac{\lambda_A v^2 v_1}{2 M_\Delta^2}, \tag{A3d}$$

$$M_h^2 = \frac{1}{2} \lambda_H v^2, \tag{A3e}$$

$$M_i^2 = m_i^2 + \frac{1}{2} (\lambda_{iH} v^2 + \lambda_{\text{in}} v_n^2), \quad i = 1, \ldots, n-1, \tag{A3f}$$

$$M_n^2 = \frac{1}{2} \lambda_n v_n^2, \tag{A3g}$$

$$M_\Delta^2 = m_\Delta^2 + \frac{1}{2} [\lambda_{n\Delta} v_n^2 + (\lambda_{H\Delta} + \lambda'_{H\Delta}) v^2]. \tag{A3h}$$

These expressions should hold in the regime, $v_i \ll v_{i+1}$, that is,

$$\varepsilon \equiv \lambda'_{i,i+1} \frac{v_{i+1}^2}{3 M_i^2} \ll 1. \tag{A4}$$

In fact, it is straightforward to show that as long as Eq. (A4) is satisfied, for any number n of scalar singlet fields, the vev of s_j, $j < n$, is simply given by

$$v_j = \prod_{k=0}^{n-j-1} \left(\frac{\lambda'_{j+k,j+k+1}}{3} \frac{v_n^2}{M_{j+k}^2} \right)^{3^k} v_n. \tag{A5}$$

If, for simplicity, one takes all $\lambda'_{ij} = \lambda'$ and $M_i = M$, then we obtain a simplified expression,

$$v_j = \left(\frac{\lambda'}{3} \frac{v_n^2}{M^2} \right)^K v_n, \quad K = (3^{n-j} - 1)/2. \tag{A6}$$

We can clearly identify the parametric suppression ε^K responsible for making $v_1 \ll v_n$. For instance, if $\varepsilon = 0.01$ and $n = 3$ we obtain $v_1 \sim 10^{-8} v_n$. Note that the expressions for the mixing angles defined in Eqs. (11) are valid for any n, and thus the phenomenological considerations regarding Higgs couplings, Majoron physics and LHC signatures will still apply.

APPENDIX B: STABILITY OF THE SCALAR POTENTIAL

Although a complete study on the stability of the scalar potential are not the main focus of this Letter, we provide here sufficient conditions for the stability. The key point is that the quartic couplings λ_A and λ'_{12} (or any $\lambda'_{i,i+1}$ in the n-step scenario) can always yield negative contributions to the potential when the values of the fields go to infinity, independently of their sign. As these couplings are the core of the generalized type II seesaw mechanism, it is important to understand how to control these contribution so that the potential is bounded from below. Although a full analysis of the stability would be very complicated, specially in the n-step scenario, we can still derive useful sufficient conditions to have stability. The idea is to split the scalar potential into pieces that will isolate each λ'_{12} or λ_A,

$$V = V_A + V_{12} + \cdots + V_0 \tag{B1}$$

and require each piece to be independently positive. For now we will focus on $n = 2$-steps and generalize the method in the end.

The first piece deals with λ_A. We define

$$V_A \equiv \lambda_{1H}(S_1^* S_1)(H^\dagger H) + \lambda_{1\Delta}\langle\Delta^\dagger\Delta\rangle(S_1^* S_1) + \lambda_{H\Delta}(H^\dagger H)\langle\Delta^\dagger\Delta\rangle + (\lambda_A H^T i\sigma_2 \Delta^\dagger H S_1^* + \text{H.c.}) \tag{B2}$$

and require it to be positive. By performing an $SU(2)$ rotation on the field one can always write [46]

$$i\sigma_2\Delta = \begin{pmatrix} a & 0 \\ 0 & be^{i\alpha} \end{pmatrix}, \qquad H = \begin{pmatrix} ce^{i\beta} \\ de^{i\gamma} \end{pmatrix}, \tag{B3}$$

and $S_i = R_i e^{i\phi_i}$. Then, it is straightforward to obtain

$$\lambda_{H\Delta} > 0, \qquad \lambda_{1\Delta} > 0, \tag{B4a}$$

$$\lambda_{1H} > 0, \qquad |\lambda_A|^2 < \lambda_{1H}\lambda_{H\Delta}. \tag{B4b}$$

Now, we handle λ'_{12} by defining

$$V_{12} \equiv \frac{\lambda_2}{4}(S_2^* S_2)^2 + \lambda_{12}(S_1^* S_1)(S_2^* S_2) - \left(\frac{2}{3}\lambda'_{12} S_1^* S_2^3 + \text{H.c.}\right), \tag{B5}$$

and requiring $V_{12} > 0$. This yields

$$\lambda_{12} > 0, \qquad \lambda_2 > 0, \qquad |\lambda'_{12}|^2 < \frac{9}{16}\lambda_{12}\lambda_2. \tag{B6}$$

We still have to deal with seven quartic couplings. First note that λ_1, λ_{2H}, and $\lambda_{2\Delta}$ need to be positive, as there is no other quartic left that can compensate for a negative contribution to the potential sourced by these couplings. The remaining parameters, λ_Δ, λ'_Δ, $\lambda_{H\Delta}$ and $\lambda'_{H\Delta}$, essentially define a usual type II seesaw potential and the stability conditions for that case are known [46]. The requirements for these seven quartics can be summarized as

$$\text{(i)}\ \lambda_H > 0, \quad \lambda_1 > 0, \quad \lambda_{2H} > 0, \quad \lambda_{2\Delta} > 0, \tag{B7a}$$

$$\text{(ii)}\ \lambda_\Delta + \lambda'_\Delta > 0, \qquad 2\lambda_\Delta + \lambda'_\Delta > 0, \tag{B7b}$$

$$\text{(iii)}\ 2\lambda'_{H\Delta} + \sqrt{\lambda_H(\lambda_\Delta + \lambda'_\Delta)} > 0, \tag{B7c}$$

$$\text{(iv)}\ 2\lambda'_{H\Delta}\sqrt{\lambda_\Delta + \lambda'_\Delta} + (2\lambda_\Delta + \lambda'_\Delta)\sqrt{\lambda_H} > 0. \tag{B7d}$$

We emphasize that if inequalities (B4), (B6), and (B7) are all satisfied, then the potential is stable.

The generalization to more n-steps is now straightforward. By defining

$$V_{i,i+1} \equiv \frac{\lambda_{i+1}}{4}(S_{i+1}^* S_{i+1})^2 + \lambda_{i,i+1}(S_i^* S_i)(S_{i+1}^* S_{i+1}) - \left(\frac{2}{3}\lambda'_{i,i+1} S_i^* S_{i+1}^3 + \text{H.c.}\right), \tag{B8}$$

and requiring $V_{i,i+1} > 0$ we obtain

$$\lambda_{i,i+1} > 0, \quad \lambda_{i+1} > 0, \quad |\lambda'_{i,i+1}|^2 < \frac{9}{16}\lambda_{i,i+1}\lambda_{i+1} \tag{B9}$$

for $i = 1 \ldots n-1$. Again, there are no quartic couplings left to compensate for λ_{iH} or $\lambda_{i\Delta}$, which demands

$$\lambda_i > 0, \qquad \lambda_{iH} > 0, \qquad \lambda_{i\Delta} > 0, \qquad i = 1 \ldots n. \tag{B10}$$

These conditions are by no means necessary, but only sufficient for having stability in the n-step realization. More general conditions may be obtained with the techniques of Ref. [47].

APPENDIX C: PARTIAL WIDTHS

We present in this Appendix the partial widths for the novel decay channels of some of the extra scalars in the generalized type II seesaw framework. In the case of δ, we will have three new channels: $\delta \to hs_1$, $\delta \to hhs_1$, and $\delta \to hs_1s_1$. As the latter is suppressed by v_1^2, we will safely neglect it in the remainder. The partial widths for the first two channels are

$$\Gamma(\delta \to hs_1) \simeq \frac{v^2}{1024\pi M_\Delta}(8\lambda_A\cos(2\theta_{\delta 1}) - \lambda_{1H}\sin(2\theta_{\delta 1}))^2,$$

$$\Gamma(\delta \to hhs_1) \simeq \frac{M_\Delta}{8192\pi^3}(2\lambda_A\cos(2\theta_{\delta 1}) - \lambda_{1H}\sin(2\theta_{\delta 1}))^2,$$

where we have neglected the phase space factor by assuming $M_1 + 2M_h \ll M_\Delta$. The phase space for 2-body decay can easily be incorporated by multiplying the partial width by

$$\bar{\beta}_{\delta \to hs_1} \equiv \sqrt{1 - \frac{2(M_1^2 + M_h^2)}{M_\Delta^2} + \frac{(M_1^2 - M_h^2)^2}{M_\Delta^4}}. \tag{C1}$$

The decay width ratios with respect to the leptonic channel, $\delta \to \nu\nu + \bar{\nu}\bar{\nu}$, are approximately given by

$$\frac{\Gamma[\delta \to h s_1]}{\Gamma[\delta \to \nu\nu + \bar{\nu}\bar{\nu}]} \simeq \lambda_A^2 \frac{v_\Delta^2}{\sum_i m_{\nu_i}^2} \frac{v^2}{M_\delta^2},$$

$$\frac{\Gamma[\delta \to hh s_1]}{\Gamma[\delta \to \nu\nu + \bar{\nu}\bar{\nu}]} \simeq \frac{\lambda_A^2}{512\pi^2} \frac{v_\Delta^2}{\sum_i m_{\nu_i}^2}.$$

In the case of the single-charged scalar δ^+, the additional channel $\delta^+ \to W^+ s_1$ is the most relevant. Its decay width is given by

$$\Gamma[\delta^+ \to W^+ s_1] = \cos^2\eta \frac{\sin^2(\theta_{\delta 1})}{8\pi} \frac{M_{\delta^+}^3}{v^2} \bar{\beta}_{\delta^+}^3,$$

with

$$\cos^2\eta \equiv 1 - \frac{2 v_\Delta^2}{v^2},$$

$$\bar{\beta}_{\delta^+} \equiv \sqrt{1 - \frac{2(M_1^2 + M_W^2)}{M_{\delta^+}^2} + \frac{(M_1^2 - M_W^2)^2}{M_{\delta^+}^4}}.$$

The ratio with the leptonic channel is approximately

$$\frac{\Gamma[\delta^+ \to W^+ s_1]}{\Gamma[\delta^+ \to \ell^+ \nu]} \simeq 2 \sin^2(\theta_{\delta 1}) \frac{v_\Delta^2}{m_{\nu_i}^2} \frac{M_{\delta^+}^2}{v^2}.$$

For s_1, we have the decay into charged fermions and neutrinos

$$\Gamma[s_1 \to f\bar{f}] = N_c \frac{M_1}{8\pi} \frac{m_f^2}{v^2} \sin^2\theta_{h1} \bar{\beta}_1, \quad \text{(C2a)}$$

$$\Gamma[s_1 \to \nu\nu] = \frac{M_1}{16\pi} \frac{m_\nu^2}{v_\Delta^2} \sin^2\theta_{\delta 1}, \quad \text{(C2b)}$$

where N_c is the number of colors and $\bar{\beta}_1 \equiv (1 - 4 m_f^2 / M_1^2)^{3/2}$. We do not present the analytic expressions for the new 3-body decay channel $\delta^{++} \to W^+ W^+ s_1$, as it is not particularly illuminating.

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