

Right handed neutrinos, TeV scale BSM neutral Higgs boson, and FIMP dark matter in an EFT framework

Geneviève Bélanger^{1,*} Sarif Khan,^{2,†} Rojalin Padhan^{3,4,‡} Manimala Mitra,^{3,4,§} and Sujay Shil^{3,4,||}

¹*Université Grenoble Alpes, USMB, CNRS, LAPTh, F-74000 Annecy, France*

²*Institut für Theoretische Physik, Georg-August-Universität Göttingen, Friedrich-Hund-Platz 1, Göttingen D-37077, Germany*

³*Institute of Physics, Sachivalaya Marg, Bhubaneswar, Odisha 751005, India*

⁴*Homi Bhabha National Institute, BARC Training School Complex, Anushakti Nagar, Mumbai 400094, India*



(Received 25 April 2021; accepted 25 August 2021; published 30 September 2021)

We consider an effective field theory framework with three Standard Model (SM) gauge singlet right-handed neutrinos, and an additional SM gauge singlet scalar field. The framework successfully generates eV masses of light neutrinos via the seesaw mechanism, and accommodates a feebly interacting massive particle (FIMP) as a dark matter candidate. Two of the gauge singlet neutrinos participate in neutrino mass generation, while the third gauge singlet neutrino is a FIMP dark matter. We explore the correlation between the v_{ev} (vacuum expectation value) of the gauge singlet scalar field which translates as mass of the BSM Higgs, and the mass of dark matter, which arises due to the relic density constraint. We furthermore explore the constraints from the light neutrino masses in this setup. We chose the gauge singlet BSM Higgs in this framework in the TeV scale. We perform a detailed collider analysis to explore the discovery prospect of the TeV scale BSM Higgs through its difatjet signature, at a future pp collider which can operate with $\sqrt{s} = 100$ TeV c.m. energy.

DOI: [10.1103/PhysRevD.104.055047](https://doi.org/10.1103/PhysRevD.104.055047)

I. INTRODUCTION

The Standard Model (SM) of particle physics, despite its accurate predictions, suffers from a few serious deficits. Two of the most serious drawbacks emerge from the observation of light neutrino masses and their mixings, and the precise measurement of dark matter (DM) relic abundance in the Universe. A number of neutrino oscillation experiments have confirmed that the solar and atmospheric neutrino mass splittings are $\Delta m_{21}^2 \sim 10^{-5} \text{ eV}^2$, $|\Delta m_{13}^2| \sim 10^{-3} \text{ eV}^2$, the Pontecorvo-Maki-Nakagawa-Sakata (PMNS) mixing angles [1,2] are $\theta_{12} \sim 33^\circ$, $\theta_{23} \sim 49^\circ$, and $\theta_{13} \sim 8^\circ$ [3]. The light neutrinos being electromagnetic charge neutral can be Majorana particles. One of the profound mechanisms to generate Majorana masses of light neutrinos is the seesaw mechanism, where tiny eV masses of the SM neutrinos

are generated from a lepton number violating (LNV) $d = 5$ operator [4,5] through electroweak symmetry breaking. Among the different UV completed theories that generate this operator, type-I seesaw [6–9] is possibly the most economic one, where particle contents of the SM are extended to include gauge singlet right-handed neutrinos (RHNS).

Different models have been postulated, where the gauge singlet fermion state can act as a DM candidate [10–12]. A number of proposed models accommodate DM as a weakly interacting massive particle (WIMP), which is in thermal equilibrium with the rest of the plasma. However the null results from various direct detection experiments cause serious tension for the WIMP paradigm, and therefore motivate us to explore alternate DM hypothesis. One of such well-motivated mechanisms is the freeze-in [13,14] production of DM. In this scenario, the DM has feeble interactions with the bath particles, and hence is referred to as a feebly interacting massive particle (FIMP). The suppressed interaction naturally explains the nonobservation of any direct detection signal. Moreover, because of the very suppressed interaction, the FIMP never attains thermal equilibrium with the SM bath. In this scenario, DM is produced from the decay and/or annihilation of the SM and beyond the Standard Model (BSM) particles which are in thermal equilibrium [13]. FIMP DM has been explored in

*belanger@lapth.cnrs.fr

†sarif.khan@uni-goettingen.de

‡rojalin.p@iopb.res.in

§manimala@iopb.res.in

||sujayshil1@gmail.com

Published by the American Physical Society under the terms of the [Creative Commons Attribution 4.0 International license](https://creativecommons.org/licenses/by/4.0/). Further distribution of this work must maintain attribution to the author(s) and the published article's title, journal citation, and DOI. Funded by SCOAP³.

different contexts, see [15–19] for FIMP DM in the $B - L$ model, [20–24] for EFT descriptions, and [25] for discussion on all nonrenormalizable operators up to dimension-eight. Specific LHC signatures of FIMP DM have been investigated in [26–29] and others [30,31].

In this work we propose an effective field theory setup which includes a FIMP DM and explains the origin of light neutrino masses. The framework contains, in addition to SM particles, three RHN states N and one BSM scalar field χ . Two of the RHN states participate in the seesaw mechanism while the third RHN is the FIMP DM. In our model, due to a discrete symmetry, the DM is completely stable. One of the specificities of our model is that the usual renormalizable Dirac mass term for light neutrino mass generation is absent and is only generated via an effective $d = 5$ operator $L\Phi N\chi/\Lambda$. Due to other sets of $d = 5$ operators involving DM and scalars ($NN\chi^2/\Lambda$, $NN\Phi\Phi/\Lambda$), the DM is mainly produced from the decay of scalars. Annihilation processes involving scalars or SM gauge bosons and fermions can also significantly contribute to DM production. The relative importance of decay and annihilation processes for DM production strongly depends on the assumption on the reheating temperature of the early Universe. We consider three different scenarios, *Scenario-I-III*; for the first two only $d = 5$ operators are responsible for both DM production in the early Universe, and generation of its mass, while in *Scenario-III* we add a bare mass term for the RHN DM and the other two RHN states. *Scenario-I* is a subset of *Scenario-II* where some of the operators are neglected for simplicity.

We find that for the first two cases, a strong correlation exists between the vev of χ and mass of DM, that emerges from the relic density constraint, while for the latter the correlation is somewhat relaxed. Demanding a TeV scale vev of χ and a TeV scale heavy Higgs H_2 which offers a better discovery prospect of this model at the collider, a lighter KeV scale DM is in agreement with the relic density constraint for *Scenario-I-II*. For *Scenario-III* we find that a much heavier DM with GeV scale mass is also consistent with a TeV scale vev of χ , and in turn a TeV scale or lighter BSM Higgs. We furthermore study the impact of the eV scale light neutrino mass constraint for these different scenarios. Using micrOMEGAs5.0 [32], we perform a scan of all the relevant parameters such as, vev of χ , mass of DM, reheating temperature, and show the variation of relic density.

Finally, we explore the collider signature of the BSM Higgs with TeV scale mass which actively participates in DM production. For this, we consider a future pp collider that can operate with c.m. energy $\sqrt{s} = 100$ TeV. We consider the decay of the BSM Higgs into two SM Higgs, followed by subsequent decays of the SM Higgs into $b\bar{b}$ states. For the TeV scale BSM Higgs, the produced SM Higgs is highly boosted, thereby giving rise to collimated decay products. We therefore study difatjet final state as our

model signature. We consider a number of possible SM backgrounds including QCD, WW/ZZ , $W + j$, $Z + j$, and $t\bar{t}$ which can mimic the signal. By judiciously applying selection cuts, we evaluate the discovery prospect of the BSM Higgs. We find that a 3σ significance can be achieved for a 1.1 TeV BSM scalar with 30 ab^{-1} luminosity for a large SM and BSM Higgs mixing angle.

The paper is organized as follows. In Sec. II we describe the model and discuss associated DM phenomenology assuming three different scenarios *Scenario I-III*, where DM is produced from the decay of the SM and BSM Higgs. In Sec. III we discuss the contributions from both the decay and annihilation processes to the relic abundance and show the variation of DM relic density with respect to various parameters such as the mass of DM, the vev of the scalar field, and the reheating temperature. We perform the collider analysis of the BSM Higgs in the difatjet channel in Sec. IV. Finally, we conclude and summarize our findings in Sec. V.

II. THE MODEL

We consider an effective field theory framework with RHNs and one BSM scalar field, χ , where we consider operators up to mass dimension $d = 5$. In addition to the SM particles, the model therefore contains three SM gauge singlet RHNs denoted as $N_{1,2,3}$, and one SM gauge singlet real scalar field χ . The two RHNs $N_{1,2}$ generate eV Majorana masses of the SM neutrinos via the seesaw mechanism, while the state N_3 is a FIMP DM. The generic Yukawa Lagrangian with N , χ and the SM Higgs field Φ has the following form,

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & M_{Bij} N_i^T C^{-1} N_j + \tilde{Y}_{ij} \bar{L}_i \tilde{\Phi} N_j + \tilde{Z}_{ij} N_i^T C^{-1} N_j \chi \\ & + \frac{c_{ij}}{\Lambda} N_i^T C^{-1} N_j \chi^2 + \frac{c'_{ij}}{\Lambda} N_i^T C^{-1} N_j \Phi^\dagger \Phi \\ & + \frac{Y_{ij}}{\Lambda} \bar{L}_i \tilde{\Phi} N_j \chi + \text{H.c.}, \end{aligned} \quad (2.1)$$

where $\tilde{\Phi} = i\sigma_2 \Phi^*$ and M_B is the bare mass term of the RHNs. Other terms are the Yukawa interaction terms with couplings \tilde{Y}_{ij} , \tilde{Z}_{ij} , c_{ij} , Y_{ij} , and c'_{ij} , where $i, j = 1, 2, 3$ are the generation indices. The parameter Λ is the cutoff scale of this theory. In our subsequent discussions we do not consider \tilde{Y} , \tilde{Z} terms separately. These interaction terms can be obtained from $LN\Phi\chi$ and $NN\chi^2$ operators via the vev of χ . A successful realization of the fermion state N as a FIMP DM demands the coupling \tilde{Z} , \tilde{Y} to be very tiny. This can naturally be obtained, if these terms are generated from $NN\chi^2$ and $LN\Phi\chi$ operators, which feature the $\frac{1}{\Lambda}$ suppression factor. Additionally, we also note that by imposing a Z_2 symmetry under which $\chi \rightarrow -\chi$, $N_i \rightarrow -N_i$, and all other SM fields are invariant, the \tilde{Y} , \tilde{Z} terms can be completely prohibited. We impose such a symmetry, hence our Lagrangian is

$$\begin{aligned} \mathcal{L}_{\text{eff}} = & M_B N_i^T C^{-1} N_i + \frac{c_{ij}}{\Lambda} N_i^T C^{-1} N_j \chi^2 \\ & + \frac{c'_{ij}}{\Lambda} N_i^T C^{-1} N_j \Phi^\dagger \Phi + \frac{Y_{ij}}{\Lambda} \bar{L}_i \tilde{\Phi} N_j \chi + \text{H.c.}, \end{aligned} \quad (2.2)$$

which only contains $d = 5$ operators as interaction terms of RHNs. For simplicity we consider the Yukawa coupling matrices c , c' , and the bare mass matrix M_B to be diagonal. As advertised before, among the N_i states, N_3 is DM. Therefore, the Yukawa matrix is required to have the following structure

$$Y = \begin{pmatrix} Y_\nu^{11} & Y_\nu^{12} & \epsilon \\ Y_\nu^{21} & Y_\nu^{22} & \epsilon \\ Y_\nu^{31} & Y_\nu^{32} & \epsilon \end{pmatrix}. \quad (2.3)$$

In the above, we consider all Y_ν^{ij} ($i = 1, 2, 3$ and $j = 1, 2$) to be equal, while ϵ is required to satisfy the hierarchy $\epsilon \ll Y_\nu^{ij}$. The requirement of stability of DM over the age of the Universe forces the parameter ϵ to be orders of magnitude smaller than the other Yukawa couplings of the matrix Y . This scenario has been explored in a minimal type-I seesaw framework in [19]. Note that the DM state N_3 can be made completely stable by imposing an additional \mathcal{Z}_2 symmetry, in which N_3 has odd charge, and all other fields are evenly charged. This forbids the mixing between N_3 and light neutrino i.e., $\epsilon = 0$. In this study we furthermore consider such a \mathcal{Z}_2 symmetry thereby making the DM state N_3 completely stable. The interaction terms proportional to ϵ are hence absent in our case. In principle, the $d = 5$ Weinberg operator, $\frac{\kappa}{\Lambda} (\bar{L} \tilde{\Phi})(\tilde{\Phi}^T L^c)$, is allowed in our framework. This operator contributes to the light neutrino mass matrix. However, the Weinberg operator does not have any impact on the DM phenomenology that we discuss in the following. Thus we set the coupling κ to be zero such that its contribution to the light neutrino mass will also be forbidden.

Scalar potential—As stated above, the model also contains a gauge singlet scalar field χ . In addition to the Yukawa Lagrangian, the scalar field χ also interacts with the SM Higgs doublet field Φ via the scalar potential,

$$\begin{aligned} V(\chi, \Phi) = & M_\Phi^2 \Phi^\dagger \Phi + m_\chi^2 \chi^2 + \lambda_1 (\Phi^\dagger \Phi)^2 \\ & + \lambda_2 \chi^4 + \lambda_3 (\Phi^\dagger \Phi) \chi^2. \end{aligned} \quad (2.4)$$

The $d = 5$ terms $\frac{1}{\Lambda} (\Phi^\dagger \Phi)^2 \chi$, and $\frac{1}{\Lambda} \chi^5$, as well as $d = 3$ term $\Phi^\dagger \Phi \chi$ are disallowed by the above mentioned \mathcal{Z}_2 symmetry. Therefore the scalar potential contains only renormalizable terms up to $d = 5$. The spontaneous symmetry breaking (SSB) in this model is similar to the SM extension with an additional singlet scalar, which has been widely discussed in the literature [33,34]. In order for the potential to be bounded from below, the couplings $\lambda_{1,2,3}$ should satisfy,

$$\begin{aligned} 4\lambda_1 \lambda_2 - \lambda_3^2 &> 0, \\ \lambda_{1,2} &> 0. \end{aligned} \quad (2.5)$$

We denote the vevs of Φ and χ by v_Φ and v_χ , respectively. After minimizing the potential $V(\chi, \Phi)$, with respect to both the vevs, we obtain

$$v_\Phi^2 = \frac{4\lambda_2 M_\Phi^2 - 2\lambda_3 m_\chi^2}{\lambda_3^2 - 4\lambda_1 \lambda_2}, \quad (2.6)$$

$$v_\chi^2 = \frac{4\lambda_1 m_\chi^2 - 2\lambda_3 M_\Phi^2}{\lambda_3^2 - 4\lambda_1 \lambda_2}. \quad (2.7)$$

The λ_3 -term in the potential enables mixing between χ and Φ states. We denote the neutral Higgs component in the Φ multiplet as H . The mass matrix between the two Higgs bosons in the basis (H, χ) is given by

$$\mathcal{M}(H, \chi) = 2 \begin{pmatrix} \lambda_1 v_\Phi^2 & \lambda_3 v_\Phi v_\chi / 2 \\ \lambda_3 v_\Phi v_\chi / 2 & \lambda_2 v_\chi^2 \end{pmatrix}. \quad (2.8)$$

The mass eigenstates (H_1, H_2) are related to the (H, χ) states as

$$\begin{pmatrix} H_1 \\ H_2 \end{pmatrix} = \begin{pmatrix} \cos \theta & -\sin \theta \\ \sin \theta & \cos \theta \end{pmatrix} \begin{pmatrix} H \\ \chi \end{pmatrix}, \quad (2.9)$$

The mixing angle θ satisfies

$$\tan 2\theta = \frac{\lambda_3 v_\chi v_\Phi}{(\lambda_2 v_\chi^2 - \lambda_1 v_\Phi^2)}. \quad (2.10)$$

We denote the masses of the physical Higgs bosons as M_{H_1} and M_{H_2} ,

$$\begin{aligned} M_{H_1}^2 = & \lambda_1 v_\Phi^2 + \lambda_2 v_\chi^2 - \sqrt{(\lambda_1 v_\Phi^2 - \lambda_2 v_\chi^2)^2 + \lambda_3^2 v_\chi^2 v_\Phi^2}, \\ M_{H_2}^2 = & \lambda_1 v_\Phi^2 + \lambda_2 v_\chi^2 + \sqrt{(\lambda_1 v_\Phi^2 - \lambda_2 v_\chi^2)^2 + \lambda_3^2 v_\chi^2 v_\Phi^2}. \end{aligned} \quad (2.11)$$

Among the two Higgs states $m_{H_1}^2 < m_{H_2}^2$, i.e., H_1 acts as the lightest state. In our subsequent discussion, we consider that H_1 is SM-like Higgs with mass $M_{H_1} \sim 125$ GeV. The interactions of H_1 and H_2 with the fermions and gauge bosons are given in the Appendix. In this work, we consider that the BSM Higgs H_2 has a mass $M_{H_2} \sim \text{TeV}$, or lower, and has a substantial mixing with the SM-like Higgs state H_1 . This large mixing facilitates the production of the BSM Higgs at colliders, which will be discussed in Sec. IV.

FIMP dark matter—As discussed above, we consider that the RHN state N_3 is a FIMP DM. The state N_3 , being gauge singlet, only interacts via Yukawa interactions $N_3 N_3 \chi^2 / N_3 N_3 \Phi \Phi$. Therefore, the production of N_3 occurs primarily from the scalar states. In particular, the dominant contribution arises from the decay of the BSM Higgs for a low reheating temperature $T_R < 10^5$ GeV. A number of annihilation channels, involving the SM/BSM Higgs and gauge boson also contribute to the relic density. For high reheating temperatures, the gauge boson annihilation channels give dominant contributions, even larger than the decay contribution. The contributions from Higgs annihilation channels for a higher reheating temperature are also significantly large. In our discussion, we consider that the FIMP DM is lighter than the Higgs states $H_{1,2}$, such that, the decay of $H_{1,2}$ into N_3 state is open. The different channels that lead to the DM production are

- (a) Decay channels: the Higgs decay $H_{1,2} \rightarrow N_3 N_3$ generate the relic abundance.
- (b) Annihilation channels: the $2 \rightarrow 2$ annihilation channels, such as, $WW/ZZ \rightarrow N_3 N_3$, $H_1 H_1 \rightarrow N_3 N_3$, $H_2 H_2 \rightarrow N_3 N_3$, $H_1 H_2 \rightarrow N_3 N_3$ contribute to the production of N_3 . We also consider annihilation of other SM particles such as, b quark.

In the subsequent discussion, we consider three different scenarios *Scenario I-III*, where we only consider the decay contribution of the SM and BSM Higgs. As stated above, this can be justified for a lower reheating temperature, for which the annihilation processes give negligible contributions and DM production is primarily governed by the decay of H_1 and H_2 . Among the three scenarios, in *Scenario-I* and *II*, we consider that the bare mass terms of $N_{1,2,3}$ states are zero. In this simplistic scenario the $d = 5$ operator determines both the relic abundance, as well as DM mass, thereby leading to a tight correlation between the mass of DM and v_{ev} of χ . In *Scenario-III* we allow a nonzero bare mass term that significantly alters the phenomenology. We analyze the constraints from DM relic density, and neutrino mass generation. We discuss the annihilation contributions in Sec. III where we depart from the assumption of a low reheating temperature.

A. Scenario-I

The RHN states N_i interact with the scalar field χ , and the Higgs doublet Φ via the following Lagrangian

$$\mathcal{L}_{\text{eff}} = \frac{c_{ij}}{\Lambda} N_i^T C^{-1} N_j \chi^2 + \frac{Y_{ij}}{\Lambda} \bar{L}_i \Phi N_j \chi + \text{H.c.} \quad (2.12)$$

In the above, c_{ij} and Y_{ij} are the Yukawa couplings, and Λ is the cutoff scale of this theory. As discussed in the previous section, we choose to work with a basis in which the Yukawa coupling c_{ij} is diagonal. The above Lagrangian, after electroweak symmetry breaking, generates the

following bilinear terms involving the light neutrinos, and RHNs (ν, N),

$$\mathcal{L}_{\text{eff}} = \frac{c_{ii}}{\Lambda} N_i^T C^{-1} N_i v_\chi^2 + \frac{Y_{ij}}{\Lambda} \bar{\nu}_i N_j v_\Phi v_\chi + \text{H.c.} \quad (2.13)$$

The Λ suppressed $d = 5$ term $NN\chi^2$ in Eq. (2.12) gives a natural explanation of the small interaction strength of the FIMP DM N_3 with all other SM (and BSM) particles. As we consider the DM to be completely stable, therefore, only the $N_{1,2}$ states participate in light neutrino mass generation. Below, we analyze the contributions of $N_{1,2}$ in light neutrino mass, and the constraint from relic density.

1. Neutrino masses

The light neutrino masses will be generated due to the seesaw mechanism, where two RHN states $N_{1,2}$ participate. In our case, $Y_{i3} = 0$ and the Dirac mass matrix M_D effectively reduces to a matrix of dimension 3×2 . The Majorana mass matrix involving $N_{1,2}$ states is a 2×2 matrix. We denote the Dirac mass matrix by M_D and the Majorana mass matrix of N_1, N_2 states by M_R , where

$$(M_D)_{\gamma\alpha} = \frac{Y_{\gamma\alpha}}{\Lambda} v_\Phi v_\chi, \quad (M_R)_{\alpha\beta} = \frac{c_{\alpha\beta}}{\Lambda} v_\chi^2 \\ (\alpha, \beta = 1, 2, \gamma = 1, 2, 3). \quad (2.14)$$

In the basis $\psi = (\nu_l, N_{R_1}^c, N_{R_2}^c)^T$, the neutral lepton mass matrix becomes

$$M_\nu = \begin{pmatrix} 0 & M_D \\ M_D^T & M_R \end{pmatrix}. \quad (2.15)$$

The seesaw approximation $M_R > M_D$ translates into the hierarchy between the two v_{ev} s $v_\chi > v_\Phi$. The light neutrino and heavy Majorana mass matrix are given by

$$m_\nu = -M_D M_R^{-1} M_D^T, \quad M_N \sim M_R. \quad (2.16)$$

These can further be rewritten as

$$m_\nu \sim -\frac{v_\Phi^2}{\Lambda} Y C^{-1} Y^T, \quad M_N \sim \frac{c}{\Lambda} v_\chi^2. \quad (2.17)$$

In the above, $M_N = \text{diag}(M_{N_1}, M_{N_2})$ are the physical masses of the $N_{1,2}$ states, respectively. For simplicity we consider $M_{N_1} = M_{N_2}$ (i.e., $c_{11} = c_{22}$) in all of the scenarios, *Scenario-I-III*. In terms of light neutrino masses, M_{N_1}, M_{N_2} become

$$M_{N_{1,2}} \sim \frac{v_\Phi^2 v_\chi^2}{\Lambda^2} (Y Y^T m_\nu^{-1}) \text{ GeV}. \quad (2.18)$$

The active-sterile mixing matrix V is related with Dirac mass matrix M_D and Majorana mass matrix M_N as

$$V \sim M_D M_N^{-1} = \frac{v_\Phi}{v_\chi} Y c^{-1}. \quad (2.19)$$

Next we discuss the relic abundance of N_3 , where we consider the decays of $H_{1,2}$ as the primary production mode.

2. Dark matter phenomenology

In general, both decay and annihilation processes can produce N_3 . However, for the low reheating temperature that we consider in this section, the production of N_3 occurs primarily from the decay of the scalar states. In our discussion, we consider that the DM is lighter than the Higgs states $H_{1,2}$, such that, the decay of $H_{1,2}$ into N_3 state is kinematically open. For illustrative purposes, in this section we consider the mass of the BSM Higgs $M_{H_2} = 250$ GeV. We have verified that for larger TeV scale masses, such as $M_{H_2} = 1.1$ TeV, the result presented in this section remains very similar. A more extensive investigation of the dependence of the relic density on the mass of the BSM Higgs is deferred to Sec. III. The values of the BSM Higgs mass and mixing that we adopt to study DM production are consistent with the collider searches, which we will discuss in Sec. IV.

First note that the decay of the SM Higgs H_1 into N_3 state is non-negligible, only if the mixing between $H_{1,2}$ is sizeable. For the mixing $\theta \simeq 0$ the DM will be produced from the decay of $\chi \approx H_2$. From Eq. (2.12) and using Eq. (2.9), the interaction Lagrangian of N_3 with the SM and BSM Higgs reduces to

$$\mathcal{L}_{N_3} = \frac{c_{33} v_\chi}{\Lambda} \bar{N}_3^c N_3 (-H_1 \sin \theta + H_2 \cos \theta) + \text{H.c.} \quad (2.20)$$

We define $\tilde{\lambda}_{1,2}$ as the couplings of N_3 with the $H_{1,2}$ states, respectively¹

$$\tilde{\lambda}_1 = -\frac{c_{33} v_\chi}{\Lambda} \sin \theta, \quad \tilde{\lambda}_2 = \frac{c_{33} v_\chi}{\Lambda} \cos \theta. \quad (2.21)$$

Since the bare mass term is zero, the mass of the DM state N_3 in this case is generated from the $c_{33} N_3 N_3 \chi^2 / \Lambda$ term once the χ state acquires vev v_χ . The mass of N_3 is therefore given by

$$M_{N_3} = \frac{v_\chi^2 c_{33}}{\Lambda}, \quad (2.22)$$

and the couplings $\tilde{\lambda}_{1,2}$ can be expressed in terms of M_{N_3} ,

$$\tilde{\lambda}_1 = -\frac{2M_{N_3}}{v_\chi} \sin \theta, \quad \tilde{\lambda}_2 = \frac{2M_{N_3}}{v_\chi} \cos \theta. \quad (2.23)$$

Note that the interaction of H_2 with DM is governed by $\cos \theta$. The LHC Higgs signal strength measurements dictate $\sin \theta < 0.36$ i.e., $\cos \theta \simeq 1$ [35]. Therefore, for similar masses of $H_{1,2}$, the BSM Higgs state H_2 primarily governs N_3 production due to a higher coupling strength $\tilde{\lambda}_2 \gg \tilde{\lambda}_1$. However, for significantly heavier H_2 , H_1 contribution in the DM relic density will be larger than from H_2 .

As we are considering the decay contribution, the relic density of the FIMP DM can be expressed as [13]

$$\Omega_{N_3} h^2 = \frac{2.18 \times 10^{27}}{g_s \sqrt{g_\rho}} M_{N_3} \sum_{i=1}^2 \frac{g_{H_i} \Gamma_{H_i}}{M_{H_i}^2}. \quad (2.24)$$

In the above, g_{H_i} is the degrees of freedom of the decaying particle, $g_{s,\rho} \simeq 103.857$ are the degrees of freedom of the Universe related to entropy and matter. The partial decay widths Γ_{H_i} for $H_i \rightarrow N_3 N_3$ are

$$\Gamma_{H_1} = \frac{\tilde{\lambda}_1^2}{16\pi} M_{H_1}, \quad \Gamma_{H_2} = \frac{\tilde{\lambda}_2^2}{16\pi} M_{H_2}. \quad (2.25)$$

The measured relic abundance is $\Omega h^2 = 0.1199 \pm 0.0012$ at 68% C.L [36]. Using the above equation and equating Eq. (2.24) with the central value of the observed relic density, we obtain the constraints on the couplings as

$$\tilde{\lambda}_i = 1.66 \times 10^{-12} \sqrt{\frac{M_{H_i}}{M_{N_3}}}, \quad (2.26)$$

where we assume that the DM is entirely produced from either of the two Higgs states, $H_{1,2}$. From Eq. (2.23), the couplings $\tilde{\lambda}_{1,2}$ depend on the mass of the DM, the vev v_χ of the BSM Higgs, and the mixing angle between two Higgs states $H_{1,2}$. Hence, using Eqs. (2.23) and (2.25), in Eq. (2.24) and equating Eq. (2.24) with the observed relic density, we obtain a correlation between the vev of χ , DM mass, and other physical parameters of this model, which are the Higgs mixing angle and mass of the Higgs. Taking into account both the $H_{1,2} \rightarrow N_3 N_3$ production modes, we find that, for a FIMP DM with mass M_{N_3} , the required value of v_χ has to satisfy the following constraint,

$$v_\chi = 1.22 \times 10^{12} M_{N_3}^{3/2} \left(\frac{\sin^2 \theta}{M_{H_1}} + \frac{\cos^2 \theta}{M_{H_2}} \right)^{1/2} \\ \sim 10^{12} M_{N_3}^{3/2} \left(\frac{\theta^2}{M_{H_1}} + \frac{1}{M_{H_2}} \right)^{1/2}, \quad (2.27)$$

where in the right-hand side we assume small values of θ and took $\cos \theta \approx 1$, $\sin \theta \approx \theta$. The strong correlation between v_χ and M_{N_3} emerges, as both the DM mass and

¹We follow this notation for the rest of our discussion.

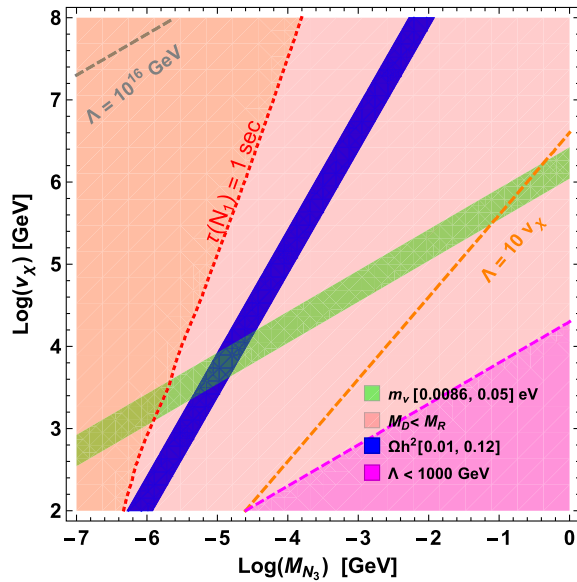


FIG. 1. This plot represents constraints on M_{N_3} and v_χ for *Scenario-I*. The blue band corresponds to the variation of the relic density from decay in between 0.01 and 0.12, where the latter satisfies the experimental constraint [36]. The green region corresponds to light neutrino mass in between 0.0086 eV to 0.05 eV. The red line corresponds to lifetime of $N_{1,2}$ as 1 second. The orange line correspond to $\Lambda > 10.0v_\chi$. Here we assume the mass of the BSM Higgs as $M_{H_2} = 250$ GeV.

its production are governed by the same $N_3 N_3 \chi^2$ operator in the Lagrangian. In Fig. 1 we show this correlation. Before presenting the discussion on Fig. 1, we note that,

(a) The contributions of $H_1 \rightarrow N_3 N_3$ and $H_2 \rightarrow N_3 N_3$ processes to the relic density are

$$(\Omega_{N_3} h^2)_{H_1} \sim \frac{\sin^2 \theta M_{N_3}^3}{v_\chi^2 M_{H_1}}, \quad (\Omega_{N_3} h^2)_{H_2} \sim \frac{\cos^2 \theta M_{N_3}^3}{v_\chi^2 M_{H_2}}. \quad (2.28)$$

Among these two, since the relic density from H_2 decay is proportional to $1/M_{H_2}$, therefore, for a very high mass of the BSM Higgs state, its contribution can be subleading. On the other hand, for a much smaller value of $\sin \theta$, the contribution from H_1 decay can also be subleading.

(b) The ratio between the two contributions is

$$\frac{(\Omega h^2)_{H_1}}{(\Omega h^2)_{H_2}} = \frac{M_{H_2}}{M_{H_1}} \left(\frac{\sin^2 \theta}{\cos^2 \theta} \right), \quad (2.29)$$

Therefore, both the contributions can be comparable if the Higgs and BSM Higgs mixing angle θ satisfies $\theta^2 \simeq 1.25 \times 10^2 / M_{H_2}$. Here we assume a small mixing angle, hence $\sin \theta \sim \theta$.

In Fig. 1, we show the constraint on v_χ and on the DM mass M_{N_3} that arises from requiring the relic density to lie in the range $0.01 < \Omega h^2 < 0.12$ (blue region), namely we allow for N_3 to account for 10%–100% of the DM abundance. Here we include both $H_1 \rightarrow N_3 N_3$ and $H_2 \rightarrow N_3 N_3$ contributions to the relic density even though the H_2 contribution dominates as we fixed $M_{H_2} = 250$ GeV. For this figure, we use the benchmark parameters given in Table I. Additionally, we also show the constraint from eV light neutrino mass in the same plot. For simplicity here and in other figures as well, we consider the light neutrino mass matrix m_ν (also M_D and M_R) as a parameter, and impose the neutrino mass constraint. Hence, y denotes the Dirac Yukawa coupling parameter in Table I. The green shaded region is compatible with eV light neutrino masses,² $0.0086 \text{ eV} < m_\nu < 0.05 \text{ eV}$, while the seesaw approximation $M_R > M_D$ is satisfied in the entire plot. We further note that, following Eq. (2.22), for higher v_χ and lower M_{N_3} the cutoff scale Λ increases. The brown dashed line in the top left corner denotes $\Lambda = 10^{16}$ GeV. We also show the line corresponding to $\Lambda = 10v_\chi$ by an orange dashed line. The magenta dashed line, assuming $\Lambda > 1$ TeV rules out the region with large M_{N_3} and low v_χ (magenta shaded region). We also check that in Fig. 1, the entire parameter space is consistent with the current experimental limits on lepton flavor violating (LFV) decays, such as $\mu \rightarrow e\gamma$, $\mu \rightarrow 3e$, and $\mu - e$ conversion [37,38]. In most of the parameter space the predicted values of the LFV decay rates are orders of magnitude smaller than the sensitivity reach of future experiments. This statement is valid for *Scenario-II*, and *III* as well.

The constraint from the relic density depends on the Yukawa coupling c_{33} , which has been rewritten in terms of M_{N_3} . However, the constraint from eV light neutrino masses depend on other parameters, such as, $M_{N_{1,2}}$, and hence the couplings c_{11} , c_{22} as well as the Dirac Yukawa y . As can be seen from the figure, to satisfy the observed DM relic density, the required value v_χ increases with DM mass M_{N_3} . For GeV scale M_{N_3} , one needs $v_\chi > 10^8$ GeV. This naturally leads to a very heavy BSM Higgs with mass³ $M_{H_2} \sim v_\chi > 10^8$ GeV for the quartic scalar coupling $\lambda_2 \sim 1$. This very heavy BSM Higgs does not have any detection prospect at colliders. Contrary to that, the coupling λ_2 needs to be extremely tiny ($\lambda_2 < 10^{-12}$) to accommodate $M_{H_2} \sim \mathcal{O}(100)$ GeV, which has better discovery prospects at the ongoing and future colliders. This

²Two of the RHN states $N_{1,2}$ will participate in neutrino mass generation. Hence, the lightest neutrino mass $m_1/m_3 = 0$ depends on normal/inverted mass hierarchy in the light neutrino sector. We therefore vary m_ν^2 in between solar and atmospheric mass-square splittings, where we consider $\Delta m_{21}^2 = 7.42 \times 10^{-5} \text{ eV}$ and $|\Delta m_{13}^2| = 2.517 \times 10^{-3} \text{ eV}$ [3].

³ M_{H_2} can not be much larger than v_χ due to the perturbativity bound of λ_2 .

TABLE I. Parameters relevant for *Scenario-I* (Fig. 1) and *Scenario-II* (Fig. 4).

	M_{H_2}	$\sin\theta$	y	c_{11} (c'_{11})	c_{33} (c'_{33})	$M_{N_{1,2}}$
<i>Scenario-I</i>	250 GeV	0.1	10^{-4}	1 (0)	2.5×10^{-6} (0)	$4 \times 10^5 M_{N_3}$
<i>Scenario-II</i>	250 GeV	0.1	10^{-4}	1 (1)	2.5×10^{-6} (2.5×10^{-6})	$4 \times 10^5 M_{N_3}$

unnatural fine tuning relaxes if the DM mass is $M_{N_3} \sim \text{KeV}$. As can be seen from the figure, DM N_3 with few KeV mass is consistent with a $v_\chi \sim \text{TeV}$ and $\lambda_2 \sim \mathcal{O}(0.1)$ and hence $M_{H_2} \sim \mathcal{O}(100)$ GeV. In conclusion, in *Scenario-I* we find that relic density constraint prefers a KeV scale DM and a TeV scale v_χ to naturally accommodate a BSM Higgs at the TeV scale or below.

3. N_1 and N_2 lifetime

Before concluding the section we also discuss the lifetime of $N_{1,2}$. For the range of the relevant parameters that we consider in Fig. 1, the mass of the RHN states $N_{1,2}$ vary from 10^{-2} GeV – 10^5 GeV, while the mixing V ranges from $V \sim 10^{-4} - 10^{-10}$. Here, for simplicity, we consider V as a parameter. For large mixing, $N_{1,2}$ states will thermalize and their decays would be constrained from the big bang nucleosynthesis (BBN). While a detailed evaluation of the BBN bound is beyond the scope of this present paper, we however show the lifetime contour in Fig. 1 that corresponds to $\tau(N_{1,2}) \sim 1$ sec. The two RHN states $N_{1,2}$ decay to various final states via their mixing V with the active neutrinos. For masses much smaller than the pion mass, the decay mode would be $\nu\gamma$ and $\nu\nu\nu$. The decay width and lifetime for this mass range are

$$\begin{aligned} \tau_{N_{1,2}}^{-1} &= \Gamma_{N_{1,2}} \simeq \frac{G_F^2 M_{N_{1,2}}^5}{96\pi^3} V^2 \\ &= 5.16 \times 10^{-24} \left(\frac{M_{N_{1,2}}}{\text{keV}} \right)^5 \left(\frac{V}{10^{-7}} \right) s^{-1}. \end{aligned} \quad (2.30)$$

For larger mass range $M_{N_{1,2}} > m_\pi + m_e$, additional decay modes $N \rightarrow l\pi^\pm$, $N \rightarrow lB^\pm/K^\pm$, and others will be open. For even an higher mass range $M_{N_{1,2}} > M_W, M_Z, M_{H_{1,2}}$, the two body modes $N \rightarrow lW, \nu Z, \nu H_{1,2}$ will be open. The expressions for these decay widths are

$$\Gamma_{N_{1,2}}(N \rightarrow l\pi) \simeq \frac{G_F^2 M_{N_{1,2}}^3}{96\pi^3} V^2, \quad (2.31)$$

$$\Gamma_{N_{1,2}}(N \rightarrow lW) \simeq \frac{g_F^2 M_{N_{1,2}}^3}{64\pi M_W^2} V^2 \left(1 - \frac{M_Z^2}{M_N^2} \right)^2 \left(1 + 2 \frac{M_W^2}{M_N^2} \right), \quad (2.32)$$

$$\Gamma_{N_{1,2}}(N \rightarrow \nu Z) \simeq \frac{g_F^2 M_{N_{1,2}}^3}{128\pi M_W^2} V^2 \left(1 - \frac{M_Z^2}{M_N^2} \right)^2 \left(1 + 2 \frac{M_Z^2}{M_N^2} \right), \quad (2.33)$$

$$\Gamma_{N_{1,2}}(N \rightarrow \nu H_{1,2}) \simeq \frac{g_F^2 M_{N_{1,2}}^3}{128\pi M_W^2} V^2 \left(1 - \frac{M_{H_{1,2}}^2}{M_N^2} \right)^2. \quad (2.34)$$

We evaluate the lifetime of $N_{1,2}$ assuming $M_{N_1} = M_{N_2}$ and show the contour of $\tau(N_{1,2}) = 1$ sec in Fig. 1 by the red line. Part of the region in the left side of the red line can be constrained from BBN as $N_{1,2}$ thermalize, and the decay of $N_{1,2}$ happens after $\tau(N_{1,2}) = 1$ sec. We estimate that for the region of Fig. 1 in agreement with both relic density and light neutrino mass, for which $M_{N_3} \sim 10\text{--}30$ KeV and $v_\chi \sim 4\text{--}10$ TeV, the cutoff scale $\Lambda \sim 10^7$ GeV, and the mixing angle $V \sim 10^{-6}$. Thus, $\tau(N_{1,2}) < 1$ sec, see Eq. (2.30), and the decay of $N_{1,2}$ in the early Universe occurs before BBN.

B. *Scenario-II*

We consider that in addition to the $N^T C^{-1} N \chi^2$ term, the Yukawa Lagrangian contains the term $N^T C^{-1} N \Phi^\dagger \Phi$. This is a more generic choice, as *Scenario-I* can be realized as only a special case of *Scenario-II* with $c' = 0$. The Lagrangian has the following terms

$$\begin{aligned} \mathcal{L}_{\text{eff}} &= \frac{c_{ij}}{\Lambda} N_i^T C^{-1} N_j \chi^2 + \frac{c'_{ij}}{\Lambda} N_i^T C^{-1} N_j \Phi^\dagger \Phi \\ &\quad + \frac{Y_{ij}}{\Lambda} \bar{L}_i \tilde{\Phi} N_j \chi + \text{H.c.} \end{aligned} \quad (2.35)$$

In this scenario, the RHN neutrino masses get contributions from both the $NN\chi^2$ and $NN\Phi^\dagger\Phi$ terms. As before, we consider c, c' to be diagonal matrix. The mass matrix of the two RHNs is

$$(M_R)_{\alpha\beta} = \frac{c_{\alpha\beta}}{\Lambda} v_\chi^2 + \frac{c'_{\alpha\beta}}{\Lambda} v_\Phi^2 \quad (\alpha, \beta = 1, 2). \quad (2.36)$$

The DM has a mass

$$M_{N_3} = \frac{c_{33}}{\Lambda} v_\chi^2 + \frac{c'_{33}}{\Lambda} v_\Phi^2. \quad (2.37)$$

The Dirac mass matrix has the same expression as in the previous section, Eq. (2.14), and the physical mass matrix of $N_{1,2}$ follows $M_N \sim M_R$. With the seesaw condition $M_R > M_D$, the light neutrino mass matrix has a similar expression as Eq. (2.16). Below, we consider the couplings $c \sim c' \simeq c_0$. Therefore, the light neutrino mass matrix

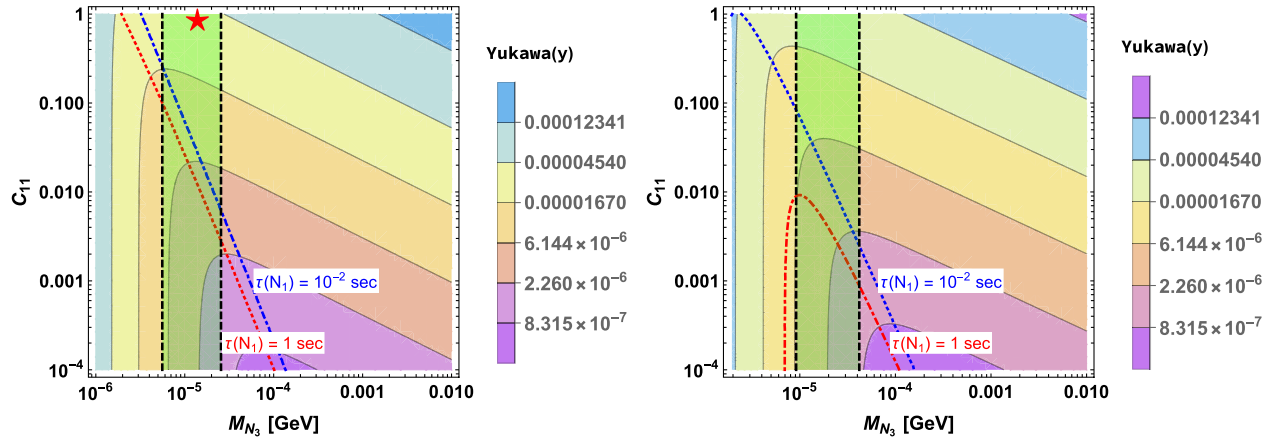


FIG. 2. Left and Right panel: the figures correspond to *scenario-IIa* for two different Higgs masses $M_{H_2} = 250$ GeV and 1.1 TeV, respectively. The color bar indicates the variation of Dirac Yukawa coupling y with respect to the variation of mass of FIMP DM M_{N_3} and the Yukawa coupling c_{11} . The green band indicates variation of v_{χ} in between 1 TeV to 10 TeV (from left to right). See the text for additional details.

receives a correction of $\mathcal{O}(\frac{v_{\Phi}^2}{v_{\chi}^2})$. The light neutrino and heavy RHN mass matrix have the following form

$$\begin{aligned} m_{\nu} &\sim -Y \frac{1}{c_0 \Lambda} v_{\Phi}^2 Y^T \left(1 - \frac{v_{\Phi}^2}{v_{\chi}^2} \right), \\ M_N &\sim \frac{c_0}{\Lambda} v_{\chi}^2 \left(1 + \frac{v_{\Phi}^2}{v_{\chi}^2} \right). \end{aligned} \quad (2.38)$$

Similar to the previous scenario, the RHN N_3 in this case is the FIMP DM. The particle is primarily produced from the two Higgs states $H_{1,2}$. The couplings of N_3 with $H_{1,2}$ states have the following form

$$\tilde{\lambda}_1 : -\frac{2v_{\chi}c_{33}}{\Lambda} \sin \theta + \frac{2v_{\Phi}c'_{33}}{\Lambda} \cos \theta, \quad (2.39)$$

$$\tilde{\lambda}_2 : \frac{2v_{\chi}c_{33}}{\Lambda} \cos \theta + \frac{2v_{\Phi}c'_{33}}{\Lambda} \sin \theta. \quad (2.40)$$

We first discuss two extreme scenarios:

- The $N_3 N_3 H_1$ coupling $\tilde{\lambda}_1$ is zero, i.e., the FIMP is produced only from H_2 decay.
- The $N_3 N_3 H_2$ coupling $\tilde{\lambda}_2$ is zero, i.e., the FIMP is produced from H_1 decay.

For the subsequent discussions, we consider the couplings c and c' independently.

(a) In the first scenario, the DM is entirely produced from the Higgs state H_2 . Imposing $\tilde{\lambda}_1 = 0$ in Eq. (2.39) leads to

$$\tan \theta = \frac{v_{\Phi}c'_{33}}{v_{\chi}c_{33}}. \quad (2.41)$$

Using Eq. (2.41), $\tilde{\lambda}_2$ can be simplified as

$$\tilde{\lambda}_2 = \frac{2M_{N_3}}{(v_{\chi} \cos \theta + v_{\Phi} \sin \theta)}. \quad (2.42)$$

Using the above coupling in Eq. (2.26) we obtain the constraint on the vev v_{χ} from $\Omega h^2 = 0.12$,

$$v_{\chi} = \frac{1.22 \times 10^{12} M_{N_3}^{3/2}}{M_{H_2}^{1/2} \cos \theta} - \frac{v_{\Phi} \sin \theta}{\cos \theta}. \quad (2.43)$$

Written in this way, the relic density constraint does not directly depend on the interaction coupling c_{33} of the Yukawa Lagrangian; the dependency is only via M_{N_3} . Rather, the vev v_{χ} depends on the mass of the DM, Higgs mass, and the SM-BSM Higgs mixing angle θ . The neutrino mass constraint, as we will derive, would be highly dependent on additional parameters. The constraint on the Dirac Yukawa is the same as Eq. (2.18), The mass of N_1 state (i.e., M_{N_1}) however gets additional contribution $c'_{11} v_{\Phi}^2 / \Lambda$ due to the $N_3 N_3 \Phi \Phi$ term

$$M_{N_1} = (c_{11} v_{\chi}^2 + c'_{11} v_{\Phi}^2) / \Lambda. \quad (2.44)$$

The cutoff scale Λ can be written in terms of the mass of the DM,

$$\Lambda = \frac{c_{33}}{M_{N_3}} (v_{\chi}^2 + \beta_3 v_{\Phi}^2), \quad (2.45)$$

where $\beta_3 = c'_{33}/c_{33}$. We combine different constraints from Eqs. (2.18), (2.41), (2.43), (2.44), and (2.45) in Fig. 2, where we show the variation of the Dirac coupling y in M_{N_3} - c_{11} plane. For this, we choose a light neutrino mass $m_{\nu} = 0.05$ eV, a Higgs mixing $\sin \theta = 0.1$, $M_{H_2} = 250$ GeV (left panel), and 1.1 TeV (right panel). Moreover

we assume $c'_{11} = 1$ and $c'_{33} = 2.5 \times 10^{-6}$. We have checked that for this choice of parameters the coupling c_{33} which dictates M_{N_3} , is perturbative in the entire region. We also display the region where v_χ is in the range 1 TeV–10 TeV. In Fig. 2 this is shown as the vertical green band as v_χ depends on M_{N_3} but not on c_{11} . The cutoff scale also increases with M_{N_3} ; we checked (using Eqs. (2.41) and (2.43) that at the boundary of the green band $\Lambda = (1, 2) \times 10^6$ GeV for left panel, and $\Lambda = (7, 14) \times 10^5$ GeV for the right panel.

It is evident from Fig. 2, that the choice of a large DM mass, M_{N_3} , together with a larger c_{11} demands a larger coupling y after imposing the light neutrino mass and relic density constraints. In the entire region the seesaw condition $M_R > M_D$ is satisfied. The red and blue lines represent the lifetime of $N_{1,2}$ as 1 sec and 10^{-2} sec, respectively for the left and for the right panel. The region enclosed by the red dashed line in the right plot corresponds to $\tau(N_1) > 1$ sec.

(b) The other scenario is where DM is produced from the SM Higgs. This can be realized for a suppressed $N_3 N_3 H_2$ coupling, we will consider the limit where this coupling is zero leading to the following constraint,

$$\tan \theta = -\frac{c_{33} v_\chi}{c'_{33} v_\Phi}. \quad (2.46)$$

Using Eq. (2.46), $\tilde{\lambda}_1$ can be simplified to the form

$$\tilde{\lambda}_1 = \frac{-2M_{N_3}}{(v_\chi \sin \theta - v_\Phi \cos \theta)}, \quad (2.47)$$

Using above coupling in Eq. (2.26) we get

$$v_\chi = \pm \frac{1.22 \times 10^{12} M_{N_3}^{3/2}}{M_{H_1}^{1/2} \sin \theta} + \frac{v_\Phi \cos \theta}{\sin \theta}. \quad (2.48)$$

In the above, the + and – sign correspond to $(v_\chi \sin \theta - v_\Phi \cos \theta) > 0$ and < 0 , respectively.

The constraint on the Dirac Yukawa in this case remains as in Eq. (2.18), where the cutoff scale Λ and the mass of the DM are related by Eq. (2.45). In Fig. 3 we plot the Dirac Yukawa as a function of c_{11} and M_{N_3} . As before we consider the parameter $\sin \theta = 0.1$. Additionally, we consider $c'_{11} = 1$, $c'_{33} = -10^{-6}$. The Yukawa coupling c_{33} varies with M_{N_3} , and is perturbative in the entire range of M_{N_3} . The seesaw condition $M_D < M_R$ is satisfied in the entire parameter space. The green band bounded by black dashed lines represent the variation of v_χ (Λ) between 3 TeV (2×10^4 GeV) and 10 TeV (5×10^4 GeV), from left to right. The lifetime of $N_{1,2}$ is less than 1 sec in the entire range. For illustration, the blue and red dashed lines

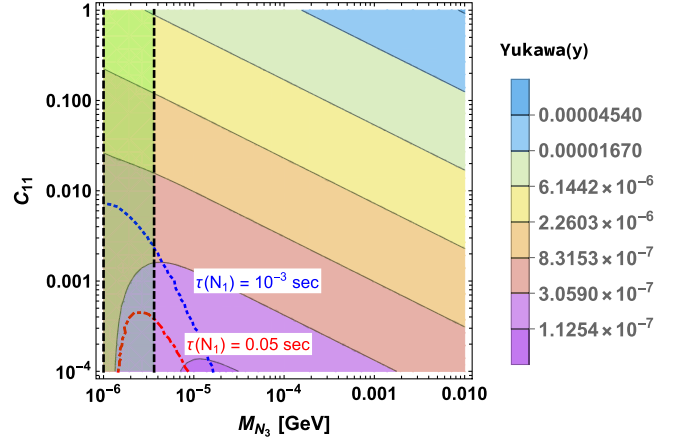


FIG. 3. The figure corresponds to *Scenario-IIB*, and represents the variation of the Yukawa coupling y w.r.t the variation of the DM mass M_{N_3} and the Yukawa coupling c_{11} . Contours of $\tau(N_{1,2}) = 0.001$ (blue) and 0.05 (red) are displayed. The vertical black lines correspond to $v_\chi = 3$ TeV (left) and 10 TeV (right).

indicate lifetime of $N_{1,2}$ as 0.001 sec, and 0.05 sec, respectively.

We also consider the generic scenario where both the $H_1 \rightarrow N_3 N_3$ and $H_2 \rightarrow N_3 N_3$ contribute to the relic density. In Fig. 4, we show different constraints in the $v_\chi - M_{N_3}$ plane. The blue band represents the total contribution from $H_1 \rightarrow N_3 N_3$ and $H_2 \rightarrow N_3 N_3$ which varies in the mentioned range. The green band represents the constraint from light neutrino mass. While the relic density constraint does not depend on the Yukawa y , the latter depends on few additional parameters. See Table. I for the details of the input parameters. Similar to *Scenario-I*, we represent

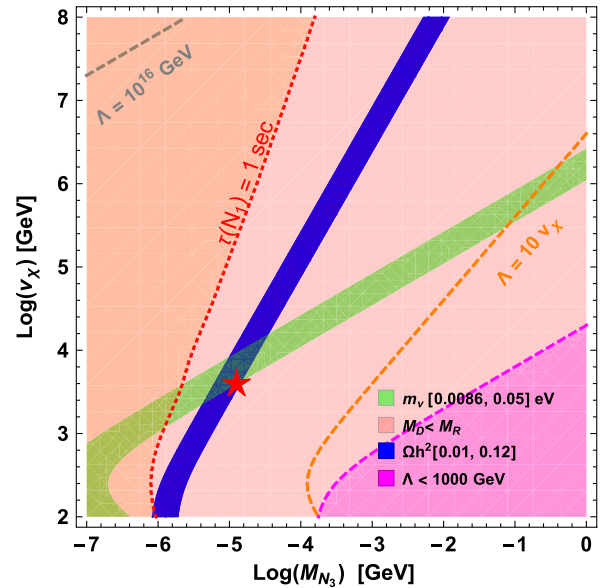


FIG. 4. Similar as the Fig. 1, but for *Scenario-II*. See text for more details.

$\tau(N_{1,2}) = 1$ sec by red line, cutoff scale $\Lambda = 10v_\chi$ by orange line. The point represented by a red star mark in this plot, corresponds to the star point shown in the left panel of Fig. 2, representing the same benchmark point. Similar to the previous scenario *Scenario-I*, a higher vev v_χ is required to satisfy relic abundance for a heavier DM mass. This happens, as for both these two scenarios, the DM mass is governed by the vev v_χ , which also governs the $H_{1,2} \rightarrow N_3 N_3$ coupling, and hence the DM production. Therefore, a TeV scale v_χ together with a TeV scale BSM Higgs with mass M_{H_2} demand that the DM mass in this case can be at most $M_{N_3} \sim \text{few KeV}$. We will see in the next section, how addition of a bare mass term in the Lagrangian relaxes this strong correlation.

C. Scenario-III

The DM phenomenology changes if a bare mass term for the FIMP DM N_3 state is being added to the Lagrangian. In this case, the tight correlation between the DM mass and vev of the gauge singlet scalar relaxes. Adding a bare mass term M_B for $N_{1,2,3}$, the Lagrangian has the following form,

$$\mathcal{L}_{\text{eff}} = \frac{c_{ij}}{\Lambda} N_i^T C^{-1} N_j \chi^2 + \frac{c'_{ij}}{\Lambda} N_i^T C^{-1} N_j \Phi^\dagger \Phi + \frac{Y_{ij}}{\Lambda} \bar{L}_i N_j H \chi + M_B N_i^T C^{-1} N_i + \text{H.c.} \quad (2.49)$$

In the above, M_B is a diagonal mass matrix $M_B = (M_{B_1}, M_{B_2}, M_{B_3})$, that represents the bare mass term for RHNs. The RHN mass matrix of $N_{1,2}$ can be written as follows,

$$(M_R)_{\alpha\beta} = \frac{c_{\alpha\beta} v_\chi^2}{\Lambda} + \frac{c'_{\alpha\beta} v_\phi^2}{\Lambda} + (M_B)_{\alpha\beta} \quad (\alpha, \beta = 1, 2). \quad (2.50)$$

Hence, the DM mass is

$$M_{N_3} = \frac{c_{33} v_\chi^2}{\Lambda} + \frac{c'_{33} v_\phi^2}{\Lambda} + M_{B_3}. \quad (2.51)$$

Note that, for c_{33} and c'_{33} of $\mathcal{O}(1)$, the mass of N_3 can primarily be governed by M_{B_3} , and the other two terms can provide subdominant contributions. The Dirac mass matrix has the same form as in previous scenarios, Eq. (2.14), and as before, $M_N \sim M_R$. The light neutrino mass matrix has the following expression,

$$m_\nu \sim -\frac{1}{\Lambda} \frac{Y v_\phi^2 v_\chi^2 Y^T}{c_{11} v_\chi^2 + c'_{11} v_\phi^2 + \Lambda M_{B_1}}, \quad (2.52)$$

where we again consider $M_{N_1} = M_{N_2}$.

TABLE II. The parameters relevant for *Scenario-III*, Fig. 5.

<i>Scenario-III</i>					
M_{H_2}	$\sin \theta$	y	$c_{33} (= c'_{33})$	$M_{N_3} - M_{B_3}$	M_{N_1}
250 GeV	0.1	1	10^{-4}	10^{-8} GeV	$4M_{N_3}$

In this scenario, the DM can be produced from $H_{1,2} \rightarrow N_3 N_3$ decays. The interaction vertex of N_3 with $H_{1,2}$ states have the same form as given in Eq. (2.39) and Eq. (2.40). As in previous scenarios we consider two extreme cases where a) $\tilde{\lambda}_1 = 0$ and the DM is produced from H_2 decay. b) $\tilde{\lambda}_2 = 0$ and the DM is produced from H_1 decay.

(a) The condition $\tilde{\lambda}_1 = 0$ together with Eq. (2.41) leads to the $N_3 N_3 H_2$ coupling

$$\tilde{\lambda}_2 = \frac{2(M_{N_3} - M_{B_3})}{(v_\chi \cos \theta + v_\Phi \sin \theta)}, \quad (2.53)$$

and to the constraint on the vev v_χ

$$v_\chi = \frac{1.22 \times 10^{12} (M_{N_3} - M_{B_3})^{3/2}}{M_{H_2}^{1/2} \cos \theta} - \frac{v_\Phi \sin \theta}{\cos \theta}. \quad (2.54)$$

(b) The condition $\tilde{\lambda}_2 = 0$ together with Eq. (2.46) leads to the $N_3 N_3 H_2$ coupling

$$\tilde{\lambda}_1 = -\frac{2(M_{N_3} - M_{B_3})}{(v_\chi \sin \theta - v_\Phi \cos \theta)}, \quad (2.55)$$

and to the constraint on the vev v_χ

$$v_\chi = \pm \frac{1.22 \times 10^{12} (M_{N_3} - M_{B_3}) M_{N_3}^{1/2}}{M_{H_1}^{1/2} \sin \theta} + \frac{v_\Phi \cos \theta}{\sin \theta}. \quad (2.56)$$

For both the scenarios *a* and *b*, it is evident that the factor $M_{N_3} - M_{B_3}$ in the numerator of Eq. (2.54) and Eq. (2.56) will have impact on the tight correlation between v_χ and M_{N_3} , found in *Scenario-I-II*.

In general both $H_1 \rightarrow N_3 N_3$ and $H_2 \rightarrow N_3 N_3$ can contribute to the relic density. With the choice of input parameters given in Table II, the constraints on M_{N_3} and v_χ are shown in Fig. 5. The pink color shaded area indicates that the seesaw approximation $M_R > M_D$ is satisfied. In the blue band, DM relic density varies in the range $0.01 < \Omega h^2 < 0.12$ from left to right. The green band represents the constraint from light neutrino mass while the red dashed line corresponds to the contour $\tau(N_{1,2}) = 1$ sec. The horizontal lines represent the cutoff scale $\Lambda = 10^{10}$ and 10^{12} GeV. From Fig. 5 it can be seen that the

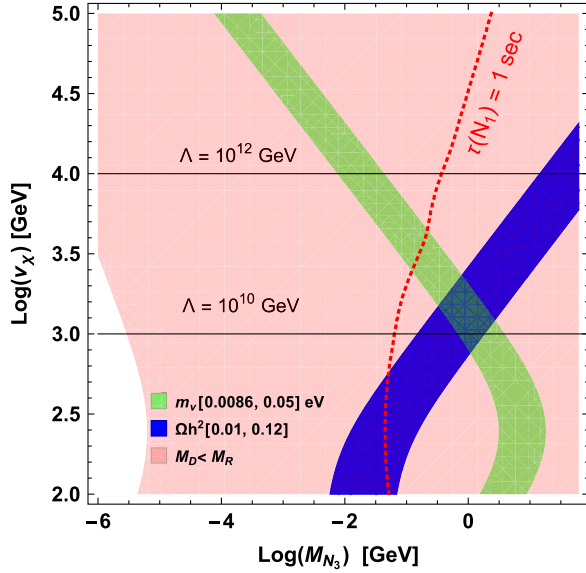


FIG. 5. *Scenario-III*: Constraints from the DM relic density (blue band) and eV scale light neutrino mass (green band) in the M_{N_3} - v_χ plane. The values of other free parameters are listed in Table II. Here we include the decay of $H_{1,2}$ for DM production.

region compatible with both the relic density and the neutrino mass constraints corresponds to $v_\chi \sim \mathcal{O}(\text{TeV})$ and $M_{N_3} \sim \mathcal{O}(\text{GeV})$. In this scenario it is therefore natural to have a BSM Higgs at the TeV scale for a coupling $\lambda_2 \sim \mathcal{O}(1)$. In Sec. IV we will consider this scenario and explore the collider signature of a TeV scale BSM Higgs.

III. DM PRODUCTION: DECAY VS ANNIHILATION

In the previous sections we considered only decay contributions of the SM and BSM Higgs in the relic density. This is justified for a not too large reheating temperature T_R [13]. In this section we deviate from this assumption and allow for a high T_R . Thus we obtain larger contributions from annihilation channels as well. To determine the comoving number density of N_3 , we need to solve the following Boltzmann equation,

$$\begin{aligned} \frac{dY_{N_3}}{dz} = & \frac{2M_{pl}}{1.66M_{H_2}^2} \frac{z\sqrt{g_*(z)}}{g_*(z)} \left[\sum_{i=1,2} \langle \Gamma_{H_i \rightarrow N_3 N_3} \rangle (Y_{H_i}^{eq} - Y_{N_3}) \right] \\ & + \frac{4\pi^2 M_{pl} M_{H_2} \sqrt{g_*(z)}}{45 \cdot 1.66 z^2} \\ & \times \left[\sum_{x=W,Z,h_1,h_2,f} \langle \sigma_{v_\chi \bar{x} \rightarrow N_3 N_3} \rangle (Y_x^{eq 2} - Y_{N_3}^2) \right], \quad (3.1) \end{aligned}$$

where $z = M_{H_2}/T$ and $g_*(z)$ depends on $g_s(z)$, $g_\rho(z)$ in the following way [39], $\sqrt{g_*(z)} = \frac{g_s(z)}{\sqrt{g_\rho(z)}} \left(1 - \frac{1}{3} \frac{d \ln g_s(z)}{d \ln z}\right)$. Y_A is the comoving number density of A , M_{pl} is the Planck mass

and the quantity inside $\langle \dots \rangle$ corresponds to the thermal average of decay rate and annihilation cross sections. From the comoving number density obtained by solving the Eq. (3.1), one can determine the relic density of DM by following the expression,

$$\Omega_{N_3} h^2 = 2.755 \times 10^8 \left(\frac{M_{N_3}}{\text{GeV}} \right) Y_{N_3}(T_0). \quad (3.2)$$

We perform our numerical simulation using micrOMEGAs5.0 [32] after implementing the model file in FeynRules [40], where we use appropriate quantum statistics for bosons and fermions.

In the left and right panel of Fig. 6, we show the variation of the DM relic density with its mass for *Scenario-III* and *Scenario-II*, respectively. We keep the other parameters fixed as shown in the caption. For these figures, we choose a high value of the reheating temperature, $T_R = 10^9$ GeV. From both the figures, one can see that the DM relic density increases with its mass M_{N_3} as also evident from Eq. (3.2). In the left panel, there is a sharp fall in the H_1 decay contribution at $M_{N_3} = \frac{M_{H_1}}{2}$ when the H_1 decay to DM is not kinematically allowed. Note that for large T_R the annihilation contribution is larger than the decay contribution by several orders of magnitude. This occurs as relic density for annihilation (from $WW/ZZ \rightarrow N_3 N_3$, $H_i H_j \rightarrow N_3 N_3$) is proportional to the value of the reheating temperature as shown in the Appendix, (see Eq. (A9)).

We also show the variation of relic density with the reheating temperature T_R (shown in the left panel), and the vev v_χ (shown in the right panel) in Fig. 7. Moreover, the individual contributions from decay and annihilation processes to the total relic density have also been shown in both the panels. In the left panel, we can see that the decay contribution does not depend on the reheating temperature, T_R , but the annihilation contribution strongly depends on the reheating temperature when $T_R \gtrsim 10^6$ GeV. As we have shown in Eq. (A9) of the Appendix, and also have been discussed in [13], the annihilation contribution depends on the reheating temperature and linearly grows with it. This is visible in Fig. 7, where for large T_R annihilation contributions increase. This occurs because of the presence of $d = 5$ $H_i H_j N_3 N_3$ ($i, j = 1, 2$) operators. We also show in the Appendix that $WW/ZZ \rightarrow N_3 N_3$ exhibits similar feature. Furthermore, as evident from Eq. (A10), the contribution from the annihilation process $f\bar{f} \rightarrow N_3 N_3$ is dominated by the low scale physics. For a very high T_R this contribution becomes almost independent of the variation of T_R . In the right panel of Fig. 7, the variation of relic density with the vev v_χ has been shown. One can see that both decay and annihilation contributions fall linearly with the increase of the v_χ . This can be explained easily, as the coupling of DM with the Higgs

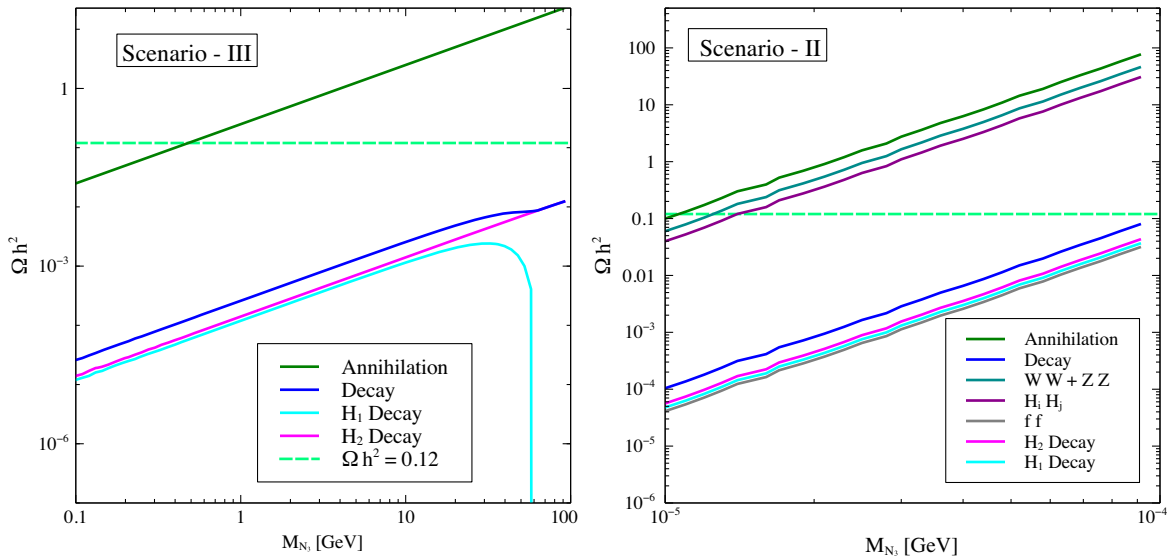


FIG. 6. Left (right) panel: Ωh^2 vs the mass of DM for *Scenario-III* (*Scenario-II*) with $\sin\theta = 0.1$, $\Lambda = 10^{11}$ GeV, $M_{H_2} = 1100$ GeV, $M_{N_{1,2}} = 5M_{N_3}$, $M_{N_3} - M_{B_3} = 10^{-8}$ GeV, $v_\chi = 3$ TeV (left panel), $M_{B_3} = 0$, $v_\chi = 15$ TeV (right panel). The green-dashed horizontal line represents $\Omega h^2 = 0.12$ [36].

is inversely proportional to the square of the vev *i.e.*, $\tilde{\lambda}_{1,2} \propto \frac{1}{v_\chi^2}$. In addition, the $WW \rightarrow N_3 N_3$, $ZZ \rightarrow N_3 N_3$ contributions also depend on the coupling $\tilde{\lambda}_{1,2}$. Moreover the contact interaction $H_i H_i \rightarrow N_3 N_3$ depends on c_{33}/Λ , which also varies as $1/v_\chi$. Therefore, as the vev increases, the relevant coupling becomes smaller, resulting in a reduced production of DM.

A. The parameter space of *Scenario-III*

In the previous section, we have evaluated the decay and annihilation contributions to the relic density for specific benchmark points. In this section we vary all the free parameters of *Scenario-III* in a wide range and present the results in the form of scatter plots. The model parameters are varied in the following range,

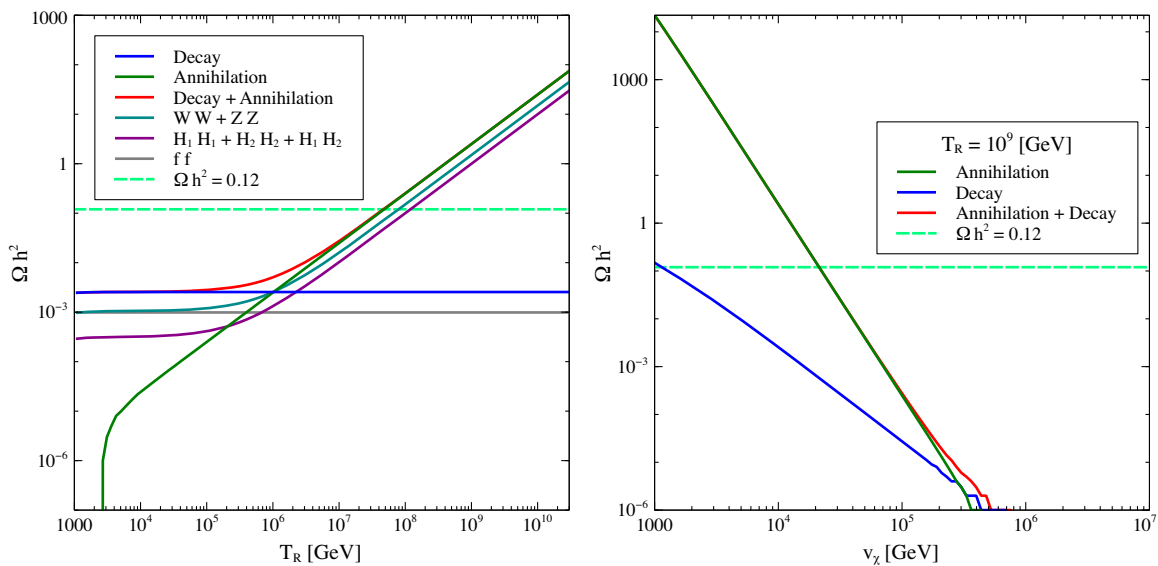


FIG. 7. Left panel: Relic density vs reheating temperature for *Scenario-III*. Right panel: variation of relic density vs vev of the singlet scalar for the same scenario. Other parameters kept fixed at $\sin\theta = 0.1$, $\Lambda = 10^{11}$ GeV, $M_{H_2} = 1100$ GeV, $M_{N_{1,2}} = 5M_{N_3}$, $M_{N_3} - M_{B_3} = 10^{-8}$ GeV and $v_\chi = 3000$ GeV. The green dashed horizontal line represents the experimental constraint on DM relic density [36].

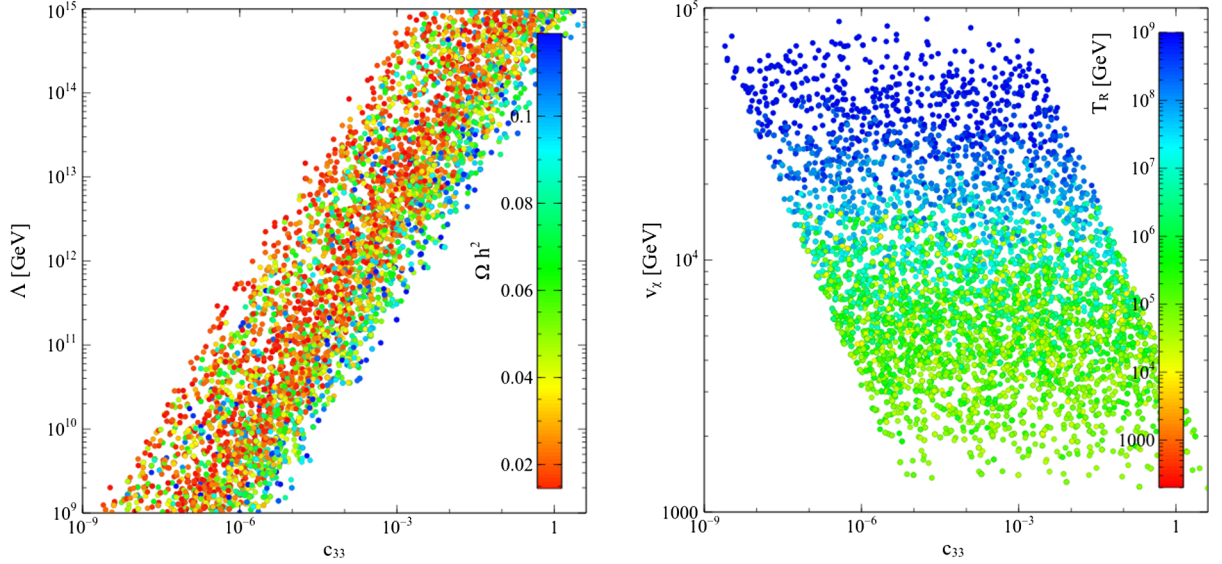


FIG. 8. Left panel: variation of relic density of N_3 in the c_{33} - Λ plane. Right panel: variation of T_R in the c_{33} - v_χ plane. We consider the same constraint as in the left panel.

$$\begin{aligned}
 200 \text{ GeV} < M_{H_2} < 3000 \text{ GeV}, \\
 10 \text{ GeV} < M_{N_3} < 100 \text{ GeV}, \\
 10^{-3} < \theta < 10^{-1}, \\
 1000 \text{ GeV} < v_\chi < 10000 \text{ GeV}, \\
 200 \text{ GeV} < T_R < 10^9 \text{ GeV}, \\
 10^9 \text{ GeV} < \Lambda < 10^{14} \text{ GeV}.
 \end{aligned} \tag{3.3}$$

To accommodate H_2 at the TeV scale together with $v_\chi \sim \text{TeV}$, the bare mass term of N_3 has to dominate its physical mass M_{N_3} , thus we impose $M_{N_3} - M_{B_3} \sim \mathcal{O}(10^{-8}) \text{ GeV}$. In our scan we require that N_3 contribute

to at least 10% of the total DM, thus we impose that its relic density falls within the range

$$0.01 < \Omega h^2 < 0.1211. \tag{3.4}$$

In Figs. 8 and 9 we display the allowed parameter space after taking into account the constraint from Eq. (3.4). The entire range of M_{N_3} mentioned above can satisfy Eq. (3.4) with the variation of other model parameters. We did not find any strong correlation between M_{N_3} and other model parameters, and hence we do not present any scatter plot for M_{N_3} .

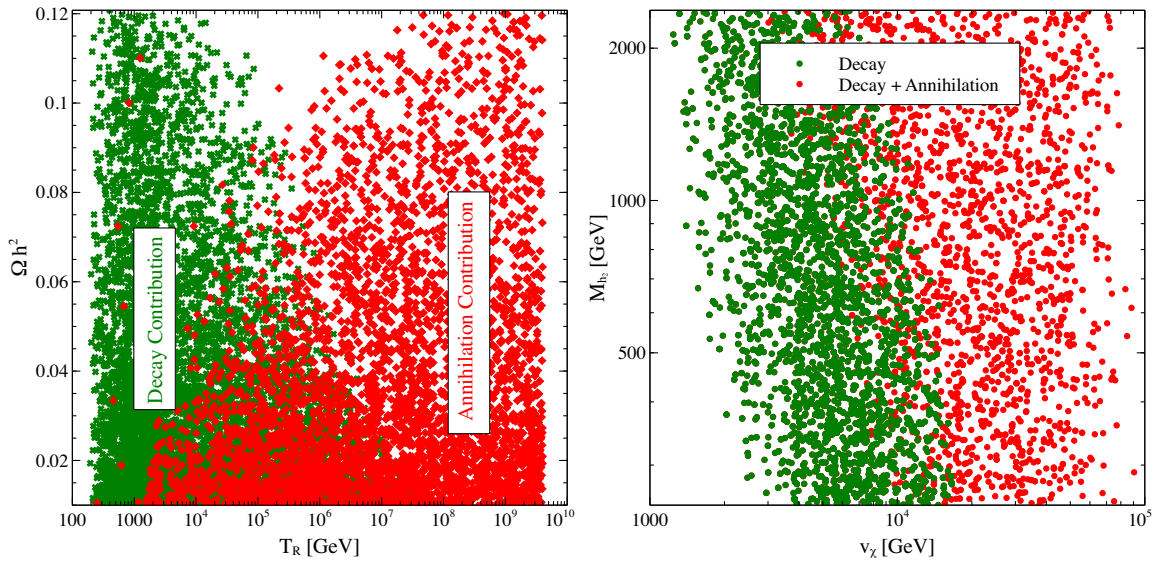


FIG. 9. Left panel: Variation of relic density of N_3 vs T_R . Right panel: scatter plot in the v_χ - M_{H_2} plane.

In the left panel of Fig. 8, we show the variation of relic density (in color bar) in the c_{33} - Λ plane after satisfying Eq. (3.4). For *Scenario-III*, we can express the coupling c_{33} in terms of the cutoff scale Λ and the $vevs$ v_ϕ, v_χ in the following way,

$$c_{33} = \frac{\Lambda(M_{N_3} - M_{B_3})}{v_\chi^2 + \beta_3 v_\phi^2}. \quad (3.5)$$

As evident from the above expression, there is a linear relation between c_{33} and Λ . Therefore, as we increase Λ , c_{33} also increases. This is clearly visible from the figure shown in the left panel. The blue scattered points satisfy the experimentally measured DM relic density constraint [36]. In the right panel of Fig. 8, we show the points that satisfy Eq. (3.4) in the v_χ - c_{33} plane. As expected from Eq. (3.5), since $v_\chi \gg \beta_3 v_\phi$, c_{33} is inversely proportional to v_χ^2 . Furthermore, right panel of Fig. 8 shows that T_R increases linearly with v_χ . This can be understood as follows. We have seen in Fig. 7 that the relic density (when dominated by the annihilation contribution) increases with the reheating temperature at large T_R , and decreases as $1/v_\chi$. Therefore, for a given value of the DM relic density, higher values of T_R will be associated with larger values of v_χ . The yellow points are not clearly visible, as they have been covered by the green points.

In the left panel of Fig. 9, we show the variation of the decay and annihilation contributions to the relic density with T_R . For the discussion on the dependency of the relic density on T_R , see Sec. V. In the right panel of the same plot, we show the relation between v_χ and M_{H_2} . As evident from the left plot, for reheating temperature $T_R \sim 10^5$ GeV, the decay and annihilation contributions are equal, while for $T_R > 10^6$ - 10^7 GeV, the annihilation contributions dominate. Lower than $T_R \sim 10^5$ GeV, decay contribution to the relic density dominates. This occurs as the annihilation contribution is directly proportional to the reheating temperature as has been explained before. In generating the scatter plots both for Fig. 8 and Fig. 9, we assume that the reheating temperature is greater than the masses of all the particles.

In the right panel of Fig. 9, we show the points in the v_χ - M_{H_2} plane which satisfy Eq. (3.4). We represent the decay contribution by green points and the total contribution by red points. For the decay contribution, there exists an inverse correlation between v_χ and M_{H_2} . The $H_2 N_3 N_3$ coupling for this case takes the following form,

$$\tilde{\lambda}_2 = \frac{2(M_{N_3} - M_{B_3}) \times (v_\chi \cos \theta + \beta_3 v_\phi \sin \theta)}{v_\chi^2 + \beta_3 v_\phi^2}. \quad (3.6)$$

The above equation together with Eq. (2.28) imply that the DM relic density decreases with the increase of both v_χ and M_{H_2} . Therefore, for a given value of the DM relic density,

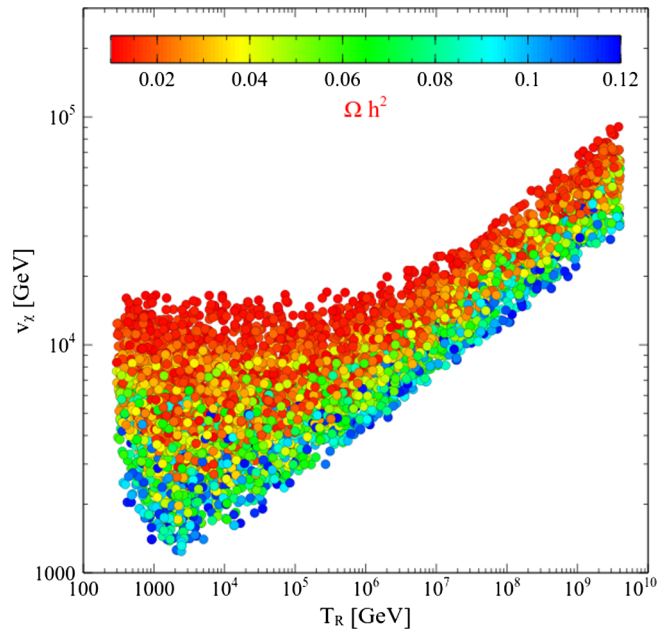


FIG. 10. Variation of relic density with respect to variation of v_χ and T_R . The points satisfy the mentioned DM relic density range.

higher values of v_χ will be associated with smaller values of M_{H_2} . For annihilation processes the correlation between v_χ and M_{H_2} is somewhat mild, as the additional parameter T_R plays a significant role in annihilation processes, and we have varied T_R in a wide range mentioned in Eq. (3.3).

Finally in Fig. 10, we show the variation of the relic density with respect to the variation of vev v_χ and the reheating temperature T_R . As we have discussed, the coupling strength to produce DM varies inversely with the vev , therefore for a smaller vev , the coupling c_{33} increases. This results in a higher value of the Ωh^2 which is represented by the blue points. We can also see for $T_R > 10^6$ GeV there exists a sharp correlation between T_R and v_χ , which is consistent with the right panel of Fig. 8.

IV. COLLIDER SIGNATURE OF H_2

Other than the SM Higgs, the model also contains a neutral BSM Higgs, which can be probed at collider experiments. A number of LHC measurements constrain the presence of such a heavy Higgs, and its mixing with the SM 125 GeV Higgs [35]. For the collider analysis we consider BSM Higgs in the TeV mass range. We pursue the study for *Scenario-III*. The model signature we study will remain the same for *Scenario-I*, and *II* as well, as the signature does not depend on DM mass. The LHC searches that constrain the BSM Higgs and its mixing are: a) The SM Higgs signal strength measurement, and b) heavy Higgs searches.

a) Higgs signal strength measurement constrains the mixing of the SM and BSM Higgs. To evaluate this, we adopt [35]. The signal strength of SM Higgs is given as

TABLE III. LHC constraints on BSM Higgs production in GF and VBF channel. We follow [41] for ZZ channel (ATLAS search-139 fb⁻¹), [42] for $VV(WW + ZZ)$ channel (ATLAS search-139 fb⁻¹), [43] for H_1H_1 (ATLAS search-36.1 fb⁻¹), and [45] for $H_1H_1 + jj$ channel (ATLAS search-126 fb⁻¹). For H_1H_1 channel, the limits from CMS search-35.9 fb⁻¹ [44] are mentioned in bracket, that we obtained after dividing $\sigma(pp \rightarrow H_2 \rightarrow H_1H_1 \rightarrow 4b)$ by $\text{Br}^2(H_1 \rightarrow b\bar{b})$.

σ (pb)	$M_{H_2} = 250$ GeV	$M_{H_2} = 500$ GeV	$M_{H_2} = 1100$ GeV
$\sigma(pp \rightarrow H_2 \rightarrow ZZ)$	<0.12	<0.02	<0.005
$\sigma(pp \rightarrow H_2 \rightarrow VV)$	<1.6 ($M_{H_2} = 300$ GeV)	<0.2	<0.008
$\sigma(pp \rightarrow H_2 \rightarrow H_1H_1)$	<0.8 (3)	<0.15 (0.25)	<0.03 (0.04)
$\sigma(pp \rightarrow H_2jj \rightarrow H_1H_1jj)$	<1	<0.05	<0.004 ($M_{H_2} = 1$ TeV)

$$\mu_{h \rightarrow xx} = \frac{\sigma_{H_1} \text{Br}(H_1 \rightarrow xx)}{\sigma_{H_1}^{\text{SM}} \text{Br}^{\text{SM}}(H_1 \rightarrow xx)}. \quad (4.1)$$

In the above, $H_1 \rightarrow xx$ represents any channel where SM Higgs can decay. In our model one of the additional final states in which SM Higgs can decay is to the DM pair ($H_1 \rightarrow N_3N_3$). However, the branching ratio of this channel is very small; $\text{Br}(H_1 \rightarrow N_3N_3) < 10^{-8}$. The other modes such as $H_1 \rightarrow \nu N_{1,2}/N_1N_1/N_2N_2$ also have small branching ratios. Therefore, for all practical purposes, these channels can be neglected. The branching ratio of H_1 decaying to any SM final state is hence almost identical to the branching ratio in the SM, i.e., $\text{Br}(H_1 \rightarrow xx) \sim \text{Br}^{\text{SM}}(H_1 \rightarrow xx)$. The production cross section of $pp \rightarrow H_1$ becomes $\sigma_{H_1} = \cos^2 \theta \sigma_{H_1}^{\text{SM}}$. Therefore, we find that the above Higgs signal strength expression takes a very simplified form,

$$\mu_{H_1 \rightarrow xx} \sim \cos^2 \theta. \quad (4.2)$$

The $\sqrt{s} = 13$ TeV LHC measurements of the Higgs signal strength in combined channel dictates $\mu = 1.17 \pm 0.1$ [35]. We find, allowing a 3σ deviation around the best fit value of $\mu = 1.17$, that the Higgs mixing angle $\sin \theta < 0.36$. In our subsequent analysis, we consider both a larger value of $\sin \theta = 0.34$ and a small value, $\sin \theta = 0.1$.

b) A number of LHC searches constrain the production of heavy scalar resonance and its decay into various SM final states. As discussed in the previous sections, one of the main production channels of the FIMP DM is the decay of H_2 to two N_3 states, which is dominant for a low reheating temperature. However, due to the negligible branching ratio $\text{Br}(H_2/H_1 \rightarrow N_3N_3) < 10^{-8}$, this channel is not constrained by the LHC searches for the invisible decays of Higgs. The main decay channels of the BSM Higgs include $H_2 \rightarrow WW, ZZ, ZZ, b\bar{b}, \tau^+\tau^-, H_1H_1$ channels. A number of CMS and ATLAS searches constrain the production cross section of the BSM Higgs in gluon fusion (GF), or vector boson fusion (VBF) channels folded with the branching ratio of the H_2 in the above mentioned modes. In Table. III, we outline the most sensitive searches for a neutral Higgs at the LHC, where we quote the limits on $\sigma \times \text{Br}$ for few illustrative mass

points. We consider the searches $pp \rightarrow H_2 \rightarrow ZZ$ [41], $pp \rightarrow H_2 \rightarrow W^+W^- + ZZ$ [42], $pp \rightarrow H_2 \rightarrow H_1H_1 \rightarrow 4b$ [43,44], and $pp \rightarrow H_2 + jj \rightarrow H_1H_1 + jj$ [45]. We find that among them $pp \rightarrow H_2 \rightarrow ZZ$ [41] is the most constraining; in particular this channel does not allow lighter masses, $M_{H_2} \sim 200$ GeV, for larger value of $\sin \theta = 0.34$. Note that this mixing angle is allowed by Higgs signal strength measurements. Such values are marginally allowed by the search $pp \rightarrow H_2 \rightarrow VV$. On the other hand, the large mixing angle, $\sin \theta \sim 0.34$ is allowed by all the above mentioned searches when $M_{H_2} \sim 1$ TeV. The mixing angle $\sin \theta = 0.1$ which we consider for the DM analysis in the previous section is allowed for the entire range $200 \text{ GeV} < M_{H_2} < \text{few TeV}$. We have cross checked our results with the results obtained from HiggsBounds [46].

Other than the WW, ZZ channels, one of the spectacular signatures of a heavy BSM scalar is the di-Higgs signal. The di-Higgs channel is, in particular, important to probe Higgs trilinear coupling. Any deviation from the SM prediction will indicate new physics. Di-Higgs production in SM, which is nonresonant, has extensively been studied at the LHC. There are different studies that analyzed $b\bar{b}\gamma\gamma$ [47–50], $b\bar{b}W^+W^-$ [51,52], $b\bar{b}\tau^+\tau^-$ [51,53,54], $b\bar{b}b\bar{b}$ [54–57], and $b\bar{b} + E_T$ [58] final states. With a heavy BSM Higgs which couples to two SM Higgs, the di-Higgs production cross section becomes large. The studies [59–68] focus on resonant production of a BSM Higgs and its decay to di-Higgs. Both resonant and nonresonant di-Higgs production processes have been extensively explored by CMS and ATLAS [43,69–78]. Most of the above mentioned studies focused on the resolved final states with isolated final state leptons, jets, and photons. However, for a very heavy mass of the BSM Higgs, the produced SM Higgs will be boosted, leading to collimated decay products. Note that, the di-Higgs production from a heavy Higgs with few TeV mass is less favourable at $\sqrt{s} = 13$ TeV LHC due to smaller production cross section. For nonresonant di-Higgs production at the proposed $\sqrt{s} = 100$ TeV LHC, see [79–83].

We instead focus on resonant di-Higgs production from the decay of a heavy BSM Higgs at 100 TeV collider, decaying into two SM Higgs. The branching ratio of this

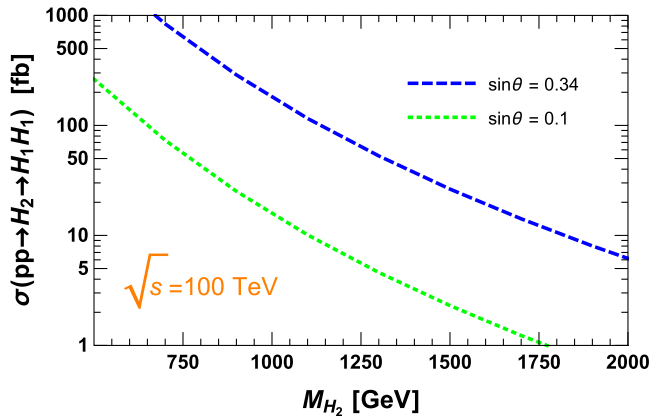


FIG. 11. Production cross section of $pp \rightarrow H_2 \rightarrow H_1 H_1$ vs mass of H_2 for different values of the Higgs mixing angle.

channel $\text{Br}(H_2 \rightarrow H_1 H_1) \sim 25\%$ for H_2 mass around 1 TeV [84]. We analyze the di-Higgs channel with subsequent decay of H_1 to $b\bar{b}$. We assume $M_{H_2} > 1\text{ TeV}$, for which the two Higgs bosons produced from H_2 are moderately boosted, leading to a peak in ΔR separation between the two b quarks as $\Delta R(b, \bar{b}) \lesssim 0.4$. Instead of a resolved analysis with four or more number of isolated b jets in the final state, we perform an analysis where we adopt a large radius jet as the jet description, which is effective in suppressing a number of SM backgrounds. Therefore, our model signature is

$$pp \rightarrow H_2 \rightarrow H_1 H_1 \rightarrow 2j_{\text{fat}}. \quad (4.3)$$

where, each of the fatjet j_{fat} contains two b quarks appearing from Higgs decay.

To evaluate the signature, we implement the Lagrangian of this model in FeynRules(v2.3) [40] to create the UFO [85] model files. The Event generator MADGRAPH5_aMC@NLO (v2.6) [86] is used to generate both the signal and the background events at leading order. Generated events are passed through PYTHIA8 [87] to perform showering and hadronization. Detector effects are simulated using DELPHES(v3.4.1) [88]. We use FastJet [89] for the clustering of fatjets and consider Cambridge-Aachen [90,91] algorithm, with radius parameter $R = 1.0$.

In Fig. 11, we show the production cross section of $pp \rightarrow H_2 \rightarrow H_1 H_1$ at $\sqrt{s} = 100\text{ TeV}$ as a function of the mass of H_2 for $\sin\theta = 0.1, 0.34$. There are a number of SM backgrounds, that can mimic the signal. This include both QCD and electroweak processes. The QCD is generated by combining $b\bar{b}b\bar{b}$ and $b\bar{b}jj$ final state. The other backgrounds which include electroweak coupling are ditop ($t\bar{t}$), diboson (WW and ZZ), Wj and Zj . Here we consider full hadronic decays of top quark, W and Z boson. At the generator level, we implement these following cuts on background samples:

- The transverse momentum of the partons: $p_T^j > 20\text{ GeV}$, $p_T^b > 15\text{ GeV}$
- The pseudo-rapidity of the partons: $|\eta_j| < 5.0$, $|\eta_b| < 3.0$
- The separation between partons: $\Delta R_{jj} > 0.4$, $\Delta R_{bb} > 0.2$, $\Delta R_{bj} > 0.4$
- Invariant mass of the two b quark: $m_{bb} > 30\text{ GeV}$
- The scalar sum of the transverse momentum of all the hadronic particles H_T of the background: $H_T > 800\text{ GeV}$

The third cut on ΔR is to avoid any divergence, that may arise from the QCD samples. The distribution of H_T is shown in Fig. 12. The cut on H_T ensures the sufficiently

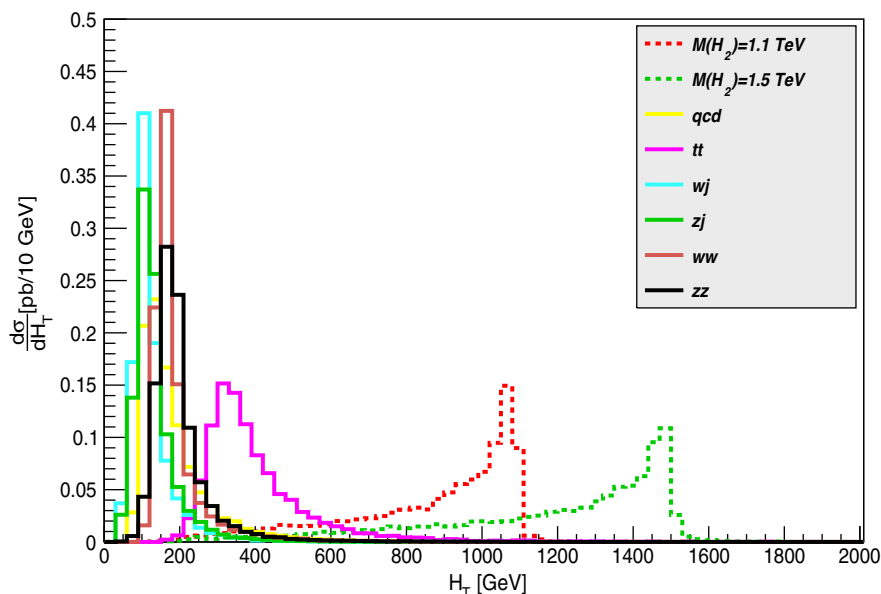


FIG. 12. H_T distribution for signal and background samples.

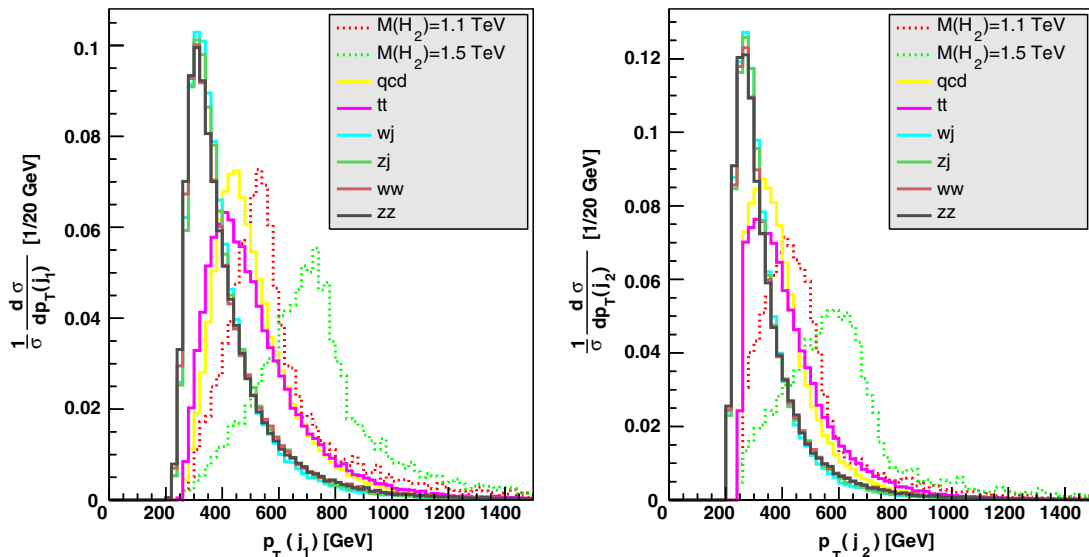


FIG. 13. p_T distribution of leading fatjet j_1 (Left panel) and subleading fatjet j_2 (right panel). We also show the p_T distributions of the fatjets arising from background samples.

large background population in the desired region, where the signal populates. We do not take the SM di-Higgs channel into consideration, as we find that after the H_T cut, the di-Higgs channel including the $h \rightarrow b\bar{b}$ branching ratio, only gives $\sigma < 2.5$ fb cross section, which is suppressed compared to other backgrounds. The signal as compared to background shows distinct features in the distributions of different kinematic variables. In Fig. 13, we show p_T distributions of the two fatjets, for two mass points of the BSM Higgs $M_{H_2} = 1.1, 1.5$ TeV. As is clearly seen in the figure the p_T distributions of the two leading jets for the signal and backgrounds are not very well separated. The peak of the p_T distributions for the first and the second jets

occurs at a relatively high values of p_T as compared to the background. Therefore, to reduce the background without affecting the signal we demand a higher value of p_T on leading and subleading fatjets as a cut, which are $p_T(j_1) > 250$ GeV and $p_T(j_2) > 250$ GeV.

In Fig. 14, the leading and subleading fatjet masses have been displayed. For the signal the two jets are produced from the decay of the SM-like Higgs, hence the peaks for the distribution of $M(j_1)$ and $M(j_2)$ occur around the Higgs mass ($M_{H_1} \sim 125$ GeV). Since these fatjets are formed after showering and hadronization of $b\bar{b}$ pair, therefore we expect a nontrivial two prong substructure inside each of the fatjets. We use the Soft Drop algorithm [92] which uses

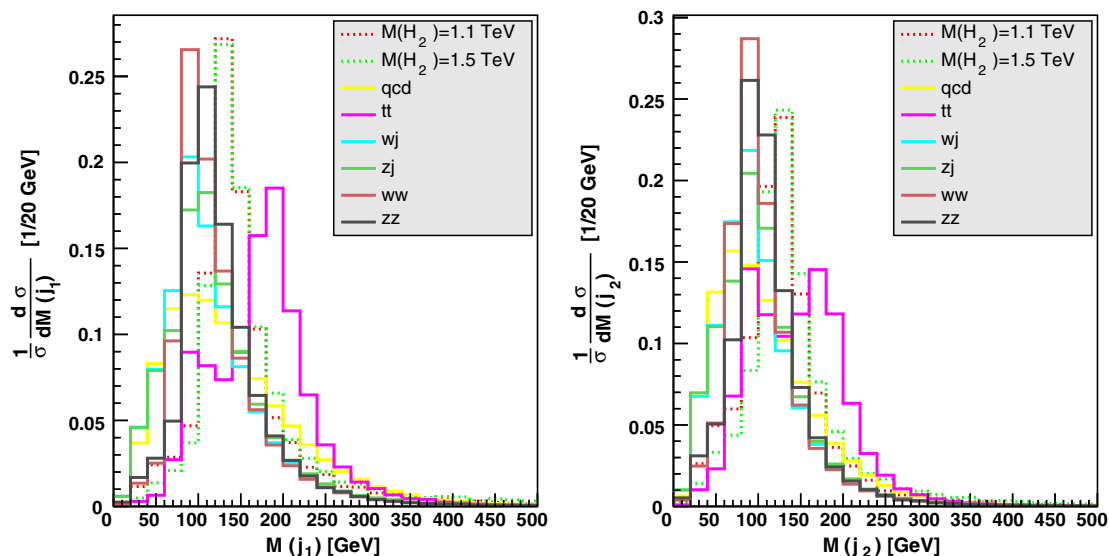


FIG. 14. Left panel: invariant mass distribution of the leading fatjet j_1 . Right panel: the same for the subleading fatjet j_2 .

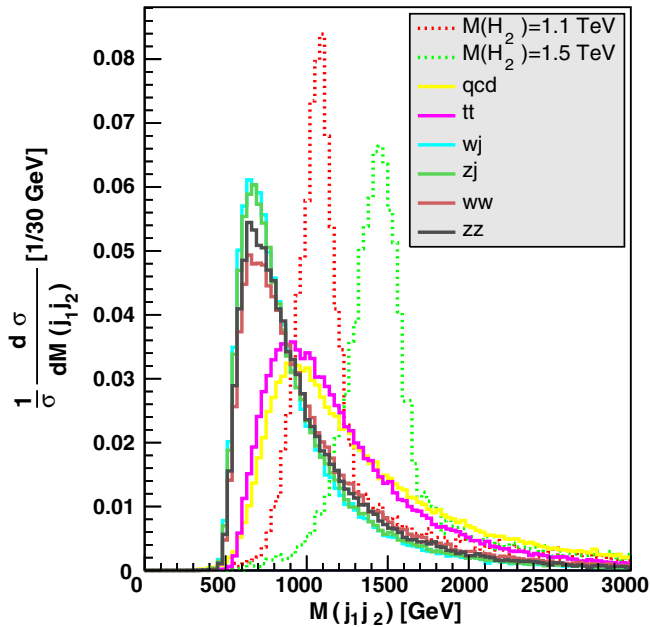


FIG. 15. Invariant mass distribution of the fatjet pair.

the condition $\frac{\min(p_T(\text{subject1}), p_T(\text{subject2}))}{(p_T(\text{subject1}) + p_T(\text{subject2}))} > 0.1$ to determine whether subjects are created from Higgs decays. All subjects which satisfy this condition are qualified as the subjects originating from Higgs decay. In Delphes a subject can not be tagged as b -jet. We implement a naive b -jet tagging for subjects in our analysis. We use B -hadron to tag the subject originating from the b quark. We consider b -tag efficiency for the subject is 80% with mistag efficiency 1%. In Fig. 15, we show the invariant mass distribution of the fatjet pairs,

TABLE IV. Background cross section before and after the cuts.

BG	σ^{partonic} [fb]	σ^{analysis} [fb] ($950 < M(j_1 j_2) < 1250$)	σ^{analysis} [fb] ($1350 < M(j_1 j_2) < 1650$)
QCD	4.1479×10^7	1789.9	211
$t\bar{t}$	7.603×10^4	0.03	9.9×10^{-5}
Wj	5.311×10^4	0.0018	0.43
Zj	5.89×10^4	1.9	0.006
WW	1.2815×10^2	2×10^{-6}	3×10^{-7}
ZZ	3.614×10^1	1.3×10^{-6}	5×10^{-8}

TABLE V. Signal ($pp \rightarrow 2j_{\text{fat}}$) and background cross sections after different selection cuts at $\sqrt{s} = 100$ TeV for $\sin\theta = 0.34(0.1)$.

	$M_{H_2} = 1.1$ TeV		$M_{H_2} = 1.5$ TeV	
	σ^s [fb]	σ^b [fb]	σ^s [fb]	σ^b [fb]
Before cut	36.22 (3.13)	4.17×10^7	8.64 (0.75)	4.17×10^7
After cut	0.745 (0.064)	1791.9	0.19 (0.016)	211.43
$\frac{\sigma^s \sqrt{\mathcal{L}}}{\sqrt{\sigma^s + \sigma^b}}, \mathcal{L} = 30 \text{ ab}^{-1}$	3.05 (0.26)		2.26 (0.19)	

the signal peaks around M_{H_2} . We use these features to reduce the backgrounds while not reducing too much the signal. Therefore our selection cuts are the following:

- c_1 : We demand at least two fatjets in the final state, $N_j \geq 2$.
- c_2 : Bound on the leading and subleading fatjets p_T are $p_T(j_1) \geq 250$ GeV and $p_T(j_2) \geq 250$ GeV.
- c_3 : The mass of leading and subleading fatjets must be within 20 GeV of the SM Higgs mass, $|M_{H_1} - M_{j_{1,2}}| \leq 20$ GeV.
- c_4 : The invariant mass of the two fatjets will deviate at most by 150 GeV from the BSM Higgs mass, $|M_{H_2} - M(j_1 j_2)| \leq 150$.
- c_5 : Pseudorapidity separation between j_1 and j_2 , $|\Delta\eta(j_1 j_2)| \leq 1.5$.
- c_6 : The leading and subleading fatjets must contain at least two subjects.
- c_7 : For the leading and subleading fatjets, each of the fatjets will contain two b -tagged subjects.

We show the partonic cross sections of different SM backgrounds in Table. IV. As can be seen, the main background is QCD, with a cross-section $\sigma_{\text{QCD}} \sim 10^7$ fb at the partonic level. The other backgrounds, such as $t\bar{t}$, Wj , Zj have a cross section $\sigma \sim 10^4$ fb. In the third and fourth column, we show the cross sections of the backgrounds after implementing all the cuts $c_1 - c_7$.

We also checked that for a resolved analysis with standard set of cuts that: a) number of b jet ≥ 4 , b) $p_T(b) > 30$ GeV, c) invariant mass of Higgs $m_{bb} = 125 \pm 20$ GeV, d) b -tagging efficiency same as the fatjet analysis, and e) invariant mass of $4b$ jet similar to the fatjet analysis, the QCD cross section is large $\sigma = 4878$ fb (2795 fb) for $M_{H_2} = 1.1$ TeV (1.5 TeV). For the resolved analysis, additional background contributions such as $W + 2j$, $Z + 2j$, jbb , and others are also relevant.

In Table V, we show the signal cross sections before and after applying the cuts $c_1 - c_7$. We consider two illustrative BSM Higgs masses $M_{H_2} = 1.1, 1.5$ TeV. The background in this table corresponds to the QCD background, shown in Table. IV, as this is the major background. As can be seen from the table the background is huge as compared to the signal. Applying the cuts however allows to improve the significance of the signal. The main remaining background is QCD . We find that for $M_{H_2} = 1.1$ TeV, one can achieve 3σ significance of the signal over background for 30 ab^{-1}

luminosity. For higher-mass values of the BSM Higgs, significance reduces. The significance of the signal can be improved over the background if one uses multivariate analysis and neural network methods. Additional final states such as $b\bar{b}\gamma\gamma/b\bar{b}\tau^+\tau^-$ are expected to give better significance, as these are clean channels. Detailed evaluation of the discovery prospect of all these channels is beyond the scope of this paper, and we will present this elsewhere.

V. CONCLUSION

In this work, we adopt an effective field theory framework that contains RHN and one SM gauge singlet scalar. Our model accommodates a FIMP DM candidate and explains the observed eV masses of light neutrinos, where SM neutrinos acquire their masses via the seesaw mechanism. Three gauge singlet RHN states $N_{1,2,3}$ and one gauge singlet real scalar χ are present in addition to the SM fields. The two RHN states $N_{1,2}$ participate in light neutrino mass generation and N_3 is the DM. There is a sizeable mixing between the BSM scalar and the SM Higgs, that offers better detection prospect of the BSM Higgs at colliders.

The FIMP DM candidate N_3 in our model interacts with the SM and BSM scalars only via effective $d = 5$ Yukawa interaction. The $d = 5$ term generates the trilinear interaction term responsible for decay once χ and Φ acquire vev . Hence FIMP DM can be produced from the decay of the SM and BSM Higgs. Annihilation of scalars and other SM particles can also lead to DM production. However, for a low reheating temperature, the decay contribution dominates. In our analysis, we therefore first consider a low reheating temperature and analyze only decay contributions. In *Scenario-I and II*, where there is no bare mass term of $N_{1,2,3}$ being added, both the DM mass and its interaction with other particles depend on the same operator. Therefore, the relic density constraint leads to a strong correlation between the vev of BSM scalar (v_χ) and the mass of DM (M_{N_3}). Keeping other parameters fixed, the required value of v_χ to satisfy the observed relic density increases with the mass of DM. The same v_χ also primarily governs the BSM Higgs mass M_{H_2} . Since in our model SM and BSM Higgs mixing can be sizeable, a TeV scale or lighter H_2 therefore has a better discovery prospect at collider as compared to a very heavy H_2 . We find that, for the TeV scale v_χ which is a natural choice for TeV scale or lower BSM Higgs states, the DM relic density constraint is satisfied only if its mass is in the KeV range. We also consider another scenario (*Scenario-III*) where we accommodate a bare mass term of the RHN states. We find that in this case, the tight correlation between vev of χ and mass of DM is somewhat relaxed, and a GeV scale DM is possible to accommodate with a TeV scale BSM Higgs/ vev v_χ .

We also consider a variation of the reheating temperature T_R and study the different annihilation channels. For a high reheating temperature we consider both the decay and annihilation contributions in relic density, where the latter dominates the relic abundance. A number of annihilation channels $WW/ZZ/H_iH_i \rightarrow N_3N_3$ can give significant contributions. In our analysis, we show the variation of relic density with respect to various parameters, such as v_χ , the mass of DM, and the reheating temperature. We find that the relic density increases with the mass of DM, and T_R (for gauge boson and scalar annihilation only), and decreases for higher $vevs$ of the BSM scalar. Assuming BSM Higgs varying in $\mathcal{O}(\text{TeV})$ range, we vary these parameters in a wide range and show the variation of relic density as scatter plot.

Finally, we explore the collider signature of the TeV scale BSM scalar at the 100 TeV future pp machine. We consider the production of the BSM scalar which has sizeable mixing with the SM Higgs, and its decay to a pair of SM Higgs states. We further consider the decay of the SM Higgs to $b\bar{b}$ states. For a TeV scale heavy Higgs, the SM Higgs is rather moderately boosted leading to collimated decay products. We consider difatjet final states as our model signature. We perform a detailed analysis considering several backgrounds, such as, QCD, $t\bar{t}$, WW/ZZ , $W + 1j$, $Z + 1j$. Following a cut based analysis we find that a 3σ significance can be achieved for a 1.1 TeV BSM scalar with 30 ab^{-1} luminosity. Thus, the difatjet channel which is sensitive to the trilinear Higgs coupling $H_2H_1H_1$ is a complementary probe for the heavy BSM Higgs, in addition to other channels, such as $pp \rightarrow H_2 \rightarrow WW/ZZ$.

ACKNOWLEDGMENTS

G. B. and M. M. acknowledge the support from the Indo-French Centre for the Promotion of Advanced Research (CEFIPRA/IFCPAR, Grant No 6304-2). M. M. thanks DST INSPIRE Faculty research Grant No. (IFA-14-PH-99). This work used the Scientific Compute Cluster at GWDG, the joint data center of Max Planck Society for the Advancement of Science (MPG) and University of Göttingen. S. S. and R. P. acknowledge the support of the SAMKHYA: High Performance Computing Facility provided by IOPB. M. M., R. P., and S. S. thank Dr. Shankha Banerjee for useful discussions on di-Higgs searches.

APPENDIX

In this section we discuss various expressions of the annihilation contributions to relic density.

1. Analytical expression of relevant cross sections

We list the cross sections for the $AB \rightarrow N_3N_3$ processes where A, B are any SM particles contributing to DM production in the freeze-in mechanism.

(a) $WW \rightarrow N_3 N_3$:

$$\begin{aligned}
g_{H_1 W^+ W^-} &= \frac{M_W e \cos \theta}{s_w}, \\
g_{H_2 W^+ W^-} &= \frac{M_W e \sin \theta}{s_w}, \\
A_{WW} &= \frac{g_{H_1 W^+ W^-} g_{H_1 N_3 N_3}}{(s - M_{H_1}^2) + iM_{H_1} \Gamma_{H_1}} + \frac{g_{H_2 W^+ W^-} g_{H_2 N_3 N_3}}{(s - M_{H_2}^2) + iM_{H_2} \Gamma_{H_2}}, \\
M_{WW} &= \frac{2}{9} \left(1 + \frac{(s - 2M_W^2)^2}{8M_W^4} \right) (s - 4M_{N_3}^2) |A_{WW}|^2, \\
\sigma_{WW \rightarrow N_3 N_3} &= \frac{1}{16\pi s} \sqrt{\frac{s - 4M_{N_3}^2}{s - 4M_W^2}} M_{WW}.
\end{aligned} \tag{A1}$$

(b) $ZZ \rightarrow N_3 N_3$:

$$\begin{aligned}
g_{H_1 ZZ} &= \frac{M_W e \cos \theta}{c_w^2 s_w}, \\
g_{H_2 ZZ} &= \frac{M_W e \sin \theta}{c_w^2 s_w}, \\
A_{ZZ} &= \frac{g_{H_1 ZZ} g_{H_1 N_3 N_3}}{(s - M_{H_1}^2) + iM_{H_1} \Gamma_{H_1}} + \frac{g_{H_2 ZZ} g_{H_2 N_3 N_3}}{(s - M_{H_2}^2) + iM_{H_2} \Gamma_{H_2}}, \\
M_{ZZ} &= \frac{2}{9} \left(1 + \frac{(s - 2M_W^2)^2}{8M_W^4} \right) (s - 4M_{N_3}^2) |A_{ZZ}|^2, \\
\sigma_{ZZ \rightarrow N_3 N_3} &= \frac{1}{32\pi s} \sqrt{\frac{s - 4M_{N_3}^2}{s - 4M_Z^2}} M_{ZZ}.
\end{aligned} \tag{A2}$$

(c) $f\bar{f} \rightarrow N_3 N_3$:

$$\begin{aligned}
g_{H_1 f\bar{f}} &= -\frac{eM_f \cos \theta}{2M_W}, \\
g_{H_2 f\bar{f}} &= -\frac{eM_f \sin \theta}{2M_W}, \\
A_{ff} &= \frac{g_{H_1 f\bar{f}} g_{H_1 N_3 N_3}}{(s - M_{H_1}^2) + iM_{H_1} \Gamma_{H_1}} + \frac{g_{H_2 f\bar{f}} g_{H_2 N_3 N_3}}{(s - M_{H_2}^2) + iM_{H_2} \Gamma_{H_2}}, \\
M_{ff} &= \frac{2}{n_c} (s - 4M_f^2) (s - 4M_{N_3}^2) |A_{ff}|^2, \\
\sigma_{f\bar{f} \rightarrow N_3 N_3} &= \frac{1}{64\pi s} \sqrt{\frac{s - 4M_{N_3}^2}{s - 4M_f^2}} M_{ff}.
\end{aligned} \tag{A3}$$

(d) $H_i H_j \rightarrow N_3 N_3$ ($i, j = 1, 2$):

$$\begin{aligned}
g_{H_1 H_1 H_1} &= -3[2v\lambda_{H_1} \cos^3 \theta + 2v_\chi \lambda_{H_2} \sin^3 \theta + \lambda_{H_1 H_2} \sin \theta \cos \theta (v \sin \theta + v_\chi \cos \theta)], \\
g_{H_1 H_1 H_2} &= [6v\lambda_{H_1} \cos^2 \theta \sin \theta - 6v_\chi \lambda_{H_2} \sin^2 \theta \cos \theta - (2 - 3\sin^2 \theta)v\lambda_{H_1 H_2} \sin \theta - (1 - 3\sin^2 \theta)v_\chi \lambda_{H_1 H_2} \cos \theta],
\end{aligned} \tag{A4}$$

$$\begin{aligned}
g_{H_2 H_2 H_2} &= 3[2v\lambda_{H_1} \sin^3\theta - 2v_\chi\lambda_{H_2} \cos^3\theta + \lambda_{H_1 H_2} \sin\theta \cos\theta(v \cos\theta - v_\chi \sin\theta)], \\
g_{H_2 H_2 H_1} &= -[6v\lambda_{H_1} \sin^2\theta \cos\theta + 6v_\chi\lambda_{H_2} \cos^2\theta \sin\theta - (2 - 3\sin^2\theta)v_\chi\lambda_{H_1 H_2} \sin\theta \\
&\quad + (1 - 3\sin^2\theta)v\lambda_{H_1 H_2} \cos\theta], \\
g_{H_1 H_1 N_3 N_3} &= \left(\frac{c'_{33}}{\Lambda} \cos^2\theta + \frac{c_{33}}{\Lambda} \sin^2\theta\right), \\
g_{H_2 H_2 N_3 N_3} &= \left(\frac{c'_{33}}{\Lambda} \sin^2\theta + \frac{c_{33}}{\Lambda} \cos^2\theta\right) \\
g_{H_1 H_2 N_3 N_3} &= \sin\alpha \cos\alpha \left(\frac{c'_{33}}{\Lambda} - \frac{c_{33}}{\Lambda}\right), \\
M_{H_i H_j} &= \frac{g_{H_i H_j H_1} g_{H_1 N_3 N_3}}{(s - M_{H_1}^2) + iM_{H_1} \Gamma_{H_1}} + \frac{g_{H_i H_j H_2} g_{H_2 N_3 N_3}}{(s - M_{H_2}^2) + iM_{H_2} \Gamma_{H_2}} + g_{H_i H_j N_3 N_3}, \\
\sigma_{H_i H_j \rightarrow N_3 N_3} &= \frac{1}{16\pi s} \sqrt{\frac{s(s - 4M_{N_3}^2)}{(s - (M_{H_i} + M_{H_j})^2)(s - (M_{H_i} - M_{H_j})^2)}} |M_{H_i H_j}|^2. \tag{A5}
\end{aligned}$$

As given in [13], we can approximately calculate the DM contribution analytically for higher value of the reheating temperature by solving the following Boltzmann equation,

$$\frac{dn_{N_3}}{dt} + 3n_{N_3}H \simeq \frac{T}{1024} \int ds d\Omega \sqrt{s} |M|_{AB \rightarrow N_3 N_3}^2 K_1 \left(\frac{\sqrt{s}}{T}\right). \tag{A6}$$

We parametrize the amplitude for the $AB \rightarrow N_3 N_3$ process to be proportional to some power of center of mass energy at very high temperature i.e., $|M|_{AB \rightarrow N_3 N_3}^2 = \alpha s^n$ where α is a constant which depends on the couplings and n is a rational number. After substituting the above amplitude and focusing on the dependence of the comoving number density on temperature, we obtain

$$\frac{dY_{UV}}{dT} = \kappa T^{2n-2}, \tag{A7}$$

where Y_{UV} is the co-moving number density of DM and κ is a constant. In the present work the amplitude varies in the following way at large s ,

$$\begin{aligned}
|M|^2 &= \alpha_{AB} s \quad \text{where } A, B = W, Z, H_1, H_2 \\
&= \alpha_{ff} \quad \text{where } f \text{ is SM fermion.} \tag{A8}
\end{aligned}$$

Therefore, for gauge bosons and Higgses we can easily show that

$$Y_{UV} = \kappa T_R, \tag{A9}$$

and for the fermion it will be

$$Y_{UV} = \kappa \left(\frac{1}{T_0} - \frac{1}{T_R}\right). \tag{A10}$$

Finally, we can conclude that for the Higgs bosons and gauge boson the relic density contribution increases linearly with T_R , whereas for fermions the contribution does not change significantly with the T_R for high value of T_R .

-
- [1] Z. Maki, M. Nakagawa, and S. Sakata, Remarks on the unified model of elementary particles, *Prog. Theor. Phys.* **28**, 870 (1962).
[2] B. Pontecorvo, Neutrino experiments and the problem of conservation of leptonic charge, *Sov. Phys. JETP* **26**, 984 (1968).
[3] I. Esteban, M. C. Gonzalez-Garcia, M. Maltoni, T. Schwetz, and A. Zhou, The fate of hints: Updated global analysis of

- three-flavor neutrino oscillations, *J. High Energy Phys.* **09** (2020) 178.
[4] S. Weinberg, Baryon and Lepton Nonconserving Processes, *Phys. Rev. Lett.* **43**, 1566 (1979).
[5] F. Wilczek and A. Zee, Operator Analysis of Nucleon Decay, *Phys. Rev. Lett.* **43**, 1571 (1979).
[6] P. Minkowski, $\mu \rightarrow e\gamma$ at a rate of one out of 10^9 muon decays? *Phys. Lett.* **67B**, 421 (1977).

- [7] R. N. Mohapatra and G. Senjanovic, Neutrino Mass and Spontaneous Parity Violation, *Phys. Rev. Lett.* **44**, 912 (1980).
- [8] T. Yanagida, Horizontal symmetry and masses of neutrinos, *Conf. Proc.* **C7902131**, 95 (1979).
- [9] M. Gell-Mann, P. Ramond, and R. Slansky, Complex spinors and unified theories, *Conf. Proc.* **C790927**, 315 (1979).
- [10] Y. G. Kim and K. Y. Lee, The Minimal model of fermionic dark matter, *Phys. Rev. D* **75**, 115012 (2007).
- [11] Y. G. Kim, K. Y. Lee, and S. Shin, Singlet fermionic dark matter, *J. High Energy Phys.* **05** (2008) 100.
- [12] L. Lopez-Honorez, T. Schwetz, and J. Zupan, Higgs portal, fermionic dark matter, and a Standard Model like Higgs at 125 GeV, *Phys. Lett. B* **716**, 179 (2012).
- [13] L. J. Hall, K. Jedamzik, J. March-Russell, and S. M. West, Freeze-in production of FIMP dark matter, *J. High Energy Phys.* **03** (2010) 080.
- [14] J. McDonald, Thermally Generated Gauge Singlet Scalars as Selfinteracting Dark Matter, *Phys. Rev. Lett.* **88**, 091304 (2002).
- [15] A. Biswas and A. Gupta, Freeze-in production of sterile neutrino dark matter in $U(1)_{B-L}$ model, *J. Cosmol. Astropart. Phys.* **09** (2016) 044; **05**(2017) (A)01.
- [16] V. De Romeri, D. Karamitros, O. Lebedev, and T. Toma, Neutrino dark matter and the Higgs portal: Improved freeze-in analysis, *J. High Energy Phys.* **10** (2020) 137.
- [17] P. Bandyopadhyay, M. Mitra, and A. Roy, Relativistic freeze-in with scalar dark matter in a gauged $B-L$ model and electroweak symmetry breaking, *J. High Energy Phys.* **05** (2021) 150.
- [18] M. Lucente, Freeze-in dark matter within the Seesaw mechanism, [arXiv:2103.03253](https://arxiv.org/abs/2103.03253).
- [19] A. Datta, R. Roshan, and A. Sil, Imprint of seesaw on FIMP dark matter and baryon asymmetry, [arXiv:2104.02030](https://arxiv.org/abs/2104.02030).
- [20] F. Elahi, C. Kolda, and J. Unwin, Ultraviolet freeze-in, *J. High Energy Phys.* **03** (2015) 048.
- [21] S.-L. Chen and Z. Kang, On UltraViolet Freeze-in Dark Matter during Reheating, *J. Cosmol. Astropart. Phys.* **05** (2018) 036.
- [22] A. Biswas, S. Ganguly, and S. Roy, Fermionic dark matter via UV and IR freeze-in and its possible X-ray signature, *J. Cosmol. Astropart. Phys.* **03** (2020) 043.
- [23] N. Bernal, J. Rubio, and H. Veermäe, Boosting ultraviolet freeze-in in NO models, *J. Cosmol. Astropart. Phys.* **06** (2020) 047.
- [24] S. Bhattacharya and J. Wudka, Effective theories with dark matter applications, [arXiv:2104.01788](https://arxiv.org/abs/2104.01788).
- [25] B. Barman, D. Borah, and R. Roshan, Effective Theory of Freeze-in Dark Matter, *J. Cosmol. Astropart. Phys.* **11** (2020) 021.
- [26] R. T. Co, F. D’Eramo, L. J. Hall, and D. Pappadopulo, Freeze-in dark matter with displaced signatures at colliders, *J. Cosmol. Astropart. Phys.* **12** (2015) 024.
- [27] G. Bélanger *et al.*, LHC-friendly minimal freeze-in models, *J. High Energy Phys.* **02** (2019) 186.
- [28] L. Calibbi, F. D’Eramo, S. Junius, L. Lopez-Honorez, and A. Mariotti, Displaced new physics at colliders and the early universe before its first second, *J. High Energy Phys.* **05** (2021) 234.
- [29] J. M. No, P. Tunney, and B. Zaldivar, Probing dark matter freeze-in with long-lived particle signatures: MATHUSLA, HL-LHC and FCC-hh, *J. High Energy Phys.* **03** (2020) 022.
- [30] G. Bélanger, C. Delaunay, A. Pukhov, and B. Zaldivar, Dark matter abundance from the sequential freeze-in mechanism, *Phys. Rev. D* **102**, 035017 (2020).
- [31] E. Molinaro, C. E. Yaguna, and O. Zapata, FIMP realization of the scotogenic model, *J. Cosmol. Astropart. Phys.* **07** (2014) 015.
- [32] G. Bélanger, F. Boudjema, A. Goudelis, A. Pukhov, and B. Zaldivar, micROMEGAS5.0: Freeze-in, *Comput. Phys. Commun.* **231**, 173 (2018).
- [33] C. P. Burgess, M. Pospelov, and T. ter Veldhuis, The Minimal model of nonbaryonic dark matter: A singlet scalar, *Nucl. Phys.* **B619**, 709 (2001).
- [34] V. Barger, P. Langacker, M. McCaskey, M. J. Ramsey-Musolf, and G. Shaughnessy, LHC Phenomenology of an Extended Standard Model with a Real Scalar Singlet, *Phys. Rev. D* **77**, 035005 (2008).
- [35] A. M. Sirunyan *et al.* (CMS Collaboration), Combined measurements of Higgs boson couplings in proton–proton collisions at $\sqrt{s} = 13$ TeV, *Eur. Phys. J. C* **79**, 421 (2019).
- [36] N. Aghanim *et al.* (Planck Collaboration), Planck 2018 results. VI. Cosmological parameters, *Astron. Astrophys.* **641**, A6 (2020).
- [37] A. Ilakovac and A. Pilaftsis, Flavor violating charged lepton decays in seesaw-type models, *Nucl. Phys.* **B437**, 491 (1995).
- [38] R. Alonso, M. Dhen, M. B. Gavela, and T. Hambye, Muon conversion to electron in nuclei in type-I seesaw models, *J. High Energy Phys.* **01** (2013) 118.
- [39] P. Gondolo and G. Gelmini, Cosmic abundances of stable particles: Improved analysis, *Nucl. Phys.* **B360**, 145 (1991).
- [40] A. Alloul, N. D. Christensen, C. Degrande, C. Duhr, and B. Fuks, FeynRules2.0 - A complete toolbox for tree-level phenomenology, *Comput. Phys. Commun.* **185**, 2250 (2014).
- [41] G. Aad *et al.* (ATLAS Collaboration), Search for heavy resonances decaying into a pair of Z bosons in the $\ell^+\ell^-\ell^+\ell^-$ and $\ell^+\ell^-\nu\bar{\nu}$ final states using 139 fb^{-1} of proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Eur. Phys. J. C* **81**, 332 (2021).
- [42] G. Aad *et al.* (ATLAS Collaboration), Search for heavy diboson resonances in semileptonic final states in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Eur. Phys. J. C* **80**, 1165 (2020).
- [43] G. Aad *et al.* (ATLAS Collaboration), Combination of searches for Higgs boson pairs in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Phys. Lett. B* **800**, 135103 (2020).
- [44] A. M. Sirunyan *et al.* (CMS Collaboration), Search for resonant pair production of Higgs bosons decaying to bottom quark-antiquark pairs in proton-proton collisions at 13 TeV, *J. High Energy Phys.* **08** (2018) 152.
- [45] G. Aad *et al.* (ATLAS Collaboration), Search for the $HH \rightarrow b\bar{b}b\bar{b}$ process via vector-boson fusion production using proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *J. High Energy Phys.* **07** (2020) 108; Erratum, *J. High Energy Phys.* **01** (2021) 145.

- [46] P. Bechtle, D. Dercks, S. Heinemeyer, T. Klingl, T. Stefaniak, G. Weiglein, and J. Wittbrodt, HiggsBounds-5: Testing Higgs sectors in the LHC 13 TeV era, *Eur. Phys. J. C* **80**, 1211 (2020).
- [47] U. Baur, T. Plehn, and D. L. Rainwater, Determining the Higgs boson selfcoupling at hadron colliders, *Phys. Rev. D* **67**, 033003 (2003).
- [48] U. Baur, T. Plehn, and D. L. Rainwater, Probing the Higgs selfcoupling at hadron colliders using rare decays, *Phys. Rev. D* **69**, 053004 (2004).
- [49] A. Azatov, R. Contino, G. Panico, and M. Son, Effective field theory analysis of double Higgs boson production via gluon fusion, *Phys. Rev. D* **92**, 035001 (2015).
- [50] F. Kling, T. Plehn, and P. Schichtel, Maximizing the significance in Higgs boson pair analyses, *Phys. Rev. D* **95**, 035026 (2017).
- [51] M. J. Dolan, C. Englert, and M. Spannowsky, Higgs self-coupling measurements at the LHC, *J. High Energy Phys.* **10** (2012) 112.
- [52] A. Papaefstathiou, L. L. Yang, and J. Zurita, Higgs boson pair production at the LHC in the $b\bar{b}W^+W^-$ channel, *Phys. Rev. D* **87**, 011301 (2013).
- [53] A. J. Barr, M. J. Dolan, C. Englert, and M. Spannowsky, Di-Higgs final states augMT2ed - selecting hh events at the high luminosity LHC, *Phys. Lett. B* **728**, 308 (2014).
- [54] S. Banerjee, C. Englert, M. L. Mangano, M. Selvaggi, and M. Spannowsky, $hh + \text{jet}$ production at 100 TeV, *Eur. Phys. J. C* **78**, 322 (2018).
- [55] D. E. Ferreira de Lima, A. Papaefstathiou, and M. Spannowsky, Standard model Higgs boson pair production in the $(b\bar{b})(b\bar{b})$ final state, *J. High Energy Phys.* **08** (2014) 030.
- [56] D. Wardrope, E. Jansen, N. Konstantinidis, B. Cooper, R. Falla, and N. Norjoharuddeen, Non-resonant Higgs-pair production in the $b\bar{b}b\bar{b}$ final state at the LHC, *Eur. Phys. J. C* **75**, 219 (2015).
- [57] J. K. Behr, D. Bortoletto, J. A. Frost, N. P. Hartland, C. Issever, and J. Rojo, Boosting Higgs pair production in the $b\bar{b}b\bar{b}$ final state with multivariate techniques, *Eur. Phys. J. C* **76**, 386 (2016).
- [58] S. Banerjee, B. Batell, and M. Spannowsky, Invisible decays in Higgs boson pair production, *Phys. Rev. D* **95**, 035009 (2017).
- [59] M. J. Dolan, C. Englert, and M. Spannowsky, New physics in LHC Higgs boson pair production, *Phys. Rev. D* **87**, 055002 (2013).
- [60] J. M. No and M. Ramsey-Musolf, Probing the Higgs Portal at the LHC Through Resonant di-Higgs Production, *Phys. Rev. D* **89**, 095031 (2014).
- [61] V. Martín Lozano, J. M. Moreno, and C. B. Park, Resonant Higgs boson pair production in the $hh \rightarrow b\bar{b}WW \rightarrow b\bar{b}\ell^+\nu\ell^-\bar{\nu}$ decay channel, *J. High Energy Phys.* **08** (2015) 004.
- [62] T. Huang, J. No, L. Pernié, M. Ramsey-Musolf, A. Safonov, M. Spannowsky, and P. Winslow, Resonant di-Higgs boson production in the $b\bar{b}WW$ channel: Probing the electroweak phase transition at the LHC, *Phys. Rev. D* **96**, 035007 (2017).
- [63] J. Ren, R.-Q. Xiao, M. Zhou, Y. Fang, H.-J. He, and W. Yao, LHC search of new Higgs boson via resonant Di-Higgs production with decays into $4W$, *J. High Energy Phys.* **06** (2018) 090.
- [64] P. Basler, S. Dawson, C. Englert, and M. Mühlleitner, Di-Higgs boson peaks and top valleys: Interference effects in Higgs sector extensions, *Phys. Rev. D* **101**, 015019 (2020).
- [65] H.-L. Li, M. Ramsey-Musolf, and S. Willocq, Probing a scalar singlet-catalyzed electroweak phase transition with resonant di-Higgs boson production in the $4b$ channel, *Phys. Rev. D* **100**, 075035 (2019).
- [66] A. Adhikary, S. Banerjee, R. Kumar Barman, and B. Bhattacharjee, Resonant heavy Higgs searches at the HL-LHC, *J. High Energy Phys.* **09** (2019) 068.
- [67] L.-C. Lü, C. Du, Y. Fang, H.-J. He, and H. Zhang, Searching heavier Higgs boson via di-Higgs production at LHC Run-2, *Phys. Lett. B* **755**, 509 (2016).
- [68] A. Alves, D. Gonçalves, T. Ghosh, H.-K. Guo, and K. Sinha, Di-Higgs production in the $4b$ channel and gravitational wave complementarity, *J. High Energy Phys.* **03** (2020) 053.
- [69] M. Aaboud *et al.* (ATLAS Collaboration), Search for pair production of Higgs bosons in the $b\bar{b}b\bar{b}$ final state using proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Phys. Rev. D* **94**, 052002 (2016).
- [70] V. Khachatryan *et al.* (CMS Collaboration), Search for heavy resonances decaying to two Higgs bosons in final states containing four b quarks, *Eur. Phys. J. C* **76**, 371 (2016).
- [71] V. Khachatryan *et al.* (CMS Collaboration), Search for two Higgs bosons in final states containing two photons and two bottom quarks in proton-proton collisions at 8 TeV, *Phys. Rev. D* **94**, 052012 (2016).
- [72] G. Aad *et al.* (ATLAS Collaboration), Searches for Higgs boson pair production in the $hh \rightarrow b\bar{b}\tau\tau, \gamma\gamma WW^*, \gamma\gamma b\bar{b}, b\bar{b}b\bar{b}$ channels with the ATLAS detector, *Phys. Rev. D* **92**, 092004 (2015).
- [73] M. Aaboud *et al.* (ATLAS Collaboration), Search for pair production of Higgs bosons in the $b\bar{b}b\bar{b}$ final state using proton-proton collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *J. High Energy Phys.* **01** (2019) 030.
- [74] M. Aaboud *et al.* (ATLAS Collaboration), Search for Higgs boson pair production in the $b\bar{b}WW^*$ decay mode at $\sqrt{s} = 13$ TeV with the ATLAS detector, *J. High Energy Phys.* **04** (2019) 092.
- [75] M. Aaboud *et al.* (ATLAS Collaboration), Search for Resonant and Non-Resonant Higgs Boson Pair Production in the $b\bar{b}\tau^+\tau^-$ Decay Channel in pp Collisions at $\sqrt{s} = 13$ TeV with the ATLAS Detector, *Phys. Rev. Lett.* **121**, 191801 (2018); Erratum, *Phys. Rev. Lett.* **122**, 089901 (2019).
- [76] M. Aaboud *et al.* (ATLAS Collaboration), Search for Higgs boson pair production in the $WW^{(*)}WW^{(*)}$ decay channel using ATLAS data recorded at $\sqrt{s} = 13$ TeV, *J. High Energy Phys.* **05** (2019) 124.
- [77] M. Aaboud *et al.* (ATLAS Collaboration), Search for Higgs boson pair production in the $\gamma\gamma b\bar{b}$ final state with 13 TeV pp collision data collected by the ATLAS experiment, *J. High Energy Phys.* **11** (2018) 040.
- [78] M. Aaboud *et al.* (ATLAS Collaboration), Search for Higgs boson pair production in the $\gamma\gamma WW^*$ channel using pp

- collision data recorded at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Eur. Phys. J. C* **78**, 1007 (2018).
- [79] Q.-H. Cao, G. Li, B. Yan, D.-M. Zhang, and H. Zhang, Double Higgs production at the 14 TeV LHC and a 100 TeV pp collider, *Phys. Rev. D* **96**, 095031 (2017).
- [80] J. Chang, K. Cheung, J. S. Lee, C.-T. Lu, and J. Park, Higgs-boson-pair production from gluon fusion at the HL-LHC and HL-100 TeV hadron collider, *Phys. Rev. D* **100**, 096001 (2019).
- [81] M. L. Mangano, G. Ortona, and M. Selvaggi, Measuring the Higgs self-coupling via Higgs-pair production at a 100 TeV p - p collider, *Eur. Phys. J. C* **80**, 1030 (2020).
- [82] J. Park, J. Chang, K. Cheung, and J. S. Lee, Measuring the trilinear Higgs boson self-coupling at the 100 TeV hadron collider via multivariate analysis, *Phys. Rev. D* **102**, 073002 (2020).
- [83] S. Borowka, C. Duhr, F. Maltoni, D. Pagani, A. Shivaji, and X. Zhao, Probing the scalar potential via double Higgs boson production at hadron colliders, *J. High Energy Phys.* **04** (2019) 016.
- [84] S. Banerjee, M. Mitra, and M. Spannowsky, Searching for a heavy Higgs boson in a Higgs-portal B-L model, *Phys. Rev. D* **92**, 055013 (2015).
- [85] C. Degrande, C. Duhr, B. Fuks, D. Grellscheid, O. Mattelaer, and T. Reiter, UFO—The universal FeynRules output, *Comput. Phys. Commun.* **183**, 1201 (2012).
- [86] J. Alwall, R. Frederix, S. Frixione, V. Hirschi, F. Maltoni, O. Mattelaer, H. S. Shao, T. Stelzer, P. Torrielli, and M. Zaro, The automated computation of tree-level and next-to-leading order differential cross sections, and their matching to parton shower simulations, *J. High Energy Phys.* **07** (2014) 079.
- [87] T. Sjöstrand, S. Ask, J. R. Christiansen, R. Corke, N. Desai, P. Ilten, S. Mrenna, S. Prestel, C. O. Rasmussen, and P. Z. Skands, An introduction to PYTHIA8.2, *Comput. Phys. Commun.* **191**, 159 (2015).
- [88] J. de Favereau, C. Delaere, P. Demin, A. Giammanco, V. Lemaître, A. Mertens, and M. Selvaggi (DELPHES 3 Collaboration), DELPHES3, A modular framework for fast simulation of a generic collider experiment, *J. High Energy Phys.* **02** (2014) 057.
- [89] M. Cacciari, G. P. Salam, and G. Soyez, FastJet user manual, *Eur. Phys. J. C* **72**, 1896 (2012).
- [90] Y. L. Dokshitzer, G. Leder, S. Moretti, and B. Webber, Better jet clustering algorithms, *J. High Energy Phys.* **08** (1997) 001.
- [91] M. Wobisch and T. Wengler, Hadronization corrections to jet cross-sections in deep inelastic scattering, in *Workshop on Monte Carlo Generators for HERA Physics (Plenary Starting Meeting)* (1998), pp. 270–279 [arXiv:hep-ph/9907280].
- [92] A. J. Larkoski, S. Marzani, G. Soyez, and J. Thaler, Soft drop, *J. High Energy Phys.* **05** (2014) 146.