Cosmological Singlet Diagnostics of Neutrinophilic Dark Matter

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The standard model (SM) of particle physics is extended by adding a fundamental dark sector containing a dark matter singlet, which is coupled to the visible sector via a gauge invariant Yukawa portal and a sterile neutrino bridge. The cases of weakly interacting matter particles and feebly interacting ones are both considered. Even after the Yukawa portal is closed, the sterile neutrino bridge keeps the communication between the dark and visible sectors open, allowing for a prolonged contact between dark matter and standard model neutrinos. Such a dark sector is UV-complete and satisfies all particle physics and cosmological constraints. In particular, it leads to an acceptable formation of dark matter substructures with reasonable mass scales for the typical first protohalos.

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I.  INTRODUCTION

Although a lot is known about the properties that dark matter should possess, its exact nature is still far from being understood. The two most popular approaches are weakly interacting massive particles and warm dark matter, i.e. sterile neutrinos. Even though both of them satisfy all the required conditions for large structures, difficulties arise when considering small-scale structure formation. This leads in particular to the “missing satellite”, “cusp vs. core” and “too big to fail” problems. While warm dark matter is able to solve the issues with masses in the keV range, cold dark matter needs kinetic decouplings after big bang nucleosynthesis. Some models of cold dark matter with late decoupling between dark matter particles and the standard model are already present: dark matter with anomalous magnetic moments, leptophilic models and dark photon models. To the best of our knowledge, none of them is able to address all enduring small-scale problems and simultaneously provide UV-completeness, zero dark photon/photonic one-loop mixing, which is strongly constrained, compatibility with big bang nucleosynthesis and respect the standard model symmetries. This might be a signal that no such cold dark matter model could exist. In this work we consider a model, which behaves like cold dark matter for large structures and like warm dark matter at small scales.

This paper is organized as follows. We begin in Section II by adding a dark sector to the standard model. We describe the particle content of this dark sector and the relevant couplings and interactions. The following section is devoted to computing the cross-sections and decay rates that will play a role for determining the cosmological observables of the model. In Section IV we consider the constraints coming from particle physics and cosmology and note how they restrict the parameter set. This allows us to compute the most important cosmological observables, like relic abundance, decoupling temperature and damping scales, in Section V. By comparing these quantities with recent experimental results we give the mass spectrum of the theory in Section VI. We summarize the results and conclude in Section VII.

II.  SINGLETS IN THE DARK SECTOR

In this section, the dark sector and its coupling to the visible sector are described in detail. The particle content of the dark sector and its justification is as follows: Dark matter is assumed to be a SM-singlet represented by a Majorana fermion F with mass parameter m. We choose F to be a Majorana fermion, instead of a Dirac one, just to have a more minimal model; but, in principle, both Majorana and Dirac fermions admit similar properties during the thermal evolution of the theory and thus would be equally good candidates.

In order to couple the newly introduced F particles to the standard model, we postulate that they serve as a source for two scalar bosons S and X, with mass parameters M and M′ respectively. These scalars mediate a gauge invariant Yukawa portal between the standard model and the dark sector. Hence, the Yukawa portal is non-minimal with respect to the bosonic particle content. This non-minimality offers the possibility of introducing a mass hierarchy between the bosonic fields relative to the mass scale of the dark matter particles F, which is desirable in order to facilitate distinct decoupling scenarios as explained below.

Furthermore, we postulate that the dark matter particles solely couple to neutrinos at tree level, and not to charged leptons. This poses a potential problem, because it violates the gauge group of the left-handed fermions. For this reason, we introduce a further SM-singlet: A sterile neutrino with mass parameter M_R generated through some dark Higgs model with vacuum expectation value v_4, so that M_R ∼ v_4^2/M_F. Here F is some heavier fermion field, such that before the spontaneous symmetry breaking the fermionic current is con-
served. The sterile neutrino is represented by a right-handed fermion $n$, which is coupled to the SM-neutrinos in the usual manifest gauge invariant way introduced in $\Phi$ with coupling $y \equiv Y/v$ and $v \equiv \sqrt{2}\langle H \rangle$.

The theory should contain a Majorana mass $m_L$ for the SM-neutrino to accommodate the final observed neutrino mass, $m_\nu$, after a hybrid-type seesaw (type I+II). After diagonalizing the neutrino mass matrix, we obtain the Majorana mass eigenstates defined by the corresponding eigenstates of the $p^2$-Casimir. In abuse of notation we call them $n$ and $\nu$ with masses $M_R$ and $m_L - Y^2/2M_R$ respectively, assuming that $m_L, Y \ll M_R$. Hence, $S$ and $X$ are coupled to the sterile neutrino $\nu$ as well. Integrating out the sterile neutrino induces an effective coupling between the massive bosons of the dark sector and the SM-neutrinos. For simplicity, we include in this work only one lepton generation, the generalization to all generations is straightforward; moreover, we ignore the presence of the PMNS matrix. Let us stress already that the sterile neutrino bridge described above constitutes the longest-lived communication channel between the dark matter and the SM-neutrinos.

At this point, we include in the spectrum a new scalar field $\Phi$ subject to some $\mathbb{Z}_2$ symmetry, which we call $\mathcal{O}$, solely coupled to the singlet $S$ quadratically with coupling strength $-x < 0$. Such a scalar potential enables a first order phase transition at times after big bang nucleosynthesis and changes dynamically the on-shell mass of the $S$-boson significantly.

For ease of notation, the dark matter and the sterile neutrino are combined in a formal doublet $N \equiv (F,n)^T$. In addition, the bosons of the theory are combined in a formal triplet $B \equiv (S,\Phi,X)^T$. For the purpose of this work, we decided to remain agnostic about the origin of the initial boson and dark matter masses and, although achievable, a self-consistent completion of the dark sector is certainly beyond the scope of this paper.

After having presented the complete particle spectrum which will be considered throughout this work, we now explain in more detail the interactions between these newly introduced particles and the standard model.

The interaction between the doublet $N$ and the triplet $B$ is described by the coupling strength $g_{\pm} = g_\pm \sigma_+ + g_- \sigma_-$, where $A \in \{S,\Phi,X\}$ and $\sigma_\pm$ are the Pauli projectors acting on $N$. For the sake of the minimalism of the model, we assume that $g_\pm = g_\mp = 0$ and we define for convenience $\lambda \equiv g_\chi$ and $g \equiv g_\nu$.

After the spontaneous symmetry breaking in the dark sector, which occurs at a temperature $T_\chi$, the propagating massive field of $\Phi$, let us call it axion $a$, acquires a mass $m_a = \sqrt{2B}$. The singlet boson $S$ admits an observed mass of $M_{\text{obs}} = \sqrt{M^2 + xT^2}$, with $T$ being the quartic self-coupling strength of the scalar $\Phi$. This allows us to write the critical temperature explicitly as $T_c \approx m_a^2/\sqrt{M_{\text{obs}}} [9]$, assuming that $M^2/xT^2 \ll 1$. Note that both contributions to $M_{\text{obs}}, M^2$ and $xT^2$, are always positive.

A minimal UV-complete dark sector including the particle content discussed above is given explicitly by

$$\mathcal{L}_{\text{ds}} = \bar{N}K_{\ell}N + B^T K_B B - V(B) \; ,$$

(1)

where $2K_{\ell} \equiv i\bar{\ell}I + gB - M_{\ell}$, $2K_B \equiv -I - M_R^2$, $M_{\ell} = \text{diag}(m^{(c)}_\ell, M_R^{(c)})$ and $M_B = \text{diag}(M,0,M')$. All mass parameters are positive. Here the label $(c)$ on a mass indicates that it is a Majorana mass term.

The communication between this dark sector and the standard model is achieved through the gauge invariant Yukawa portal discussed above and the following sterile neutrino bridge: $\mathcal{L}_{\text{nh}} = igL\sigma_2 H^* n \pm h.c.$, with $y$ denoting the Yukawa coupling mentioned before. The dark sector induces a dominant effective coupling $g_\nu$ between the $S$ field and the $\nu$'s after the electroweak phase transition. This is summarized in the interaction term $g_\nu S\bar{\nu}/2$, where $(g_\nu)^{1/2} = (g_\bar{\nu})^{1/2} (\delta r + \bar{\delta} r^2)$, using the abbreviation $\delta r \equiv Y/\sqrt{2M_R}$. This mechanism accommodates naturally effective microcharges between the SM-neutrinos and the $S$-bosons.

At energies below the rest mass of the heavy sterile neutrino $n$ and the massive scalar field $X$, these fields can be integrated out, resulting in the following effective fermion sector

$$\mathcal{L}_{\text{eff}} \supset \mathcal{L}_{\text{eff}} = \bar{F}K_F F + \bar{\nu}K_{\nu}\nu \; \text{ with} \; \mathcal{L}_{\text{eff}}$$

(2)

$$2K_F \equiv i\bar{\ell} + g_\ell S - m \; \text{and} \;$$

$$2K_{\nu} \equiv i\bar{\nu} + g_\nu S + \sqrt{\delta r}/S^2 - m_\nu \; .$$

(3)

Before concluding this section, we present the proposed mass spectrum. We assume the following mass hierarchy in the dark sector:

$$M' \gg M_R > m, m_h \gg M_{\text{obs}} \gg M, m_\nu, m_a \; .$$

(4)

Such an arrangement allows various decays, like for example $n \rightarrow h\nu$, $B \rightarrow \nu\nu$, $X \rightarrow FF$, and keeps $F$ stable without requiring an additional symmetry. Note that the relation of $m_a$ with the other masses appears as a natural consequence of requiring a consistent thermal evolution of the theory. Furthermore, $M$ can be thought as a $t$-channel regulator in order to simplify the calculations.

III. CROSS-SECTIONS & DECAY RATES

This section is devoted to the presentation of cross-sections and decay rates relevant for particle physics and structure formation processes in the primordial Universe, in which the dark sector $\Phi$ participates.

It is useful to introduce the dark sector Fermi constant of the theory of neutrino scattering $G_D/\sqrt{2} \equiv g_\ell g_\nu/AM^2$. Annihilation into singlets $S$ implies a sink term in the dark matter abundance that is characterized by the $p$-wave cross-section

$$\langle v_{\text{rel}}^2 (\sigma_{\text{ann}}) \rangle = \frac{\pi \alpha^2}{2m^2} \langle v_{\text{rel}}^2 \rangle \; ,$$

(5)
where \( \langle \ldots \rangle \) denotes the thermal average using relative velocities, \( \alpha = \left( \frac{2}{g} \right) \frac{1}{\pi T} \) and \( T \) is the average photon temperature. In principle, if \( m > M_R \), \( F \)'s annihilate to \( n \)'s; however, this possibility is excluded by the assumed mass hierarchy (4).

Finally, the dark matter abundance is sourced by \( X \) decays with a rate

\[
\Gamma_{X \rightarrow FF} = \frac{\lambda^2}{16\pi} M'.
\tag{6}
\]

This concludes the brief presentation of the dominant sinks and sources for the dark matter abundance in the temperature regime \( m \gg T \gg M \).

Elastic scatterings involving dark matter particles determine their kinetic properties and allow to quantify their deviation from local thermal equilibrium in terms of viscous processes sourcing entropy production. As mentioned above, the longest running channel in this respect is the one with SM-neutrinos as scattering partners. At temperatures \( M \gg T \gg m_\nu \), and at lowest order in the Fermi constant, the averaged momentum transfer elastic cross-section, defined by \( \sigma_T \equiv \int d\Omega (1 - \cos \theta) d\sigma_{el}/d\Omega \), is given by

\[
\langle v_{rel}\sigma_T \rangle = \frac{80\zeta(5) M^2 T^2}{\pi\zeta(3)} \left( \frac{T}{T_\nu} \right)^2
\tag{7}
\]

after applying the equipartition theorem. Note that in the above expression all neutrino masses have been neglected. \( T_\nu \) denotes the neutrino temperature. The elastic scattering of \( S \) particles to neutrinos is extremely suppressed. Indeed, although one can integrate out the sterile neutrino field to obtain \( g_\nu \) for the physical neutrinos, the internal 2-point functions include sterile neutrino propagators, which are not bare propagators but must be perturbatively corrected with internal neutrinos. These rates are suppressed by a factor of \( g_\nu^4 \) for the IR dominant part.

The numerical solutions in various regimes of the momentum transfer cross-section for fermions can be found in (3) and for the classical limit, \( m v/M \gg 1 \), we cite semianalytic formulas for the corresponding total cross-sections in different kinematical regimes:

\[
\sigma_T \approx \begin{cases} 
\frac{4\pi}{M^2} \beta^2 \ln \left( 1 + \beta^{-1} \right) & , \quad \beta \in (0, 10^{-1}) \\
\frac{8\pi}{M^2} \left( 1 + 1.5 \beta^{1/2} \right) & , \quad \beta \in [10^{-1}, 10^3] \\
\frac{8\pi}{M^2} \left( \ln \beta + 1 - \frac{1}{\beta^{1/2}} \right)^2 & , \quad \beta \in (10^3, \infty)
\end{cases}
\tag{8}
\]

where \( \beta \equiv 2 v M/(m_\nu v_{rel}) \). This dependence seems to be of vital importance in order to resolve small-scale anomalies that are typically present when structures form in non-self interacting dark matter (non-SIDM). Self-interactions among the dark matter particles can be mediated via \( X \) as well. In this case, however, there is no significant Sommerfeld enhancement.

**IV. PARAMETER CONSTRAINTS**

We proceed to the discussion of the constraints on the dark sector arising from particle physics and cosmology.

**A. Constraints from particle physics**

The dark matter particles \( F \) are assumed to be stable and produced in local thermal equilibrium. Then partial-wave unitarity of the scattering matrix bounds the annihilation cross section in the primordial Universe, which in turn bounds the relic abundance and the universal mass of the dark matter (10), \( m < 0.03 \) MeV. The effective number of neutrino generations as measured by the Planck satellite implies a lower mass bound \( m > 0.1 \) MeV. Note that this bound holds precisely for the dark matter introduced in (1), since it is kept in local thermal equilibrium for an extended period via elastic scatterings with SM-neutrinos. A priori there are no such bounds on the sterile neutrino mass \( M_R \) or on the masses \( M, M' \), other than the imposed mass hierarchy (4). Lower bounds on these mass parameters do, however, arise in scenarios where the sterile neutrinos and the dark bosons are kept in local thermal equilibrium with the SM-neutrinos after their respective kinetic decoupling from other SM-species. An extensive list of a variety of constraints on leptonic-boson interactions can be found in (11).

Let us turn to the particle physics constraints on the dark sector couplings. The couplings combined in \( g \) are assumed to respect the perturbative domain. In addition the coupling in the Yukawa bridge is bounded by \( Y^2/2M \sim m_L \lessgtr 5 \) GeV (12). Accordingly, in this work we set the upper bounds \( \alpha, \tilde{\alpha} \leq 1/2 \), which in turn imply an upper bound on \( m \approx 1.5 \) TeV, as will be shown below. Note that the dark sector (11) does not facilitate a coupling between the dark bosons and charged leptons. Therefore, \( g_\nu \) is only subject to restrictions arising from the three-body decays of \( Z^0, W^\pm \) and from the \( K^\pm \) decays, although such constraints are less tight than those coming from cosmology. Furthermore, we anticipate that such constraints should not be as strong as the light scalar-neutrino interaction bounds (13), due to the absence of a longitudinal contact term.

Consider the invisible three-body decay \( Z^0 \rightarrow \nu \nu S \) with rate \( \Gamma_{Z^0\rightarrow\nu\nu S} \approx 0.18 \) GeV \( g_\nu^2(N_\nu/3) \) for an observed \( S \)-boson mass parameter of \( O(10) \) MeV. This experimental error on \( Z^0 \) decays is 0.0023 GeV (14). This allows to constrain the effective neutrino coupling by \( g_\nu \leq 0.12 \sqrt{N_\nu} \). The result does not depend strongly on \( M_{obs} \). Note that if \( S \) decays sufficiently fast, then the four-body decay \( Z^0 \rightarrow 4\nu \) is dominant, which lightens the previous constraint. The dark sector (11) allows the decay \( W \rightarrow v e S \) as well. The experimental error on the decay of charged electroweak gauge bosons is 0.042 GeV (15). For \( M_{obs} \sim O(10) \) MeV this implies \( g_\nu \lesssim O(1) \), which is not a significant constraint in our context. Fi-
nally, the constraint on the effective neutrino coupling arising from $K^\pm$ decays is more restrictive. In addition to the standard decay $K \to \nu \mu$, the channel $K \to \nu\mu S$ is available as well. For masses $M$ of order $O(1)$ MeV and energies between 165.5 MeV and 205.5 MeV for the outgoing muon, the relative rate is almost constant and can be explicitly written \cite{13} as

$$\frac{\Gamma_{K \to \nu\mu S}}{\Gamma_{K \to \nu\mu}} = 7.4 \times 10^{-4} g_\nu^2.$$ \hspace{1cm} (9)

The experimental bound on this ratio is $3.6 \times 10^{-6}$, implying a bound on the effective neutrino coupling $g_\nu \lesssim 7 \times 10^{-2}$.

Note that, contrary to the situation analyzed in \cite{13}, there is no constraint from the elastic $\nu e \to \nu e$ scattering, since the dark singlets in \cite{14} are decoupled from charged leptons at tree level. Additional bounds on $g_\nu$ due to lepton number violation and mesonic decays are computed in \cite{10}. However, they are less tight than the cosmological constraints, which will be discussed in the next section. Since $g_\nu \lesssim 0.12/\sqrt{N_L}$ and $\Gamma_{\nu e} \sim N_L g_\nu^2$, the extremal result for $\Gamma_{\nu e}$ does not depend on the actual number of neutrinos coupled effectively to the singlet boson; however, if $N_S$ denotes the number of identical $S$-bosons which are present, then $g_\nu$ scales with $1/\sqrt{N_S}$. Further bounds on self-annihilation thermal averaged cross-sections of $F$'s to neutrinos can be found in \cite{17}.

B. Constraints from cosmology

In this part, we study the constraints on the couplings $x, g_\nu$ and $\lambda$, which arise from astrophysical observables.

The desired value of $M_{\text{obs}}$ lies in the sub-GeV/multiple-MeV interval in order to solve the small-scale issues. Furthermore, $M$ should lie in the (sub-)keV range to enable a late kinetic decoupling. Therefore, we discuss only the case, where the phase transition in the dark sector takes place after the kinetic decoupling of the $F$ particles from the neutrinos; otherwise, the relevant interactions are highly suppressed. If $\phi$ breaks after $T_{\text{bd}} \sim O(10^2)$ eV, the $S$ particles will annihilate rapidly to $a$ particles, as long as $m_a \ll M_{\text{obs}}$. We demand that at the phase transition $S$ and $\Phi$ are in local thermal equilibrium. This condition implies the following constraint for the coupling strength between $S$ and $\Phi$

$$x \gtrsim 10^{-13} \left( \frac{T_c}{100 \text{eV}} \right).$$ \hspace{1cm} (10)

If the mass parameters $m_a$ and $M$ of the axion $a$ and the $S$ bosons respectively were below $O(1)$ MeV, the presence of these particles in the primordial plasma would modify big bang nucleosynthesis. Furthermore, if these mass parameters were close to the temperature of the primordial plasma at recombination, the cosmic microwave background radiation would carry a fingerprint of the dark sector. However, if the light particles of the dark sector and the neutrinos freeze-in after the end of big bang nucleosynthesis, at around 30 keV, the impact is minimized. This leads to a conservative upper bound for the effective neutrino coupling $g_\nu^2 \lesssim 10^{-12}$ using the IR-dominant part of the elastic cross-sections in the absence of a fundamental cubic scalar interaction in the dark sector. Furthermore, an important bound of $g_\nu \gtrsim 1.6 \times 10^{-6}$MeV/$M_{\text{obs}}$ arises for $M_{\text{obs}} \sim O(1)$ MeV, as computed from the luminosity and deleptonization arguments regarding the observation of SN1987A in \cite{18}. In this model values of $M_{obs} \approx 50$ MeV alleviate the cusp vs. core and the too big to fail problems together with a suitable set of $\{m_a, g_\nu\}$. Therefore, the final constraint on $g_\nu$ reads $10^{-15} \lesssim g_\nu^2 \lesssim 10^{-12}$. Taking the saturation limit of this relation, the effective neutrino coupling can be eliminated as a free parameter; hence, we consider a typical value of $g_\nu^2 = 10^{-13}$ in this work.

For simplicity we assume that the light dark sector particles are in local thermal equilibrium after big bang nucleosynthesis and before the phase transition takes place. After the condensation of $\phi$, the dark axion $a$ emerges as a massive propagating degree of freedom. The metastable massive axions contribute to the energy density of the universe by a factor of $\Omega_a h^2 \approx m_a/300$ eV. Nevertheless, the condition that $\Omega_a h^2$ be small enough is easily fulfilled, since $m_a$ lies well below the eV-scale.

Since the scalar field couples directly only to the $S$-boson, we turn our attention to the properties of the interactions between $S$, $\Phi$ and the standard model. To test the above construction, we calculate the deviation of the effective neutrino degrees of freedom, which parametrizes the energy density of the universe, in different periods of its evolution. This is encapsulated in the following definition, assuming that the neutrinos already decoupled at $T_{\nu D} = 2.3$ MeV \cite{19},

$$\Delta N_{\nu|\text{BBN}} = N_{\nu} \frac{\rho_{S+\Phi}}{\rho_\nu}.$$ \hspace{1cm} (11)

$\rho_\nu$ is the equilibrium energy density of the $i$-th particle species with initial conditions defined at $T_{\nu D}$ and including all available degrees of freedom. This number is important because it parameterizes the cosmic energy budget and can therefore be probed with high precision. Let us introduce the abbreviation $\varepsilon_{T_{\nu D}} \equiv (T_S/T_{\nu D})^3$.

Examining the parameter space we find that a common value of $\varepsilon \sim 0.1$ is compatible with big bang nucleosynthesis \cite{20} and cosmic microwave background \cite{21} 1σ measurements. Indeed one obtains $\Delta N_{\nu|\text{BBN}} \approx 0.05$ and $\Delta N_{\nu|\text{CMB}} \approx -0.02$ assuming that $T_c > T_{\nu D}$. This may also explain the recent tension about the decrease of the deviation of effective neutrino number from BBN to CMB-based measurements. A difference $\Delta N_{\nu|\text{CMB}} - \Delta N_{\nu|\text{BBN}} < 0$ is possible in the underlying theory. However, if $T_c < T_{\nu D}$ and at the same time all light degrees of freedom are not in local thermal equilibrium with neutrinos at the recombination period, then $\Delta N_{\nu|\text{CMB}} \approx 0.05$; therefore, for optimal big bang nucleosynthesis/cosmic microwave background compat-
bility the phase transition is preferred to happen before the recombination period with all light particles in local thermal equilibrium. Moreover, we make the bound less tight by demanding that $N_{\text{eff}} + \Delta N_{\text{eff}} \lesssim 3.5$ \cite{21}. This sets an upper bound to the number of singlet bosons, $N_S \lesssim 8$.

It is crucial that the $S$-boson should decouple from the primordial plasma early enough, when all or almost all the degrees of freedom of the standard model are accessible. Quantitatively this means that the $S$-boson decoupling from the rest plasma should take place before the QCD phase transition, $T_{\text{freeze-out}} > T_{\text{QCD}}$, to obtain an allowed value of $\varepsilon \lesssim 0.2$. This corresponds to a lower bound for $M_R$ with respect to the couplings of the theory. We obtain up to $\mathcal{O}(1)$ factors,

$$\left( \frac{g_*}{10^{-6}} \right) \left( \frac{\bar{g}_s}{1} \right) \left( \frac{\text{TeV}}{M_R} \right) \lesssim \mathcal{O}(1).$$

(12)

In other words, light dark scalar bosons went through kinetic decoupling from the SM-species at temperatures $T_{SD} > T_{\nu D}$ due to their IR-suppressed part of the cross-section for elastic scatterings with SM-neutrinos. For instance, $T_{SD} = \mathcal{O}(10^6) T_{\nu D}$ for $M_R = \mathcal{O}(10)$ TeV, $\bar{g}_s = 0.05$ and $g^2 \approx 10^{-13}$. In this example kinetic decoupling from the SM-species happens at $T_{SD} \approx 0.5$ TeV.

The above discussion does not involve masses larger than $T_{\nu D}$ and as longs as $\tau_S \ll \tau_{\text{BBN}}$ no particular constraints arise. It turns out that the kinetic properties of the dark sector depend on the mass of the light dark bosons relative to the neutrino decoupling temperature, as will be worked out in great detail in the next section. Anticipating distinct kinetic regimes, the dark sector will be referred to as model $DS_{\text{FIMP}}$ if $M \leq \mathcal{O}(1)$ keV $\ll T_{\nu D}$, and as model $DS_{\text{WIMP}}$ if $M > T_{\nu D}$. Let us briefly turn to $DS_{\text{WIMP}}$ and consider the modification of the effective neutrino degrees of freedom in this case,

$$\Delta N_{\text{eff}} = \frac{60}{7\pi^4} \int_{x_{\text{BBN}}}^{\infty} dz \frac{z^2 (z^2 - x_{\text{BBN}})}{\exp(z) - 1} ,$$

(13)

where $x_{\text{BBN}} = M/T_{\text{BBN}}$. As an example, for $M = 10.5$ MeV we obtain $N_{\text{eff}} = 3.15$, which matches perfectly with the value inferred from the Planck measurements \cite{21}. At the $2\sigma$-level, these also impose a lower bound for the mass of the light dark bosons of $M > 3.5$ MeV.

In other words, the light dark scalar bosons in $DS_{\text{WIMP}}$ are roughly six orders of magnitude heavier compared to those in the $DS_{\text{FIMP}}$ model. In the next section it will be shown how this renders $DS_{\text{WIMP}}$ less attractive from a phenomenological point of view.

Finally, the feebly interacting massive boson $X$ should have a decay channel into dark matter flavors after chemical decoupling at $x_f$, and this channel should close before the epoch of big bang nucleosynthesis. Using the corresponding decay rate \cite{14}, the following interval for the coupling of the massive to the dark matter can be inferred,

$$4 \times 10^{-5} \left( \frac{T_{\nu D}}{2.3 \text{ MeV}} \right) \left( \frac{g_*}{10.75} \right)^{1/4} \lesssim \left( \frac{\lambda}{3 \times 10^{-16}} \right) \left( \frac{M}{\text{PeV}} \right)^{1/2} \lesssim \left( \frac{30}{x_f} \right) \left( \frac{g_*}{108.75} \right)^{1/4} \left( \frac{m}{\text{TeV}} \right) ,$$

(14)

where $g_*$ counts the effective degrees of freedom in the primordial plasma. This concludes the discussion on coupling constraints.

V. COSMOLOGICAL OBSERVABLES

In this section we calculate the relic abundance of the dark matter candidates in the spectrum of \cite{11}, their kinetic decoupling temperature, both for $DS_{\text{FIMP}}$ and $DS_{\text{WIMP}}$, and evaluate the characteristic damping scales implied by kinetic decoupling and free streaming of the dark matter candidates together with their cross sections for self-interaction. Collisional and collision-free damping impact structure formation and determine the properties of the typical first protohalos. Finally, we investigate whether the dark sector \cite{11} allows to resolve challenges for $\Lambda$CDM posed by the formation of cosmic structure as it is observed. In this paper, we neglect a possible second period of dark matter annihilation and the possible formation of dark matter bound states \cite{2}. Alternatively, one can consider out of equilibrium dark matter production from late decays. We are only interested in investigating the possibility of the existence of a simple field theoretical model, which alleviates the large and small structure problems; furthermore, we assume zero dark matter asymmetry. We note that all temperatures are given in the photon plasma frame, unless stated differently.

A. Dark matter relic abundance

Let us first assume that the $F$ fields are the dominant dark matter population, with annihilations into light singlet bosons as the dominant sink for this population towards chemical decoupling. $\alpha$ is fixed by requiring that the observed relic density is correctly retrieved in the unbroken phase. The defining equation of the dark matter distribution function $f(p(t))$ per degree of freedom of the underlying field in a Friedmann-Robertson-Walker
universe yields
\[(L - C)[f](p) = 0 \quad (15)\]
where \(L := E(\partial_t - H \cdot \nabla_p)\) is the Liouville operator, which gives the change of \(f\) with respect to an affine parameter along a geodesic and \(C\) is the collision term. We can simplify the Boltzmann equation by taking all the participating particles up to the dark matter ones to admit equilibrium thermal distributions. The Boltzmann equation for non-relativistic chemical decoupling leads to the change of the number density \(n\) of dark matter particles per entropy density \(s\), denoted by \(Y\), with respect to the temperature \(T = m/x\) according to the Riccati-type equation,
\[
\frac{dY_F}{dx} = -\frac{E}{x^2} \left( Y_F^2 - Y_{eq}^2 \right), \quad (16)
\]
where the efficiency is given by \(E = s(x)\langle v_{rel} \sigma_{ann} \rangle / H(x) \rangle \big|_{x=1}\), and \(\ell = 3\) for \(p\)-wave annihilations, approximating massless final states, more precisely \(M/(m/x) \ll 1\). We define the chemical freeze-out value \(x_t\) by demanding that
\[
\langle v_{rel} \sigma_{ann} \rangle n_{eq}^F(x_t) \frac{1}{H(x_t)}. \quad (17)
\]
The exact solution of this equation is given in terms of the Lambert-W function, but a sufficiently good approximation can be found iteratively. Already after one iteration:
\[
x_t \approx \ln(C) - \frac{\frac{4}{5}\ln(\ln(C)) + \frac{1}{5}}{\ln^{-1}(C)},
\]
\[
C = \sqrt{\frac{45}{32}} \sqrt{\frac{g_F}{\langle v_{rel} \sigma_{ann} \rangle |_{x=1}}} \frac{M_{Pl} m}{5}.
\quad (18)
\]
This assumes that the number of relativistic degrees of freedom varies slowly with the temperature towards chemical decoupling. Inserting (18) into (16), the differential equation for \(Y_F\) can be solved numerically for the dark flavor coupling parameter \(\alpha\). Requiring that \(n_F(x_t) \gg n_{eq}^F(\infty)\), and in the absence of feebly couplings, we obtain for the relic abundance of the dark matter population as a function of \(m\) and \(\alpha\),
\[
\Omega_F h^2 \approx 0.12 \left( \frac{\alpha}{0.1} \right)^{-2} \left( \frac{m}{1 \text{ TeV}} \right)^2.
\quad (19)
\]
The self-annihilation cross section \(\sigma_{ann}\) is well below the experimental sensitivity \(\frac{1}{10}\). Note that (19) has been obtained assuming \(g \gg g_s\) in order to neglect annihilations into SM-neutrinos. We also neglected the Sommerfeld effect, which gives corrections of \(O(1)\) to the above result.

Let us stress that the constraints and the relic abundance (19) fix \(g\) completely. However, in order to resolve the missing satellite problem, typically larger values of \(\alpha\) are needed for such masses; as a consequence, a different (primary or subsidiary) mechanism to populate the dark matter is required, which is still compatible with the framework above. The only option left in the dark sector is the out of equilibrium decay of the heavy scalar \(X\). The number of \(X\) particles per entropy varies with the temperature \(T = M/x\) as described by
\[
\frac{dY_X}{dx} \approx \frac{3}{8\pi^2} \sqrt{\frac{5}{\pi}} \frac{M_{Pl} \Gamma_X \rightarrow p F}{g_{_M} \ell^3/2(T_{_F}) M_{_F}^2} x^3 K_1(x), \quad (20)
\]
where \(K_1\) is the modified Bessel function of the second kind with index one. The \(\ell = -3\) behaviour is a characteristic property of models with feebly interacting massive particles. From this we find for the \(F\) relic abundance,
\[
\Omega_F h^2 \approx 0.12 \left( \frac{\lambda}{8 \times 10^{-11}} \right)^2 \left( \frac{114}{g_{_M}(T_{_F})} \right)^{3/2} \times \left( \frac{m}{1 \text{ TeV}} \right) \left( \frac{10^5 \text{ PeV}}{M'} \right).
\quad (21)
\]
It is very promising that this scenario is perfectly compatible with the constraints on the \(\lambda\) coupling \(\left[14\right]\). Concluding, in the presence of multiple heavy bosons, smaller couplings are needed, which allow smaller dark matter masses for higher values of \(M'\) due to \(\left[14\right]\). The same can be achieved after using Dirac fermions for \(F\) and/or upgrading \(X\) to a vector boson.

**B. Decoupling temperatures**

As calculated above, chemical decoupling of the dark particles happens long before the SM-neutrinos decouple. However, this is not the case for the cease of efficient momentum exchange: The kinetic coupling of dark matter to neutrinos is active much longer. In order to have a phenomenologically viable scenario, kinetic decoupling should be scheduled for times long after big bang nucleosynthesis, but still before the 100 eV epoch, due to the constraints from the Lyman-\(\alpha\) measurements \(\left[23\right]\).

A rough but sufficiently accurate estimate for the temperature \(T_{_{Ly}}\) of the typical last elastic contact between the dark matter and the SM-neutrinos, which represents the longest active communication channel, is obtained by equating the corresponding elastic scattering rate \(\Gamma_{el} \equiv \langle v_{rel} \sigma_{el}\rangle n_{eq}^\nu\) with the Hubble expansion rate. For example, assuming three generations of SM-neutrinos and \(M = 0(1)\) keV, which fixes \(g_s = 0(10^{-7})\), we find \(T_{_{Ly}} \approx T_{eq}\), where \(T_{eq}\) denotes the temperature of matter-radiation equality. This characterizes the typical spatial hypersurface from which dark matter enters the free streaming regime.

The temperature \(T_{_{Ly}}\) at which the dark sector kinetically decouples from the standard model, when the elastic interactions between the dark matter particles and the neutrino generations cease to sustain local thermal equilibrium \(\left[24\right]\). In model \(D_{\text{SM}}\) it depends on the parameters: \(\{ g_X, M, M, g_s\} \). Hence, the temperature at kinetic decoupling is not completely fixed by the relic dark matter abundance. Moreover, these parameters allow to
accommodate other constraints, like the one on $\langle \sigma_T/m \rangle$, to address the cusp versus core problem.

We note that $T_{kd}$ is not very sensitive to small changes of $M \sim O(1)$ keV, but it changes rapidly with respect to deviations of the neutrino effective coupling. Numerically, using the rate of the averaged momentum transfer elastic cross-section as in [11] and equating it to the Hubble rate, we obtain $T_{kd} \lesssim 267$ eV and at the same time the variable set $\{g, m, M, g_\nu\}$ fulfills all previous constraints. If a larger value of $\alpha$ is needed, then the FIMP scenario comes into play and cures this issue. An elegant instantaneous kinetic decoupling can be by achieved setting the critical temperature equal to the desired $T_{kd}$.

C. Damping Scales

Eventually, around a temperature $T_{kd}$, elastic scattering processes between dark matter and SM-neutrinos happen too infrequently to sustain local thermal equilibrium. Since the neutrino abundance is much larger than the dark matter one, every non-relativistic dark matter particle will interact a huge amount of times with the relativistic neutrinos until the time of kinetic decoupling. On the other hand, hardly any neutrinos will ever interact with a dark matter particle. Around $10^{10}$ interactions between dark matter and neutrinos are needed for a $O(1)$ momentum transfer. This is easily achieved since the abundances of the two types of particles are so different from each other and since there is a long enough period of time until kinetic decoupling. For this reason, any structure smaller than $M_d = (4\pi/3)\rho_{dm}(T_{kd})/H^3(T_{kd})$ will be destroyed before $T_{kd}$. We find the following estimate of the characteristic damping mass:

$$M_d \approx 2.7 \times 10^8 \frac{g_s(T_{kd})}{g_{s,0}^{1/2}(T_{kd})} \left( \frac{\text{keV}}{T_{kd}} \right)^3 M_\odot ,$$

where $g_s(T)$ and $g_{s,0}$ denote the effective number of relativistic degrees of freedom contributing to the entropy and energy density at temperature $T$ respectively.

A relevant independent quantity is the free-streaming length of dark matter particles [22]. However, we expect that the acoustic damping is much more efficient than free streaming for very late decoupling temperatures, since the latter has a $1/T_{kd}^{3/2}$ dependence in contrast to the $1/T_{kd}$ dependence of $M_d$. The mass of the smallest possible protohalo, that could be formed, is found by taking the largest of the above two masses. A possible solution to small-scale abundance problems of CDM cosmology, i.e. the missing satellite problem, can be found by after suppressing the power spectrum at scales as large as that of dwarf galaxies [24], which is provided by damping masses of order $\log_{10} \left( M_d/M_\odot \right) \gtrsim 9$ and not larger than $5 \times 10^{10} M_\odot$ as stated in [26, 11, 27] and [23].

On the other hand, the mass scale of acoustic damping [28] is fixed at $10^{10} M_\odot$. In model $D S_{WIMP}$, the cut-off masses are much smaller, namely we approximately obtain $M_d DS_{WIMP} \approx 5 \times 10^3 M_\odot$. These values are far from promising. Although they allow protohalo masses that are not excluded by current collider and direct search constraints, they do not provide a solution to the missing satellite problem. We note that for neutralinos in the MSSM the predicted masses are around the Earth mass [29]. Such masses are well below the masses that the current numerical simulations studying the large scale structures can resolve [2].

Concluding, we notice that the model $DS_{WIMP}$ is ruled out by astrophysics, if one assumes that the astrophysical constraint of order 9 is actually a bound. This does not alleviate the enduring missing satellite problem respecting the Planck measurements as stated in [24]; in other words, at least $M \sim O(1)$ keV is needed in order to generate masses near $10^{10} M_\odot$. The FIMP scenario makes almost the whole parameter set $\{g, m\}$ accessible, hence one can always obtain $M_d DS_{WIMP} = 10^9 M_\odot$. On the contrary, a dominant WIMP production mechanism fails to deliver the necessary values of $g$. The missing satellite problem is easily solved for the WIMP case, by considering for example $m = 300$ GeV, $M = 0.1$ keV, $g_\nu^2 = 10^{-13}$ and $\alpha \approx 0.1$ or $m = 1$ TeV, $M = 10$ eV, $g_\nu^2 = 10^{-13}$ and $\alpha \approx 1/\pi$. All effective neutrino couplings are in perfect accordance with [13]. Alternatively, setting $T_c = 100$ eV, we assure that an instantaneous late kinetic decoupling at the desired temperature $T_c \approx T_{kd}$ takes place.

D. Effects of SIDM cross-sections on dwarf galaxies

The final test of the dark matter models is given here: We investigate whether this parameter set supports noticeable effects on dwarf galaxies. As we mentioned before, the SIDM elastic cross-section in galaxy clusters should lie within a constrained range of values. Quantitatively, the cross-section for elastic scatterings between the dark matter particles should lie in the interval $\langle \sigma_T/m \rangle_{\text{therm}} \in [0.1, 10]$ cm$^2$g$^{-1}$ [5] to resolve the cusp vs. core problem as shown in [29] where the indicated average is taken with respect to a Maxwell-Boltzmann distribution with $v_{\text{therm}} = O(10^{-4})$ as the most probable velocity.

Exploring the parameter space of this family of models, we find that the values of SIDM elastic cross-sections in model $DS_{WIMP}$ are of order $\sim 0.2 - 0.4$ cm$^2$g$^{-1}$. Therefore, after setting $x, \lambda$ and $\ell$ to zero without loss of generality, no further modification of the dark sector is needed. However, $O(1)$ keV mediators need either $\mu$-dark matter charges or multi-TeV values of $m$. Both possibilities are strictly excluded, as found previously, due to the previous constraints, thermal evolution or perturbation theory. Therefore, in model $D S_{\nu WIMP}$, as we mentioned before, we make use of the scalar potential, which should admit a phase transition at $T_c \lesssim T_{kd}$. This does not have an effect on the damping masses, but cures the SIDM values automatically. For the relevant subspace of parameters, the obtained values can stay always inside the constrained interval as one sees in [4]. The SIDM cross-
sections between $0.1 - 10 \text{ cm}^2 \text{g}^{-1}$ are easily accessible for a range value of sub-TeV/TeV dark matter resonances together with sub-GeV mediators. These results are in line with the maximum circular velocities observed and the values affecting dwarf galaxy scales \[3\]. Exemplary, for the previous parameter set considering the FIMP scenario with $m = 0.3 \text{ TeV}$, a value of $M_{\text{dark}} = 50 \text{ MeV}$ leads to the thermal averaged SIDM cross-section values of $\sim 0.3 - 0.6 \text{ cm}^2 \text{g}^{-1}$.

VI. THE UNDERLYING SPECTRUM

In the present section we examine the resulting spectrum of the theory, which is able to provide solutions to the \( \Lambda \)CDM problems. Besides the stable $F$, both the bosons and the sterile neutrino decay. In addition, for perturbative couplings $g_x/5 = y = 1/10$ and $g_s^2 = 10^{-13}$, the sterile neutrino admits a mass $M_R \approx 20 \text{ TeV}$ and $\delta \nu \approx 10^{-3}$. For this parameter set one finds that the $S$-singlet decouples from the SM-neutrino plasma at very high temperatures, around $m, M_R$, due to the tiny value of the effective neutrino coupling $g_\nu$; at that time almost all the degrees of freedom of the standard model are relativistic. This means that the $B$-triplet freezes-out early enough and it does not leave much time for them to achieve local thermal equilibrium with neutrinos, which relaxes the strong constraints on $M_{\text{dark}}$ from Plank measurements \[3\]. This is a unique feature of this family of models, since the $S$-mode tree level interactions include 2-point functions of $F$’s or $n$’s, which at the zero momentum limit are severely suppressed due to the large masses, $m$ and $M_R$. Furthermore, the axion mass is of order $m_a \sim \mu \text{eV}$.

Recently, the IceCube collaboration published results about extraterrestrial ultra-energetic neutrinos, i.e. Bert, Ernie \[20\] and Big Bird with approximate energies in the PeV-regime. We could easily accommodate these observations in $D_{\text{S,TIME}}$: $X$ admits a mass of $M' = 1 \text{ PeV}$ together with a minimum DM mass of $m = 300 \text{ GeV}$; additional FIMP sources, i.e. more than one heavy $X$-boson, allow lower values of $m$ due to the constraints on the $\lambda$-coupling \[14\]. We stress out that this parameter set still resolves all three sub-scale structure problems simultaneously after $T_c$, delivering damping masses and SIDM cross-sections of the desired order, together with the predicted value of $\Omega_{\text{CDM}}$.

VII. CONCLUSION

In this paper we started with an overview of the dark matter problem. We asked, whether it is possible to write down a minimal QFT as a standard model UV-complete extension, that respects all cosmological and particle physics constraints and alleviates the small-structure problems. We therefore introduced a singlet family of purely neutrinoophilic dark matter models, in order to satisfy experimental and theoretical constraints. The proposed dark sector interactions work as a fundamental extension of the standard model, while respecting the SM-symmetries. Afterwards, we computed astrophysically important observables (cross-sections and decay rates) in this framework. We presented and explained a list of constraints on the parameter space emerging from the particle physics and cosmological nature of the models. Finally, we estimated the chemical and kinetic decoupling and last scattering temperatures and we determined the relevant smallest protohalo masses. These masses are far above the Earth mass produced by the usual WIMP models. The derived $M_d$ may provide a solution to the missing satellite problem and the SIDM cross-sections could be able to alleviate the cuspy profile and massive subhalos issues of CDM simultaneously.

It is worth noticing, that in this theory the strongly constrained $\text{SU}(2)_L$ violating $\nu - S$ coupling arises naturally as an effective interaction. Furthermore, the case of a simple dominant WIMP production mechanism does not solve all the problems simultaneously and is ruled out, as we showed previously, at least in the perturbative regime of $\alpha$. Nevertheless, the FIMP scenario is able to alleviate all small-scale issues of the CDM paradigm. In addition, both dark matter production schemes provide excellent big bang nucleosynthesis/cosmic microwave background compatibility. Therefore, \( \Lambda \)CDM models are proven hard to die and seem more attractive in this neutrinoophilic perspective.

Nature supplies us with great evidence, that different forces exist for different purposes: The strong force keeps the nucleons in the nuclei together and, at the same time, the weak force makes our world as it looks, since it defines the abundances of primordial elements; moreover, the electromagnetic interactions are always present to ensure the stability of matter. To the best of our knowledge, there is no single mediator for all these forces and therefore, it is natural to proceed similarly when extending the standard model and introduce new fields to solve the various problems at hand independently. In our case, the sterile neutrino works as the gauge invariant bridge between the standard model and the dark sector connecting the known neutrino species with the stable dark matter candidate $F$ and giving rise to a particle physics-compatible effective coupling. Then the large-scale property of the correct relic abundance could be obtained also from the out of equilibrium decays of the heavy bosons $X$. The missing satellite problem is alleviated by the light singlet $S$ and its late decoupling properties, the cusp vs. core and the too big to fail issues are solved after the symmetry $\Phi$ breaks and the condensation of the scalar field $\Phi$ takes place. The simplicity and straightforwardness of this minimal model is intriguing. In future works we hope to return to this model, investigate the possibility of a more minimal construction and consider different spectrum hierarchies.
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