Small-scale structure from neutron dark decay

Georgios K. Karananas and Alexis Kassiteridis

Arnold Sommerfeld Center, Ludwig-Maximilians-Universität München, Theresienstraße 37, 80333 München, Germany
E-mail: georgios.karananas@physik.uni-muenchen.de, a.kassiteridis@physik.uni-muenchen.de

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Abstract. It was recently proposed that the disagreement in the experimental measurements of the lifetime of the neutron might be eradicated if the neutron decays to particles responsible for the dark matter in the Universe. In this paper we construct a prototype self-interacting dark matter model which, apart from reproducing the correct relic abundance, resolves all small-scale problems of the ΛCDM paradigm. The theory is compatible with the present cosmological observations and astrophysical bounds.

Keywords: dark matter theory, particle physics - cosmology connection

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1 Introduction

Since the 30’s, when Zwicky correctly claimed the existence of dark matter (DM) [1, 2], a lot of things have changed in the DM community. If one assumes that DM is made of particles with a certain mass, then there are different approaches to describe our universe: thermally produced Weakly-Interacting Massive Particles (WIMPs), Feebly-Interacting Massive Particles (FIMPs) or even warm DM (WDM), see for instance [3–5]. Towards the end of the 20th century, the sole purpose of a DM candidate was to solve the large-structure problems. However, critical questions have arisen regarding the small-structure formation: the missing satellite problem, the cuspy profiles of dark galaxies in the ΛCDM cosmology and the fact that these galaxies are too big to fail in producing luminous content; for a thorough review and possible solutions of these problems see [7–9]. The DM model-building “industry” is usually dedicated to the pursuit of theories which produce the right DM relic density, while solving all small-scale problems in the local group.

The idea that DM can be naturally accommodated within a “hidden world” that communicates with the Standard Model (SM) of particle physics via a neutron portal, is certainly not a new one. Actually, this possibility was first spelled out, in the context of braneworld scenarios, by Dvali and Gabadadze [10]. Nevertheless, owing to the fact that the aforementioned portal is omnipresent in theories with dark sectors (DS) [10], there is a number of interesting implications. For instance, in [11], it was shown that if the hidden sector involves a big number of SM copies, then the DM abundance can be naturally produced, since the particle content of the other copies are excellent DM candidates [12]. Recently, the idea that DM might be due to the interaction of the neutron with a DS has been revived and studied.

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1 WDM models are usually able to solve the abundance problem of dwarf galaxies by employing particles with masses in the keV range [6]. On the other hand, cold DM (CDM) theories require late kinetic decouplings $T \sim \text{keV}$, in order to provide a possible solution [6–8].
in great detail in [13]; various aspects of this proposal have been addressed in a number of follow-up works [14–25].

A new dark decay channel for the neutron can resolve the experimental discrepancy — which is more than 4σ — between the bottle and beam measurements of the neutron lifetime. The neutron lifetime inferred from the β-decays of free neutrons in a beam is found to be \( \tau_{n}^{\beta^{-}} \approx 888 \text{ s} \) [26]. At the same time, the direct counting of the remaining neutrons in the bottle experiment leads to \( \tau_{n} \approx 880 \text{ s} \) [27]. If yet another channel for neutron decay is present, then the correct lifetime is the one measured in the bottle experiment, meaning that its decay rate is

\[
\Gamma \equiv \frac{1}{\tau_{n}} = \Gamma_{\beta^{-}} + \Gamma_{DS},
\]

where \( \Gamma_{\beta^{-}} = 1/\tau_{n}^{\beta^{-}} \), while \( \Gamma_{DS} \) corresponds to the decay channel into the DS particles. It is straightforward to show that [13]

\[
\Gamma_{DS} \approx 10^{-5} \text{ s}^{-1}.
\]

Kinematically, the neutron decay to the DS is allowed if the sum of the masses of its decay products, \( \sum m \), is smaller than the mass of the neutron, \( m_{n} = 939.565 \text{ MeV} \). On the other hand, in order to be compatible with all present experimental constraints due to the stability of the proton, as well as nuclear transitions between stable nuclei, \( \sum m \) must be bigger than 937.900 MeV [13, 28]. Let us note that, in principle, the hidden modes could decay back into (SM) protons, therefore, we consider only

\[
\sum m < m_{p} + m_{e} + \text{min}[m_{\nu}],
\]

where \( \text{min}[m_{\nu}] \) denotes the smallest of the neutrino masses. In other words, only a very small window for the allowed mass-sum of the DM particles is open

\[
937.900 \text{ MeV} < \sum m < 938.783 \text{ MeV}.
\]

The main aim of this paper is to show that the products of the possible hidden decays of the neutron — that could successfully explain its lifetime anomaly — not only give rise to the DM relic abundance, but might be able to simultaneously provide solutions to the enduring small-scale structure problems of the ΛCDM cosmology. As a proof of concept, we construct a self-interacting dark matter (SIDM) model [29], which is compatible with all astrophysical and particle physics constraints and respects all known symmetries of the SM. To the best of our knowledge, up to now there has not appeared a theory in which the products of the hidden neutron decay are able to alleviate all the small-scale problems.

In our setup, the present DM abundance is assumed to consist of a Dirac fermion with mass in the GeV range — charged under a hidden abelian gauge symmetry — that plays the role of the stable DM candidate. Both the SIDM virtue of the model, as well as its late kinetic decoupling, is due to the secluded interaction mediated by a vector. Note that this is the only choice of fields that allows the theory to accommodate repulsive self-interactions, see for example [8, 21, 30]. It should be stressed that, contrary to what stated in [31], Yukawa self-interactions mediated by a scalar are always attractive, irrespective of whether the involved fermions are Dirac or Majorana.\

To ensure the compatibility of the model with the various constraints, it is nevertheless necessary to introduce more degrees of freedom. These comprise a heavy Dirac fermion whose

\footnote{We thank Jonathan Cornell for bringing this to our attention, which in turn led to the revision of our paper.}
mass needs to be of the order of TeV, that opens the hidden decay channel of the neutron and at the same time modifies the WIMP production of the DM particles, as well as a very light complex scalar field. Both of them should admit a baryon number, while the scalar is also charged under the hidden symmetry.

This paper is organized as follows. In section 2, we introduce the model and discuss its field content. In section 3, we compute certain important observables (decay rates and cross sections) and we give a complete list of constraints stemming from particle physics and cosmological considerations. In section 4, we study the thermal evolution of the theory by computing the present-day DM abundance and the kinetic decoupling temperature. In section 5, we examine how the small-scale problems are solved successfully in this context. Our conclusions can be found in section 6.

2 The effective theory for neutron dark decays

The proposed model is an effective theory which is valid below some cutoff scale \( \Lambda \) and connects the neutron hidden decay with an invisible sector that plays the role of DM. The most general color-singlet operators that open up such a channel read

\[
\frac{1}{\Lambda^{2+N}} \Psi_N P_s d u^T C P_s d + \text{h.c.},
\]

with \( \Psi_N \) some fermionic operator (fundamental or composite) that carries unit baryon number. Here, \( P_s \) and \( C \) denote the chiral projection and charge conjugation operators in the spinor-space, respectively, and \( s \) is running over the chiralities. To keep the notation simple, we have suppressed the SU(3)-color labels.

The thermal evolution of the theory takes place at temperatures \( T \ll \Lambda \), thus the communication of the SM with the dark sector is due to the lowest-order effective neutron portal (corresponding to \( N = 0 \)), i.e.

\[
\frac{1}{\Lambda^2} \Psi_0 P_s d u^T C P_s d + \text{h.c.}.
\]

Here, \( \Psi_0 \) is a Dirac fermion which is a singlet under the SM gauge group, but charged under the (global) baryon symmetry; in what follows we will call it heavy baryon. For our considerations, higher order terms can be safely neglected at these energies.\(^3\)

As it will become clear in a while, in order to have a phenomenologically viable DM model, the spectrum of the hidden sector, in addition to \( \Psi_0 \), should comprise: \( i \) a stable Dirac fermion \( F \) with mass \( m_F \lesssim \mathcal{O}(\text{GeV}) \), which is a singlet under the SM gauge group and does not carry baryon number, it is, however, charged under a hidden local abelian symmetry with charge \( g \); \( ii \) a complex scalar field \( \Sigma \) with mass \( m_\Sigma \sim \mathcal{O}(\text{eV}) \), unit baryon number and secluded charge \(-g\), which we call scalar baryon; \( iii \) a real vector field \( \sigma \) with mass \( m_\sigma \sim \mathcal{O}(\text{keV}) \) and no baryon number, acting as the force mediator of the hidden gauge symmetry.

In this work we are agnostic as to how the participating fields acquire their masses. It should be stated clearly that we are interested in working out a purely phenomenological model for the dark decay of the neutron, which successfully addresses the DM problems. There is no a priori reason for the mass hierarchy to be chosen like this, besides fulfilling

\(^3\)Note that a possible ultraviolet completion of the effective interaction (2.2) can certainly be found. For instance, one can introduce a colored scalar field in an appropriate representation of the SM gauge group (see e.g. [13, 28]), and mass of the order of \( \Lambda \).
the various observational constraints. For instance, the scalar baryon $\Sigma$ could, in principle, communicate directly with the SM sector through a quartic interaction with the Higgs field. However, the corresponding coupling should be finetuned to be extremely small, even for a feebly production DM mechanism; otherwise, the desired value for $m_\Sigma$ will not be generated. In any case, such interactions are irrelevant regarding the thermal evolution of the proposed model, so we will not discuss them further in the present paper.

The cosmologically relevant interaction terms appearing in the hidden sector are the following (see also eq. (B26) in [6])

\begin{equation}
L_{\text{hidden}} \supset -\lambda \left( \Sigma \bar{\Psi}_0 F + \text{h.c.} \right) + g \sigma^\mu \bar{F} \gamma_\mu F - ig \sigma^\mu \Sigma^* \partial_\mu \Sigma,
\end{equation}

with $\lambda$ and $g$ dimensionless couplings and $\Sigma^* \partial_\mu \Sigma \equiv \Sigma^* \partial_\mu \Sigma - \Sigma \partial_\mu \Sigma^*$.

Let us now turn to the justification of the terms appearing in (2.3). The first one, apart from opening up the desired dark decay channel for the neutron (if of course the mass spectrum allows it), it is also responsible for the freeze-out annihilation of the DM particles. The second, enables a SIDM scenario and gives rise to possible infrared-dominant interactions in the DM sector of the theory. Such interactions are vital for an acceptable small-scale structure formation [8]. Finally, the third term in (2.3), makes $\sigma$ potentially unstable, since they could decay mainly in scalar baryons. In other words, the scalar baryons serve as way-out particles and a possible overclosure of the universe at late times is avoided.\(^4\) For later convenience, the decay width is given by

\begin{equation}
\Gamma_{\sigma \rightarrow 2\Sigma} \approx \frac{\alpha'}{12} m_\sigma
\end{equation}

with $\alpha' = g^2/4\pi$, and assuming $m_\sigma \gg m_\Sigma$. In principle the scalar baryon can directly couple to the SM-Higgs field through a gauge invariant quadratic-scalar interaction with some coupling $\lambda_\Sigma$. However, $m_\Sigma$ should lie at the eV or sub-eV in order to avoid a possible overclosure; this leads to extremely small values for the corresponding coupling $\lambda_\Sigma$, making such interaction terms cosmologically irrelevant (even for a FIMP scenario).

The model under consideration admits all the properties of a prototype SIDM theory, due to the presence of the mediator. The SIDM cross sections per dark matter mass are repulsive [8, 30] and are strongly velocity-dependent when $2\alpha' m_\sigma/m_F v_{\text{rel}}^2 \lesssim 10^3$ [8]. This fact is in principle essential for addressing the small-scale problems that are present when matter structures form in non-interacting (purely WIMP) models. Furthermore, it is this very SIDM interaction that enables the observed neutron star formation and makes the proposed theory compatible with the latest observations.

3 Constraints from particle physics and cosmology

In this section we discuss the possible constraints on the couplings and masses of the effective theory. For an extensive list of constraints on interactions between SM-particles and DM, see [33].

\(^4\)In principle, one could also allow an additional effective coupling between $\sigma$ and the SM-neutrinos $\nu$ of the form $\sigma^\mu \bar{\nu} \gamma_\mu \gamma_5 \nu$; such couplings are discussed in [32].
3.1 Particle and astroparticle physics

The motivation behind this model is purely phenomenological and aims in connecting the existence of a hidden decay channel for the neutron to a dark sector which addresses all the DM problems. As we already mentioned, for the dark neutron decay channel to be available, we have to require that

\[ \sum m \equiv m_F + m_\Sigma \] should lie in the allowed mass interval (1.4), and that

\[ m_{\Psi_0} \gg m_F + m_\Sigma. \] (3.1)

The validity of the effective field theory approach dictates that

\[ m_{\Psi_0} < \Lambda. \] (3.2)

At temperatures below the quark confinement scale \( \Lambda_{\text{QCD}} \approx 200 \text{ MeV} \), the neutron portal (2.2) boils down to

\[ \frac{f_n}{\Lambda^2} \bar{\Psi}_0 n + \text{h.c.}, \] (3.3)

where \( f_n \approx 10^{-2} \text{ GeV}^3 \) is the neutron decay constant [34]. After diagonalizing the mass matrix and upon integrating out the heavy \( \Psi_0 \), the hidden decay channel arises from the following effective Yukawa term

\[ y \Sigma \bar{n} F + \text{h.c.}, \] (3.4)

where the dimensionless coupling is defined as \( y \equiv \lambda f_n / \Lambda^2 m_{\Psi_0} \), and in abuse of language \( n \) and \( F \) correspond to the mass eigenstates with eigenvalues \( m_n \) and \( m_F \), respectively.

From the above interaction term, the invisible two-body decay width of the neutron can be easily calculated. At the lowest order in the coupling constant \( y \), one finds

\[ \Gamma_{\text{DS}} \approx \frac{y^2}{8\pi m_n^3} \left[ (m_n + m_F)^2 - m_\Sigma^2 \right] \left[ (m_n^2 - m_F^2 - m_\Sigma^2)^2 - 4m_F^2m_\Sigma^2 \right]^{1/2}. \] (3.5)

In order to reproduce the desired value (1.2), the coupling constant should be \( y \sim \mathcal{O}(10^{-13}) \). This means that \( y \) is actually completely fixed via the neutron anomaly and is not a free parameter. However, if \( y \) were to be complex, then this effective interaction could in principle modify the electric dipole moment of the neutron. Nevertheless, due to the tiny value of \( y \), the corresponding correction should not lead to additional constraints.

The DM particles \( F \) are assumed to be stable and produced in local thermal equilibrium (LTE). Furthermore, possible upper bounds on the dark matter mass [35] are irrelevant due to \( m_F \sim \mathcal{O}(1 \text{ GeV}) \). Lower bounds on the boson mass are not valid since the dark radiation (or in other words the \( \sigma \) and \( \Sigma \) particles) is not in LTE with the SM modes during the typical neutrino decoupling at \( T_{\nu \text{D}} = 2.3 \text{ MeV} \) [36]. The low-temperature annihilation of \( F \) into scalar baryons implies that \( m_F > m_\Sigma \).

The presence of the effective interaction (3.3), between \( F \) and the neutrons leads to a 2-to-2 interaction with spin-independent cross section

\[ \sigma_{nF \rightarrow nF} \sim \frac{\Gamma_{\text{DS}}^2}{(v m_n)^4} \sim \mathcal{O}(10^{-69}) \text{cm}^2, \] (3.6)

where the characteristic velocity of the scattering is \( v \sim \mathcal{O}(10^{-3}) \) in units of speed of light [37]. Similar considerations apply also to the spin-dependent cross sections. These are totally in line with the LUX and XENON direct detection measurements for WIMP/neutron elastic scattering [37–40].
Finally, since the dark sector does not break the baryon symmetry, constraints inferred from baryon-violating processes such as the $(\Delta B = 2) n - \bar{n}$ oscillations [41], and $2n \rightarrow 2\pi$ dinucleon decays into pions [42], are not applicable here. If, on the other hand, $\Sigma$ were real, then the relevant parameter space of the proposed theory would be excluded. It should be stressed at this point that the $\Psi_0$ is uncharged under the hidden local symmetry, which is of vital importance in order to obtain a less constrained parameter space.

### 3.2 Cosmology

In order to solve the small-scale issues of the $\Lambda$CDM paradigm, $m_\sigma$ should lie in the sub-MeV region. Such values give rise to interesting effects at the dwarf and cluster galaxy scales and could allow a late kinetic decoupling of $F$ from the plasma. However, much lower values of $m_\sigma$, i.e. sub-keV masses tend to oversolve the small-scale problems of $\Lambda$CDM cosmology leading to unacceptably large protohalo masses, while the corresponding large SIDM cross sections destroy any core structures at dwarf and cluster scales [7].

Since $m_\sigma < \mathcal{O}(\text{MeV})$, these modes are in LTE and ultra-relativistic during the Big Bang Nucleosynthesis (BBN). In principle they could modify the processes that take place at this period, and consequently affect the energy density of the universe after the neutrino decoupling. This is encapsulated in the deviation of the effective neutrino degrees of freedom $\Delta N_{\text{eff}}$. For the proposed theory, we find that the dark sector decouples from the SM plasma before the annihilation of the bottom quark. This leads to an entropy dilution $(T_{\text{DS}}/T_\nu)^3 |_{T_{\nu,D}} \sim 0.2$, in line with BBN [43] and CMB [44] $1\sigma$-measurements. Such a value corresponds to $\Delta N_{\text{eff}}|_{\text{BBN}} \approx 0.35$, while during recombination $\Delta N_{\text{eff}}|_{\text{CMB}} \approx 0.15$, for $m_\sigma \sim \text{keV}$ and $m_\Sigma \sim \text{eV}$, which are perfectly compatible with all present measurements.

At the same time, the scalar baryons comprise less than 2% of the total CDM abundance. Moreover, this result may also explain the recent tension about the difference between $\Delta N_{\text{eff}}$ at BBN and CMB periods. In our model, $\Delta N_{\text{eff}}|_{\text{CMB}} - \Delta N_{\text{eff}}|_{\text{BBN}} < 0$, appears naturally, since the force mediators become nonrelativistic after the BBN period but long before the era of recombination. Note that keeping $\Sigma$ massless or $m_\Sigma \lesssim 0.3 \text{ eV}$, leads to an increase of $\Delta N_{\text{eff}} = 0.49$, which nevertheless is still $2\sigma$-compatible.

In order to obtain the thermal evolution of the theory, where all participating modes were in LTE at previous times and still be compatible with the effective description we proposed, we should demand that the heavy baryons be in LTE with the light quarks at temperatures below the cutoff $\Lambda$. It turns out that this is always the case as long as $\Lambda \ll M_{\text{Pl}} = 1.22 \times 10^{19} \text{ GeV}$ and $y \sim \mathcal{O}(10^{-13})$, which directly relates a hidden neutron decay to a non-zero relic abundance. Meanwhile, the decoupling of the DS from the SM plasma should take place well before the QCD phase transition after which the relativistic degrees of freedom reduce drastically. Otherwise, the extracted value of the entropy ratio would lead to unacceptably large deviations of effective neutrino degrees of freedom, excluded from Planck satellite measurements. The above reasoning yields $1 \text{ TeV} \lesssim \Lambda \ll M_{\text{Pl}}$.

Recent constraints [18–21], about the hidden sector communicating with the neutron, are inferred from the observed masses of neutron stars assuming WIMP theories. If the repulsive (vector-mediated) self-interactions between $F$ are strong inside the star, then the equation of state is modified and the bounds might be evaded [18]. Since the theory under consideration has all the virtues of a SIDM scenario, the observations of neutron stars with size of the order of $2M_\odot$, can in principle be accommodated. More precisely, the coupling constant $\alpha'$ should lie in the milli-regime, while the mediators should admit keV masses,
yielding [18, 21]

$$\alpha' \gtrsim \mathcal{O}(10^{-11}) \left( \frac{m_\sigma}{\text{keV}} \right)^2 . \quad (3.7)$$

In addition, for such values of $m_\sigma \sim \mathcal{O}(\text{keV})$ and $\alpha' \sim \mathcal{O}(10^{-4})$, and if the DM mass lies well below the GeV regime, then in principle dark halos could appear around neutron stars, as shown and explained in [45]. Such non-compact objects imply further constraints on the repulsive DM self-interactions. As explained previously, the DM mass is almost fixed and lies in the small interval (1.4), namely $m_F \sim \mathcal{O}(\text{GeV})$; this relaxes the proposed dark halo constraints on $\alpha'$ for this particular $m_\sigma$.

4 Thermal evolution of the theory

In this part of our work we study the thermal history of the proposed theory. We calculate the dark matter relic abundance, which consists of the $F$-modes after they chemically decouple from the plasma and we examine the possibility of a late kinetic decoupling regime. All temperatures that appear here are given in the reference frame of the photons, unless stated otherwise.

4.1 The present-day relic density

The fact that $m_F \sim \mathcal{O}(\text{GeV})$ and $m_\sigma \sim \mathcal{O}(\text{keV})$, implies that the SIDM couplings are much smaller than the usual WIMP couplings. Consequently, it is safe to assume that $\lambda \gg g$. In turn, for temperatures $T < m_F$, this leads into the rapid dominant annihilation of $F$ into the scalar baryons. This process is mediated through the heavy baryon $\Psi_0$, as long as $m_\Sigma \ll m_F$. The corresponding thermal cross section reads

$$\langle v_{\text{rel}} \sigma_{\text{ann}} \rangle_{F \rightarrow \Sigma} = \frac{\pi \alpha^2}{4m_\Psi^2} \langle v_{\text{rel}}^2 \rangle , \quad (4.1)$$

where $\langle \ldots \rangle$ denotes the thermal average using the relative velocities $v_{\text{rel}}$, and for later convenience we introduced $\alpha \equiv \lambda^2/4\pi$. Assuming that the $F$’s dominate the DM population, their annihilation into $\Sigma$ reduces their number. This fixes the coupling constant $\alpha$ for a given set of masses reproducing the measured DM relic density $\Omega_{\text{CDM}}$.

In what follows, we undertake the approach outlined in [32] and [46]. The Boltzmann equation for the distribution function $f_F(t, p)$ of the DM particles $F$ in a Friedmann-Lemaître-Robertson-Walker universe reads

$$\frac{\partial f_F}{\partial t} - 3H_0 p \frac{\partial f_F}{\partial p} = \frac{1}{2p^0} C , \quad (4.2)$$

where $H$ is the Hubble parameter and $C$ the collision integral, whose formal expression can be found for instance in [47] and [6]. Note that the collision term can be simplified by assuming that the final states follow an equilibrium thermal distribution. As customary, it is quite convenient to eliminate the Hubble parameter from (4.2), by working in terms of the number density per comoving volume, and study its evolution with respect to $m_F/T$.

The resulting differential equation is solved numerically by setting the initial conditions at the chemical freeze-out [32, 46, 48]. This yields the present-day relic abundance of $F$ as a function of $m_\Psi_0$ and the coupling constant $\alpha$

$$\Omega_F h^2 \approx \frac{0.12}{2} \left( \frac{\alpha}{0.1} \right)^2 \left( \frac{m_\Psi_0}{0.8 \text{TeV}} \right)^2 , \quad (4.3)$$
with $h \approx 0.67$, the reduced Hubble constant [44]. One notices immediately the absence of $m_F$ in the DM relic density at the first order approximation. This is an aftermath of the presence of the much heavier field $\Psi_0$ in the spectrum of the theory, which acts as the mediator during the annihilation process.

4.2 Kinetic decoupling at the keV-scale

Up to this point, we have shown that the theory is capable of explaining the anomaly of the neutron lifetime, and at the same time producing the desired relic abundance of DM. However, the solution to the small-scale “crisis” lies also in the elastic scattering between $F$ and the scalar baryons in the dark sector. The kinetic decoupling temperature, henceforth $T_{kd}$, describes the moment when these processes cease to sustain LTE [49].

The phenomenologically interesting case is when the kinetic decoupling takes place long after the BBN period ($T_{BBN} \sim \text{few MeV}$), in order to suppress the structure formation at scales similar to those of dwarf galaxies. For a novel approach of the thermal kinetic decoupling regime see [49], and for a detailed analysis of theories that enable such a late decoupling see [6].

We assume that $F$ at high temperatures are in LTE with the scalar baryons. As the temperature falls below $m_\sigma$, these modes annihilate and decay into scalar baryons. To estimate $T_{kd}$, we have to equate the Hubble parameter $H(T)$ to the elastic scattering rate $\Gamma_{el}(T)$, which is given by

$$\Gamma_{el}(T) \approx \frac{2}{3\pi^2 m_F} \int_0^\infty dE f_\Sigma(E) \frac{\partial}{\partial E} \left(E^4 \sigma_{el}\right).$$

(4.4)

As usual, $f_\Sigma(E)$ is the equilibrium distribution function per d.o.f. of the scalar baryons at temperature $T$, $E$ its energy, and $\sigma_{el}$ the momentum-transfer elastic cross section, given in [6].

At temperatures $T \ll m_F$, and at the lowest order in perturbation theory, one obtains for the IR-dominant expression

$$\sigma_{el} \approx \frac{4\pi\alpha'^2}{E^2} \left[ \log \left( \frac{4E^2 + m_\sigma^2}{m_\sigma^2} \right) - \frac{E^2}{E^2 + \frac{1}{4}m_\sigma^2} \right].$$

(4.5)

Note that as the temperature falls below $m_\sigma$, the elastic scattering becomes dominant. Here, we ignored the optical term due to the decay of the vector field since $\Gamma_{\sigma \to 2\Sigma} \ll m_\sigma$, see (2.4).

Before concluding this section, let us compute the kinetic decoupling temperature for certain values of the relevant parameters. For example, considering the benchmark point $m_\sigma \approx 10\ \text{keV}$ and milli-charges of $\alpha' \approx 3 \times 10^{-8}$ (cf. equation (3.7)) we find

$$T_{kd} \approx 0.45\ \text{keV}.$$  

(4.6)

In principle, much smaller values of $m_\sigma$ or/larger charges, give rise to lower decoupling temperatures leading to unacceptable results as explained previously.

5 The small-scale structures

The aim of this section is to show that the products of possible hidden decays of the neutron can solve the DM small-scale structure problems. Therefore, we discuss the implications on the small-scale structure formation due to the late kinetic decoupling, together with the SIDM nature of the model. Our analysis is based on [32] and [46].
5.1 The protohalo mass: the abundance of satellite galaxies

The kinetic decoupling is of great importance for the small-scale structures: the efficient momentum exchange damps the perturbations and determines the masses of the first gravitationally-bound objects (protohalos) of DM particles [33, 50]. This process terminates after $T_{kd}$. Technically, the fluctuations are exponentially damped with a characteristic mass scale given by

$$M_{damp} = \frac{4\pi \rho_m(T_{kd})}{3H^3(T_{kd})},$$

with $\rho_m(T_{kd})$ the matter density at the time of the kinetic decoupling. In other words, the dark matter substructures cannot form within the Hubble volume at kinetic decoupling, since the scalar baryon abundance within this volume suffices to keep $F$ in approximate LTE. As an example, masses of order $M_{damp} \sim 10^8 M_\odot$ correspond to $T_{kd}$ of the order of keV.5

The small-scale problems of the ΛCDM cosmology can be addressed more efficiently after suppressing the linear power spectrum at scales similar to those of dwarf galaxies [7, 9, 51]. This means that damping masses between $10^8–10^9$ solar masses are required. Such cut-offs are generated when the positrons kinetically decouple from their scattering partners at $T_{kd} \sim$ keV, as stated in [8, 52]. On the other hand, the cutoff masses should not be larger than the bounds set by Lyman–α measurements [33, 51, 53–55].

The proposed theory provides naturally $T_{kd} \sim \mathcal{O}(m_\sigma)$, and therefore cutoff masses of the desired order are easily accessible, while respecting all present constraints. For comparison, we note that CDM theories, where no late kinetic decoupling is present, usually predict masses between the Earth mass [50] and below the mass of Sagittarius A* [6].

5.2 Self-interaction scattering in the dwarf and cluster scales

In this last part of the present work, we study the impact to the small-scale structure formation at the non-linear regime due to the protohalo masses and the SIDM properties of the theory. In other words, we examine whether this model is able to resolve successfully the enduring small-scale problems of the ΛCDM paradigm: the cusp vs. core and the too big to fail problems. Moreover, if one considers also the missing satellite issue as an actual problem, then it is already solved due to the produced damping masses discussed previously. However, we tend to understand the solution of this problem more as a virtue of the underlying theory helping to address the cusp vs. core and the too big to fail problems more easily.

It is well known that a natural solution to both of these problems can be found in the context of SIDM [8]. For this to be possible, the elastic cross-section per unit mass should at least be $\langle \sigma_T/m \rangle_{therm} \sim 1\ cm^2\ g^{-1}$, where the thermal average is taken with respect to a Maxwell-Boltzmann distribution, with $v_{therm}$ as the most probable velocity. However, due to the typical velocities of the DM particles on cluster scales, it turns out that $\langle \sigma_T/m \rangle_{therm} \sim 0.1\ cm^2\ g^{-1}$ [8, 56], which implies that the SIDM elastic cross section should have a (mild) velocity dependence.6 If we ignore the damping of the power spectrum due to the kinetic decoupling, then the above values seem to be able to resolve the cusp vs. core and the too big to fail problems as shown in [57, 58] and explained in [8]. On the other hand, due to the combination of the SIDM effects together with the late kinetic decoupling, one has to make sure that the aforementioned problems are not actually over-solved.

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5 We do not consider the free-streaming impact due to the much heavier DM particles appearing in this theory as compared to the usual WDM mass interval.

6 This is also the case for neutron-proton scattering in the SM.
For the model that we discussed here, it is easy to obtain $\langle \sigma_T/m \rangle_{\text{therm}} \sim (0.1-10) \text{cm}^2 \text{g}^{-1} \text{keV}$, as long as $m_\sigma \sim \mathcal{O}(\text{keV})$ and $\alpha' \sim 10^{-8}$. For instance, we present the following benchmark point: for thermally produced $F$ with $m_F = 937.9 \text{ MeV}$, $m_\Psi_0 = 0.8 \text{ TeV}$, $\Lambda \approx 14 \text{ TeV}$ and $m_\sigma = 10 \text{ keV}$, a late kinetic decoupling takes place at $T_{kd} \approx 0.45 \text{ keV}$ (see eq. (4.6)). At the same time, the thermally-averaged SIDM elastic cross sections are similar to the ones of the tuned framework ETHOS-4 [7]. This model is perfectly compatible with the latest cosmological data, alleviates the missing satellite issue, and solves the too big to fail and the cusp vs. core problems as well [52], without over-solving them.

6 Conclusions

It was recently suggested that the large discrepancy in the measured neutron lifetime from the different experiments may be due to the presence of a hidden sector, which sources the present-day DM relic density. However, apart from reproducing the appropriate $\Omega_{DM}$, a theory describing DM should also address the enduring small-scale problems of the $\Lambda$CDM cosmology, while staying in line with all recent experiments.

The main aim of this paper was to provide an existence proof of a theory in which the possible dark decay products of the neutron solve the DM problems in the local group. More precisely, they admit the observed DM relic density and at the same time address the enduring small-scale structure problems of the $\Lambda$CDM paradigm, while staying compatible with all astrophysical and particle physics constraints.

In our context, a late decoupling takes place naturally for mediator masses at the keV scale. Furthermore, it is this very mediator that enables magnitudes of the corresponding SIDM cross sections. These are not only necessary in order to solve the enduring small-scale problems of the concordance cosmology, but also to generate sufficiently heavy neutron stars compatible with all observations, while reducing the recent tension about the effective neutrino degrees of freedom. The fact that the parameter space favors masses of order $\mathcal{O}(1-10) \text{ keV}$, might be directly related to experimental evidence: a few years ago a mysterious 3.55 keV photon ray was measured [60, 61], indicating the possibility of the existence of lighter modes in the Universe.

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