Heavy dark matter, neutrino masses, and Higgs naturalness from a strongly interacting hidden sector

Mayumi Aoki,^{1,*} Vedran Brdar^{2,†} and Jisuke Kubo^{2,3,‡}

¹*Institute for Theoretical Physics, Kanazawa University, Kanazawa 920-1192, Japan*

²*Max-Planck-Institut für Kernphysik, 69117 Heidelberg, Germany*

³*Department of Physics, University of Toyama, 3190 Gofuku, Toyama 930-8555, Japan*

(Received 17 July 2020; accepted 11 August 2020; published 27 August 2020)

We consider the extension of the Standard Model (SM) with a strongly interacting QCD-like hidden sector, at least two generations of right-handed neutrinos, and one scalar singlet. Once the scalar singlet obtains a nonzero vacuum expectation value, active neutrino masses are generated through a type-I seesaw mechanism. Simultaneously, the electroweak scale is generated through the radiative corrections involving these massive fermions. This is the essence of the scenario that is known as the “neutrino option” for which the successful masses of right-handed neutrinos are in the range 10^7 – 10^8 GeV. The main goal of this work is to scrutinize the potential to accommodate dark matter in such a realization. The dark matter candidates are Nambu-Goldstone bosons which appear due to the dynamical breaking of the hidden chiral symmetry. The mass spectrum studied in this work is such that masses of Nambu-Goldstone bosons and the singlet scalar exceed those of right-handed neutrinos. Having the masses of all relevant particles several orders of magnitude above $\mathcal{O}(\text{TeV})$, the freeze-out of dark matter is not achievable, and, hence, we turn to alternative scenarios, namely, freeze-in. The Nambu-Goldstone bosons can interact with particles that are not in the SM but, however, have non-negligible abundance through their not-too-small couplings with the SM. Utilizing this, we demonstrate that the dark matter in the model is successfully produced at a temperature scale where the right-handed neutrinos are still stable. We note that the lepton number asymmetry sufficient for the generation of observable baryon asymmetry of the Universe can be produced in right-handed neutrino decays. Hence, we infer that the model has the potential to simultaneously address several of the most relevant puzzles in contemporary high-energy physics.

DOI: [10.1103/PhysRevD.102.035026](https://doi.org/10.1103/PhysRevD.102.035026)

I. INTRODUCTION

Despite being a great success, the Standard Model (SM) has several shortcomings. It predicts all three species of neutrinos to be massless, in contrast to the observation from neutrino oscillation experiments. It also does not contain a particle that can be a viable dark matter (DM) candidate as well as a successful mechanism for generation of baryon asymmetry of the Universe. Many of the proposed theories beyond the Standard Model (BSM) predict degrees of freedom at very high scales which could, for instance, address the aforementioned issues. However, such an approach has its own difficulties: The increment of the

mass scale at which heavy particles reside leads to the larger Higgs mass correction through processes at loop level, and this leads to a feature dubbed the hierarchy problem. In this paper, we will try to address most of the aforementioned issues simultaneously, within a framework that is a well-motivated SM extension. The core is connection between the generation of light neutrino masses and electroweak scale via processes involving right-handed neutrinos. The light neutrino masses induced via a type-I seesaw [1–4] are

$$m_\nu \simeq \frac{y_\nu^2 v_h^2}{m_N}, \quad (1)$$

where y_ν represents lepton portal Yukawa coupling, $v_h = 246$ GeV, and m_N stands for the right-handed neutrino mass. If the right-handed neutrinos are heavy, they can significantly contribute to the correction to the Higgs mass term, $\mathcal{L}_{\text{SM}} \subseteq -\mu_H^2 H^\dagger H$, in higher orders in perturbation theory, where H is the SM Higgs doublet. At one loop, the correction (finite term) is given by [5–8]

*mayumi.aoki@staff.kanazawa-u.ac.jp

†vbrdar@mpi-hd.mpg.de

‡kubo@mpi-hd.mpg.de

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$$|\Delta\mu_H^2| \sim \frac{y_\nu^2 m_N^2}{4\pi^2}, \quad (2)$$

where $\sim$ stands to indicate that the correction depends on the renormalization scale at which it is evaluated. The idea of the “neutrino option” [9] is based on the assumption that Eq. (2) is the main source for the Higgs mass term, i.e., $\mu_H^2 \simeq \Delta\mu_H^2$. Simultaneous solution of Eq. (1) and $\mu_H^2 \simeq \Delta\mu_H^2$ reveals m_N between 10 and 100 PeV, with $y_\nu \lesssim 10^{-4}$ [9] (see also Refs. [10,11]). The neutrino option thus establishes a link between the heavy right-handed neutrinos and the electroweak scale.

We are naturally led to the desire to embed the neutrino option into a classically scale-invariant theory [12], because the Higgs mass term, which is the only scale-symmetry-violating term in the SM Lagrangian, is assumed to be absent or extremely suppressed for the neutrino option to be sensible: We would like to understand the origin of the electroweak scale as a consequence of spontaneous (dynamical) scale symmetry breaking [13–15]. To be complete, we also have to understand the origin of the right-handed neutrino mass m_N in this manner. The simplest way is to realize $m_N \sim \langle S \rangle$ [15], where S stands for a SM singlet real scalar field. That is, the Majorana mass term for the right-handed neutrino N_R is replaced by the Yukawa interaction between S and N_R [15]. Along this line, a conformal UV completion of the neutrino option has been performed in Ref. [16]: The mass squared correction $\Delta\mu_H^2$ transmutes indeed into a radiative correction to the dimensionless coupling λ_{HS} before the spontaneous scale symmetry breaking [16], where λ_{HS} is the coupling for $S^2 H^\dagger H$, and it is assumed that its one-loop correction $\sim y_\nu^2 y_M^2 / 16\pi^2$ is the main source for this quartic coupling.

The main goal of this work is to successfully embed DM while preserving the aforementioned neutrino option property. To this end, we will introduce a strongly interacting QCD-like hidden sector, in which the mass scale is generated in a nonperturbative way via condensation that breaks chiral symmetry dynamically [17–19]. Such a scenario is in contrast with the aforementioned conformal UV completion of the neutrino option [16,20], where the scale is generated perturbatively *à la* Coleman and Weinberg [13]. Realistic models employing QCD-like strong dynamics have been proposed in Refs. [21–26]. A crucial observation in the presence of QCD-like hidden sector is that it features the appearance of (quasi-)Nambu-Goldstone (NG) bosons, hidden mesons, arising through dynamical chiral symmetry breaking and that due to their stability they are good DM candidates [21–23,26–28].

We would like to bring up at this stage that the use of Eq. (2) as the main source for the Higgs mass term in a scotogenic model, coupled with a QCD-like hidden sector, has been proposed in Ref. [11], where the DM candidate sits in an additionally introduced inert Higgs doublet [29]. Furthermore, the authors of Ref. [30] have recently shown

that the neutrino option can be related to both DM and baryon asymmetry of the Universe through inflation. In this work, we will take a complementary path and demonstrate successful embedding of the hidden mesons as DM in the model that does not rely on inflaton decay, although features related to primordial physics will be discussed primarily in the context of the QCD-like hidden sector that we expect to be rather cold. This is a new aspect compared to the previous models considered in Refs. [21–28].

While we will not study leptogenesis [31,32] for the purpose of generating successful baryon asymmetry of the Universe, we point out that our DM production is fully consistent with resonant leptogenesis [33–35] that was previously shown to be successful for the neutrino option [36,37]. Namely, DM production precedes right-handed neutrino decays, and in addition any relevant washout processes involving particles from a hidden sector are suppressed by the heavy mass scale.

The paper is organized as follows. In Sec. II, we introduce the set of considered BSM fields and interactions. In Sec. II A, we present techniques employed for the analysis of the strongly interacting sector, namely, the Nambu–Jona-Lasinio (NJL) description. Furthermore, in Secs. II B and II C, we elaborate on mass spectrum and DM interactions, respectively. Section III is reserved for our DM analysis: We start by discussing the thermodynamics of the hidden sector and then move on toward relevant Boltzmann equations and calculations of thermally averaged quantities (Sec. III A). Finally, in Sec. III B, we demonstrate success of the DM analysis and present results chiefly in the form of numerical scans. In Sec. IV, we conclude.

II. THE MODEL

As we have announced in Sec. I, we consider the model that is scale invariant at high energies. Instead of the Coleman-Weinberg mechanism [13] for symmetry breaking, we employ a strongly interacting, QCD-like hidden sector to generate a mass scale [21–23,26,27,38,39]. In contrast to the realization in Ref. [16], we introduce only one singlet scalar, S , which is coupled to right-handed neutrinos but also interacts with the vectorlike hidden fermions ψ_i ($i = 1, \dots, n_f$) belonging to the fundamental representation of $SU(n_c)_H$. The relevant part of the Lagrangian reads

$$\begin{aligned} \mathcal{L} \supseteq & \frac{1}{2} \partial_\mu S \partial^\mu S + \frac{i}{2} \bar{N}_R \not{\partial} N_R - \frac{1}{2} \text{Tr} F^2 \\ & + \text{Tr} \bar{\psi} (i \gamma^\mu \partial_\mu + g_H \gamma^\mu G_\mu - \mathbf{y} S) \psi \\ & - \frac{1}{2} y_M S N_R^T C N_R - \left(y_\nu \bar{L} \tilde{H} \frac{1}{2} (1 + \gamma_5) N_R + \text{H.c.} \right) \\ & - V_{\text{tree}}(H, S). \end{aligned} \quad (3)$$

Here, H ($\tilde{H} = i\sigma_2 H^*$) and L are the SM Higgs and lepton doublet, respectively, and N_R denotes the

right-handed neutrinos ($C\bar{N}_R^T = N_R$), whereas G_μ is the gauge field for the hidden QCD ($F^{\mu\nu}$ denotes the corresponding field strength tensor). The generation indices of the SM sector are suppressed in Eq. (3). Strictly speaking, the Yukawa couplings y_ν and y_M should be matrices in the generation space. However, we will not consider the flavor structure, but instead y_ν and y_M will be representative real numbers. This is justifiable, because the detailed matrix structure of these couplings is not relevant for the dark matter analysis that will be performed in this work. The $\mathbf{y}$ is an $n_f \times n_f$ Yukawa coupling matrix which can be taken as a diagonal matrix without loss of generality, i.e., $\mathbf{y} = \text{diag}(y_1, \dots, y_{n_f})$, where the diagonal entries y_i are assumed to be positive. The tree-level scalar potential V_{tree} in Eq. (3) reads

$$V_{\text{tree}}(H, S) = \lambda_H (H^\dagger H)^2 + \frac{1}{4} \lambda_S S^4 + \frac{1}{2} \lambda_{\text{HS}} S^2 (H^\dagger H), \quad (4)$$

where $H^T = (H^+, (h + iG^0)\sqrt{2})$ is the SM Higgs doublet field, with H^+ and G^0 as the would-be NG fields, and the (tree-level) stability conditions for the scalar potential are given by

$$\lambda_H > 0, \quad \lambda_S > 0, \quad 2\sqrt{\lambda_H \lambda_S} + \lambda_{\text{HS}} > 0. \quad (5)$$

In the strongly interacting hidden sector, described by the non-Abelian gauge theory based on $SU(n_c)_H$, the $SU(n_c)_H$ invariant chiral bilinear dynamically forms a $U(n_f)$ invariant condensate $\langle \bar{\psi}_i \psi_j \rangle \propto \delta_{ij}$, and at the same time the real scalar singlet S acquires a vacuum expectation value (VEV), $v_S = \langle S \rangle$. Sequentially, the right-handed neutrinos become massive with a Majorana mass $m_N = y_M v_S$; this allows for the generation of the loop-induced Higgs mass term that is at the core of the neutrino option idea [9].

Since the Higgs mass m_h , the VEV of the Higgs field h (denoted as v_h), and also the mass scale of light neutrinos are experimentally fixed, the parameter space in the model is strongly constrained. Namely, as already introduced in Eqs. (1) and (2), m_h and v_h fix the value of the product $y_\nu^2 m_N^2$, while the seesaw mechanism constrains y_ν^2 / m_N^2 . It turns out that the successful parameter space is between 10 and 100 PeV for m_N , while $y_\nu \lesssim 10^{-4}$. Without loss of generality, in our analysis we will stick to the value of $m_N = 5 \times 10^7$ GeV and the corresponding value of y_ν with which the scale of light neutrinos is fixed to $m_\nu \simeq 0.1$ eV.

Following Refs. [23, 27, 28], we consider $n_f = n_c = 3$. In this case, the hidden chiral symmetry $SU(3)_L \times SU(3)_R$ is dynamically broken to its diagonal subgroup $SU(3)_V$ by the nonzero chiral condensate, which implies the existence of eight NG bosons. Strictly speaking, they are quasi-NG bosons, because the Yukawa coupling $\mathbf{y} S \bar{\psi} \psi$ breaks explicitly the chiral symmetry, such that the NG bosons acquire a mass and as we will show in this work can become a cold DM candidate; these fields are stable due to the remnant

unbroken $SU(3)_V$ flavor group or its subgroup (depending on the choice of y_i). To simplify the situation, we assume in the following discussion that $y = y_1 = y_2 = y_3$ so that the unbroken flavor group is $SU(3)_V$.

Consequently, the free parameters of the model are

$$g_H, \quad \lambda_S, \quad \lambda_{\text{HS}}, \quad y_M, \quad \text{and} \quad y. \quad (6)$$

The Higgs portal λ_{HS} has to be tiny, because such coupling contributes to the Higgs mass term even at tree level as $\lambda_{\text{HS}} v_S^2$, and the idea is to generate the scale through radiative threshold corrections involving right-handed neutrinos instead of employing the scalar sector for that purpose. We require $\lambda_{\text{HS}} v_S^2 \lesssim m_h^2$, which reads

$$\lambda_{\text{HS}} \lesssim 6 \times 10^{-12} y_M^2, \quad (7)$$

for considered right-handed neutrino scale. Instead of setting this coupling around the value given in Eq. (7), we will conservatively study the $\lambda_{\text{HS}} \rightarrow 0$ limit, or, in other words, we will not rely on the smallness of this parameter to contribute to the generation of DM abundance.

The relevant system to investigate DM contains the NG boson fields ϕ^a ($a = 1, \dots, 8$) which are DM candidates in our model, S and N_R . We denote the mass of ϕ^a by m_{DM} and that of S by $m_S \simeq 3\lambda_S v_S$. The dynamics of the system is of a nonperturbative nature. In the following subsections, we will be using an effective theory, the Nambu–Jona-Lasinio theory, to compute the chiral condensate and m_{DM} and also to describe the interactions between ϕ^a and S . The couplings relevant for our DM analysis in Sec. III are y , y_M , and λ_S .

A. Nambu–Jona-Lasinio description

In order to analyze the strongly interacting hidden sector described by

$$\mathcal{L}_H = -\frac{1}{2} \text{Tr} F^2 + \text{Tr} \bar{\psi} (i\gamma^\mu \partial_\mu + g_H \gamma^\mu G_\mu - yS) \psi,$$

$$\text{with } y = y_1 = y_2 = y_3 \quad \text{and} \quad n_c = n_f = 3, \quad (8)$$

we replace the Lagrangian $\mathcal{L}_H$ (8) by the Nambu–Jona-Lasinio (NJL) Lagrangian that serves as an effective Lagrangian for the dynamical chiral symmetry breaking [17, 18]. It reads

$$\begin{aligned} \mathcal{L}_{\text{NJL}} = & \text{Tr} \bar{\psi} (i\gamma^\mu \partial_\mu - yS) \psi + 2G \text{Tr} \Phi^\dagger \Phi \\ & + G_D (\det \Phi + \text{H.c.}), \end{aligned} \quad (9)$$

where

$$\Phi_{ij} = \bar{\psi}_i (1 - \gamma_5) \psi_j = \frac{1}{2} \sum_{a=0}^8 \lambda_{ji}^a [\bar{\psi} \lambda^a (1 - \gamma_5) \psi], \quad (10)$$

with $\lambda^0 = \sqrt{2/3} \mathbf{1}$ and λ^a ($a = 1, \dots, 8$) that are Gell-Mann matrices. The dimensionful parameters G and G_D have

canonical dimensions of -2 and -5 , respectively. In order to deal with the nonrenormalizable Lagrangian (9), we work in the self-consistent-mean-field (SCMF) approximation of Refs. [40,41]. The mean fields σ and ϕ_a ($a = 0, \dots, 8$) are defined in the ‘‘Bardeen-Cooper-Schrieffer’’ (BCS) vacuum as

$$\sigma\delta_{ij} = -4G\langle\bar{\psi}_i\psi_j\rangle, \quad \phi_a = -2iG\langle\bar{\psi}\gamma_5\lambda^a\psi\rangle, \quad (11)$$

where the CP -even mean fields corresponding to the nondiagonal elements of $\langle\bar{\psi}_i\psi_j\rangle$ are suppressed; they do not play any role in our analysis. The NJL Lagrangian $\mathcal{L}_{\text{NJL}}$ is split into two parts as $\mathcal{L}_{\text{NJL}} = \mathcal{L}_{\text{MFA}} + \mathcal{L}_I$, where $\mathcal{L}_I$ is normal ordered with respect to the BCS vacuum, i.e., $\langle 0|\mathcal{L}_I|0\rangle = 0$, while $\mathcal{L}_{\text{MFA}}$ is computed in the SCMF approximation:

$$\begin{aligned} \mathcal{L}_{\text{MFA}} = & \text{Tr}\bar{\psi}(i\partial - M)\psi - i\text{Tr}\bar{\psi}\gamma_5\phi\psi - \frac{1}{8G}\left(3\sigma^2 + 2\sum_{a=1}^8\phi_a\phi_a\right) \\ & + \frac{G_D}{8G^2}\left(-\text{Tr}\bar{\psi}\phi^2\psi + \sum_{a=1}^8\phi_a\phi_a\text{Tr}\bar{\psi}\psi + i\sigma\text{Tr}\bar{\psi}\gamma_5\phi\psi + \frac{\sigma^3}{2G} + \frac{\sigma}{2G}\sum_{a=1}^8(\phi_a)^2\right) \end{aligned} \quad (12)$$

with $\phi = \sum_{a=1}^8\phi_a\lambda^a$. Here ϕ_0 has been suppressed¹ and M is given by

$$M(S, \sigma) = \sigma + yS - \frac{G_D}{8G^2}\sigma^2. \quad (13)$$

The one-loop effective potential obtained from $\mathcal{L}_{\text{MFA}}$ (12) can be obtained by integrating out the hidden fermions:

$$V_{\text{NJL}}(S, \sigma) = \frac{3}{8G}\sigma^2 - \frac{G_D}{16G^3}\sigma^3 - 3n_c I_0(M, \Lambda_H). \quad (14)$$

Here the function I_0 is given by

$$\begin{aligned} I_0(M, \Lambda) = & \frac{1}{16\pi^2}\left[\Lambda^4\ln\left(1 + \frac{M^2}{\Lambda^2}\right) \right. \\ & \left. - M^4\ln\left(1 + \frac{\Lambda^2}{M^2}\right) + \Lambda^2 M^2\right], \end{aligned} \quad (15)$$

with a four-dimensional momentum cutoff Λ ; we denote the cutoff in the hidden sector by Λ_H . For a certain interval of the dimensionless parameters $G^{1/2}\Lambda_H$ and $(-G_D)^{1/5}\Lambda_H$, we have $v_S = \langle\sigma\rangle \neq 0$ and $v_\sigma = \langle S\rangle \neq 0$ [23,27,28]. In view of Eq. (11), it is implied that the dynamics of the hidden sector creates a nonvanishing chiral condensate $\langle 0|\bar{\psi}\psi|0\rangle \neq 0$. The actual value of Λ_H can be indirectly fixed by the right-handed neutrino mass, because v_S is fixed from $m_N = y_M v_S$ for a given m_N and y_M : λ_{HS} is no longer an active portal coupling for the neutrino option and, hence, does not influence the value of Λ_H . Furthermore, one can see that the potential $V_{\text{NJL}}(S, \sigma)$ is asymmetric in σ by inspecting the last term in the NJL Lagrangian (9) as well as

the constituent mass M (13); due to the latter, chiral phase transition can become of first order.

It is noted that the mean fields σ and ϕ_a are non-propagating classical fields at the tree level. Therefore, their kinetic terms are generated by integrating out the hidden fermions at the one-loop level, which will be seen in Sec. II B where two-point functions are calculated.

The NJL parameters for the hidden QCD sector are obtained by scaling up the values of G , G_D , and the cutoff Λ from QCD hadron physics. Following Refs. [23,27,28], we employ the following dimensionless combinations:

$$G^{1/2}\Lambda_H = 1.82, \quad (-G_D)^{1/5}\Lambda_H = 2.29, \quad (16)$$

which are satisfied for the SM hadrons; hence, we assume that these relations remain unchanged for $\Lambda_H \gg 200$ MeV.

B. Mass spectrum

Once the VEVs $\langle\sigma\rangle$ and $\langle S\rangle$ (before the electroweak symmetry breaking) are obtained, the scalar mass spectrum can be calculated from the corresponding two-point functions at one loop with the hidden fermions. The VEVs can be computed from

$$V_{\text{S+NJL}}(S, \sigma) = \frac{1}{4}\lambda_S S^4 + V_{\text{NJL}}, \quad (17)$$

where V_{NJL} is given in Eq. (14). The CP -even scalars σ and S mix with each other; we will, however, neglect the mixing and use the tree-level mass for S :

$$m_S^2 = 3\lambda_S v_S^2 \quad (v_S = \langle S\rangle). \quad (18)$$

To obtain the mass of DM candidate(s), we have to compute one-loop diagrams of Fig. 1 to obtain [23]

¹Because of the chiral $U(1)$ anomaly, ϕ_0 is not a NG boson and is not stable.

$$\Gamma_{\text{DM}}(p^2) = -\frac{1}{2G} + \frac{G_D \langle \sigma \rangle}{8G^3} + \left(1 - \frac{G_D \langle \sigma \rangle}{8G^2}\right)^2 2n_c I_{\phi^2}(p^2, \langle M \rangle, \Lambda_H) + \frac{G_D}{G^2} n_c I_V(\langle M \rangle, \Lambda_H), \quad (19)$$

where

$$\langle M \rangle = \langle \sigma \rangle + y \langle S \rangle - \frac{G_D}{8G^2} \langle \sigma \rangle^2, \quad (20)$$

and the loop functions are given by

$$I_V(m, \Lambda) = \int_{\Lambda} \frac{d^4 k}{i(2\pi)^4} \frac{m}{(k^2 - m^2)} = -\frac{1}{16\pi^2} m \left[\Lambda^2 - m^2 \ln \left(1 + \frac{\Lambda^2}{m^2} \right) \right], \quad (21)$$

$$\begin{aligned} I_{\phi^2}(p^2, m, \Lambda) &= \int_{\Lambda} \frac{d^4 k}{i(2\pi)^4} \frac{\text{Tr}(\not{k} - \not{p} + m) \gamma_5 (\not{k} + m) \gamma_5}{((k-p)^2 - m^2)(k^2 - m^2)} \\ &= \frac{1}{4\pi^2} \left[\Lambda^2 + (p^2/2 - m^2) \ln(1 + \Lambda^2/m^2) - \sqrt{(4m^2 - p^2)p^2} \arctan \frac{1}{\sqrt{(4m^2/p^2 - 1)}} \right. \\ &\quad \left. + \frac{(2\Lambda^2 + 4m^2 - p^2)}{\sqrt{(4\Lambda^2/p^2 + 4m^2/p^2 - 1)}} \arctan \frac{1}{\sqrt{(4\Lambda^2/p^2 + 4m^2/p^2 - 1)}} \right]. \end{aligned} \quad (22)$$

Then we can calculate the DM mass from

$$\Gamma_{\text{DM}}(m_{\text{DM}}^2) = 0 \quad (23)$$

and the wave function renormalization constant Z from

$$Z^{-1} = \left. \frac{d\Gamma_{\text{DM}}(p^2)}{dp^2} \right|_{p^2=m_{\text{DM}}^2}. \quad (24)$$

C. Dark matter coupling with S

The diagrams in Figs. 2 and 3 show dark matter interactions with the singlet S . We set the external momenta to zero in order to obtain their local form (see Ref. [28]):

$$\mathcal{L}_{\phi S} = \frac{1}{4} G_{\phi^2 S^2} \phi_a \phi_a S^2 + \frac{1}{2} G_{\phi^2 S} \phi_a \phi_a S, \quad (25)$$

where

$$G_{\phi^2 S^2} = -Z n_c \left(1 - \frac{G_D}{8G^2} \langle \sigma \rangle \right)^2 y^2 [2I_{\phi^2 S^2}^A(\langle M \rangle, \Lambda_H) + I_{\phi^2 S^2}^B(\langle M \rangle, \Lambda_H)] - \frac{1}{2} Z n_c \left(\frac{G_D}{8G^2} \right) y^2 I_{\phi^2 S^2}^C(\langle M \rangle, \Lambda_H), \quad (26)$$

$$G_{\phi^2 S} = -2Z n_c \left(1 - \frac{G_D}{8G^2} \langle \sigma \rangle \right)^2 y I_{\phi^2 S}^A(\langle M \rangle, \Lambda_H) - 2Z n_c \left(\frac{G_D}{8G^2} \right) y I_{\phi^2 S}^B(\langle M \rangle, \Lambda_H), \quad (27)$$

and $\langle M \rangle$ and Z are given in Eqs. (20) and (24), respectively. The loop functions are given by

$$\begin{aligned} I_{\phi^2 S^2}^A(m, \Lambda) &= (-1) \int_{\Lambda} \frac{d^4 l}{i(2\pi)^4} \frac{\text{Tr}(\not{l} + m) \gamma_5 (\not{l} - \not{p}' + m) (\not{l} + \not{p} - \not{k} + m) (\not{l} + \not{p} + m) \gamma_5}{(l^2 - m^2)((l-p')^2 - m^2)((l+p-k)^2 - m^2)((l+p)^2 - m^2)} \Big|_{p=p'=k=0} \\ &\quad + (k \leftrightarrow k') = -\frac{1}{2\pi^2} [2 - \ln(\Lambda^2/m^2) + O(m^2/\Lambda^2)], \end{aligned} \quad (28)$$

$$\begin{aligned} I_{\phi^2 S^2}^B(m, \Lambda) &= (-1) \int_{\Lambda} \frac{d^4 l}{i(2\pi)^4} \frac{\text{Tr}(\not{l} + m) (\not{l} + \not{k}' + m) \gamma_5 (\not{l} + \not{p} - \not{k} + m) (\not{l} + \not{p} + m) \gamma_5}{(l^2 - m^2)((l+k')^2 - m^2)((l+p-k)^2 - m^2)((l+p)^2 - m^2)} \Big|_{p=k=k'=0} \\ &\quad + (k \leftrightarrow k') = -\frac{1}{2\pi^2} [-1 + \ln(\Lambda^2/m^2) + O(m^2/\Lambda^2)], \end{aligned} \quad (29)$$

$$I_{\phi^2 S^2}^C(m, \Lambda) = (-1) \int_{\Lambda} \frac{d^4 l}{i(2\pi)^4} \frac{\text{Tr}(\not{l} - \not{k}' + m)(\not{l} + m)(\not{l} + \not{k} + m)}{((l+k)^2 - m^2)(l^2 - m^2)((l-k')^2 - m^2)} \Big|_{k'=k=0} + (k \leftrightarrow k')$$

$$= \frac{m}{2\pi^2} [5 - 3 \ln(\Lambda^2/m^2) + O(m^2/\Lambda^2)], \quad (30)$$

$$I_{\phi^2 S}^{2A}(m, \Lambda) = (-1) \int_{\Lambda} \frac{d^4 l}{i(2\pi)^4} \frac{\text{Tr}(\not{l} + \not{p} + m)\gamma_5(\not{l} + m)\gamma_5(\not{l} - \not{p}' + m)}{((l+p)^2 - m^2)(l^2 - m^2)((l-p')^2 - m^2)} \Big|_{p=p'=0} + (p \leftrightarrow p')$$

$$= \frac{m}{4\pi^2} [-1 + \ln(\Lambda^2/m^2) + O(m^2/\Lambda^2)], \quad (31)$$

$$I_{\phi^2 S}^B(m, \Lambda) = (-1) \int_{\Lambda} \frac{d^4 l}{i(2\pi)^4} \frac{\text{Tr}(\not{l} + \not{p} + \not{p}' + m)(\not{l} + m)}{((l+p+p')^2 - m^2)(l^2 - m^2)} \Big|_{p=p'=0}$$

$$= \frac{1}{4\pi^2} [\Lambda^2 + 2m^2 - 3m^2 \ln(\Lambda^2/m^2) + m^2 O(m^2/\Lambda^2)]. \quad (32)$$

III. DARK MATTER AND FREEZE-IN

The quasi-NG bosons in the hidden sector are our DM candidates primarily due to their stability stemming from the unbroken $SU(3)_V$ flavor symmetry. However, they are very heavy as shown in the left panel in Fig. 4; there the masses of m_{DM} (solid lines) and m_S (dashed lines) are shown with respect to the absolute scale of the hidden sector that is fixed by $v_S = m_N/y_M = 5 \times 10^7/y_M$ GeV. As we see from this panel, m_{DM} is larger than 2.5×10^6 GeV for $y \gtrsim 10^{-6}$, $\lambda_S \gtrsim 10^{-5}$, and $y_M \lesssim 0.1$. In our

FIG. 1. One-loop diagrams for Γ_{DM} .

FIG. 2. One-loop diagrams for $G_{\phi^2 S}$.

FIG. 3. One-loop diagrams for $G_{\phi^2 S}$.

scans we will focus on the mass spectrum $m_{\text{DM}} \simeq m_S > m_N$.

Clearly, such NG bosons are too heavy for a conventional freeze-out at the measured value of DM number density [42] (for freeze-out production in the presence of a hidden strongly interacting sector, see [43,44]). The situation at hand is actually very similar to that considered in Refs. [45,46]: X (Y) in Ref. [46] would correspond to our ϕ (S). However, our S is not “highly decoupled” from the SM sector; it decays with the lifetime $\tau_S \lesssim 4 \times 10^{-24}$ s for $y_M, \lambda_S > 10^{-5}$, and this should be contrasted with the lifetime of $Y \sim \mathcal{O}(1)$ s [45,46]. The bottom line is that our scenario is not compatible with either conventional freeze-out or its derivatives. Hence, we will therefore explore the possibility of the freeze-in mechanism [47]. Since our dark sector is at masses $\gtrsim 10^7$ GeV, the temperatures where the freeze-in is most efficient will be of similar order.

An important question to ask is whether the hidden sector is thermalized. If so, then it undergoes the chiral phase transition at the critical temperature T_C , which is about $10^{-1} \times \Lambda_H$ [23,48]. In passing, we mention that $T_C \sim 0.1$ GeV in QCD with nearly massless quarks [49–53]. In the right panel in Fig. 4, we show Λ_H in the same parameter space as for the left panel discussed above. We see that, in contrast to m_{DM} , the smaller y is, the larger Λ_H becomes; it is larger than $\sim 5 \times 10^7$ GeV for $y \lesssim 0.1$, $\lambda_S \gtrsim 10^{-5}$, and $y_M \lesssim 0.1$. So, the critical temperature T_C would be higher than $\sim 5 \times 10^6$ GeV. If the dark sector temperature reaches T_C during the evolution of the Universe, the hidden sector undergoes a phase transition and NB Goldstones appear in the theory. Notice that at $T > T_C$ such particles are not even present in the model.

If, on the other hand, the hidden sector is not thermalized above T_C , there will be no chiral phase transition; the chiral symmetry is dynamically broken all the way to the present, and the hidden sector is in the confining phase all the time. In this case, our DM candidate is present from the

FIG. 4. Left: m_{DM} and m_S shown against the Yukawa coupling y , where the masses are normalized by $m_N/y_M = v_S$. The solid and dashed lines stand for $m_{\text{DM}}/(m_N/y_M)$ and $m_S/(m_N/y_M)$, respectively. Right: Λ_H/v_S against the Yukawa coupling y . In both panels, we show curves corresponding to $\lambda_S = 10^{-1}$, 10^{-3} , and 10^{-5} .

beginning. Its number density would be zero if the hidden sector was completely disconnected with the SM, because the SM sector could not heat up the hidden sector in such a situation.² We recall that there is no direct contact between the hidden sector and the SM sector; they are connected only indirectly via mediator S , which has contact with the SM sector through the Higgs portal interaction λ_{HS} , which we previously argued to be very tiny.

An indirect contact of S with the SM sector is established by the Yukawa interaction y_M , while the right-handed neutrino N_R is in direct contact with the SM thermal bath due to the Yukawa interaction y_ν which is fixed to $\mathcal{O}(10^{-4})$ in order to reproduce light neutrino masses. Therefore, the connection of the hidden sector to the SM sector is considerably suppressed. In the absence of inflation-scale

effects, let us further mention that the thermalization of S would yield the thermalization of the hidden sector. We have, however, found that this is not satisfied for the considered range of parameters as will be shown in Sec. III B. Hence, in what follows, we proceed with the assumption that the hidden sector is cold and, in particular, not reheated during inflation.

A. Boltzmann equations and thermally averaged cross sections

The relevant physical degrees of freedom to be considered are ϕ , S , and N_R as well as the SM fields which are in thermal equilibrium. The set of Boltzmann equations to be solved reads

$$\begin{aligned} \frac{dY_\phi}{dx} = & - \left[\left(\frac{\pi g_*}{45} \right)^{1/2} \frac{\mu M_{\text{PL}}}{x^2} \right] \left\{ \langle \sigma(\phi\phi; H^\dagger H) v \rangle (Y_\phi Y_\phi - \bar{Y}_\phi \bar{Y}_\phi) \right. \\ & \left. + \langle \sigma(\phi\phi; SS) v \rangle \left(Y_\phi Y_\phi - \frac{Y_S Y_S}{\bar{Y}_S \bar{Y}_S} \bar{Y}_\phi \bar{Y}_\phi \right) + \langle \sigma(\phi\phi; N_R N_R) v \rangle \left(Y_\phi Y_\phi - \frac{Y_{N_R} Y_{N_R}}{\bar{Y}_{N_R} \bar{Y}_{N_R}} \bar{Y}_\phi \bar{Y}_\phi \right) \right\}, \end{aligned} \quad (33)$$

$$\begin{aligned} \frac{dY_S}{dx} = & - \left[\left(\frac{\pi g_*}{45} \right)^{1/2} \frac{\mu M_{\text{PL}}}{x^2} \right] \left\{ \langle \sigma(SS; H^\dagger H) v \rangle (Y_S Y_S - \bar{Y}_S \bar{Y}_S) \right. \\ & \left. + \langle \sigma(SS; N_R N_R) v \rangle \left(Y_S Y_S - \frac{Y_{N_R} Y_{N_R}}{\bar{Y}_{N_R} \bar{Y}_{N_R}} \bar{Y}_S \bar{Y}_S \right) - 8 \langle \sigma(\phi\phi; SS) v \rangle \left(Y_\phi Y_\phi - \frac{Y_S Y_S}{\bar{Y}_S \bar{Y}_S} \bar{Y}_\phi \bar{Y}_\phi \right) \right\} \\ & - \left[\left(\frac{90}{8\pi^3 g_*} \right)^{1/2} \frac{x M_{\text{PL}}}{\mu^2} \right] \left\{ \langle \Gamma(S; N_R N_R) \rangle \left(Y_S - \frac{Y_{N_R} Y_{N_R}}{\bar{Y}_{N_R} \bar{Y}_{N_R}} \bar{Y}_S \right) + \langle \Gamma(S; H^\dagger H) \rangle (Y_S - \bar{Y}_S) \right\}, \end{aligned} \quad (34)$$

²Assuming that there exists no other mechanism that can heat up the hidden sector. This should be contrasted to the assumption of Ref. [46] that the SM and hidden sectors are both thermally populated during postinflation reheating and they maintain separate temperatures.

FIG. 5. Annihilation diagrams for ϕ , where filled circle indicates the effective coupling (26) for the upper-left diagram and (27) for the other diagrams. They contribute to $\sigma(\phi\phi; SS)$ (upper), $\sigma(\phi\phi; N_R N_R)$ (lower left), and $\sigma(\phi\phi; H^\dagger H)$ (lower right), respectively.

$$\begin{aligned}
 \frac{dY_{N_R}}{dx} = & - \left[\left(\frac{\pi g_*}{45} \right)^{1/2} \frac{\mu M_{\text{PL}}}{x^2} \right] \left\{ (\langle \sigma(N_R N_R; \tilde{H}^\dagger \tilde{H}) v \rangle + \langle \sigma(N_R N_R; \bar{L} L) v \rangle) \right. \\
 & \times (Y_{N_R} Y_{N_R} - \bar{Y}_{N_R} \bar{Y}_{N_R}) - (8/3) \langle \sigma(\phi\phi; N_R N_R) v \rangle \left(Y_\phi Y_\phi - \frac{Y_{N_R} Y_{N_R}}{\bar{Y}_{N_R} \bar{Y}_{N_R}} \bar{Y}_\phi \bar{Y}_\phi \right) - (1/3) \langle \sigma(SS; N_R N_R) v \rangle \\
 & \times \left(Y_S Y_S - \frac{Y_{N_R} Y_{N_R}}{\bar{Y}_{N_R} \bar{Y}_{N_R}} \bar{Y}_S \bar{Y}_S \right) \left. \right\} - \left[\left(\frac{90}{8\pi^3 g_*} \right)^{1/2} \frac{x M_{\text{PL}}}{\mu^2} \right] \left\{ -(1/3) \langle \Gamma(S; N_R N_R) \rangle \left(Y_S - \frac{Y_{N_R} Y_{N_R}}{\bar{Y}_{N_R} \bar{Y}_{N_R}} \bar{Y}_S \right) \right. \\
 & \left. + \langle \Gamma(N_R; LH) \rangle (Y_{N_R} - \bar{Y}_{N_R}) \right\}, \tag{35}
 \end{aligned}$$

where M_{PL} is the nonreduced Planck mass, g_* is the number of degrees of freedom fixed to 106.75, $1/\mu = 1/m_{\text{DM}} + 1/m_S + 1/m_N$, $x = \mu/T$, and we assume the spectrum $m_{\text{DM}} > m_S > m_N$. Here Y_ϕ is the number density of one ϕ divided by entropy, dubbed yield in what follows [$Y_\phi = Y_{\phi_1} = \dots = Y_{\phi_8}$ because of the unbroken $SU(3)_V$]. Similarly, Y_{N_R} is the yield of one N_R ; we also have $Y_{N_R} = Y_{N_{R1}} = Y_{N_{R2}} = Y_{N_{R3}}$. Note, however, that the internal degrees of freedom for N_R are counted as two. The yields denoted with the bar represent values in equilibrium. When solving the Boltzmann equations, we apply the initial condition where all yields are equal to zero at the starting temperature that exceeds DM mass.

The Feynman diagrams for all relevant terms in Eq. (35) are shown in Figs. 5–9.

The thermally averaged cross section is defined as [54]

$$\begin{aligned}
 \langle \sigma v \rangle &= \frac{\int d^3 p_1 d^3 p_2 (\sigma v) \exp(-E_1/T) \exp(-E_2/T)}{\int d^3 p_1 d^3 p_2 \exp(-E_1/T) \exp(-E_2/T)} \\
 &= \frac{\int d\Omega_1 d\Omega_2 dE_1 dE_2 p_1 p_2 [E_1 E_2 (\sigma v)] \exp(-E_1/T) \exp(-E_2/T)}{(4\pi \int dE_1 p_1 E_1 \exp(-E_1/T))^2}, \tag{36}
 \end{aligned}$$

FIG. 6. Annihilation diagrams for S , which contribute to $\sigma(SS; N_R N_R)$ (left and middle) and $\sigma(SS; H^\dagger H)$ (right), respectively.

FIG. 7. Annihilation diagrams for N_R , which contribute to $\sigma(N_R N_R; \tilde{H}^\dagger \tilde{H})$ (upper) and $\sigma(N_R N_R; \bar{L} L)$ (lower), respectively.

FIG. 8. Diagrams for the S decay into $N_R N_R$ (left) and $H^\dagger H$ (right).

while the thermally averaged decay rate reads

$$\langle \Gamma \rangle = \frac{\int d^3 p \Gamma \exp(-E/T)}{\int d^3 p \exp(-E/T)} = \frac{\int dE p E \Gamma \exp(-E/T)}{\int dE p E \exp(-E/T)}. \quad (37)$$

FIG. 9. Diagrams for the N_R decay into $\tilde{H} \bar{L}$ (left) and $\tilde{H}^\dagger L$ (right).

Here, E_1 and E_2 are the energies of the annihilating particles, and E is the energy of the decaying particle. Note that $E_1 E_2(\sigma v)$ can be expressed in a covariant form as [54]

$$E_1 E_2(\sigma v) = \sigma F(m), \quad \text{where } F(m) = \sigma \sqrt{s(s - 4m^2)}/2, \quad (38)$$

and s is the Mandelstam variable in units of energy squared. Since $\sigma F(m)$ is a function of s , dubbed $\tilde{\sigma}(s)$, the integration (36) cannot be carried out analytically, in general. However, if $\tilde{\sigma}(s)$ is a constant independent of s , for instance, we can perform the integration analytically. So, to proceed, we approximate the integral by replacing s by $4m_{\text{DM}}^2$ in $\tilde{\sigma}(s)$ for the DM annihilation and also by $4m_S^2$ for the S annihilation. In this way, the high-temperature behavior of the thermally averaged cross section, $\langle \sigma v \rangle \propto 1/T^2$, can be partly included in the s -wave contribution. This procedure greatly simplifies numerical evaluation of Eq. (35); note that we have checked its robustness by making an exact evaluation of Eq. (36) for one of the processes and found minor differences at temperatures where the DM production occurs. As argued above, this simplification now allows us to perform the thermal averaging analytically, and here are the computed expressions³:

$$\langle \sigma(\phi\phi; SS)v \rangle \simeq \frac{[G_{\phi^2 S^2} + 6\lambda_S v_S G_{\phi^2 S} \Delta(4m_{\text{DM}}^2)]^2}{64\pi m_{\text{DM}}^2} \left(1 - \frac{m_S^2}{m_{\text{DM}}^2}\right)^{1/2} \times \left[\frac{K_1(m_{\text{DM}}/T)}{K_2(m_{\text{DM}}/T)}\right]^2, \quad (39)$$

$$\langle \sigma(\phi\phi; H^\dagger H)v \rangle \simeq \frac{[v_S \lambda_{HS} G_{\phi^2 S} \Delta(4m_{\text{DM}}^2)]^2}{16\pi m_{\text{DM}}^2} \left[\frac{K_1(m_{\text{DM}}/T)}{K_2(m_{\text{DM}}/T)}\right]^2, \quad (40)$$

$$\langle \sigma(\phi\phi; N_R N_R)v \rangle \simeq 3 \times \frac{[G_{\phi^2 S y_M} \Delta(4m_{\text{DM}}^2)]^2}{8\pi} \left(1 - \frac{m_N^2}{m_{\text{DM}}^2}\right)^{3/2} \left[\frac{K_1(m_{\text{DM}}/T)}{K_2(m_{\text{DM}}/T)}\right]^2, \quad (41)$$

$$\langle \sigma(SS; N_R N_R)v \rangle \simeq \frac{3}{2\pi m_S^4} \{ (v_S \lambda_{S y_M})^2 - 4v_S \lambda_{S y_M}^3 m_N + 4y_M^4 m_N^2 \} \times \left(1 - \frac{m_N^2}{m_S^2}\right)^{3/2} \left[\frac{K_1(m_S/T)}{K_2(m_S/T)}\right]^2, \quad (42)$$

³We neglect below $\sigma(N_R N_R; \tilde{H}^\dagger \tilde{H})$ and $\sigma(N_R N_R; \bar{L} L)$; their contribution is proportional to $y_\nu^4 = 10^{-14}$, while $\Gamma(N_R; LH) \propto y_\nu^2 = 10^{-7}$.

$$\langle \sigma(SS; H^\dagger H) v \rangle = \frac{\lambda_{\text{HS}}^2}{16\pi m_S^2} \left[\frac{K_1(m_S/T)}{K_2(m_S/T)} \right]^2, \quad (43)$$

$$\langle \Gamma(S; H^\dagger H) \rangle = \frac{(v_S \lambda_{\text{HS}})^2}{8\pi m_S} \left[\frac{K_1(m_S/T)}{K_2(m_S/T)} \right], \quad (44)$$

$$\langle \Gamma(S; N_R N_R) \rangle = 3 \times \frac{y_M^2 m_S}{16\pi} \left(1 - \frac{4m_N^2}{m_S^2} \right)^{3/2} \left[\frac{K_1(m_S/T)}{K_2(m_S/T)} \right], \quad (45)$$

$$\langle \Gamma(N_R; LH) \rangle = \frac{m_\nu m_N^2}{8\pi v_h^2} \left[\frac{K_1(m_N/T)}{K_2(m_N/T)} \right] (v_h = 246 \text{ GeV}), \quad (46)$$

where $\Delta(4m_{\text{DM}}^2) = (4m_{\text{DM}}^2 - m_S^2)^{-1}$. In the calculation we have used

$$\int_m dE (E^2 - m^2)^{1/2} \exp(-E/T) = m T K_1(m/T), \quad (47)$$

$$\int_m dE (E^2 - m^2)^{1/2} E \exp(-E/T) = m^2 T K_2(m/T). \quad (48)$$

As we see from the above expressions, the approximate thermally averaged cross section can be obtained from the corresponding s -wave cross section by multiplying it with $[K_1(m/T)/K_2(m/T)]^2$; the final expression approaches 1 as T goes to zero and can be approximated as $m^2/4T^2$ for large T . We also note that factors $[K_1(m/T)/K_2(m/T)]$ appear in the thermally averaged decay rate for which we made an exact calculation without approximations.

$$\begin{aligned} m_N &= 5 \times 10^7 \text{ GeV}, & m_\nu &= 0.1 \text{ eV}, & y_\nu &= \sqrt{m_\nu m_N}/v_h = 2.874 \times 10^{-4}, \\ y_M &= 7 \times 10^{-4}, & y &= 1.6 \times 10^{-4}, & \lambda_S &= 2.8 \times 10^{-4}, & \lambda_{\text{HS}} &= 0, \\ v_S &= m_N/y_M = (5/7) \times 10^{11} \text{ GeV}, \end{aligned} \quad (49)$$

which gives

$$\begin{aligned} \Lambda_H &= 2.60 \times 10^{11} \text{ GeV}, & G_{\phi^2 S^2} &= -1.39 \times 10^{-8}, & G_{\phi^2 S} &= -4.16 \times 10^5 \text{ GeV}, \\ m_{\text{DM}} &= 3.44 \times 10^9 \text{ GeV}, & m_S &= 2.07 \times 10^9 \text{ GeV}, & v_\sigma &= 4.17 \times 10^{10} \text{ GeV}. \end{aligned} \quad (50)$$

The VEVs v_S and v_σ are obtained from the potential $V_{\text{S+NJL}}$ (17), where V_{NJL} is given in Eq. (14). Since their scale is fixed by Λ_H [see also Eq. (16)], we vary Λ_H to obtain $v_S = (5/7) \times 10^{11} \text{ GeV}$. Using these values for Eq. (20), we obtain the constituent mass $\langle M(S, \sigma) \rangle = 4.64 \times 10^{10} \text{ GeV}$. The DM and S masses are calculated from Eqs. (18) and (23), respectively, where the wave function renormalization constant is $Z = 3.98$. The effective couplings $G_{\phi^2 S^2}$ and $G_{\phi^2 S}$ are defined in Eq. (25) and

FIG. 10. Temperature dependence of ϕ , N_R , and S yields for the benchmark point given in Eq. (49). Corresponding equilibrium values are shown as dashed lines.

B. Results

In this section, we will explicitly demonstrate the success of DM production in the model. Before presenting the results of a scan, we illustrate our findings by showing a benchmark point for which $\Omega h^2 \simeq 0.12$ and the temperature-dependent yields of all relevant particles involved show a typical behavior. The temperature dependence of ϕ , N_R , and S yields is shown in Fig. 10 for a representative benchmark point

can be computed from Eqs. (26) and (27), where we use the approximated expressions of the loop functions [see Eqs. (28)–(32)].

In Fig. 10, one can observe that across the relevant temperature range Y_{N_R} gradually approaches the thermal equilibrium value. At smaller temperatures than those shown, N_R would decay into SM particles through a lepton portal; it corresponds to the temperature range at which DM is already produced. Such decays, while irrelevant for DM

FIG. 11. Results of our numerical scans; in each of the panels, we fix y_M . It reads 7×10^{-4} (left), 2×10^{-4} (middle), and 1.38×10^{-3} (right). The colors indicate the change of the relic density as a function of y and λ_S , while the black lines indicate several specific values of Ωh^2 . In particular, each panel features a 0.12 line that matches observation. The red star in the left panel indicates our benchmark point (see Fig. 10).

production, can, however, play a role in producing the required amount of baryon asymmetry of the Universe, as shown in Refs. [36,37]. S abundance increases and reaches maximum slightly before it decays away. Since we take very tiny portal λ_{HS} and, in addition, S is typically lighter than ϕ , the only available decay channel is into a pair of right-handed neutrinos (see Fig. 8). Most importantly, we observe how DM candidate ϕ neatly undergoes freeze-in; its abundance follows the characteristic curve for such nonthermal DM production. In particular, the strongest production occurs at the temperature of the order of the DM mass scale, i.e., $T \simeq m_{DM}$. We have checked that the dominant and, in fact, only relevant production arises from the $\phi\phi \leftrightarrow N_R N_R$ process, leaving $\phi\phi \leftrightarrow SS$ subdominant. Notice that the situation is very similar but still a little bit different with respect to the conventional freeze-in [47]: Namely, instead of producing DM through feeble interactions with SM particles that have a thermal abundance, in this case production occurs through interactions with BSM particles that happen to have non-negligible abundance due to rather strong interactions with the thermal bath (decays and inverse decays of N_R). We have also checked that the DM relic abundance Ωh^2 does not depend on the starting temperature T_0 (at which all yields are set equal to zero) as long as $T_0 > m_{DM}$ is satisfied.

In Fig. 11, we present the results of the parameter scans. As already elaborated in Sec. II, there are essentially only three relevant parameters (if m_N is fixed): y_M , λ_S , and y . In the left panel, we fix y_M to the value corresponding to our benchmark point; then $\Omega h^2 = 0.12$ is expected for values of y and λ_S given in Eq. (49) (our benchmark point is shown as a red star in the left panel). In the middle and right panels, we show the cases when y_M is fixed to 2×10^{-4} and 1.38×10^{-3} , respectively. In these two figures, we demonstrate that, for a given value of y_M , DM abundance can be relatively independent on one of the remaining two parameters: In the middle (right) panel, Ωh^2 is only mildly

dependent on λ_S (y). These are limiting cases; typically, relic abundance is a nontrivial function of both parameters and clearly also strongly dependent on y_M ; a rather mild change of this parameter yields very different values of y and λ_S for which the observed value of DM relic abundance is obtained.

Overall, we have shown that the successful production of DM is achievable in the vast portion of parameter space without employing any tuning.

IV. SUMMARY AND CONCLUSIONS

The neutrino option [9] establishes a link between the heavy right-handed neutrinos N_R and the electroweak scale. We have considered a scale-invariant realization of the neutrino option and studied the possibility of incorporating DM into the model. In contrast to Refs. [16,20], where such a realization is achieved perturbatively *à la* Coleman and Weinberg [13], we have coupled a strongly interacting QCD-like sector via a real SM singlet scalar S to the SM with three right-handed neutrinos introduced. The scale invariance is dynamically broken by the chiral condensate in the hidden sector, and the scale generated in this way is transmitted via S to the SM sector. The dynamical chiral symmetry breaking in the hidden sector produces NG bosons in the same way as in QCD. They are massive, because the Yukawa coupling of S with the hidden chiral fermions ψ explicitly breaks the chiral symmetry. Moreover, they are stable due to the unbroken vectorlike flavor symmetry [$SU(3)_V$ in our case] and, therefore, can be DM candidates. The main task of this work has been to show that they can indeed be a realistic DM with the abundance that matches observations.

Unlike previously studied models of a similar kind [21–23,26–28], in which the quartic coupling $\lambda_{HS} S^2 H^\dagger H$ is the portal for the transmitted scale, we have assumed that the Yukawa coupling $y_M S N_R^T C N_R$ is the main portal, because the neutrino option scenario would be spoiled if

(the tree-level or bare) λ_{HS} is not sufficiently suppressed. This has a considerable influence on the scale Λ_H of the hidden sector: Since the neutrino option works if the right-handed neutrino mass $m_N = y_M \langle S \rangle$ is in the range $10^7 - 10^8$ GeV, Λ_H cannot be of $\mathcal{O}(\text{TeV})$ as is the case in Refs. [21–23, 26–28]. For $y_M \lesssim 0.1$, we consequently have $\langle S \rangle \gtrsim 10^8 - 10^9$ GeV, which implies a large Λ_H . More importantly, we have found that the DM mass m_{DM} is so large [several orders of magnitude above $\mathcal{O}(\text{TeV})$] that the freeze-out of DM is not achievable to obtain a consistent DM relic abundance, and, hence, we have turned to the freeze-in mechanism [47]. For the conventional freeze-in process, the DM should be sufficiently disconnected with the thermal bath. In our case, ϕ (DM), S , and N_R interplay with each other during the freeze-in process in a nontrivial way, where, due to the assumption $\lambda_{\text{HS}} \sim 0$, ϕ and S have only an indirect contact with the SM via N_R , which has a direct contact with the SM through the Dirac-Yukawa coupling. Since m_N and y_ν are basically fixed (by the neutrino option), there are only three independent parameters left: y (for $S\bar{\psi}\psi$), y_M (for $SN_R^T CN_R$), and λ_S (for S^4). We have found that there exists a sufficient parameter space in which consistent DM relic abundance can be obtained. We emphasize that the DM in the model is successfully produced at temperature scale T where the right-handed neutrinos are still stable, i.e., $T \gtrsim m_N$.

Throughout the analysis, we have assumed that the hidden sector, described by a non-Abelian gauge theory, is not thermalized. This is because we have assumed that the hidden sector can be thermalized only through a contact

with the SM particles in the thermal bath; if S is not thermalized, the hidden sector, too, is not thermalized. We have found that in the parameter space, in which the freeze-in mechanism for DM works, the singlet S is not thermalized. That is, our hidden sector is cold and dark (“cold dark sector”). It would be an interesting question whether or not our DM scenario works if the hidden sector is thermalized. This is possible if, for instance, the hidden fermions have a nonvanishing SM $U(1)_Y$ hypercharge [27, 28]. We may address this question in a future publication.

We finally note that the lepton number asymmetry sufficient for the generation of observable baryon asymmetry of the Universe can be produced in right-handed neutrino decays as previously shown in the literature [36, 37] and such production is independent of a DM production mechanism. In combination with the generation of light neutrino masses and absence of the hierarchy problem (together with potential baryon asymmetry production from N_R decays), we have, hence, demonstrated that the considered well-motivated model can successfully and simultaneously tackle some of the most relevant puzzles in high-energy physics.

ACKNOWLEDGMENTS

The work of M. A. is supported in part by the Japan Society for the Promotion of Sciences Grant-in-Aid for Scientific Research (Grants No. 17K05412 and No. 20H00160). J. K. is partially supported by the Grant-in-Aid for Scientific Research (C) from the Japan Society for Promotion of Science (Grant No. 19K03844).

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