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Scotogenic RνMDM at three-loop level

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A B S T R A C T

We propose a model in which the radiative neutrino (Rν) masses are induced by fermion quintuplet and scalar septuplet fields from the minimal-dark-matter (MDM) setup. In conjunction with the 2HDM fields, on top of which our model is built, these hypercharge zero fields and additional scalar quintuplet lead to an accidental DM-protecting Z_2 symmetry and establish the RνMDM model at the three-loop level. We assess the potential for discovery of quintuplet fermions on present and future pp colliders.

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1. Introduction

After the discovery of $\sim$125 GeV particle [1,2] looking very much like the one Higgs boson of the Standard Model (SM), searching for the dark matter (DM) is one of the main targets for the next run of the LHC. At the same time, the scotogenic models of radiative neutrino mass relate the experimental evidence of neutrino masses and the existence of the DM in the Universe [3]. The scotogenic one-loop model proposed by Ma [4] augments the SM particle content by three singlet Majorana fermions and a second scalar doublet, and remains the simplest scotogenic realization. It imposes an exactly conserved Z_2 symmetry to both eliminate the unwanted Yukawa interactions responsible for the tree-level neutrino masses and to simultaneously stabilize the DM candidates. Here we take under scrutiny the notions of (i) *loop-generated* neutrino masses and (ii) Z_2 *symmetry*, which have since then become a common theme among many studies.

A three-loop neutrino masses under consideration bear an unquestionable appeal of naturally explaining the twelve orders of magnitude hierarchy between neutrino masses and the electroweak scale. The other theme we address here is how to avoid a common use of an *ad hoc* Z_2 symmetry. In a recent proposal [5] to promote Z_2 to a local gauge $U(1)_D$ symmetry in a model with two dark scalar doublets transforming as ±1 under $U(1)_D$, a breaking of gauge symmetry provides the dynamical origin of an exact Z_2 symmetry. Alternatively, the DM protecting Z_2 symmetry may arise "accidentally", on account of the SM symmetry and a choice

of the field content.[1] A realization which is in our focus here is enabled by employing higher weak multiplets, studied within the minimal dark matter (MDM) model [8]. There, an isolated fermion quintuplet $\Sigma \sim (5, 0)$ or an isolated scalar septuplet $\chi \sim (7, 0)$ have been selected to provide a viable DM candidate. Recent minimalistic variant of MDM model, with wino-like fermion triplet [9], relies on the enforcement of accidental B–L symmetry for DM stability.

In the present study we reconsider original MDM multiplets in such a way that both fermion quintuplet and scalar septuplet have to be employed to produce the neutrino masses at three-loop level. In the original model of radiative neutrino masses with MDM (RνMDM [10]) aiming at an automatic Z_2 symmetry by employing the higher multiplets, the zero hypercharge quintuplet $\Sigma \sim (5, 0)$ has been accompanied by a hypercharge-one sextuplet $\Phi \sim (6, 1)$ to close a one-loop neutrino mass diagram. However, as observed in [11], there is an additional renormalizable quartic term in this model which violates the Z_2 symmetry and threatens the stability of the proposed DM candidate.

In the present account we attempt to restore the RνMDM idea in a three-loop variant which employs both the fermion quintuplet $\Sigma \sim (5, 0)$ and the scalar septuplet $\chi \sim (7, 0)$. This is in contrast to an early three-loop model proposed by Krauss, Nasri and Trodden (KNT) [12] followed by similar recent attempts to explain small neutrino masses by adding to the SM content only additional weak singlets [13], or by substituting a real scalar triplet

---

[1] In a recent scotogenic variant realized by a real triplet scalar field [6] there is no need to eliminate the tree-level contribution, but additional Z_2 symmetry is needed to stabilize a DM candidate. Also, an antisymmetric tensor field may be stable [7] without introducing a new Z_2 symmetry.
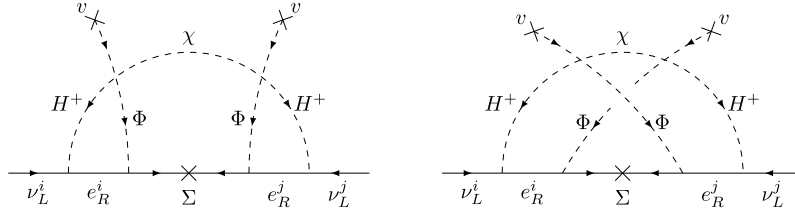
Fig. 1. The diagrams for generating tiny neutrino masses.

for a charged scalar singlet in the KNT model [14]. The KNT model has been partly self-criticized [15] because of employing unobservable singlet DM. However, recent generalizations [16–18] of the KNT three-loop topology employ non-singlet multiplets which provide charged components as their tracers. Thereby, while Ref. [16] still imposes an exact Z_2 symmetry for a real fermion triplet DM, Ref. [17] considers a real fermion quintuplet in the context of softly-broken accidental Z_2 symmetry and classifies a "tower" of model possibilities realized in Refs. [12,16,17].

Our model generalizes these studies to different three-loop topology, proposed by Aoki, Kanemura and Seto (AKS) [19,20]. It involves a second Higgs doublet and provides another well motivated scenario for a study of the 2-Higgs Doublet Model (2HDM) recently reviewed in [21]. Our novel three-loop RνMDM neutrino mass model is presented in Section 2. In Section 3 we present some phenomenological signatures of the Majorana quintuplet and the beyond SM (BSM) scalars at the LHC. We summarize our results in the concluding section.

2. Three-loop RνMDM model

In the original one-loop RνMDM model [10] the fermion quintuplet $\Sigma \sim (5, 0)$ has been proposed to provide its neutral component Σ^0 as a DM candidate on account of an accidental Z_2 symmetry. However, a scalar sextuplet field Φ which closes a neutrino mass diagram in this model also generates a quartic term [11],

$$\lambda \Phi^* \Phi^* \Phi H^* + \text{H.c.} ,$$
$$\Phi^* \Phi^* \Phi H^* = \Phi^{*iabcd} \Phi^{*pqrst} \Phi_{abpqr} H^{*n} \epsilon_{in} \epsilon_{cs} \epsilon_{dt} , \qquad (1)$$

which breaks the proposed DM-protecting discrete symmetry.

Similar DM instability controlled by a single parameter has been studied by authors of Ref. [17]. In their model with a scalar quintuplet $\Phi \sim (5, -2)$ and a charged scalar singlet $S^+ \sim (1, 2)$ there is a single Z_2-violating term

$$\lambda S^+ \Phi^* \Phi \Phi + \text{H.c.} , \quad \Phi^* \Phi \Phi = \Phi^{*abcd} \Phi_{abkl} \Phi_{cdmn} \epsilon^{km} \epsilon^{ln} , \qquad (2)$$

leading to an instability of the neutral component of $\Sigma \sim (5, 0)$. However, in the limit $\lambda \to 0$ their model has an accidental Z_2 symmetry: $(\phi, \Sigma) \to (-\phi, -\Sigma)$.

In our modified scenario, where the SM Higgs in Eq. (1) or the charged singlet scalar in Eq. (2) are effectively replaced by nonminimal scalar septuplet, the corresponding quartic-interaction term is Z_2-even and thus harmless.

2.1. Structure and role of 2HD sector

The mass matrix M_{ij}^ν of active neutrinos is generated by three-loop diagrams in Fig. 1 which belong to the AKS topology. The outer loops of these diagrams are opened by a pair of charged scalars $H^\pm$ that originates from the two Higgs doublets. This enables the conversion of an active neutrino to the SM lepton singlet e_R^i, connecting in the next step the outer neutrino lines with the inner one-loop box formed by exotic non-singlet particles.

The two Higgs doublets $H_{1,2} \sim (2, 1)$ of a generic non-supersymmetric 2HDM, on which our model is built, can be written as

$$H_1 = \begin{pmatrix} G^+ \cos\beta - H^+ \sin\beta \\ \frac{1}{\sqrt{2}} (v_1 - h \sin\alpha + H \cos\alpha + i (G \cos\beta - A \sin\beta)) \end{pmatrix},$$
$$(3)$$

$$H_2 = \begin{pmatrix} G^+ \sin\beta + H^+ \cos\beta \\ \frac{1}{\sqrt{2}} (v_2 + h \cos\alpha + H \sin\alpha + i (G \sin\beta + A \cos\beta)) \end{pmatrix},$$
$$(4)$$

where their electroweak vacuum expectation values (VEVs) define $\tan\beta \equiv v_2/v_1$. Besides physical charged scalars $H^\pm$, there are three Goldstone bosons ($G, G^\pm$) and three physical neutral scalars: two CP-even states h and H with their mixing angle α, and a CP-odd neutral scalar A.

Let us note that the VEVs v_1 and v_2 (which are related to the SM VEV $v = 246$ GeV by $v^2 = v_1^2 + v_2^2$) originate from m_{11}^2 and m_{22}^2 terms through the minimization conditions of the most general CP-conserving 2HD potential

$$V(H_1, H_2) = m_{11}^2 H_1^\dagger H_1 + m_{22}^2 H_2^\dagger H_2 - [m_{12}^2 H_1^\dagger H_2 + \text{h.c.}]$$
$$+ \frac{1}{2}\lambda_1 (H_1^\dagger H_1)^2 + \frac{1}{2}\lambda_2 (H_2^\dagger H_2)^2$$
$$+ \lambda_3 (H_1^\dagger H_1)(H_2^\dagger H_2) + \lambda_4 (H_1^\dagger H_2)(H_2^\dagger H_1)$$
$$+ \left\{ \frac{1}{2}\lambda_5 (H_1^\dagger H_2)^2 \right.$$
$$\left. + \left[\lambda_6 (H_1^\dagger H_1) + \lambda_7 (H_2^\dagger H_2) \right] H_1^\dagger H_2 + \text{h.c.} \right\}. \qquad (5)$$

Here the quartic couplings λ_1 to λ_5 can be traded for the four physical Higgs boson masses as free input parameters and the mixing parameter $\sin(\beta - \alpha)$.

The Yukawa couplings of the fermions are *a priori* free parameters, but then they lead to flavour-changing neutral currents (FCNC) mediated by 2HD scalars at the tree level. Usually they are eliminated by imposing a discrete symmetry under which $(H_1, H_2) \to (H_1, -H_2)$. This symmetry imposed on the potential (5), henceforth denoted as $\tilde{Z}_2$, is exact as long as m_{12}^2, λ_6 and λ_7 vanish. A recent detailed study within the 2HD scenario [22] shows that the exact $\tilde{Z}_2$, in an absence of the soft breaking m_{12}^2 term, does not require intervention of new physics below $\sim$10 TeV scale. Note that at this scale the exotic states of our model already entered into the play.

Out of four different ways the Higgs doublets are conventionally assigned charges under a $\tilde{Z}_2$ symmetry [23], we adopt here the "lepton-specific" (Type X or Type IV) model implemented by AKS [19,20], corresponding to Table 1. In this model H_2 couples to all quarks whereas H_1 couples to all leptons and provides the $H^+ \nu_L e_R$ coupling enhanced by $\tan\beta$.

Table 1
Charge assignment under an automatic Z_2 and an imposed $\tilde{Z}_2$ symmetry in the model.

	Q_i	u_{iR}	d_{iR}	L_{iL}	e_{iR}	H_1	H_2	Φ	χ	Σ_α
Z_2 accidental	$+$	$+$	$+$	$+$	$+$	$+$	$+$	$-$	$-$	$-$
$\tilde{Z}_2$ exact, imposed	$+$	$-$	$-$	$+$	$+$	$+$	$-$	$+$	$-$	$+$

2.2. Exotic BSM multiplets

The model that we propose is based on the symmetry of the SM gauge group $SU(3)_C \times SU(2)_L \times U(1)_Y$, where in the following only the relevant electroweak part is explicated. In addition to usual SM fermions, we introduce three generations of exotic real fermions transforming as $\Sigma_\alpha \sim (5, 0)$, where $\alpha = 1, 2, 3$ labels generations. Also, in addition to the described scalar doublets there are two exotic scalars $\Phi \sim (5, -2)$ and $\chi \sim (7, 0)$. In Table 1 we summarize the assignment of the charges under an imposed $\tilde{Z}_2$ symmetry and by it guaranteed accidental exact Z_2 symmetry in our model. All exotic additional fields are totally symmetric tensors Σ_{abcd}, Φ_{abcd} and χ_{abcdef}, and the components of Σ_{abcd} read

$$\Sigma_{1111} = \Sigma_R^{++}$$
$$\Sigma_{1112} = \frac{1}{\sqrt{4}} \Sigma_R^+$$
$$\Sigma_{1122} = \frac{1}{\sqrt{6}} \Sigma_R^0$$
$$\Sigma_{1222} = \frac{1}{\sqrt{4}} (\Sigma_L^+)^c$$
$$\Sigma_{2222} = (\Sigma_L^{++})^c \,, \tag{6}$$

and form two charged Dirac fermions and one neutral Majorana fermion

$$\Sigma^{++} = \Sigma_R^{++} + \Sigma_R^{--C} \,, \quad \Sigma^+ = \Sigma_R^+ - \Sigma_R^{-C} \,, \quad \Sigma^0 = \Sigma_R^0 + \Sigma_R^{0C} \,. \tag{7}$$

The components of two exotic scalar fields read

$$
\begin{array}{ll}
\Phi_{1111} = \phi^+ & \chi_{111111} = \chi^{+++} \\
\Phi_{1112} = \frac{1}{\sqrt{4}} \phi^0 & \chi_{211111} = \frac{1}{\sqrt{6}} \chi^{++} \\
\Phi_{1122} = \frac{1}{\sqrt{6}} \phi^- & \chi_{221111} = \frac{1}{\sqrt{15}} \chi^+ \\
\Phi_{1222} = \frac{1}{\sqrt{4}} \phi^{--} & \chi_{222111} = \frac{1}{2\sqrt{5}} \chi^0 \\
\Phi_{2222} = \phi^{---} \,, & \chi_{222211} = \frac{1}{\sqrt{15}} \chi^- \\
& \chi_{222221} = \frac{1}{\sqrt{6}} \chi^{--} \\
& \chi_{222222} = \chi^{---} \,,
\end{array}
\tag{8}
$$

where we distinguish ϕ^- and $(\phi^+)^*$. The Yukawa interaction is given by

$$\mathcal{L}_Y = -y_{e_i} \bar{L}_{iL} H_1 e_{iR} - Y_{i\alpha} \overline{(e_{iR})^c} \Phi^* \Sigma_{\alpha R} + \text{h.c.} \,, \tag{9}$$

where only the Higgs doublet H_1 in the lepton-specific 2HD model couples to SM leptons. Therefore the SM lepton mass m_e, which in "minimal" SM corresponds to the Yukawa strength $y_{e_i}^{SM} = \sqrt{2} m_{e_i}/v$, in the 2HD context reads $\sqrt{2} m_{e_i} \tan\beta/v$. The Yukawa terms involving new fields read in components as

$$
\begin{aligned}
(\overline{\Sigma_R})^{klmn} \Phi_{klmn} (e_R)^c &= \phi^{---} \overline{\Sigma_R^{--}} (e_R)^c + \phi^{--} \overline{\Sigma_R^-} (e_R)^c \\
&+ \phi^- \overline{\Sigma_R^0} (e_R)^c + \phi^0 \overline{\Sigma_R^+} (e_R)^c + \phi^+ \overline{\Sigma_R^{++}} (e_R)^c.
\end{aligned}
\tag{10}
$$

The scalar potential contains the gauge invariant pieces

$$
\begin{aligned}
V(H_1, H_2, \Phi, \chi) = {} & V(H_1, H_2) + V(\Phi) + V(\chi) \\
& + V_m(H_1, H_2, \Phi) + V_m(H_1, H_2, \chi) \\
& + V_m(\Phi, \chi) + V_m(H_1, H_2, \Phi, \chi) \,,
\end{aligned}
\tag{11}
$$

where $V(H_1, H_2)$ is given in Eq. (5). The $\tilde{Z}_2$-symmetric mixing potential $V_m(H_1, H_2, \Phi, \chi)$ provides the quartic term

$$
\begin{aligned}
V_m(H_1, H_2, \Phi, \chi) &= \kappa H_1 H_2 \Phi \chi + \text{h.c.} \,, \\
H_1 H_2 \Phi \chi &= H_{1i} H_{2j} \Phi_{klmn} \chi_{abcdfg} \epsilon^{ng} \epsilon^{ia} \epsilon^{jb} \epsilon^{kc} \epsilon^{ld} \epsilon^{mf} \,,
\end{aligned}
\tag{12}
$$

which includes the couplings needed to close our three-loop mass diagrams:

$$
\begin{aligned}
H_1 H_2 \Phi \chi \supset {} & -\frac{1}{\sqrt{6}} \chi^{--} H_1^+ H_2^0 \phi^+ + \frac{2}{\sqrt{15}} \chi^- H_1^+ H_2^0 \phi^0 \\
& - \frac{\sqrt{3}}{\sqrt{10}} \chi^0 H_1^+ H_2^0 \phi^- + \frac{2}{\sqrt{15}} \chi^+ H_1^+ H_2^0 \phi^{--} \\
& - \frac{1}{\sqrt{6}} \chi^{++} H_1^+ H_2^0 \phi^{---} + (H_1 \leftrightarrow H_2) \,.
\end{aligned}
\tag{13}
$$

These couplings enable the decays of the exotic scalars, making the scalar septuplet unstable.

Let us stress that without an important interplay of the 2HD and exotic sectors in our model there would be an additional dimension-three Z_2-noninvariant operator

$$\mu \Phi \Phi^* \chi = \mu \Phi_{ijkl} \Phi^{*iabc} \chi_{abcpqr} \epsilon^{jp} \epsilon^{kq} \epsilon^{lr} \,. \tag{14}$$

This coupling, which could make the scalar septuplet DM candidate χ^0 unstable already at the tree level and which also threatens a stability of the fermion quintuplet, would be a single Z_2-noninvariant term[2] in our model. However, the operator in (14) is forbidden by the $\tilde{Z}_2$ symmetry enforced on the 2HD sector and mandatory for a whole model, with parities explicated in Table 1.

2.3. Three-loop-induced neutrino mass

The neutrino mass term

$$\mathcal{L}^{\text{eff}} = \overline{\nu_{Li}^c} M_{ij} \nu_{Lj} \tag{15}$$

is generated by the three-loop diagrams in Fig. 1. They include five diagrams corresponding to five pairs of (χ, Σ) fields propagating in the inner loop in Fig. 1,

$$
\begin{aligned}
& (\chi^{--}, \Sigma^{++}), \ (\chi^-, \Sigma^+), \ (\chi^0, \Sigma^0), \\
& (\chi^+, (\Sigma^+)^c), \ (\chi^{++}, (\Sigma^{++})^c) \,.
\end{aligned}
\tag{16}
$$

They correspond to five components of the Φ field in the loop (ϕ^+ to ϕ^{---}) building the quartic vertices which result from the following substitution in five terms listed in Eq. (13):

$$H_1^+ H_2^0 + H_2^+ H_1^0 \to v \cos 2\beta \, H^+ \,. \tag{17}$$

In conjunction with appropriate Yukawa couplings, these quartic couplings lead to radiatively generated lepton number breaking Majorana neutrino masses.

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[2] Another potentially dangerous trilinear term χ^3 is Bose-forbidden for a septuplet field χ.
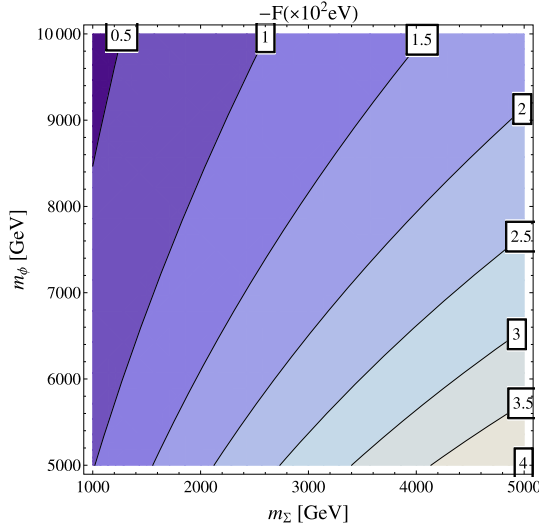
Fig. 2. The contour plot of the loop-integral function $-F(\times 10^2$ eV) that practically does not change for $M_{H^\pm}$ between 100 GeV to 1 TeV and $m_\chi \geq m_\Phi$ in 5 to 20 TeV range.

When we neglect the mass differences within Φ, χ and Σ_α multiplets, the mass matrix M_{ij}^ν for active neutrinos keeps the AKS form [20]

$$M_{ij} = \sum_{\alpha=1}^{3} C_{ij}^\alpha \, F(m_{H^\pm}, m_\Phi, m_\chi, m_{\Sigma_\alpha}), \qquad (18)$$

where the coefficient C_{ij}^α comprises the vertex coupling strengths

$$C_{ij}^\alpha = \frac{7}{3} \kappa^2 \tan^2 \beta \cos^2 2\beta \, y_{e_i}^{SM} Y_i^\alpha y_{e_j}^{SM} Y_j^\alpha, \qquad (19)$$

and the loop integral is represented by function F,

$$F(m_{H^\pm}, m_\Phi, m_\chi, m_{\Sigma_\alpha})$$
$$= \left(\frac{1}{16\pi^2}\right)^3 \frac{(-m_{\Sigma_\alpha})}{m_{\Sigma_\alpha}^2 - m_\chi^2} \frac{v^2}{m_{H^\pm}^4}$$
$$\times \int_0^\infty x dx \left\{ B_1(-x, m_{H^\pm}, m_\Phi) - B_1(-x, 0, m_\Phi) \right\}^2$$
$$\times \left(\frac{m_{\Sigma_\alpha}^2}{x + m_{\Sigma_\alpha}^2} - \frac{m_\chi^2}{x + m_\chi^2} \right),$$

where B_1 denotes the Passarino–Veltman function for one-loop integrals [24].

The magnitude of the integral function F in Fig. 2 is practically insensitive on the mass of the charged Higgs boson $H^\pm$ in the 100 GeV to 1 TeV range discussed in Sect. 3.3.2. Its magnitude is plotted as a function of m_Σ for the values of $m_\Sigma \leq m_\Phi$. It is also rather insensitive on the value of m_χ, which we take $m_\chi \geq m_\Phi$. The magnitude of F which is of order 10^2 eV in the wide range of the parameter space reproduces the neutrino masses with the coefficient $C_{ij}^\alpha \leq 10^{-4}$ that is easily achieved without fine tuning.

3. Phenomenology of BSM states

There are two major phenomenological issues generally related to the models for radiative neutrino masses: (i) an enlarged Higgs sector with charged scalar bosons; (ii) the TeV-scale right-handed
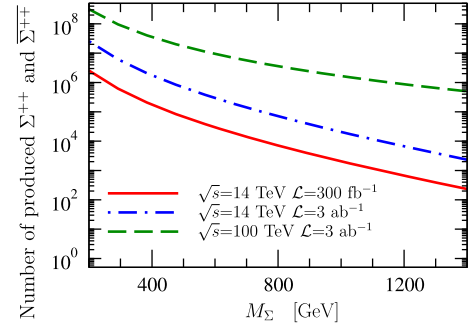
Fig. 3. Number of Σ^{++} and $\overline{\Sigma^{++}}$ particles produced for three characteristic collider setups, in dependence on the heavy lepton mass M_Σ.

neutrinos with Majorana masses, as possible DM candidates. In order to account for all of the observed DM abundance, the mass m_Σ of MDM Majorana quintuplet [8] approaches 10 TeV range. Accordingly, the Majorana quintuplet may be studied at the LHC only by relaxing this condition like in presence of non-thermal production mechanisms, in non-standard cosmological scenarios or if it accounts only for a fraction of the DM abundance.

3.1. DM candidate at the LHC

The present study of a Majorana quintuplet on colliders can be compared with a previous one [25] of the Majorana triplet employed in Type-III seesaw. Let us note that the quintuplet in presence of a half integer scalar quartet [26] $\Phi \sim (4, 1)$ leads to neutrino masses both at tree level[3] and one-loop level. It has been shown in [26] that it is possible to falsify this tree-level option at the LHC by mere non-observation of related light quintuplet states.

Here we are reassessing the LHC production of a quintuplet fermions studied previously [26] for the 2011/2012 run of the LHC. Now we consider it at a design 14 TeV energy both for the design luminosity of 300 fb^{-1} and for high luminosity (HL–LHC) of 3 ab^{-1}, as well as for a futuristic 100 TeV pp collider with 3 ab^{-1}.

Quintuplet lepton pairs in proton–proton collisions are mostly produced via Drell–Yan process, mediated by neutral and charged gauge bosons. Detailed expressions for cross-sections are given in [26]. In particular, for $M_\Sigma = 400$ GeV and 300 fb^{-1} of integrated luminosity at $\sqrt{s} = 14$ TeV, the LHC should produce about $1.8 \cdot 10^5$ doubly-charged Σ^{++} or $\overline{\Sigma^{++}}$ fermions, and $2.2 \cdot 10^5$ $\Sigma - \overline{\Sigma}$ pairs in total. In Fig. 3 we plot the expected number of produced Σ^{++} and $\overline{\Sigma^{++}}$ particles for three characteristic collider setups in dependence on the heavy lepton mass M_Σ. There would be roughly the same number of singly-charged Σ fermions produced, and half as many neutral ones.

The mass differences within the multiplet are of the order of few hundred MeV, so charged Σ leptons decay to neutral DM candidate by cascade radiation of soft undetected pions, e.g. $\Sigma^{++} \rightarrow \Sigma^+ \pi^+ \rightarrow \Sigma^0 \pi^+ \pi^+$.

Detailed study of collider signature of wino-like DM triplet in [9] identified disappearing tracks and monojets as most promising search channels. It turns out that regarding monojets our model is similar enough that we can, within some reasonable assumptions, directly use properly transformed results of [9]. In particular, we have monojets generated by same diagrams as in the triplet model, and, additionally, diagrams with doubly-charged leptons. Taking into account these additional contributions and different electroweak charges, in our model five times as many DM parti-

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[3] Therefore the model of Ref. [26] has been dubbed [27] the Type-V seesaw.
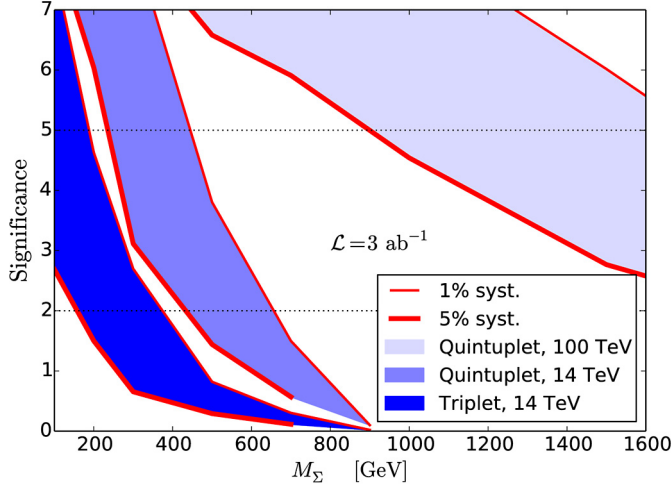
Fig. 4. Reach of monojet searches for quintuplet DM at the high-luminosity LHC and at a futuristic 100 TeV proton–proton collider. Bands correspond to variation of background systematic uncertainty from 5% to 1%. For comparison, corresponding reach of monojet searches for triplet model from [9] is also shown.

cles are produced as in the triplet model; either directly or via cascade decays mentioned above.

Since this enhancement factor turns out to be the same for various production mechanisms any kinematic cuts will cut away the same fraction of signal in both models,[4] and SM background should also be about the same. Thus we can estimate the LHC monojet search reach for our model by taking significance Z_{triplet} as defined in Eq. (3.1) of [9] and plotted in their Fig. 2, and scaling it using formula

$$Z_{\text{quintuplet}} = \frac{1}{\sqrt{\left(\frac{\sigma_{\text{triplet}}}{\sigma_{\text{quintuplet}}}\right)^2 \left(\frac{1}{Z_{\text{triplet}}^2} - \beta^2\right) + \beta^2}} , \qquad (20)$$
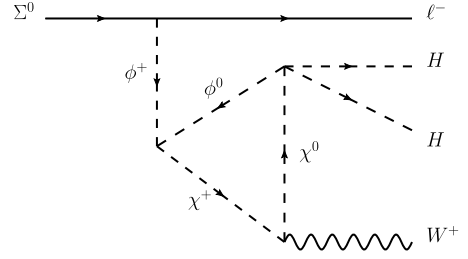
where σ is the total cross section for Drell–Yan production of Σ lepton pair with any charges, and β is the signal systematic uncertainty, which we should take to be 10%, same as in [9].

Resulting significance is plotted in Fig. 4. One notices that whereas the triplet model can only be *excluded* with 95% C.L. for $M < 350\,(150)$ GeV, the high-luminosity LHC can make a 5σ *discovery* of quintuplet model DM for $M < 450\,(250)$ GeV, with systematic uncertainty of the background of 1% (5%), respectively. Futuristic 100 TeV collider can easily extend the discovery reach beyond 1 TeV.

3.2. Majorana quintuplet in the broken Z_2 scenario

In the proposed model $\tilde{Z}_2$ symmetry, introduced in 2HDM for flavour physics reasons, implies also accidental Z_2 symmetry, thus protecting the stability of Σ^0 as a DM candidate. However, it is known that soft breaking of $\tilde{Z}_2$ in 2HDM may have some interesting features, like introducing a new source of CP violation and preserving stability and unitarity of the scalar potential. In the 2HD sector there is a potential soft m_{12} term (5), whereas the Z_2-breaking soft term (14) in the sector of exotic scalar fields

---

Fig. 5. An example of decay of the heavy lepton Σ at the loop level in the case of broken Z_2 symmetry.

would destabilize the Majorana quintuplet Σ^0 via loop-induced effects.

Let us for completeness estimate the value of the small, 'tHooft-natural coefficient carried by the term (14), so that the loop-induced decays of Σ^0 state may be long enough in comparison to the lifetime of the Universe.

For example, the loop diagram in Fig. 5 leads to the decay amplitude for $\Sigma^0 \to l_R^- W^+ H_1^0 H_2^0$. It can be compared to the decay amplitude for $\Sigma^0 \to \nu W^+ W^- H$ in Fig. 1 in [11], induced by dim-4 operator which destabilized DM candidate in one-loop RνMDM model [10]. In the present case the width for decay into the four particle final state $l_R^- W^+ H_1^0 H_2^0$ is given by

$$\Gamma \sim \frac{\kappa^2 g^2}{192\pi} \left(\frac{1}{16\pi^2}\right)^4 Y^2 \mu^2 \frac{m_\Sigma^9}{m_\Phi^8 m_W^2} . \qquad (21)$$

By assuming $\kappa = g = 0.65$ and by adopting the values $m_\Sigma = m_\Phi = 10$ TeV and $Y = 10^{-1}$, as used in [10], the lifetime of our DM candidate exceeds the age of the Universe $\approx 10^{17}$ s if the soft term coupling μ does not exceed the value of neutrino masses, $\mu < 0.1$ eV. A stronger bound $\mu < 10^{-9}$ eV can be obtained in the context of the decaying DM [28–30]. Notably, switching off the soft μ term does not affect the neutrino mass diagram but makes the Σ^0 state a viable DM candidate.

So, in order to confirm the relation of the quintuplet particles to neutrinos one has to study their decays. Let us stress that there are ample decays to purely SM final state particles in case that the quintuplets are not constrained by Z_2 symmetry, like in scenarios where they are accompanied by even-plet scalars [26,10].

Therefore a recent scrutiny [31] is welcome as a way to possibly exclude such modes and scenarios, what would be in favor of the present three-loop model in which fermion quintuplets are accompanied by odd-plet scalars. Then only less visible cascade decays of charged Σ states remain, and the search is focused on monojets or disappearing tracks like in case of wino-like particles [9]. In such scenario the Σ^0 is the DM candidate.

3.3. BSM scalars at colliders

The 2HD scalar sector has been studied in detail (for a review see [21] and [23]), independently of possible further extensions of the scalar sector. A study in the context of additional singlet scalars of AKS model faces challenging separation of signals addressed in [32]. The case of the larger multiplets of the present account is characterized by the fact that the 2HD and the exotic scalar sectors are decoupled due to the SM gauge symmetry, and thus can be treated separately.

3.3.1. Exotic quintuplet and septet scalars

Since in this work the scalar quintuplet and septuplet states are assumed to be heavier than Majorana quintuplet accounting for all DM abundance, we do not expect that they are accessible at the LHC. In particular, the collider signatures of these exotic scalars

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[4] Neutral current partonic production processes are enhanced by electric charge factors (thanks to the hypercharge being zero) like $Q(\Sigma^{++})^2 + Q(\Sigma^+)^2 = 2^2 + 1^2 = 5$, while the corresponding charged current partonic processes are enhanced by electroweak SU(2) ladder operators matrix elements $\sqrt{(T \mp T_3)(T \pm T_3 + 1)/2}$, giving $\sqrt{2}^2 + \sqrt{3}^2 = 5$ again.

are out of scope of the present paper. Still, their multiple-charged components are studied in number of papers for different reason.

As already pointed out, the sectors of small and large scalar multiplets are to large extent decoupled and can be treated separately. Here we are explicating the Z_2 symmetry breaking terms connecting small and large scalar multiplets, which are allowed by the exact $\tilde{Z}_2$ symmetry. A doublet-quintuplet mixing is represented by a dim 5 operator,

$$\frac{1}{\Lambda}\Phi\Phi^*\Phi H_{\mathbf{1}}H_{\mathbf{1}}, \quad \frac{1}{\Lambda}\Phi\Phi^*\Phi H_{\mathbf{2}}H_{\mathbf{2}}. \tag{22}$$

However, in order to enable similar doublet-septuplet mixing, we must go to a suppressed dim 7 operator

$$\frac{1}{\Lambda^3}\chi(H_{\mathbf{1}}H_{\mathbf{2}}^*)^3. \tag{23}$$

Since the natural mass scale of these fields is in a multi-TeV range, a study of 2HD states may be more promising.

3.3.2. Charged Higgs in lepton-specific 2HDM

The charged Higgs boson phenomenology for lepton-specific 2HDM at LEP was presented in [33], and for direct searches at LHC and ILC in [34]. The standard procedure [35–37] is to adjust the scenario for the neutral h state to be the SM-like Higgs boson with mass 125 GeV. In the first case corresponding to the *decoupling regime* where $\sin(\beta-\alpha) \simeq 1$, the $H^\pm$ mass $M \gg v$ may be related to large soft breaking scale of the discrete symmetry $\tilde{Z}_2$. In another case where the $H^\pm$ mass is at the scale $M \sim v$, these charged states may be probed through the Higgs diphoton decay.

Let us note that there is another three-loop model [38] which has a lepton-specific 2HDM as a part of the scalar sector extension. Therefore the Higgs phenomenology presented there applies to the charged states $H^\pm$. On the other hand for a charged Higgs heavier than top, the BaBar data based on $b \to s\gamma$ exclude a charged Higgs lighter than 380 GeV, independently of $\tan\beta$ (see Ref. [39] and references therein). In addition, in Ref. [40] the charged Higgs decay channels have been identified, which are promising for non-supersymmetric 2HD models.

4. Conclusions

We propose a model addressing in a common framework the open questions of the neutrino masses and the dark matter of the Universe. We follow the idea of the RνMDM proposed at one-loop level in [10] and look at alternative particle content that may actually realize it at three-loop level. We adopt a three-loop topology proposed by AKS [19,20] in a way that their imposed Z_2 symmetry appears accidentally due to higher representations added to the SM particle content. Thereby we keep a Z_2-even second Higgs doublet field, while the Z_2-odd weak singlet fields S, η and N_R of AKS are replaced by our higher multiplets Φ, χ and Σ, respectively.

We can compare the present study to the recent [41–43] and in particular to [16,17] which are based on modifications of KNT topology [12] and include also higher multiplets. A common feature of the models in [16,17] are super-renormalizable terms violating a desired Z_2 symmetry. In our case such unique term (14) is forbidden by exact $\tilde{Z}_2$ symmetry originating in the 2HD sector in our setup. The $\tilde{Z}_2$ charges have to be attributed also to exotic sector states as displayed in Table 1, in order to enable the neutrino mass loop diagram. The essential couplings in Eq. (13) make a neutral component of the scalar septuplet unstable so that the Majorana quintuplet $\Sigma \sim (5, 0)$ remains a single DM candidate.

There are cosmological scenarios in which this DM candidate can be within the reach of the LHC, where we take its mass as a free parameter. The studies of testability of the exotic quintuplet states at hand that are testable at the LHC can be compared to the LHC phenomenology of wino-like fermion triplet [9] belonging to recent minimalistic variant of MDM model. Their phenomenology reduces to a restricted subset of decays which are allowed by a Z_2 symmetry. Mere observation of a decays into purely SM final state particles would falsify the scenario with exact Z_2 symmetry. In the absence of such decays the proposed set of exotic BSM particles keeps its appealing features of contributing to the DM and generating three-loop suppressed neutrino masses. Using monojet searches, our DM candidate with mass $M < 450$ GeV is discoverable at the high-luminosity LHC.

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