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Discoverable Matter: an Optimist's View of Dark Matter and How to Find It

by

Robert Stephen McGehee Jr.

A dissertation submitted in partial satisfaction of the

requirements for the degree of

Doctor of Philosophy

in

Physics

in the

Graduate Division

of the

University of California, Berkeley

Committee in charge:

Professor Hitoshi Murayama, Chair

Professor Alexander Givental

Professor Yasunori Nomura

Summer 2020

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Abstract

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An abundance of evidence from diverse cosmological times and scales demonstrates that 85% of the matter in the Universe is comprised of nonluminous, non-baryonic dark matter. Discovering its fundamental nature has become one of the greatest outstanding problems in modern science. Other persistent problems in physics have lingered for decades, among them the electroweak hierarchy and origin of the baryon asymmetry. Little is known about the solutions to these problems except that they must lie beyond the Standard Model. The first half of this dissertation explores dark matter models motivated by their solution to not only the dark matter conundrum but other issues such as electroweak naturalness and baryon asymmetry. The latter half of this dissertation approaches the dark matter enigma from a different vantage point inspired by the null results at dark matter direct detection experiments. The theory community has explored alternative dark matter candidates and production mechanisms while the experimental program has made progress on larger and more sensitive experiments. In this dissertation, we take a complementary approach by investigating signals of novel dark matter models which may have been overlooked in current experiments.

For Carol.

Contents

Contents	ii
1 Motivation	1
2 SIMPs through the Axion Portal	6
2.1 Introduction	6
2.2 The Framework	7
2.3 Boltzmann Equations	9
2.4 Axions and Pions in Equilibrium	11
2.5 Axions and Photons in Equilibrium	17
2.6 Discussion	22
3 Freezing-in Twin Dark Matter	25
3.1 Introduction	25
3.2 The Mirror Twin Higgs & Cosmology	26
3.3 Kinetic Mixing & a Massive Twin Photon	28
3.4 Freezing- <i>Twin</i> Dark Matter	31
3.5 Conclusion	35
4 Asymmetric Matters from a Dark First-Order Phase Transition	36
4.1 Introduction	36
4.2 Dark-Sector Baryogenesis	37
4.3 Dark Matter and Experimental Signatures	40
4.4 Conclusion	44
5 Directly Detecting Signals from Absorption of Fermionic Dark Matter	46
5.1 Introduction	46
5.2 Signals and Operators	47
5.3 Neutral Current Signals: Nuclear Recoils	48
5.4 Charged Current Signals: Dark-Matter-Induced β^- Decays	51
5.5 Discussion	54

6	Absorption of Fermionic Dark Matter by Nuclear Targets	55
6.1	Introduction	55
6.2	UV Completions	57
6.3	Relevant Current Experiments	65
6.4	Neutral Current Nuclear Recoils	65
6.5	Charged Current: Induced β Decays	73
6.6	Charged Current: β Endpoint Shifts	83
6.7	Discussion	85
7	Summary	88
	Bibliography	89

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Chapter 1

Motivation

The long history and abundance of evidence for dark matter (DM) from diverse cosmological times and scales has made discovering its fundamental nature one of the greatest outstanding problems in modern science. Some of the earliest evidence for dark matter was inexplicable mass-to-light ratios across a range of scales [1–4]. Decades later, galactic rotation curves began suggesting that most of the matter in galaxies was mysteriously invisible [5–10]. In the past few years, *Planck* results have continued to uphold this picture. Fits for the baryon and matter densities inferred from the power spectrum of the cosmic microwave background (CMB) anisotropy, $\Omega_B h^2 = 0.022$ and $\Omega_M h^2 = 0.12$ respectively, clearly demonstrate that 85% of the matter in the Universe is comprised of nonluminous, non-baryonic *dark matter* [11]. Big Bang Nucleosynthesis (BBN) allows for an independent verification of the baryon density using the measured deuterium abundance [12–14], while the imprints of baryon acoustic oscillations on large scale structure allow for an independent verification of the matter density [15–21].

All of the above information about dark matter is granted solely through its gravitational interactions with Standard Model (SM) particles. Its non-gravitational interactions and particle description(s) remain unknown and only broadly constrained. Dark matter self interactions are bounded above from observations of galaxy cluster mergers [22–24] as well as halo shapes [25, 26]. Dark matter’s lifetime is often constrained to be much longer than the age of the Universe, depending on its decay products (see *e.g.* [27]). We also know that dark matter cannot be “hot,” where it would decouple from the early Standard Model bath while still relativistic, due to the unobserved erasure of small-scale structure (see *e.g.* [28]). If dark matter is a fermion, there exists a lower bound on its mass due to the inability to pack fermions on small scales from Pauli’s principle [29]. In sum, there are very few model-independent bounds on dark matter.

Since so little is known about dark matter, it has given particle theorists an expansive, cosmic playground to explore exciting phenomenologies. Often, particular models have been motivated not only by their solution to the dark matter conundrum, but also by their relation to other outstanding problems within the Standard Model. Weakly Interacting Massive Particles (WIMPs) have been one of the most popular classes of dark matter candidates

for decades. They generically refer to particles of roughly weak-scale masses with weak-scale interactions which realize the so-called “WIMP miracle”: they easily reproduce the observed dark matter abundance. It is no surprise that these WIMPs are found in a variety of theories which solve the electroweak hierarchy problem such as supersymmetric [30, 31] or extra-dimensional models [32, 33]. See *e.g.* [34, 35] for recent reviews of WIMPs. Another perennial dark matter favorite is the QCD axion which appears in solutions to the strong CP problem of QCD [36–41]. Other well motivated candidates include sterile neutrinos (see *e.g.* [42]) and gravitinos (see *e.g.* [43]), though both of these are increasingly constrained¹.

In a similar spirit, the next three chapters of this dissertation expound dark matter candidates which are well motivated by their connection to other problems. The models discussed in Chapters 2 through 4 are well motivated by their potential solution of many well known (and debated) small-scale structure issues (see, *e.g.* [46] for a recent review). One such issue is the core vs. cusp problem: simulations which assume cold, collisionless dark matter (CDM) often generate more cuspy dark matter distributions than those found in dwarf galaxies [47–65]. Additionally, low surface brightness galaxies continue to suggest that a cored dark matter profile may be preferred to a CDM, cuspy one [66]². Other CDM issues include the missing satellites [68–71], diversity [72–74], and too-big-to-fail problems [75–78]. All of these issues may be resolved by the particular dark matter candidates we propose in Chapters 2 through 4 which each have strong self interactions³.

The electroweak hierarchy problem has long provided theoretical motivation for the presence of new TeV-scale physics [81–87]. With the curious absence of new physics observed at the LHC, the issue of Higgs naturalness has become increasingly urgent and prompts the consideration of new approaches to this old problem. One such approach is the ‘Neutral Naturalness’ paradigm in which the states responsible for stabilizing the electroweak scale are not charged under (some of) the Standard Model gauge symmetries [88–100], thus explaining the lack of expected signposts of naturalness. The model of dark matter in Chapter 3 is itself part of a broader class of mirror twin Higgs (MTH) models which fall under this ‘Neutral Naturalness’ umbrella and solve the hierarchy problem up to cutoffs around $\Lambda \sim 5 - 10$ TeV.

Another of the greatest standing mysteries in physics is the origin of the baryon asymmetry which is puzzling because, although baryon and lepton number are individually conserved at tree level in the Standard Model, cosmological measurements observe a net baryon asymmetry. A mechanism to generate this asymmetry must contain three ingredients outlined by Sakharov: (1) C- and CP-violation, (2) baryon number violation, and (3) departure from thermal equilibrium [101]. There are two broad classes of solutions: leptogenesis [102] and

¹For sterile neutrinos, a combination of experimental and observational constraints are forcing the neutrino mixing angle to be ever smaller. For gravitinos, a combination of cosmological bounds and naturalness requirements are making it more difficult to realize concrete models (see *e.g.* [44, 45]).

²However, see [67] which demonstrates that this newly discovered Antlia 2 satellite galaxy may be reasonably described by both CDM and self interacting dark matter models.

³At their core, these issues are discrepancies between observations and simulations, which have often only included CDM. So, there has been a lot of effort to include relevant baryonic processes in these simulations to see if these issues can be resolved within a CDM framework [79, 80]

electroweak baryogenesis (see *e.g.* [103–105]). The former is well motivated by the seesaw explanation of neutrino masses [106–108], but occurs at such high temperatures in the early Universe that it is difficult to probe⁴. The latter requires only minimal additions to the Standard Model matter content [111, 112], but is often excluded by measurements of electric dipole moments [113]. There has been a recent push to devise theories with electroweak-like baryogenesis mechanisms which are not yet excluded, but are more imminently falsifiable than theories of leptogenesis [114–121]. In Chapter 4, we propose such a model of ‘dark sector baryogenesis’ which explains the origins of both the baryon asymmetry and dark matter in a minimal, testable model.

The last half of this dissertation approaches the dark matter enigma from a different vantage point. Instead of other unsolved problems motivating particular dark matter candidates, null search results at dark matter direct detection experiments motivate novel dark matter signals which may have been overlooked (see *e.g.* [122] for an excellent review of dark matter direct detection). Despite increasing sensitivities, dark matter direct detection experiments have continued to yield null results for the long sought after WIMP [123–125]. As such, both theoretical and experimental programs have moved beyond the WIMP paradigm. On the theoretical side, explorations of alternative dark matter candidates and production mechanisms have motivated dark matter masses below a GeV [126–138]. On the experimental side, progress has been made on the size frontier — where ton-scale experiments are needed to search for signals with rates consistent with current bounds. In tandem, new proposals hope to probe regions of parameter space interesting for lighter dark matter candidates [139–158]. Another strategy for progress is to explore novel dark matter direct detection signals (see *e.g.*, [159–161]). With this direction in mind in Chapters 5 and 6, we consider the novel signals coming from the absorption of a fermionic dark matter particle in a detector— a scenario in which the dark matter mass energy is available to the target (in contrast to the well studied elastic scenario where only its velocity-suppressed kinetic energy may be imparted).

In sum, Chapters 2, 3, and 4 focus on new ways to produce dark matter within motivated dark sectors while Chapters 5 and 6 concentrate on new ways to directly detect dark matter. In detail, the remaining chapters’ contents are:

Chapter 2 Dark matter could be a thermal relic comprised of strongly interacting massive particles (SIMPs), where $3 \rightarrow 2$ interactions set the relic abundance. Such interactions generically arise in theories of chiral symmetry breaking via the Wess-Zumino-Witten term. In this Chapter, we show that an axion-like particle can successfully maintain kinetic equilibrium between the dark matter and the visible sector, allowing the requisite entropy transfer that is crucial for SIMPs to be a cold dark matter candidate. Constraints on this scenario arise from beam dump and collider experiments, from the CMB, and from supernovae. We find a viable parameter space when the axion-like particle is close in mass to the SIMP dark matter, with strong-scale masses of order a few hundred MeV. Many planned experiments are set to probe the parameter space in the near future. This chapter is based on Ref. [162].

⁴For counterexamples, see [109, 110].

Chapter 3 The MTH addresses the little hierarchy problem by relating every Standard Model particle to a twin copy, but is in tension with cosmological bounds on light degrees of freedom. Asymmetric reheating has recently been proposed as a simple way to fix MTH cosmology by diluting the twin energy density. We show that this dilution sets the stage for an interesting freeze-in scenario where both the initial absence of dark sector energy and the feeble coupling to the SM are motivated for reasons unrelated to dark matter production. We give the twin photon a Stueckelberg mass and freeze-in twin electron and positron dark matter through the kinetic mixing portal. The kinetic mixing required to obtain the dark matter abundance is of the loop-suppressed order expected from infrared contributions in the MTH. This chapter is based on Ref. [163].

Chapter 4 We introduce a model for matters-genesis in which both the baryonic and dark matter asymmetries originate from a first-order phase transition in a dark sector with an $SU(3) \times SU(2) \times U(1)$ gauge group and minimal matter content. In the simplest scenario, we predict that dark matter is a dark neutron with mass either $m_n = 1.33$ GeV or $m_n = 1.58$ GeV. Alternatively, dark matter may be comprised of equal numbers of dark protons and pions. This model, in either scenario, is highly discoverable through both dark matter direct detection and dark photon search experiments. The strong dark matter self interactions may ameliorate small-scale structure problems, while the strongly first-order phase transition may be confirmed at future gravitational wave observatories. This chapter is based on Ref. [164].

Chapter 5 We present a new class of direct detection signals: absorption of fermionic dark matter. We enumerate the operators through dimension six which lead to fermionic absorption, study their direct detection prospects, and summarize additional constraints on their suppression scale. Such dark matter is inherently unstable as there is no symmetry which prevents dark matter decays. Nevertheless, we show that fermionic dark matter absorption can be observed in direct detection and neutrino experiments while ensuring consistency with the observed dark matter abundance and required lifetime. For dark matter masses well below the GeV scale, dedicated searches for these signals at current and future experiments can probe orders of magnitude of unexplored parameter space. This chapter is based on Ref. [165].

Chapter 6 Absorption of fermionic dark matter leads to a range of distinct and novel signatures at dark matter direct detection and neutrino experiments. We study the possible signals from fermionic absorption by nuclear targets, which we divide into two classes of four Fermi operators: neutral and charged current. In the neutral current signal, dark matter is absorbed by a target nucleus and a neutrino is emitted. This results in a characteristically different nuclear recoil energy spectrum from that of elastic scattering. The charged current channel leads to induced β decays in isotopes which are stable in vacuum as well as shifts of the kinematic endpoint of β spectra in unstable isotopes. To confirm the possibility of observing these signals in light of other constraints, we introduce UV completions of

example higher dimensional operators that lead to fermionic absorption signals and study their phenomenology. Most prominently, dark matter which exhibits fermionic absorption signals is necessarily unstable leading to stringent bounds from indirect detection searches. Nevertheless, we find a large viable parameter space in which dark matter is sufficiently long lived and detectable in current and future experiments. This chapter is based on Ref. [166].

Chapter 7 We conclude with a summary of the work in this dissertation and an optimistic outlook.

We sincerely hope that nothing in this dissertation will be of any use fifty years from now because dark matter will have been discovered and a generation of particle theorists will have pivoted to understanding its details in minutiae: its cosmological history, its interactions with the Standard Model, and its role in the larger dark sector. It is unfathomable that even ϵ of the content herein will still have relevance after its discovery. Nonetheless, in the following chapters, we sketch pictures of dark matter which are for now just an optimist's dream.

Chapter 2

SIMPs through the Axion Portal

2.1 Introduction

As discussed in Chapter 1, dark matter’s interactions and origin remain unknown. One possibility for its origin in the early Universe is as a thermal relic. The most well-studied thermal scenario is that of a WIMP. The number density of WIMPs is set by $2 \rightarrow 2$ annihilations of the dark matter into Standard Model particles, and the observed dark matter relic abundance is achieved when both the dark matter mass and coupling to Standard Model particles are near the scales relevant for electroweak processes.

An alternative thermal setup was proposed in Ref. [128] where $3 \rightarrow 2$ dark matter self-interactions set its abundance. In this scenario, the observed relic density indicates that the dark matter mass and self-coupling should be near the strong scale. This mechanism of strongly interacting massive particles (SIMPs) was shown to be generic in strongly coupled theories of chiral symmetry breaking, where the pions play the role of dark matter [129]. The $3 \rightarrow 2$ interactions are then sourced by the well known Wess-Zumino-Witten (WZW) action [167–169]. This provides a simple and calculable realization of the SIMP mechanism, although by no means the only one [170–180].

In addition to providing a novel thermal mechanism for explaining the dark matter abundance, SIMPs also offer a possible explanation for the issues related to small-scale structure formation discussed in Chapter 1. While many of these issues may be resolved with better understanding of astrophysical processes (for instance in Ref. [181]), it is also possible to mitigate these issues if the dark matter can self-scatter (see Ref. [46] for a recent review). The strong self-annihilations of SIMP dark matter imply that their self-scatterings are also large, such that they naturally address these small-scale puzzles [128, 129].

The $3 \rightarrow 2$ dark matter annihilations would raise the temperature of the residual dark matter due to conservation of comoving entropy. Therefore, the dark matter must be in thermal equilibrium with a heat sink, such as the Standard Model bath, until after freeze-out [128]. Otherwise, the $3 \rightarrow 2$ dark matter annihilations would cause the steady depletion of dark matter particles and heating of the remaining dark matter, a scenario referred

to as cannibalization. While cannibalization was originally proposed to provide a class of dark matter models intermediate between hot and cold dark matter, such models are not observationally-viable [131, 182]. Obtaining the observed dark matter abundance inevitably leads to an unacceptable washout of small-scale structure.

To allow for adequate thermalization between the dark matter and the Standard Model, Refs. [172, 173] explored the kinetically-mixed hidden photon portal. Here, we explore the possibility of a pseudoscalar portal using axion-like particles to accomplish the entropy transfer to photons. For brevity, we refer to axion-like particles simply as “axions” throughout this Chapter. We note that Ref. [183] also considered an axion portal, but focused on the regime where semi-annihilations set the relic abundance. In contrast, we focus on the SIMP regime where $3 \rightarrow 2$ annihilations determine the relic density. For concreteness, we will use the SIMPlest pion realization of the dark matter based on an $\text{Sp}(2N_c)$ gauge theory with four doublet Weyl fermions following Ref. [129]. $\text{Sp}(2N_c)$ gauge groups with a larger number of flavors or $\text{SU}(N_c)$ and $\text{SO}(N_c)$ gauge groups allow for semi-annihilations which can control the relic abundance, although there may still be parameter space where $3 \rightarrow 2$ annihilations determine the dark matter density.

The remainder of this Chapter is organized as follows. In Section 2.2, we describe the framework and identify the interactions responsible for setting the correct dark matter relic abundance and cooling the dark matter via the axion portal. In Section 2.3, we review and enumerate the relevant Boltzmann equations. In order to cool the dark matter effectively, the axion must be in thermal equilibrium with both the dark matter and the Standard Model. In Section 2.4, we illustrate the theoretical and empirical requirements of axion-pion thermal equilibrium, while in Section 2.5 we do the same for axion-Standard Model thermal equilibrium. Concluding remarks and discussions follow in Section 2.6.

2.2 The Framework

Our starting point is an $\text{Sp}(2N_c)$ gauge theory with $2N_f$ Weyl fermions that couple to an axion-like field a as

$$\mathcal{L}_{aq} = -\frac{1}{2}m_a^2 a^2 - \left(\frac{1}{2}m_Q e^{ia/f_{a\pi}} J^{ij} q_i q_j + \text{h.c.} \right) \quad (2.1)$$

where m_a is the axion mass, m_Q is the quark mass matrix, q_i are the confining quarks and J is the $\text{Sp}(2N_f)$ group invariant.¹ Upon dynamical chiral symmetry breaking, the ground state is expected to be given by

$$\langle q_i q_j \rangle = \mu^3 J_{ij}. \quad (2.2)$$

¹Note that in terms of a 4-component spinor $\psi^T = (q \quad q^\dagger)$, the identities $i\bar{\psi}\gamma_5\psi = -iqq + iq^\dagger q^\dagger$ and $\bar{\psi}\psi = qq + q^\dagger q^\dagger$ hold.

Any transformation by the flavor symmetry $V \in \text{SU}(N_f)$ would also be a ground state, and in general

$$\langle q_i q_j \rangle = \mu^3 (V J V^T)_{ij}. \quad (2.3)$$

Switching the description to the chiral Lagrangian, a spacetime-dependent flavor rotation gives the low-energy excitations,

$$\langle qq \rangle \rightarrow \mu^3 \Sigma, \quad \Sigma \equiv V J V^T, \quad V = \exp(i\pi/f_\pi), \quad (2.4)$$

where $\pi \equiv \pi^b T^b$, T^b are the $\text{Sp}(2N_f)$ generators and f_π is the pion decay constant. We use the normalization $\text{Tr } T^b T^c = 2\delta^{bc}$ for the generators. In terms of the pion fields,

$$\begin{aligned} -\frac{i}{2} m_Q J^{ij} q_i q_j + h.c. &\Rightarrow \frac{m_\pi^2}{6f_\pi} \text{Tr } \pi^3 + \mathcal{O}(\pi^5), \\ -\frac{1}{2} m_Q J^{ij} q_i q_j + h.c. &\Rightarrow -\frac{m_\pi^2}{4} \text{Tr } \pi^2 + \mathcal{O}(\pi^4). \end{aligned} \quad (2.5)$$

The theory has an $\text{SU}(2N_f)/\text{Sp}(2N_f)$ flavor structure, where the residual $\text{Sp}(2N_f)$ is exact due to the quark masses' proportionality to J . For $N_f \geq 2$, the fifth homotopy group of the coset space is non-vanishing and the WZW term exists [167–169],

$$\mathcal{S}_{\text{WZW}} = \frac{-iN_c}{240\pi^2} \int \text{Tr } (\Sigma^\dagger d\Sigma)^5. \quad (2.6)$$

Generally, both $a\text{Tr}(\pi^3)$ and $a^2\text{Tr}(\pi^2)$ terms can appear in the interaction Lagrangian. However, the former introduces semi-annihilations of pions into a pion and an axion which might contribute to determining the relic abundance of the dark matter. Here, we are interested in exploring the role of an axion mediator in the SIMP mechanism of $3 \rightarrow 2$ self-annihilations of pions. Consequently, we focus on an $\text{Sp}(2N_c)$ gauge theory with $N_f = 2$ fermions, where the flavor symmetry is $\text{SU}(4)/\text{Sp}(4)$ and $N_\pi = (N_f - 1)(2N_f + 1) = 5$ pions emerge. In this theory, the semi-annihilation process is absent since $\text{Tr}(\pi^3) = 0$, and pure $3 \rightarrow 2$ annihilations of pions via the WZW term are guaranteed to control the relic abundance of dark matter. For more flavors, or for other gauge groups, $3 \rightarrow 2$ annihilation may still control the relic abundance, though in a smaller region of parameter space. To leading order in pion fields, the WZW term for our choice of gauge group takes the form

$$\mathcal{L}_{\text{WZW}} = \frac{2N_c}{15\pi^2 f_\pi^5} \epsilon^{\mu\nu\rho\sigma} \text{Tr} [\pi \partial_\mu \pi \partial_\nu \pi \partial_\rho \pi \partial_\sigma \pi] = \frac{8N_c}{15\pi^2 f_\pi^5} \epsilon^{\mu\nu\rho\sigma} \epsilon_{abcde} \pi^a \partial_\mu \pi^b \partial_\nu \pi^c \partial_\rho \pi^d \partial_\sigma \pi^e. \quad (2.7)$$

The excess kinetic energy generated in the dark sector from $3 \rightarrow 2$ annihilations needs to be transferred out, which can be obtained through kinetic coupling of the pions to the axions and the axions to the Standard Model bath. Since the semi-annihilation term is absent for our flavor group of choice, the interaction Lagrangian between pions and axions is

$$\mathcal{L}_{a\pi} \supset \frac{\kappa}{4} a^2 \pi^b \pi^c \delta^{bc}. \quad (2.8)$$

If the axion coupling to the pions arises in a similar manner to what occurs in QCD, as in Eq. (2.1), the mass term for the hidden quarks q in the $\text{Sp}(2N_c)$ gauge theory gives rise to an axion potential:

$$\begin{aligned}\mathcal{L}_{a\pi} &= -\frac{1}{2}m_a^2 a^2 - \frac{1}{2}m_Q \mu^3 e^{ia/f_{a\pi}} \text{Tr} J \Sigma + \text{h.c.} \\ &= -\frac{1}{2} \left(m_a^2 a^2 + \frac{2m_\pi^2 f_\pi^2}{f_{a\pi}^2} \right) a^2 + \frac{m_\pi^2}{8f_{a\pi}^2} a^2 \text{Tr} \pi^2 + \dots\end{aligned}\quad (2.9)$$

where m_π is the pion mass. Using the normalization of $\text{Tr} \pi^2 = 2\pi^b \pi^c \delta^{bc}$, we identify the Feynman rule for the $aa\pi^b \pi^c$ vertex in Eq. (2.8) as

$$i\kappa \delta^{bc} = i \frac{m_\pi^2}{f_{a\pi}^2} \delta^{bc} . \quad (2.10)$$

Meanwhile, the interaction Lagrangian between the axions and Standard Model photons is

$$\mathcal{L}_{a\gamma} = \frac{1}{4f_{a\gamma}} a F^{\mu\nu} \tilde{F}_{\mu\nu} . \quad (2.11)$$

As long as the kinetic equilibrium between the pions and the Standard Model is maintained through the interactions of Eqs. (2.9) and (2.11), the preferred mass for the dark matter is $m_\pi \approx 300$ MeV [129] with $m_\pi \sim 2\pi f_\pi$ to set the observed relic abundance. We find below that viable ALP masses are around the same scale, 10 MeV $\lesssim m_a \lesssim 1$ GeV. Couplings that satisfy $m_\pi \sim 2\pi f_\pi$ correspond to the strongly-interacting regime of the theory, where self-interactions are important on astrophysical scales. In this regime, $\mathcal{O}(1)$ corrections to perturbative results are expected, and therefore should be thought of as a proxy for the scales involved. Phenomenologically interesting pion masses lie at the edge of perturbativity, where higher order corrections and vector meson effects can impact the range of observationally-viable pion masses [180, 184].

2.3 Boltzmann Equations

Before deriving the theoretical and empirical requirements of axion-pion and axion-Standard Model thermal equilibrium, we briefly review and enumerate the relevant Boltzmann equations. The Boltzmann equations govern the evolution of the phase space $f_X(p, t)$ for particle X ,

$$\frac{\partial f_X}{\partial t} - H \frac{p^2}{E} \frac{\partial f_X}{\partial E} = C[f_X] \quad (2.12)$$

where the left hand side is the relativistic Liouville operator in a Friedmann-Robertson-Walker spacetime and $C[f_X]$ is the collision term. In the regime where $2 \rightarrow 2$ annihilations are negligible, the relevant terms which appear in the collision term are the $3 \rightarrow 2$ interactions

which set the relic abundance, the dark matter self-interactions which give it a thermal distribution, the elastic interactions which transfer energy between the pions and axions, and decays (and inverse decays) of axions to Standard Model particles. We neglect axions converting to photons (and vice versa) via t -channel scattering off electrons, which is less efficient than decays (and inverse decays) at thermalizing the axions with the Standard Model. In the parameter space of interest (with the exception of axions in the out-of-equilibrium decay scenario), all particles will be interacting sufficiently frequently so that they have thermal distributions,

$$f_X = \frac{1}{e^{(E-\mu)/T} \pm 1}, \quad (2.13)$$

where the $+$ sign ($-$ sign) is for Fermi-Dirac (Bose-Einstein) statistics and μ is the chemical potential. At temperatures below the mass of a given particle, the effects of quantum statistics become negligible and the Maxwell-Boltzmann (MB) distribution is recovered.

To solve the Boltzmann equations, it is most useful to look at two moments of the phase space distribution, which correspond to the number density and energy density

$$n_X = g_X \int d^3p f_X \quad (2.14)$$

$$\rho_X = g_X \int d^3p E f_X, \quad (2.15)$$

where g_X is the number of degrees of freedom and $d^3p \equiv d^3p/(2\pi)^3$. The Boltzmann equations for axion-mediated SIMPs are

$$\frac{\partial n_\pi}{\partial t} + 3H n_\pi = -\langle \sigma_{3 \rightarrow 2} v^2 \rangle (n_\pi^3 - n_\pi^2 n_\pi^{\text{eq}, T_\pi}), \quad (2.16)$$

$$\frac{\partial \rho_\pi}{\partial t} + 3H (\rho_\pi + p_\pi) = -\langle \sigma_{el} v \delta E \rangle n_\pi n_a, \quad (2.17)$$

$$\frac{\partial n_a}{\partial t} + 3H n_a = -(\langle \Gamma \rangle_{T_a} n_a - \langle \Gamma \rangle_{T_{\text{SM}}} n_a^{\text{eq}, T_{\text{SM}}}), \quad (2.18)$$

$$\frac{\partial \rho_a}{\partial t} + 3H (\rho_a + p_a) = \langle \sigma_{el} v \Delta E \rangle n_\pi n_a - m_a \Gamma (n_a - n_a^{\text{eq}, T_{\text{SM}}}). \quad (2.19)$$

where p_X is the pressure density of species X (which is related to the energy density through the equation of state w_X), and $n_X^{\text{eq}, T}$ is the thermal equilibrium density for species X at temperature T . Additionally, $\langle \Gamma \rangle_T = \Gamma \langle m_a/E_a \rangle_T$ is the thermally averaged decay rate of the axion at temperature T , $\langle \sigma_{3 \rightarrow 2} v^2 \rangle$ is the thermally-averaged $3 \rightarrow 2$ cross-section of the pions for this choice of gauge group, labeled $i = 1 \dots 5$,

$$\begin{aligned} \langle \sigma_{3 \rightarrow 2} v^2 \rangle &= \frac{1}{3!2!n_\pi^3} \int \left(\prod_{i=1}^5 \frac{d^3p_i}{2E_i} \right) f_1 f_2 f_3 |\mathcal{M}_{123 \rightarrow 45}|^2 (2\pi)^4 \delta^4(p_1 + p_2 + p_3 - p_4 - p_5), \\ &= \frac{6N_c^2}{\sqrt{5}\pi^5} \frac{m_\pi^3 T_F^2}{f_\pi^{10}} \end{aligned} \quad (2.20)$$

and $n_\pi n_a \langle \sigma_{el} v \Delta E \rangle$ is the energy transfer rate between the pions and axions (with the initial and final states labeled as 1 and 2),

$$n_\pi n_a \langle \sigma_{el} v \Delta E \rangle = \int \frac{d^3 p_{\pi_1}}{2E_{\pi_1}} \frac{d^3 p_{a_1}}{2E_{a_1}} \frac{d^3 p_{\pi_2}}{2E_{\pi_2}} \frac{d^3 p_{a_2}}{2E_{a_2}} \left[(E_{\pi_1} - E_{\pi_2}) f_{\pi_1} f_{a_1} |\mathcal{M}_{\pi a \rightarrow \pi a}|^2 \right. \\ \left. \times (2\pi)^4 \delta^4(p_{\pi_1} + p_{a_1} - p_{\pi_2} - p_{a_2}) \right]. \quad (2.21)$$

For MB statistics, the equilibrium values for the number, energy, and pressure densities for a particle of mass m and temperature T with g degrees of freedom are

$$n^{\text{eq}} = g \int d^3 p f^{\text{eq}} = \frac{gm^2 T}{2\pi^2} K_2\left(\frac{m}{T}\right) \quad (2.22)$$

$$\rho^{\text{eq}} = g \int d^3 p E f^{\text{eq}} = \frac{gm^2 T}{2\pi^2} \left(m K_1\left(\frac{m}{T}\right) + 3TK_2\left(\frac{m}{T}\right) \right) \quad (2.23)$$

$$p^{\text{eq}} = g \int d^3 p \frac{p^2}{3E} f^{\text{eq}} = \frac{gm^2 T^2}{2\pi^2} K_2\left(\frac{m}{T}\right) \quad (2.24)$$

and the thermally-averaged boost factor is

$$\left\langle \frac{m}{E} \right\rangle_T = \frac{1}{n^{\text{eq}}} \int g \frac{d^3 p}{(2\pi)^3} \frac{m}{E} f^{\text{eq}} = \frac{K_1(m/T)}{K_2(m/T)}. \quad (2.25)$$

2.4 Axions and Pions in Equilibrium

Theoretical Requirements

The interaction Lagrangian in Eq. (2.9) leads to annihilations of pions into axions and to elastic scattering of axions off of pions. The SIMP mechanism requires the former process to be suppressed at the time of $3 \rightarrow 2$ freeze-out while the latter is active.

The requirement in order for the $3 \rightarrow 2$ pion self-annihilations to dominate the $2 \rightarrow 2$ annihilations of pions into axions at the time of freeze-out is

$$n_\pi \langle \sigma v \rangle_{\text{ann}} \lesssim H|_{T_F}, \quad (2.26)$$

where $\langle \sigma v \rangle_{\text{ann}}$ is the thermally averaged cross section for the annihilation process $\pi\pi \rightarrow aa$. The Hubble parameter at freeze-out is given by

$$H|_{T_F} = \sqrt{\frac{g_{*,F} \pi^2}{90}} \frac{T_F^2}{M_{\text{Pl}}}, \quad (2.27)$$

where T_F is the freeze-out temperature of $3 \rightarrow 2$ interactions (typically $T_F \sim m_\pi/20$ in the SIMP setup), M_{Pl} is the Planck mass, and $g_{*,F}$ is the effective number of relativistic degrees of freedom at the time of freeze-out. We have verified numerically that this requirement on

the annihilation rate does maintain the correct relic abundance as set by the $3 \rightarrow 2$ SIMP mechanism.

The thermally averaged annihilation cross section that appears in Eq. (2.26) can be readily calculated. For a trivial matrix element, $\mathcal{M}$, of a process $12 \rightarrow 34$ (as is relevant for the Lagrangian of Eq. (2.9)) in which all states obey Maxwell-Boltzmann statistics, the thermally averaged cross section entering the Boltzmann equations [185] is expressed in terms of

$$\gamma_{12 \rightarrow 34} = \frac{g_1 g_2 g_3 g_4 T |\mathcal{M}|^2}{2^9 \pi^5} \int_{s_{\min}}^{\infty} ds \sqrt{s} \lambda^{1/2}(\sqrt{s}, m_1, m_2) \lambda^{1/2}(\sqrt{s}, m_3, m_4) K_1(\sqrt{s}/T), \quad (2.28)$$

where g_i counts degrees of freedom for particle i ,

$$s_{\min} = \max\{(m_1 + m_2)^2, (m_3 + m_4)^2\}, \quad (2.29)$$

$$\lambda(x, y, z) \equiv (1 - (y + z)^2/x^2)(1 - (y - z)^2/x^2), \quad (2.30)$$

and K_1 is the first order modified Bessel function of the second kind. The amplitude $|\mathcal{M}|^2$ which appears is averaged over *all* degrees of freedom. For pion-axion scattering and pion annihilation to axions, the relevant amplitude is therefore

$$|\mathcal{M}|^2 \equiv \frac{1}{g_\pi} \frac{1}{g_\pi} \sum_{b,c} \kappa^2 \delta^{bc} = \frac{m_\pi^4}{g_\pi f_{a\pi}^4} \quad (2.31)$$

since the trace requires that the pions be the same. The thermally averaged cross section for the annihilation process $\pi\pi \rightarrow aa$ is

$$\langle \sigma v \rangle_{\text{ann}} = \frac{1}{2} \frac{\gamma_{\text{ann}}}{(n_\pi^{\text{eq}})^2}, \quad (2.32)$$

where n_i^{eq} denotes the number density of particle i in equilibrium,

$$n_i^{\text{eq}} = \frac{g_i}{2\pi^2} T_i m_i^2 K_2(m_i/T_i), \quad (2.33)$$

where K_2 is the second order modified Bessel function of the second kind. In Eq. (2.32), the phase-space factor of $1/2$ for identical initial particles is canceled because the number density changes by two particles per annihilation. For $m_a \ll m_\pi$, Eq. (2.32) simplifies to

$$\langle \sigma v \rangle_{\text{ann}} \approx \frac{g_a^2 m_\pi^2}{64 g_\pi \pi f_{a\pi}^4}. \quad (2.34)$$

In addition to the suppression of the $2 \rightarrow 2$ annihilations, the SIMP mechanism requires that the rate of energy transfer in the scattering process $a\pi \rightarrow a\pi$ is fast enough to successfully cool the dark matter. We require that thermal decoupling occurs after freeze-out,

$T_D < T_F$, so that the energy transfer is efficient for the entirety of the freeze-out process. In the limit of small m_a (using Bose-Einstein statistics), we follow the analytic derivation of Ref. [186] which obtains the thermal decoupling temperature T_D of the $\pi a \rightarrow \pi a$ elastic scattering process. In this range of temperatures, the pions are guaranteed to be nonrelativistic since $T_F \sim m_\pi/20$. Therefore, the energy transfer rate can be re-written as

$$n_\pi n_a \langle \sigma_{el} v \Delta E \rangle \simeq - \int \frac{d^3 p_{\pi_1}}{2E_{\pi_1}} \frac{p_{\pi_1}^2}{2m_\pi} C[f_{\pi_1}]. \quad (2.35)$$

The form of the collision term in the integrand will depend on whether the axions are relativistic or not at around the time of freeze-out.

In the regime where the axions are still relativistic at the time of freeze-out, there are well known methods for computing the collision term [187]. For pion-axion scattering in this regime, the collision term takes the form

$$C[f_{\pi_1}] = \frac{\pi g_a^2 m_\pi^6}{360 f_{a\pi}^4} \left(\frac{T_a}{m_\pi} \right)^4 \left(m_\pi T_a \nabla_{p_{\pi_1}}^2 + \vec{p}_{\pi_1} \cdot \vec{\nabla}_{p_{\pi_1}} + 3 \right) f_{\pi_1}(T_\pi). \quad (2.36)$$

Integrating over the pion phase space in Eq. (2.35) yields

$$n_\pi n_a \langle \sigma_{el} v \Delta E \rangle = \frac{\pi g_a^2 m_\pi^5}{120 f_{a\pi}^4} \left(\frac{T_a}{m_\pi} \right)^4 (T_\pi - T_a). \quad (2.37)$$

While the $3 \rightarrow 2$ is actively depleting the number density, the pions are nonrelativistic and follow MB statistics, which means that their energy density Boltzmann equation can then be expressed as

$$\frac{\partial T_\pi}{\partial T_a} = 3 \frac{T_\pi^2}{m_\pi T_a} + \frac{\pi g_a^2 m_\pi^5}{120 f_{a\pi}^4 H|_{T=m_\pi}} \frac{T_a T_\pi^2 (T_\pi - T_a)}{m_\pi^4}. \quad (2.38)$$

The first term on the right-hand side comes from the $3 \rightarrow 2$ and causes the pion temperature to increase, while the second term comes from pion-axion elastic scattering and pushes $T_\pi \rightarrow T_a$. The second term cannot keep up with the first as the temperature drops and the pions and axions decouple. The temperature of decoupling is

$$T_D \simeq m_\pi \left(\frac{\pi g_a^2 m_\pi^5}{120 f_{a\pi}^4 H|_{T=m_\pi}} \right)^{-1/4}. \quad (2.39)$$

The standard result for the collision term derived in Ref. [187] applies only when the axion is relativistic. However, as long as the momentum transfer is still small, we can still apply the same methods as Ref. [187] in deriving the collision term. We will be interested in the regime where the axion mass is still smaller than the pion mass (which kinematically enforces that the momentum transferred in a single collision is relatively small) but where the axion is sufficiently heavier than the freeze-out temperature $T_F \sim m_\pi/20$ such that it

becomes Boltzmann suppressed. Most generally, the collision term for $2 \rightarrow 2$ scattering of pions and axions can be written as

$$C = \frac{1}{2} \int \frac{d^3 p_{a_1}}{2E_{a_1}} \frac{d^3 p_{a_2}}{2E_{a_2}} \frac{d^3 p_{\pi_2}}{2E_{\pi_2}} (2\pi)^4 \delta^{(4)}(p_{a_1} + p_{\pi_1} - p_{a_2} - p_{\pi_2}) |\mathcal{M}|^2 J \quad (2.40)$$

where J is the relevant combination of phase space factors. In the regime of interest, everything is MB distributed at thermal decoupling so

$$J = e^{-E_{\pi_1}/T_\pi} e^{-E_{a_1}/T_a} - e^{-E_{\pi_2}/T_\pi} e^{-E_{a_2}/T_a}. \quad (2.41)$$

The collision term can be written as an expansion in the momentum transfer, $C = \sum C^j$ where

$$C^j = \frac{(2\pi)^4}{2j!} \int \frac{d^3 p_{a_1}}{2E_{a_1}} \frac{d^3 p_{a_2}}{2E_{a_2}} \frac{d^3 p_{\pi_2}}{2E_{\pi_2}} \left[\delta(E_{a_1} + E_{\pi_1} - E_{a_2} - E_{\pi_2}) |\mathcal{M}|^2 \right. \\ \left. \times J((\vec{p}_{a_2} - \vec{p}_{a_1}) \cdot \vec{\nabla}_{p_{\pi_2}})^j \delta^{(3)}(p_{\pi_1} - p_{\pi_2}) \right]. \quad (2.42)$$

In this expansion, C^0 vanishes simply because if the momentum transfer is zero and the number of a species does not change, the collision term is identically zero. For a contact interaction which has no angular dependence (as in the scenario we consider here, $|\mathcal{M}|^2 \sim \text{const.}$), C^1 also vanishes because the angular integral contains the integrand $(\vec{p}_{a_2} - \vec{p}_{a_1}) \cdot \vec{p}_\pi \equiv \vec{q} \cdot \vec{p}_{\pi_1}$, which is odd over the angular domain. Therefore, the leading-order term is C^2 . The momentum transfer scales like $\Delta p_a \sim (m_\pi p_a - m_a p_\pi)/(m_a + m_\pi)$, and plugging in thermal values for the typical momentum indicates that when truncated at $\mathcal{O}((\Delta p_a/p_\pi)^2)$, the expansion in is accurate at the $\sim 10\%$ level when the axion mass is $m_a \lesssim m_\pi/3$.

Following Ref. [187] and plugging in the matrix element of Eq. (2.31), the leading order piece of the collision term is then

$$C^2 = \frac{\pi m_\pi^4}{8g_\pi(2\pi)^3 f_{a\pi}^4} \int \frac{d^3 p_{a_1}}{2E_{a_1}} \int d\Omega_2 dE_{a_2} p_{a_2} \left[\frac{2(\vec{q} \cdot \vec{p}_{\pi_1})^2}{E_{\pi_1}^3} J \partial_{E_{a_2}}^2 \delta(E_{a_1} - E_{a_2}) \right. \\ \left. + \left(\frac{q^2}{E_{\pi_1}^2 T_\pi} - \frac{(\vec{q} \cdot \vec{p}_{\pi_1})^2}{E_{\pi_1}^3 T_\pi^2} - \frac{3(\vec{q} \cdot \vec{p}_{\pi_1})^2}{E_{\pi_1}^4 T_\pi} \right) J' \delta(E_{a_1} - E_{a_2}) \right. \\ \left. + \left(\frac{q^2}{E_{\pi_1}^2} J - \frac{3(\vec{q} \cdot \vec{p}_{\pi_1})^2}{E_{\pi_1}^4} J + \frac{2(\vec{q} \cdot \vec{p}_{\pi_1})^2}{E_{\pi_1}^3 T_\pi} J' \right) \partial_{E_{a_2}} \delta(E_{a_1} - E_{a_2}) \right] \\ = \frac{2m_\pi^4 e^{-E_{\pi_1}/T_\pi} e^{-m_a/T_a} (T_a - T_\pi) T_a^3 (m_a^2 + 3m_a T_a + 3T_a^2)}{g_\pi(2\pi)^3 f_{a\pi}^4 E_{\pi_1}^2 T_a T_\pi} \quad (2.43)$$

where $J' \equiv e^{-E_{\pi_2}/T_\pi} e^{-E_{a_2}/T_a}$. This feeds into the calculation of the energy transfer rate of

Eq. (2.35) under the assumption that the pions are non-relativistic,

$$\begin{aligned} n_\pi n_a \langle \sigma v \Delta E \rangle &= \frac{6e^{-m_\pi/T_\pi} e^{-m_a/T_a} (T_\pi - T_a) T_a^3 (m_a^2 + 3m_a T_a + 3T_a^2) \sqrt{\pi(m_\pi T_\pi)^5}}{g_\pi \sqrt{2} (2\pi)^5 f_{a\pi}^4 T_a T_\pi} \\ &= \frac{3m_\pi e^{-m_a/T_a} (T_\pi - T_a) T_a^2 (m_a^2 + 3m_a T_a + 3T_a^2)}{g_\pi (2\pi)^3 f_{a\pi}^4} n_\pi \end{aligned} \quad (2.44)$$

By analogy to Eq. (2.38), the term in the Boltzmann equations that equalizes temperatures between the two sectors is

$$\frac{3g_a^2 m_\pi^3 m_a^2}{(2\pi)^3 f_{a\pi}^4 H|_{T=m_\pi}} \left(\frac{T_a}{m_\pi} \right)^{-1} \left(\frac{T_\pi}{m_\pi} \right)^2 \left(\frac{T_\pi - T_a}{m_\pi} \right) e^{-m_a/T_a} \quad (2.45)$$

and decoupling happens when this is order unity. In summary, in the regime where $T_F < m_a \lesssim m_\pi$, we find that decoupling occurs when

$$\frac{3g_a^2 m_\pi m_a^2 T_D^2 e^{-m_a/T_D}}{(2\pi)^3 f_{a\pi}^4 H|_{T=m_\pi}} \sim 1, \quad T_F < m_a \lesssim m_\pi. \quad (2.46)$$

We find that the low- and high-mass axion decoupling temperatures, given by Eqs. (2.39) and (2.46), match onto each other exactly when a relative factor $\pi^4/90$ is taken into account due to the different numerical prefactors between Bose-Einstein and Maxwell-Boltzmann statistics. Additionally, we have checked numerically that the requirement on thermalization between the axions and pions, $T_D < T_F$, does keep the dark matter cool.

In addition to the above requirements, the decay constant $f_{a\pi}$ must be greater than the cutoff scale of chiral symmetry breaking. Otherwise, the description in Eq. (2.1) breaks down. We require that $f_{a\pi} \gtrsim 2\pi f_\pi$, where f_π is determined for a given m_π from the solution to the Boltzmann equation. Since the $\text{Sp}(2N_c)$ gauge theory with $N_f = 2$ we discuss here points to the strongly interacting regime where $m_\pi \sim 2\pi f_\pi$, we require that $f_{a\pi} \gtrsim m_\pi$. In practice, however, suppressing $2 \rightarrow 2$ annihilations at freeze-out is always a stronger requirement.

An additional preference, though not a requirement, comes from considering how chiral symmetry breaking contributes to the axion mass in Eq. (2.9),

$$\Delta m_a^2 = \frac{2m_\pi^2 f_\pi^2}{f_{a\pi}^2}. \quad (2.47)$$

The natural range for the axion mass-squared is therefore where $\Delta m_a^2 \lesssim m_a^2$, such that no fine tuning is present against an unspecified negative contribution, possibly from another confining gauge theory with $\theta \approx \pi$.

Satisfying the above requirements on $f_{a\pi}$ as a function of m_a for a variety of dark matter masses m_π yields the viable SIMP regions depicted in Fig. 2.1. We take $g_{*,F} = 10.75$ at freeze-out since for the dark matter masses we consider, freeze-out happens below the temperature of muon-antimuon annihilation. We learn that viable SIMP-axion thermalization is achieved over a broad range of axion masses and couplings $f_{a\pi}$.

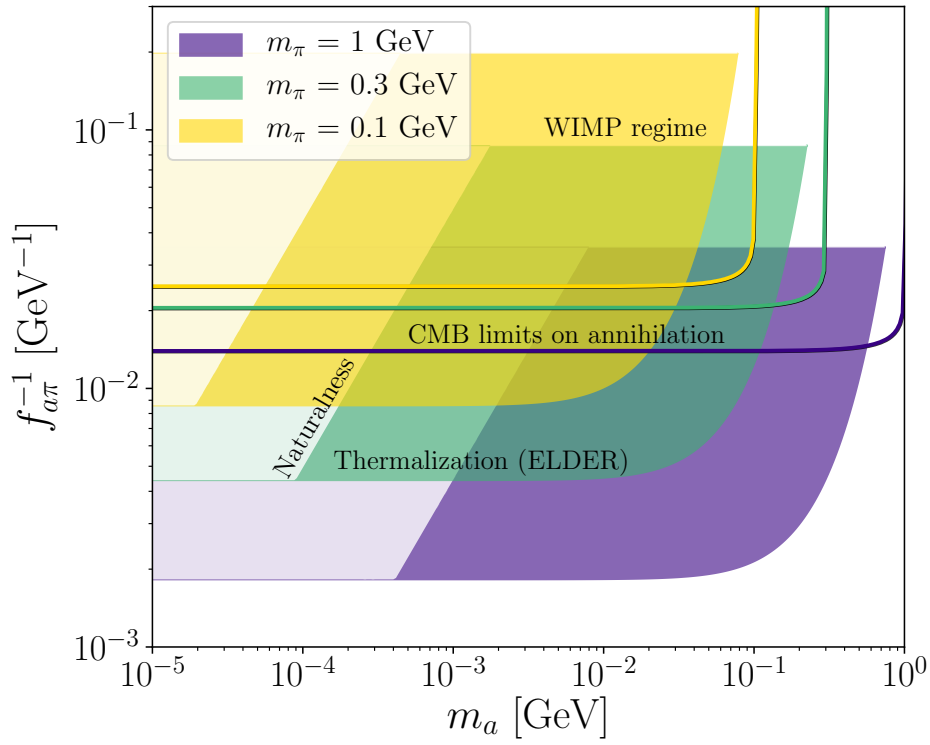


Figure 2.1: The parameter space for axions coupling to pions. The shaded regions correspond to the regions where the SIMP mechanism is theoretically viable for a given dark matter mass m_π . The boundaries of this region are set by requiring that the $3 \rightarrow 2$ rather than the $2 \rightarrow 2$ process sets the relic abundance (labeled as “WIMP regime”) and that the pions can transfer their excess heat to the axions and hence the Standard Model sector (labeled as “Thermalization”). Note that the parameter space along the “Thermalization” curve corresponds to the scenario where dark matter is an elastically decoupled relic (ELDER). We also indicate the natural mass range where the axion mass is at least as large as its contribution from chiral symmetry breaking (labeled as “Naturalness”). Also shown are the empirical upper limits on pion annihilation from energy injection into the CMB (thick solid lines labeled as “CMB limits on annihilation”).

We note that elastically decoupling relic (ELDER) dark matter [130, 186] is obtained along the thermalization curve in Fig. 2.1. For ELDER dark matter, the kinetic decoupling between the dark matter and Standard Model baths occurs before $3 \rightarrow 2$ pion self-annihilations freeze out. This causes the relic abundance to be exponentially sensitive to this elastic scattering while being relatively insensitive to the strength of the $3 \rightarrow 2$ pion process. On the thermalization curve in Fig. 2.1, the elastic scattering of pions off of axions dominates over the $3 \rightarrow 2$ pion self-annihilation process in setting the relic abundance.

Empirical Requirements

Having established the theoretical requirements on the axion-SIMP parameter space, we now move to the observational constraints coming from the CMB. In standard cold dark matter cosmology, the intergalactic medium (IGM) is almost entirely neutral after recombination and CMB photons free stream. If some fraction of the dark matter annihilates to Standard Model particles and partially ionizes the IGM, this will cause some CMB photons to re-scatter which modifies CMB anisotropies in a characteristic way. For the scenario we consider in this Chapter, the process of interest is $\pi\pi \rightarrow aa \rightarrow 4\gamma$. In the parameter space where there is sufficient thermalization between axions and the Standard Model (see Section 2.5), the decay of the intermediate-state axions happens immediately. Thus, the use of the narrow-width approximation is appropriate and the cross section for this process is set by the cross section for the annihilation process $\pi\pi \rightarrow aa$. We use limits derived in Ref. [188], which are not very sensitive to whether there are two final-state photons or four ². The resulting upper limits are shown in Fig. 2.1 as a set of thick, solid lines corresponding to the different depicted pion masses. We thus find a viable SIMP-axion parameter space below the CMB curve and above the thermalization curve.

2.5 Axions and Photons in Equilibrium

Theoretical Requirements

Fig. 2.1 presents the viable region where the SIMP and axion maintain thermal equilibrium. If the axion and Standard Model also maintain thermal equilibrium via the axion-photon coupling in Eq. (2.11), then the pions will share a temperature with the Standard Model. For most of the axion masses we consider, decays and inverse decays into Standard Model photons are the most efficient processes for kinetic equilibrium with the Standard Model at freeze-out. The rate for these decays at rest is

$$\Gamma_a = \frac{m_a^3}{64\pi f_{a\gamma}^2}. \quad (2.48)$$

In order for the pions to maintain thermal equilibrium with the Standard Model, two conditions must be satisfied: the decays and inverse decays of the axions need to thermalize the axions with the Standard Model, and the axions need to lose kinetic energy via decays faster than they gain energy from the pion $3 \rightarrow 2$ heating. We have verified numerically that as long as these conditions are satisfied at the freeze-out temperature of the pion, the relic abundance of dark matter is unaffected and the pions constitute cold dark matter.

To understand the requirement that axions maintain thermal contact with the Standard Model, we can ignore the pions and consider only the relevant Boltzmann equations for the

²This conclusion was reached after private communication with T. R. Slatyer.

axions:

$$\frac{\partial n_a}{\partial t} + 3Hn_a = -\Gamma_a m_a \left(\langle E_a^{-1} \rangle_{T_a} n_a - \langle E_a^{-1} \rangle_{T_{\text{SM}}} n_a^{\text{eq}, T_{\text{SM}}} \right) \quad (2.49)$$

$$\frac{\partial \rho_a}{\partial t} + 3H(\rho_a + P_a) = -m_a \Gamma_a (n_a - n_a^{\text{eq}, T_{\text{SM}}}) \quad (2.50)$$

where the average axion energy is $\langle E_a \rangle = \rho_a / n_a$. To make the notation less cumbersome, the label for equilibrium distributions will denote chemical equilibrium *and* kinetic equilibrium between the axions and Standard Model, *i.e.* $T_a = T_{\text{SM}} \equiv T$. With this notation, these equations can be re-expressed as

$$-T \frac{\partial n_a}{\partial T} + 3n_a = -\frac{m_a \Gamma_a}{H} n_a \left(\langle E_a^{-1} \rangle - \langle E_a^{-1} \rangle^{\text{eq}} \frac{n_a^{\text{eq}}}{n_a} \right) \equiv -\frac{m_a \Gamma_a}{H} n_a c_n \quad (2.51)$$

$$-T \frac{\partial \langle E_a \rangle n_a}{\partial T} + 3 \langle E_a \rangle n_a (1 + w_a) = -\frac{m_a \Gamma_a}{H} n_a \left(1 - \frac{n_a^{\text{eq}}}{n_a} \right) \equiv -\frac{m_a \Gamma_a}{H} n_a c_\rho, \quad (2.52)$$

where w_a is the equation of state of the axion, which is a function of time and axion temperature $w_a = w_a(T_a)$. In order to eliminate $\partial n_a / \partial T$, Eqs. (2.51) and (2.52) can be combined to give a differential equation for $\langle E_a \rangle$:

$$\frac{\partial \langle E_a \rangle}{\partial T} = 3w_a \frac{\langle E_a \rangle}{T} - \frac{m_a \Gamma_a}{TH} (\langle E_a \rangle c_n - c_\rho). \quad (2.53)$$

In the first term, the expansion is driving the change in the average energy, while in the second term, the decay is driving the average energy. One can check that the first term matches the expectations for a decoupled particle. Meanwhile, for further examination of the second term, we define the variable $\alpha = \alpha(T_a)$ such that

$$\langle E_a^{-1} \rangle \equiv \alpha / \langle E_a \rangle. \quad (2.54)$$

The value of α changes monotonically $\alpha \in [1, \pi^6 / (360 \zeta(3)^2)]$ as T_a goes from 0 to ∞ . Using this definition, we find

$$\frac{\partial \langle E_a \rangle}{\partial T} = 3w_a \frac{\langle E_a \rangle}{T} + \frac{m_a \Gamma_a}{TH} \left[(1 - \alpha) - \frac{n_a^{\text{eq}}}{n_a} \left(1 - \alpha \frac{\langle E_a \rangle}{\langle E_a \rangle^{\text{eq}}} \right) \right]. \quad (2.55)$$

The second term vanishes when the particle is in equilibrium, *i.e.*, $n_a = n_a^{\text{eq}}$ with $T_a = T$. If the particle is driven out of equilibrium (for instance by the expansion), this term will push it back into equilibrium. In order to overcome the expansion, Γ_a needs to be large enough so that the second term is larger than the first.

First we consider the case that $m_a \ll T$ such that $w_a = 1/3$ and $\alpha = \pi^6 / (360 \zeta(3)^2)$. Assuming the axion is near equilibrium and expanding Eq. (2.55) around $T_a = T$ to leading order gives

$$\frac{\partial \langle E_a \rangle}{\partial T} = \frac{\langle E_a \rangle}{T} - \frac{m_a \Gamma_a}{TH} (2\alpha - 3) \frac{(T_a - T)}{T}. \quad (2.56)$$

If the two temperatures differ, then the second term will drive the system back into equilibrium if it is comparable to the first,

$$\frac{m_a \Gamma_a}{TH} \gtrsim \frac{\langle E_a \rangle}{(2\alpha - 3)T} \simeq 4 \quad m_a \ll T. \quad (2.57)$$

Below, we find the strongest requirement on kinetic equilibrium when $m_a \gtrsim m_\pi$. For intermediate masses $T \lesssim m_a \lesssim m_\pi$, we analytically continue the condition in Eq. (2.57) and require

$$\frac{\Gamma_a m_a e^{-m_a/2T_F}}{4T_F H|_{T_F}} \gtrsim 1. \quad (2.58)$$

We have checked numerically that this requirement ensures thermal equilibrium between the axions and Standard Model for the entire region $m_a \lesssim m_\pi$ in which this is the strongest condition since the axion is relatively light and abundant.

In the regime where the axion mass is comparable to or larger than that of the pion, a second condition on their decay becomes stronger than just pure thermalization between the axions and Standard Model. The rate of kinetic energy density loss through decays for non-relativistic axions is $\Gamma_a n_a T^2/m_a$. Meanwhile, for the axions to sufficiently cool the pions, we require that $\langle \sigma v \Delta E \rangle n_a n_\pi \gtrsim m_\pi \dot{n}_\pi \sim H m_\pi^2 n_\pi / T$ at freeze-out when the pions are still barely in chemical equilibrium. Therefore, the requirement is

$$\frac{\Gamma_a T_F^2 n_a}{m_a} \gtrsim \frac{H|_{T_F} m_\pi^2 n_\pi}{T_F}. \quad (2.59)$$

This second condition requires that the kinetic energy transferred to the axions from pion $3 \rightarrow 2$ can be compensated by axion decays. This matters more for higher axion masses $m_a \gtrsim m_\pi$ since the axion number density is lower than the pion number density, so that each axion decay must transfer several pions' entropy. We have numerically checked that this condition keeps the pion, axion, and Standard Model at the same temperature in the regime $m_a \gtrsim m_\pi$.

Decays and inverse decays come into equilibrium at late times. *A priori*, this could suggest the need for Eqs. (2.58) and (2.59) to hold prior to freeze-out in order to sufficiently transfer entropy from the annihilating pions into the Standard Model particles. However, we verified numerically that this is not the case: the SIMP relic abundance is unaffected if decays and inverse decays into Standard Model particles only come into equilibrium close to the time of freeze-out.

For axions with masses below the freeze-out temperature, the scattering process $ae \rightarrow \gamma e$ is more efficient than the decays considered above. This arises because the rate of scattering is enhanced relative to decays like $\sim (T/m_a)^3$ for axions that are relativistic at the time of freeze-out, while the rate of decays is suppressed due to the boost factor of $\sim (m_a/T)$. We find that the parameter space with $m_a < T_F$ is tightly constrained for the pion masses we consider and therefore do not include the scattering process $ae \rightarrow \gamma e$ in our analytics or

numerics, since it is expected to be subdominant for axion masses with $m_a \gtrsim T_F$. Note that by including decays as the only channel to transfer entropy, we are being conservative since adding the $ae \rightarrow \gamma e$ channel would only lower the required coupling strength between axions and the Standard Model.

In Fig. 2.2, we depict the requirement on $f_{a\gamma}$ such that decays and inverse decays sufficiently transfer entropy between the sectors. Each solid curve corresponds to the lower bound on $f_{a\gamma}^{-1}$ to maintain thermal contact for a fixed pion mass. We use the full Boltzmann equations and full energy transfer rates. A crossover between two regimes occurs at $m_a \sim m_\pi$, where the lower axion number density starts to matter and Eq. (2.59) becomes a stronger condition than Eq. (2.58). In the regime where $m_a \gtrsim m_\pi$, kinetic equilibrium is maintained by axions in the exponential tail of the distribution, which causes the precipitous increase in $f_{a\gamma}^{-1}$. As is evident, kinetic equilibrium between the axion and the Standard Model through decays and inverse decays is possible over a range of axion masses.

The conditions outlined above amount to requiring that the axions and Standard Model have the same temperature at freeze-out; however, there is another possibility which we outline here. For the dark matter to cool, it only needs to transfer entropy to the axions throughout the freeze-out process. Then, instead of keeping the axions and Standard Model at the same temperature, axions can decay out of equilibrium into the Standard Model at some later time. Relative to the scenario where the axions are thermalized with the Standard Model, the pion-axion sector will be slightly hotter than the Standard Model and the relic abundance will be slightly larger for the same value of $3 \rightarrow 2$ rate. Therefore if the axions and Standard Model are not thermalized at freeze-out, the value of f_π must be increased slightly to give the right relic abundance. For sufficiently large m_a , the universe undergoes a brief matter-dominated phase where the axions dominate the energy density of the universe. When the axions decay, they reheat the Standard Model components, the universe becomes radiation-dominated again, and the pion abundance is diluted. This happens when

$$H(T_{\text{RH}}) = \sqrt{\frac{g_{*,\text{RH}}\pi^2}{90} \frac{T_{\text{RH}}^2}{M_{\text{Pl}}}} \sim \Gamma_a. \quad (2.60)$$

We require that the reheat temperature be larger than the temperature of neutrino decoupling: if the reheat temperature is lower, then the photons get preferentially heated and the effective number of relativistic neutrinos (N_{eff}) becomes smaller than allowed by observations of the CMB [189]. We take the neutrino decoupling temperature to be ~ 3 MeV [190], at which point $g_* = 10.75$. Having such a high reheat temperature also enforces that the decay products do not affect BBN. Therefore, the region in Figs. 2.2 between the solid curves and straight line labeled “Reheat before neutrino decoupling” may be viable with a slightly different $3 \rightarrow 2$ cross-section than in the standard SIMP scenario. More dedicated numerical studies are necessary in this case and will be explored in future work.

Empirical Requirements

Having established a theoretically-viable parameter space, we must check whether it is allowed by current experiments and observations. Constraints arise from early universe cosmology, astrophysical bodies, and terrestrial experiments.

Light Degrees of Freedom

If light axions are in thermal equilibrium with the Standard Model bath, a bound on their mass arises from their effect on the temperature ratio T_ν/T_γ after neutrinos have decoupled. This difference alters the effective number of neutrino species contributing to the radiation density, N_{eff} , which can be measured in the CMB by comparing the photon diffusion scale to the sound horizon scale [189, 191]. Such constraints are relevant for light particles in equilibrium with the photon or electron plasma beneath the temperature of neutrino decoupling unless the particle couples to neutrinos as well. When applicable, this bound is stronger than the BBN bound of $\sim \text{MeV}$, which comes from the fact that changes to N_{eff} change the expansion history and hence modify the abundance of the light elements. Because of this bound, only values of $m_a > 2.6 \text{ MeV}$ are shown Fig. 2.2.

Supernova 1987a

The direct coupling of the axion to photons can lead to excess emission from supernovae (SN) via the Primakoff scattering process [192]. When the coupling between axions and the Standard Model is sufficiently strong, the scattering of $e\gamma \rightarrow ea$ produces axions in the stellar medium which leads to excess cooling if the axions escape the SN. However, if the coupling is too strong, then trapping occurs via the inverse $ea \rightarrow e\gamma$ process along with axion decays, in which case the axion does not carry any energy out of the star. SN cooling primarily proceeds through neutrino emission; due to the observed neutrino signal from SN 1987A, any new SN cooling process must carry away less energy than the neutrinos, $\sim 3 \times 10^{53}$ ergs. The region of parameter space excluded by the excess cooling of SN 1987A [193] is shown in Fig. 2.2. For photon couplings that are too weak to produce significant energy loss in the supernova, there are still constraints from escaping axions decaying into an observable burst of photons [194], which we also show in Fig 2.2.

Terrestrial

The couplings between axions and Standard Model particles are constrained by terrestrial experiments. However, these constraints often come with assumptions about how the axions interact with the Standard Model. We classify constraints on the axion-photon coupling based on different assumptions about its fundamental origin, namely that the photon coupling arises from:

1. solely coupling to $U(1)_Y$;

2. solely coupling to $SU(2)_W$;
3. equal couplings to $U(1)_Y$ and $SU(2)_W$, in which case, the $aZ\gamma$ coupling vanishes.

Measurements from the LEP collider and CDF constrain the decay $Z \rightarrow \gamma\gamma(\gamma)$ [193, 195, 196] as shown in Fig. 2.2. BaBar also constrains the decay $Z \rightarrow \gamma + \text{invisible}$ [193]. In the third case above, in which the $aZ\gamma$ coupling vanishes due to equal couplings to $U(1)_Y$ and $SU(2)_W$, both of these Z decay constraints are alleviated. In their place, there is a LEP bound on $e^+e^- \rightarrow \gamma\gamma$ [196] and a BaBar bound on $e^+e^- \rightarrow \gamma + \text{inv}$ [193]. Constraints from electron beam dump experiments SLAC 137, SLAC 141 [193, 197–199], and proton beam dump experiments CHARM and NuCal [200, 201] apply for axions coupled to photons regardless of how the coupling arises. Constraints from $K_L \rightarrow \pi^0 a$ and $K^\pm \rightarrow \pi^\pm a$ with $a \rightarrow \gamma\gamma$ assume the axion couples to $SU(2)_W$. These kaon results were obtained in Ref. [202] from analyses of fixed-target kaon rare decay experiments by the E949 [203], NA62 and NA48/2 [204], and KTeV [205] collaborations.

In addition to existing constraints, we show the projected reach of several future experiments and analyses on the photon coupling to an axion, indicated by the dashed black curves in Fig. 2.2. We include the projected reach of SHiP [201, 206], NA62 [201], BaBar [202], Belle II [193, 202], SeaQuest [207] and FASER [208]. In principle there could be a constraint from a process involving the aZZ coupling for all three scenarios, though we expect it would be weaker than the constraints we present and at this time we are not aware of any existing or projected limits from such a process.

2.6 Discussion

In this Chapter, we have considered the pion realization of SIMP dark matter in strongly coupled gauge theories, and have shown that it can be realized with axions as the thermalization portal between the dark matter and Standard Model. Throughout this Chapter, we have required that all three sectors — the SIMPs, axions, and Standard Model — share the same temperature as the $3 \rightarrow 2$ annihilations freeze out. This requirement sets a target range of masses and couplings for this mechanism to be theoretically viable.

In examining the couplings between the SIMPs and axions, we have required that the coupling is strong enough to thermalize the two sectors via $2 \rightarrow 2$ scattering. At the same time, we require that the coupling not be strong enough for $2 \rightarrow 2$ annihilations to overwhelm the $3 \rightarrow 2$ process that is the hallmark of the SIMP mechanism. Combined, these requirements lead to a well-defined range of couplings between the pion dark matter and the axions such that the SIMP mechanism can work. Constraints on annihilation coming from the CMB narrow the allowed range of couplings, though a broad parameter space remains. It is possible that a future CMB spectral distortions experiment can probe this parameter space further, though exploring this possibility is beyond the scope of this Chapter.

Considering the couplings between the axions and the Standard Model, we focused on the coupling to photons. For a given pion mass, there is a range of axion masses allowed

for maintaining SIMP-axion equilibrium. Within this axion mass range, the main requirement for axion-Standard Model thermalization is that the axions decay quickly enough to successfully transfer the entropy from the pions to the Standard Model, which can easily be achieved. The relevant couplings to photons can be probed in a multitude of ways. The range of axion masses considered here are at an energy scale that is relevant for supernovae, which constrains part of the parameter space. Additionally, terrestrial beam dump and collider experiments have probed complementary parameter space. We find that the SIMP mechanism can be realized in a broad region of parameter space that is not excluded by current constraints. Several upcoming experiments are forecast to probe much of the viable parameter space that is currently allowed, providing an excellent handle for testing the framework.

There are several possible ways to extend the parameter space for axion-mediated SIMPs. Some of these possibilities are already excluded by existing limits. For instance, heavy axions which mediate the entropy transfer through off-shell interactions (both through the interactions we consider here and through the CP-violating interaction $\mathcal{L} \supset \frac{f_{\text{CPV}}}{2} a \pi^b \pi^b$) are excluded by the LEP constraint shown in Fig. 2.2. Another heavily-constrained scenario is that the axions couple to all fermions through a universal Yukawa coupling, which is almost entirely ruled out by SLAC 137, CHARM, kaon decays, B decays, supernova 1987a, BaBar and the muon anomalous magnetic moment [209–211]. A more promising possibility is that axions couple only to electrons or to charged leptons; in this case, there are known limits from SLAC 137 and the muon anomalous magnetic moment [209, 211]. In addition, the axion-electron parameter space should be constrained by limits from supernova 1987a and from loop-induced couplings to photons. However, such constraints have not been explored in the literature and will be the subject of future work. Another possibility is that the axions have a long enough lifetime that they decay out of equilibrium, dumping entropy to the Standard Model and diluting the SIMPs— this too will be explored in future work.

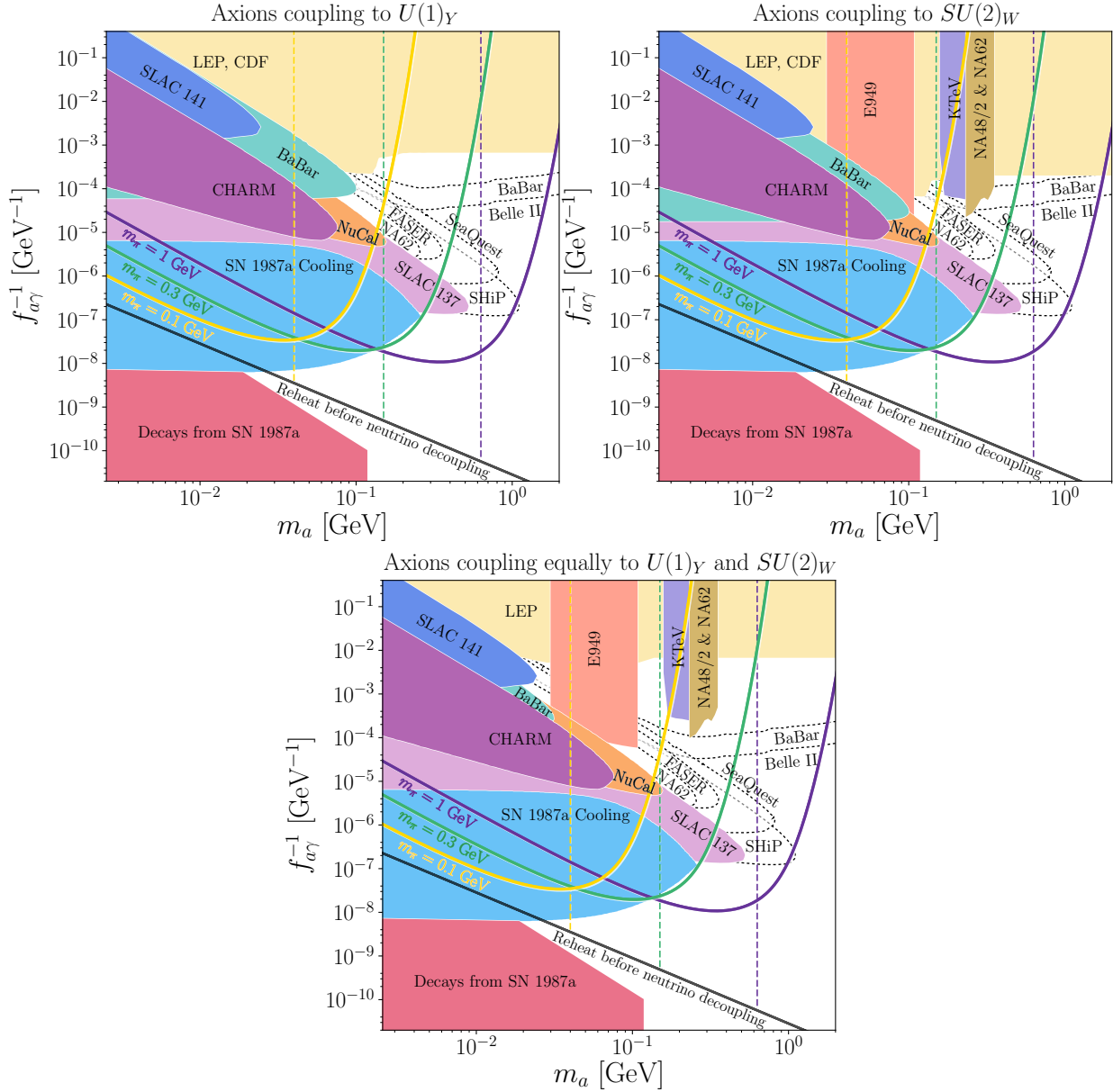


Figure 2.2: The parameter space for axions coupling to photons through $U(1)_Y$ (top left), $SU(2)_W$ (top right), and both in equal amounts (bottom). Solid curves correspond to the lower bound on the decay rate for thermalization between the axions and Standard Model for various pion masses. Vertical, dashed lines correspond to the largest axion mass allowed by CMB constraints for a given pion mass (see Fig. 2.1). Below the thermalization curves, the SIMP mechanism may still be viable down to the black solid line if the axion decays out of equilibrium and reheats above the neutrino decoupling temperature. Solid, filled regions correspond to existing constraints. Regions enclosed by dotted, black lines correspond to projected reach by proposed experiments. See the main text for more details.

Chapter 3

Freezing-in Twin Dark Matter

3.1 Introduction

As discussed in Chapter 1, the lack of new TeV-scale physics at the LHC has motivated solutions to the electroweak hierarchy problem which belong to the ‘Neutral Naturalness’ paradigm: the new particles responsible for stabilizing the electroweak scale are not charged under (some of) the Standard Model gauge symmetries [88–100]. The first and perhaps most aesthetically pleasing of these Neutral Naturalness approaches is the mirror twin Higgs (MTH) [88], which stabilizes the Higgs potential up to a cutoff $\Lambda \sim 5 - 10$ TeV. In the MTH, the states responsible for ensuring naturalness comprise another copy of the Standard Model related to ours by a $\mathbb{Z}_2$ symmetry. Since these ‘twin’ states are neutral under the Standard Model gauge group, they easily evade LHC constraints. The most serious empirical challenges to these models comes instead from cosmology. In particular, the presence of thermalized light twin neutrinos and photons [212] is significantly ruled out by BBN [213] and the CMB [11].

This cosmological tension has been alleviated with the observation that an exact $\mathbb{Z}_2$ symmetry is not truly necessary for naturalness. This was first exploited in [214], where only the third generation of Standard Model fermions was copied and twin hypercharge was not gauged. This philosophy has since been taken up by the authors of [215–218] who introduce clever $\mathbb{Z}_2$ -asymmetries into the MTH mass spectrum to alleviate cosmological problems. These models have proved to be a boon for phenomenology. Among other things, they quite generally motivate looking for Higgs decays to long-lived particles at colliders [219–226] and contain well motivated dark matter candidates [227–238].

In this Chapter, we build instead on a framework where ‘hard’ breaking of the $\mathbb{Z}_2$ is absent. In [212, 239], it was realized that late-time asymmetric reheating of the two sectors could arise naturally in these models if the spectrum were extended by a single new state. This asymmetric reheating would dilute the twin energy density and so attune the MTH with the cosmological constraints. This dilution of twin energy density to negligible levels would seem to hamper the prospect that twin states might constitute the dark matter,

and generating dark matter was left as an open question. This presents a major challenge toward making such cosmologies realistic. However, we show that asymmetric reheating perfectly sets the stage for a MTH realization of the ‘freeze-in’ mechanism for dark matter production [136, 240–246].

Freeze-in scenarios are characterized by two assumptions: 1) dark matter has a negligible density at some early time and 2) dark matter interacts with the Standard Model so feebly that it never achieves thermal equilibrium with the Standard Model bath.¹ This second assumption is motivated in part by the continued non-observation of non-gravitational dark matter-Standard Model interactions (see Chapter 1). Both assumptions stand in stark contrast to freeze-out scenarios such as the one considered in the previous chapter.

Freeze-twin dark matter is a particularly interesting freeze-in scenario because both assumptions are fulfilled for reasons orthogonal to dark matter considerations: 1) the negligible initial dark matter abundance is predicted by the asymmetric reheating already necessary to resolve the MTH cosmology, and 2) the kinetic mixing necessary to achieve the correct relic abundance is of the order expected from infrared contributions in the MTH. To allow the frozen-in twin electrons and positrons to be dark matter, we need only break the $\mathbb{Z}_2$ by a relevant operator to give a Stueckelberg mass to twin hypercharge. Additionally, the twin photon masses we consider can lead to dark matter self-interactions at the level relevant for the small-scale structure problems discussed in Chapter 1.

This Chapter is structured as follows. In Section 3.2, we review the MTH and its cosmology in models with asymmetric reheating, and in Section 3.3 we introduce our extension and include some discussion of the irreducible IR contributions to kinetic mixing in the MTH. In Section 3.4, we calculate the freeze-in yield for twin electrons and discuss the parameter space to generate dark matter and constraints thereon. We discuss future directions and conclude in Section 3.5.

3.2 The Mirror Twin Higgs & Cosmology

The mirror twin Higgs framework [88] introduces a twin sector B , which is a ‘mirror’ copy of the Standard Model sector A , related by a $\mathbb{Z}_2$ symmetry. Upgrading the $SU(2)_A \times SU(2)_B$ gauge symmetry of the scalar potential to an $SU(4)$ global symmetry adds a Higgs-portal interaction between the A and B sectors:

$$V = \lambda (|\mathcal{H}|^2 - f^2/2)^2, \quad (3.1)$$

where $\mathcal{H} = \begin{pmatrix} H_A \\ H_B \end{pmatrix}$ is a complex $SU(4)$ fundamental consisting of the A and B sector Higgses in the gauge basis. The Standard Model Higgs is to be identified as a pseudo-Goldstone mode arising from the breaking of $SU(4) \rightarrow SU(3)$ when $\mathcal{H}$ acquires a vacuum expectation value

¹We note that the feeble connection between the two sectors may originate as a small dimensionless coupling or as a small ratio of mass scales, either of which deserves some explanation.

(vev) $\langle \mathcal{H} \rangle = f/\sqrt{2}$. Despite the fact that the global $SU(4)$ is explicitly broken by the gauging of $SU(2)_A \times SU(2)_B$ subgroups, the $\mathbb{Z}_2$ is enough to ensure that the quadratically divergent part of the one-loop effective action respects the full $SU(4)$. The lightness of the Standard Model Higgs is then understood as being protected by the approximate accidental global symmetry up to the UV cutoff scale $\Lambda \lesssim 4\pi f$, at which point new physics must come in to stabilize the scale f itself.

Current measurements of the Higgs couplings at the LHC are consistent with the Standard Model, implying $f \gtrsim 3v$ [214, 247], where v is the Standard Model Higgs vev. This requires some vacuum misalignment via soft $\mathbb{Z}_2$ -breaking which leads the two $SU(2)$ doublets to get vevs $v_A \approx v/\sqrt{2}$ and $v_B \approx f/\sqrt{2}$, to lowest order. The twin spectrum is thus a factor of f/v heavier than the standard model spectrum. We refer to twin particles by their Standard Model counterparts primed with a superscript $'$, and we refer the reader to [88, 214] for further discussion of the twin Higgs mechanism.

The thermal bath history in the conventional MTH is fully dictated by the Higgs portal in Eq. (3.1) which keeps the Standard Model and twin sectors in thermal equilibrium down to temperatures $\mathcal{O}(\text{GeV})$. A detailed calculation of the decoupling process was performed in [212] by tracking the bulk heat flow between the two sectors as a function of Standard Model temperature. It was found that for the benchmark of $f/v = 4$, decoupling begins at a Standard Model temperature of $T \sim 4 \text{ GeV}$ and by $\sim 1 \text{ GeV}$, the ratio of twin-to-Standard Model temperatures may reach $\lesssim 0.1$ without rebounding. While heat flow rates become less precise below $\sim 1 \text{ GeV}$ due to uncertainties in hadronic scattering rates, especially close to color-confinement, decoupling between the two sectors is complete by then for $f/v \gtrsim 4$. For larger f/v , the decoupling begins and ends at higher temperatures.

After the two sectors decouple, chemical processes in the two sectors change their temperatures independently. The twin sector eventually cools to a slightly lower temperature than the Standard Model due to the modification of mass thresholds. However, within a standard cosmology, this still leaves far too much radiation in the twin sector to agree with BBN and CMB observations. To quantify this tension, the effective number of neutrinos, N_{eff} , is defined in

$$\rho_r = \left(1 + \frac{7}{8} \left(\frac{4}{11} \right)^{4/3} N_{\text{eff}} \right) \rho_\gamma, \quad (3.2)$$

where ρ_r is the total radiation energy density, the factors $7/8$ and $(4/11)^{4/3}$ come from Fermi-Dirac statistics and Standard Model electron-positron annihilations, and ρ_γ is the Standard Model photon energy density. The Standard Model neutrinos contribute $N_{\text{eff}} \approx 3.046$ [248, 249], with additional radiative degrees of freedom collectively contributing to $\Delta N_{\text{eff}} \equiv N_{\text{eff}} - 3.046$. In a generic MTH model, the twin neutrinos and photon contribute $\Delta N_{\text{eff}} \sim 5 - 6$ [212, 239], significantly disfavored by both BBN [213] and the more stringent *Planck* measurement, $2.99^{+0.34}_{-0.33}$ (at 95% confidence) [11].²

²Care must be taken when applying this constraint since the twin neutrinos may be semi-relativistic by matter-radiation equality [212]. But, within a standard cosmology, the MTH is unambiguously ruled out.

As mentioned above, one class of solutions to this N_{eff} problem uses hard breaking of the $\mathbb{Z}_2$ at the level of the spectra [214–218] while keeping a standard cosmology. An alternative proposal is to modify the cosmology with asymmetric reheating to dilute the energy density of twin states. For example, [239] uses late, out-of-equilibrium decays of right-handed neutrinos, while [212] uses those of a scalar singlet. These new particles respect the $\mathbb{Z}_2$, but dominantly decay to Standard Model states due to the already-present soft $\mathbb{Z}_2$ -breaking in the scalar sector. In [239], this is solely due to extra suppression by f/v -heavier mediators, while in [212], the scalar also preferentially mass-mixes with the heavier twin Higgs. [212] also presented a toy model of ‘Twinflation’, where a softly-broken $\mathbb{Z}_2$ -symmetric scalar sector may lead to inflationary reheating which asymmetrically reheats the two sectors to different temperatures. In any of these scenarios, the twin sector may be diluted to the level where it evades *Planck* bounds [11] on extra radiation, yet is potentially observable with CMB Stage IV [250].

We will stay agnostic about the particular mechanism at play, and merely assume that by $T \sim 1$ GeV, the Higgs portal interactions have become inefficient and some mechanism of asymmetric reheating has occurred such that the energy density in the twin sector has been largely depleted, $\rho_{\text{twin}} \approx 0$.³ This is consistent with the results of the decoupling calculation in [212] given the uncertainties in the rates at low temperatures, and will certainly be true once one gets down to $\text{few} \times 10^2$ MeV.

One may be concerned that there will be vestigial model-dependence from irrelevant operators induced by the asymmetric reheating mechanism which connect the two sectors. However, these operators will generally be suppressed by scales above the reheating scale, as in the example studied in [212]. Prior to asymmetric reheating, the two sectors are in thermal equilibrium anyway, so these have little effect. After the energy density in twin states has been diluted relative to that in the Standard Model states, the temperature is far below the heavy masses suppressing such irrelevant operators, and thus their effects are negligible. So we may indeed stay largely agnostic of the cosmological evolution before asymmetric reheating as well as the details of how this reheating takes place. We take the absence of twin energy density as an initial condition, but emphasize that there are external, well-motivated reasons for this to hold in twin Higgs models, as well as concrete models that predict this occurrence naturally.

3.3 Kinetic Mixing & a Massive Twin Photon

In order to arrange for freeze-in, we add to the MTH kinetic mixing between the Standard Model and twin hypercharges and a Stueckelberg mass for twin hypercharge. At low energies, these reduce to such terms for the photons instead, parametrized as

$$\mathcal{L} \supseteq \frac{\epsilon}{2} F_{\mu\nu} F'^{\mu\nu} + \frac{1}{2} m_{\gamma'}^2 A'_\mu A'^\mu. \quad (3.3)$$

³If asymmetric reheating leaves some small $\rho_{\text{twin}} > 0$, then mirror baryon asymmetry can lead to twin baryons as a small subcomponent of dark matter [251].

This gives each Standard Model particle of electric charge Q an effective twin electric charge ϵQ .⁴ The twin photon thus gives rise to a ‘kinetic mixing portal’ through which the Standard Model bath may freeze-in light twin fermions in the early universe.

The Stueckelberg mass constitutes soft $\mathbb{Z}_2$ -breaking,⁵ but has no implications for the fine-tuning of the Higgs mass since hypercharge corrections are already consistent with naturalness [214]. We will require $m_{\gamma'} > m_{e'}$, to prevent frozen-in twin electron/positron annihilations, and $m_{\gamma'} > 2m_{e'}$, to ensure that resonant production through the twin photon is kinematically accessible. Resonant production will allow a much smaller kinetic mixing to generate the correct relic abundance, thus avoiding indirect bounds from supernova cooling. We note that while taking $m_{\gamma'} \ll f$ does bear explanation, the parameter is technically natural.

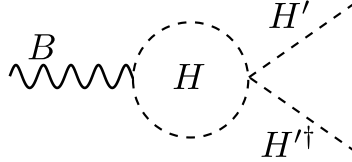
On the other hand, mixing of the twin and Standard Model $U(1)$ s preserves the symmetries of the MTH EFT, so quite generally one might expect it to be larger than that needed for freeze-in. However, it is known that in the MTH a nonzero ϵ is not generated through three loops [88]. While such a suppressed mixing is phenomenologically irrelevant for most purposes, it plays a central role in freeze-twin dark matter. Thus, we discuss here at some length the order at which it is expected in the low-energy EFT. Of course, there may always be UV contributions which set ϵ to the value needed for freeze-in. However, if the UV completion of the MTH disallows such terms - for example, via supersymmetry, an absence of fields charged under both sectors, and eventually grand unification in each sector (see e.g. [256–261])- then the natural expectation is for mixing of order these irreducible IR contributions.

To be concrete, we imagine that $\epsilon = 0$ at the UV cutoff of the MTH, $\Lambda \lesssim 4\pi f$. To find the kinetic mixing in the regime of relevance, at momenta $\mu \lesssim 1$ GeV, we must run down to this scale. As we do not have the technology to easily calculate high-loop-order diagrams, our analysis is limited to whether we can prove diagrams at some loop order are vanishing or finite, and so do not generate mixing. Thus our conclusions are strictly always ‘we know no argument that kinetic mixing of this order is not generated’, and there is always the possibility that further hidden cancellations appear. With that caveat divulged, we proceed and consider diagrammatic arguments in both the unbroken and broken phases of electroweak symmetry.

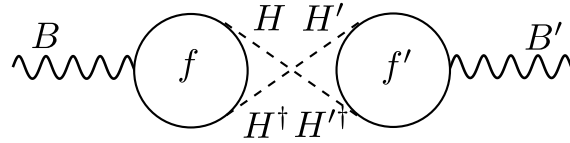
⁴Note that twin charged states do not couple to the Standard Model photon. Their coupling to the Standard Model Z boson has no impact on freeze-in at the temperatures under consideration. Furthermore, the miniscule kinetic mixing necessary for freeze-in has negligible effects at collider experiments. See Ref. [252] for details.

⁵While we are breaking the $\mathbb{Z}_2$ symmetry by a relevant operator, the extent to which a Stueckelberg mass is truly *soft* breaking is not clear. Taking solely Eq. (3.3), we would have more degrees of freedom in the twin sector than in the Standard Model, and in a given UV completion it may be difficult to isolate this $\mathbb{Z}_2$ -asymmetry from the Higgs potential. One possible fix may be to add an extremely tiny, experimentally allowed Stueckelberg mass for the Standard Model photon as well [253], though we note this may be in violation of quantum gravity [254, 255] or simply be difficult to realize in UV completions without extreme fine-tuning. We will remain agnostic about this UV issue and continue to refer to this as ‘soft breaking’, following [252].

Starting in the unbroken phase, we compute the mixing between the hypercharge gauge bosons. Two- and three-loop diagrams with Higgs loops containing one gauge vertex and one quartic insertion vanish. By charge conjugation in scalar QED, the three-leg amplitude of a gauge boson and a complex scalar pair must be antisymmetric under exchange of the scalars. However, the quartic coupling of the external legs ensures that their momenta enter symmetrically. As this holds off-shell, the presence of a loop which looks like

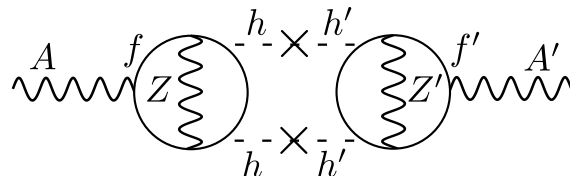


causes the diagram to vanish. However, at four loops the following diagram can be drawn which avoids this issue:



where the two hypercharges are connected by charged fermion loops in their respective sectors and the Higgs doublets' quartic interaction. This diagram contributes at least from the MTH cutoff $\Lambda \lesssim 4\pi f$ down to f , the scale at which twin and electroweak symmetries are broken. We have no argument that this vanishes nor that its unitarity cuts vanish. We thus expect a contribution to kinetic mixing of $\epsilon \sim g_1^2 c_W^2 / (4\pi)^8$, with g_1 the twin and Standard Model hypercharge coupling and $c_W = \cos \theta_W$ appearing as the contribution to the photon mixing operator. In this estimate we have omitted any logarithmic dependence on mass scales, as it is subleading.

In the broken phase, we find it easiest to perform this analysis in unitary gauge. The Higgs radial modes now mass-mix, but the emergent charge conjugation symmetries in the two QED sectors allow us to argue vanishing to higher-loop order. The implications of the formal statement of charge conjugation symmetry are subtle because we have two QED sectors, so whether charge conjugation violation is required in both sectors seems unclear. However, similarly to the above case, there is a symmetry argument which holds off-shell. The result we rely on here is that in a vector-like gauge theory, diagrams with any fermion loops with an odd number of gauge bosons cancel pairwise. Thus, each fermion loop must be sensitive to the chiral nature of the theory, so the first non-vanishing contribution is at five loops as in:



where the crosses indicate mass-mixing insertions between the two Higgs radial modes which each contribute $\sim v/f$. Thus, both the running down to low energies and the finite contributions are five-loop suppressed. From such diagrams, one expects a contribution $\epsilon \sim e^2 g_A^2 g_V^2 (v/f)^2 / (4\pi)^{10}$, where with g_V and g_A we denote the vector and axial-vector couplings of the Z , respectively. We note there are other five loop diagrams in which Higgses couple to massive vectors which are of similar size or smaller.

Depending on the relative sizes of these contributions, one then naturally expects kinetic mixing of order $\epsilon \sim 10^{-13} - 10^{-10}$. If ϵ is indeed generated at these loop-levels, then mixing on the smaller end of this range likely requires that it becomes disallowed not far above the scale f . However, we note that our ability to argue for higher-loop order vanishing in the broken versus unbroken phase is suggestive of the possibility that there may be further cancellations. We note also the possibility that these diagrams, even if nonzero, generate only higher-dimensional operators. Further investigation of the generation of kinetic mixing through a scalar portal is certainly warranted.

The diagrams which generate kinetic mixing will likely also generate higher-dimensional operators. These will be suppressed by (twin) electroweak scales and so, as discussed above for the irrelevant operators generated by the model-dependent reheating mechanism, freeze-in contributions from these operators are negligible.

3.4 Freezing-*T*win Dark Matter

As we are in the regime where freeze-in proceeds while the temperature sweeps over the mass scales in the problem, it is not precisely correct to categorize this into either “UV freeze-in” or “IR freeze-in”. Above the mass of the twin photon, freeze-in proceeds through the marginal kinetic mixing operator, and so a naive classification would say this is IR dominated. However, below the mass of the twin photon, the clearest approach is to integrate out the twin photon, to find that freeze-in then proceeds through an irrelevant, dimension-six, four-Fermi operator which is suppressed by the twin photon mass. Thus, at temperatures $T_{\text{SM}} \lesssim m_{\gamma'}$, this freeze-in looks UV dominated. This leads to the conclusion that the freeze-in rate is largest at temperatures around the mass of the twin photon. Indeed, this is generally true of freeze-in — production occurs mainly at temperatures near the largest relevant scale in the process, whether that be the largest mass out of the bath particles, mediator, and dark matter, or the starting thermal bath temperature itself in the case that one of the preceding masses is even higher.

As just argued, freeze-in production of dark matter occurs predominantly at and somewhat before $T \sim m_{\gamma'}$. We require $m_{\gamma'} \ll 1$ GeV so that most of the freeze-in yield comes from when $T < 1$ GeV, which allows us to retain ‘UV-independence’ in that we need not care about how asymmetric reheating has occurred past providing negligible density of twin states at $T = 1$ GeV. Specifically, we limit ourselves to $m_{\gamma'} < 2m_{\pi^0}$, both for this reason and to avoid uncertainties in the behavior of thermal pions during the epoch of the QCD phase transition. However, we emphasize that freeze-in will remain a viable option for producing a

twin dark matter abundance for even heavier dark photons. But the fact that the freeze-in abundance will be generated simultaneously with asymmetric reheating demands that each sort of asymmetric reheating scenario must then be treated separately. Despite the additional difficulty involved in predicting the abundance for larger twin photon masses, it would be interesting to explore this part of parameter space. In particular, it would be interesting to consider concrete scenarios with twin photons in the range of tens of GeV [262].

To calculate the relic abundance of twin electrons and positrons, we use the Boltzmann equation for the number density of e' :

$$\dot{n}_{e'} + 3Hn_{e'} = \sum_{k,l} -\langle\sigma v\rangle_{e'\bar{e}'\rightarrow kl} (n_{e'}n_{\bar{e}'} - n_{e'}^{\text{eq}}n_{\bar{e}'}^{\text{eq}}), \quad (3.4)$$

where $\langle\sigma v\rangle_{e'\bar{e}'\rightarrow kl}$ is the thermally averaged cross section for the process $e'\bar{e}' \rightarrow kl$, the sum runs over all processes with Standard Model particles in the final states and $e'\bar{e}'$ in the initial state, and $n_{e'}^{\text{eq}}$ is the equilibrium number density evaluated at temperature T . As we are in the parametric regime in which resonant production of twin electrons through the twin photon is allowed, $2 \rightarrow 2$ annihilation processes $\bar{f}f \rightarrow \gamma' \rightarrow \bar{e}'e'$, with f a charged Standard Model fermion, entirely dominate the yield.

In accordance with the freeze-in mechanism, $n_{e'}$ remains negligibly smaller than its equilibrium number density throughout the freeze-in process, and so that term is ignored. It is useful to reparametrize the abundance of e' in terms of its yield, $Y_{e'} = n_{e'}/s$ where $s = \frac{2\pi^2}{45}g_{*s}T^3$ is the entropy density in the Standard Model bath. Integrating the Boltzmann equation using standard methods, we then find the yield of e' today to be

$$Y_{e'} = \int_0^{T_i} dT \frac{(45)^{3/2}}{\sqrt{2}\pi^3\sqrt{g_{*s}g_{*s}}} \frac{M_{Pl}}{T^5} \left(\frac{1}{T} + \frac{\partial_T g_{*s}}{3g_{*s}} \right) \sum_{\bar{f}f \rightarrow \bar{e}'e'} \langle\sigma v\rangle_{\bar{f}f \rightarrow \bar{e}'e'} n_{\bar{e}'}^{\text{eq}} n_{e'}^{\text{eq}}, \quad (3.5)$$

where $T_i = 1$ GeV is the initial temperature of the Standard Model bath at which freeze-in begins in our setup, $g_*(T)$ is the number of degrees of freedom in the bath, and M_{Pl} is the reduced Planck mass. We will calculate this to an intended accuracy of 50%. To this level of accuracy, we may assume Maxwell-Boltzmann statistics to vastly simplify the calculation [263]. As a further simplification, we observe that the $\partial_T g_{*s}$ term is negligible compared to $1/T$ except possibly during the QCDPT - where uncertainties on its temperature dependence remain [264] - and so we ignore that term. The general expression for the thermally averaged cross section of the process $12 \rightarrow 34$ is then

$$\begin{aligned} \langle\sigma v\rangle n_1^{\text{eq}} n_2^{\text{eq}} &= \frac{T^4}{2^9 \pi^5 s_{34}} \int_{\text{Max}\left(\frac{m_1+m_2}{T}, \frac{m_3+m_4}{T}\right)}^{\infty} dx x^2 \sqrt{[1,2]} \sqrt{[3,4]} K_1(x) \int d(\cos\theta) |\mathcal{M}|_{12 \rightarrow 34}^2 \\ \sqrt{[i,j]} &= \sqrt{1 - \left(\frac{m_i + m_j}{xT}\right)^2} \sqrt{1 - \left(\frac{m_i - m_j}{xT}\right)^2}, \end{aligned} \quad (3.6)$$

where s_{34} is 1 if the final states are distinct and 2 if not, $x = \sqrt{s}/T$, and $|\mathcal{M}|_{12 \rightarrow 34}^2$ is the matrix element squared summed (not averaged) over all degrees of freedom. To very

good approximation, the yield results entirely from resonant production, and so we may analytically simplify the matrix element squared for $\bar{f}f \rightarrow \bar{e}'e'$ using the narrow-width approximation

$$\int d(\cos \theta) |\mathcal{M}|_{\bar{f}f \rightarrow \bar{e}'e'}^2 \approx \frac{256\pi^3 \alpha^2 \epsilon^2}{3} (2m_f^2 + m_{\gamma'}^2) \frac{(2m_{e'}^2 + m_{\gamma'}^2)}{\Gamma_{\gamma'} m_{\gamma'}^2 T} \delta(x - m_{\gamma'}/T). \quad (3.7)$$

$\Gamma_{\gamma'}$ is the total decay rate of the twin photon.

For the range of $m_{\gamma'}$ we consider, the twin photon can only decay to twin electron and positron pairs. Thus, its total decay rate is

$$\Gamma_{\gamma'} = \frac{\alpha (m_{\gamma'}^2 + 2m_{e'}^2)}{3m_{\gamma'}} \sqrt{1 - \frac{4m_{e'}^2}{m_{\gamma'}^2}}. \quad (3.8)$$

Its partial widths to Standard Model fermion pairs are suppressed by ϵ^2 , and so contribute negligibly to its total width.

The final yield of twin electrons is then

$$Y_{e'} \approx \frac{3m_{\gamma'}^2}{2\pi^2} \frac{(45)^{3/2} M_{Pl}}{\sqrt{2}\pi^3} \sum_f \int_{T_f}^{T_i} dT \Gamma_{\gamma' \rightarrow \bar{f}f} \frac{K_1(\frac{m_{\gamma'}}{T})}{\sqrt{g_*} g_{*s} T^5}, \quad (3.9)$$

where $T_f = \Lambda_{\text{QCD}}$ for quarks, $T_f = 0$ for leptons, $\Gamma_{\gamma' \rightarrow \bar{f}f}$ is the partial decay width of the twin photon to $f\bar{f}$, and the sum is over all Standard Model fermion-antifermion pairs for which $m_{\gamma'} > 2m_f$.

Since we have approximated the yield as being due entirely to on-shell production and decay of twin photons, the analytical expression for the yield in Eq. (3.9) exactly agrees with the yield from freezing-in γ' via ‘inverse decays’ $\bar{f}f \rightarrow \gamma'$, as derived in [136]. We have validated our numerical implementation of the freeze-in calculation by successfully reproducing the yield in similar cases found in [263, 265]. We have furthermore checked that reprocessing of the frozen-in dark matter [244, 266] through $e'\bar{e}' \rightarrow e'\bar{e}'e'\bar{e}'$ is negligible here,⁶ as is the depletion from $e'\bar{e}' \rightarrow \nu'\bar{\nu}'$.

An equal number of twin positrons are produced as twin electrons from the freeze-in processes. Requiring that ϵ reproduce the observed dark matter abundance today, we find

$$\epsilon = \sqrt{\frac{\Omega_\chi h^2 \rho_{\text{crit}}/h^2}{2m_{e'} \tilde{Y}_{e'} s_0}}, \quad (3.10)$$

where $\Omega_\chi h^2 \approx 0.12$, $\rho_{\text{crit}}/h^2 \approx 1.1 \times 10^{-5} \text{ GeV/cm}^3$, and $s_0 \approx 2900/\text{cm}^3$ [269]. $\tilde{Y}_{e'}$ is the total yield with the overall factor of ϵ^2 removed. This requisite kinetic mixing appears in

⁶To be conservative, we calculate the rate assuming all interactions take place at the maximum $\sqrt{s} \simeq m_{\gamma'}$ and find that it is still far below Hubble. We perform the calculation of the cross section using MadGraph [267] with a model implemented in Feynrules [268].

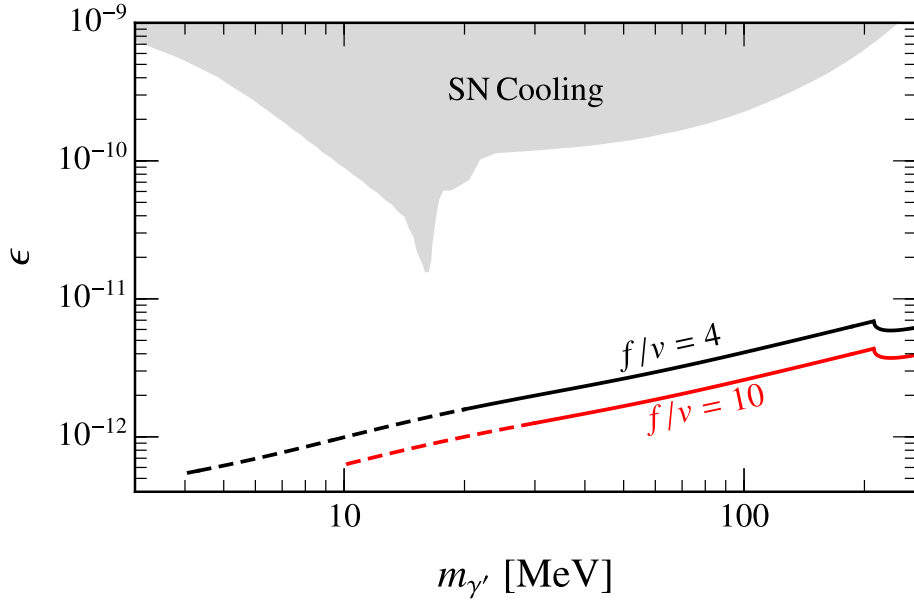


Figure 3.1: Contours in the plane of twin photon mass $m_{\gamma'}$ and kinetic mixing ϵ which freeze-in the observed dark matter abundance for two values of f/v . The dip at high masses corresponds to additional production via muon annihilations. In the dashed segments, self-interactions occur with $\sigma_{\text{elastic}}/m_{e'} \gtrsim 1 \text{ cm}^2/\text{g}$. Also included are the combined supernova cooling bounds from [270, 271].

Fig. 3.1 as a function of the twin photon mass $m_{\gamma'}$ for the two benchmark f/v values 4 and 10. In grey, we plot constraints from anomalous supernova cooling. To be conservative, we include both, slightly different bounds from [270, 271]. The dashed regions of the lines show approximately where self-interactions through Bhabha scattering are relevant in the late universe, $\sigma_{\text{elastic}}/m_{e'} \gtrsim 1 \text{ cm}^2/\text{g}$. Self-interactions much larger than this are constrained by the Bullet Cluster [22–24] among other observations. Interestingly, self-interactions of this order have been suggested to fix small-scale issues, and some claimed detections have been made as well. We refer the reader to [46] for a recent review of these issues.

As mentioned above, the level of kinetic mixing required for freeze-in is roughly of the same order as is expected from infrared contributions in the MTH. It would be interesting to develop the technology to calculate the high-loop-order diagrams at which it may be generated. In the context of a complete model of the MTH where kinetic mixing is absent in the UV, ϵ is fully calculable and depends solely on the scale at which kinetic mixing is first allowed by the symmetries. Calculating ϵ would then predict a minimal model at some $m_{\gamma'}$ to achieve the right dark matter relic abundance, making this effectively a zero-parameter extension of MTH models with asymmetric reheating. Importantly, even if ϵ is above those shown in Fig. 3.1, that would simply point to a larger value of $m_{\gamma'}$ which would suggest that the parameter point depends in more detail on the mechanism of asymmetric reheating. We

note that in the case that the infrared contributions to ϵ are below those needed here, the required kinetic mixing may instead be provided by UV contributions and the scenario is unaffected.

3.5 Conclusion

The mirror twin Higgs is perhaps the simplest avatar of the Neutral Naturalness program, which aims to address the increasingly severe little hierarchy problem. Understanding a consistent cosmological history for this model is therefore crucial, and an important step was taken in [212, 239]. As opposed to prior work, the cosmology of the MTH was remedied without hard breaking of the $\mathbb{Z}_2$ symmetry by utilizing asymmetric reheating to dilute the twin energy density. Keeping the $\mathbb{Z}_2$ as a good symmetry should simplify the task of writing high energy completions of these theories, but low-scale reheating may slightly complicate cosmology at early times. These works left as open questions how to set up cosmological epochs such as dark matter generation and baryogenesis in such models. We have here found that at least one of these questions has a natural answer.

In this Chapter, we have shown that twin electrons and positrons may be frozen-in as dark matter following asymmetric reheating in twin Higgs models. This requires extending the mirror twin Higgs minimally with a single free parameter: the twin photon mass. Freezing-in the observed dark matter abundance pins the required kinetic mixing to a level expected from infrared contributions in MTH models. In fact, the prospect of calculating the kinetic mixing generated in the MTH could make this an effectively parameter-free extension of the MTH. Compared to generic freeze-in scenarios, it is interesting in this case that the “just so” stories of feeble coupling and negligible initial density were already present for reasons entirely orthogonal to dark matter.

This minimalism in freeze-twin dark matter correlates disparate signals which would allow this model to be triangulated with relatively few indirect hints of new physics. If deviations in Higgs couplings are observed at the HL-LHC or a future lepton collider, this would determine f/v [223, 225, 272, 273], which would set the dark matter mass. An observation of anomalous cooling of a future supernova through the measurement of the neutrino ‘light curve’ might allow us to directly probe the $m_{\gamma'}, \epsilon$ curve [270, 271], though this would rely on an improved understanding of the Standard Model prediction for neutrino production.⁷ Further astrophysical evidence of dark matter self-interactions would point to a combination of f/v and $m_{\gamma'}$. All of this complementarity underscores the value of a robust experimental particle physics program whereby new physics is pursued via every imaginable channel.

⁷We thank Jae Hyeok Chang for a discussion on this point.

Chapter 4

Asymmetric Matters from a Dark First-Order Phase Transition

4.1 Introduction

As outlined in Chapter 1, the origin of the baryon asymmetry is one of the great persistent mysteries of modern physics. One of the most historically popular mechanisms to generate this asymmetry is electroweak baryogenesis [274–284], in which the departure from thermal equilibrium arises from a strongly first-order electroweak phase transition. Because the minimal Standard Model phase transition is a crossover [285–287] and CP-violation is too small [288–291], models typically introduce additional singlet scalars [112, 280] or an extra Higgs doublet [111, 277–279, 281–283, 292, 293]. However, strong constraints on Standard Model CP-violation have made these models increasingly in tension with experiment [113].

As discussed at length in Chapter 1, the Standard Model must also be extended to account for dark matter. The similarity of dark and baryon abundances has motivated studies of asymmetric dark matter, in which the baryon and dark matter asymmetries originate from the same mechanism (see the classic reviews [294–296] and references therein).

In this Chapter, we introduce a minimal model in which the baryon and dark matter asymmetries originate from electroweak-like baryogenesis in a dark sector with an $SU(3) \times SU(2) \times U(1)$ gauge group, two Higgs doublets, and one generation of Standard Model-like matter content. This Chapter builds upon recent work [297] in which the Standard Model baryon asymmetry is the result of electroweak-like baryogenesis in a dark sector with an $SU(2)$ gauge group and two “lepton” doublets. A right-handed neutrino singlet and the Standard Model electroweak sphaleron transfer the dark lepton asymmetry into a Standard Model baryon asymmetry. We show that by extending the gauge group to $SU(3) \times SU(2) \times U(1)$, one may straightforwardly obtain GeV-scale asymmetric dark hadronic dark matter. The symmetric component of the dark baryons annihilates into massive dark photons, which decay to Standard Model states via a testable kinetic mixing, leaving the asymmetric component as dark matter.

We consider two dark matter possibilities in detail. In the first, the dark neutron is the lightest baryon and comprises all of dark matter. In the second, the dark proton is the lightest baryon and acts as dark matter together with dark pions. We find both of these scenarios are testable at current and future dark photon and direct detection experiments. Because the dark matter consists of GeV-scale dark hadrons, they may also have velocity-dependent self-interactions at the correct scale to address the small-scale structure issues discussed in Chapter 1.

There is an extensive history of dark sectors with an $SU(3) \times SU(2) \times U(1)$ gauge group, such as the twin sector of the previous chapter, but particularly in the context of mirror world models (see [298, 299] for a review). Additionally, dark $SU(3)$ gauge groups are common features of baryonic dark matter and many models of asymmetric dark matter. Often, mirror asymmetric dark matter models assume high-scale leptogenesis produces the asymmetries and sometimes, that an exact $\mathbf{Z}_2$ symmetry between the standard and mirror sectors exists (see [300–302] for recent interesting examples).

The idea that the Standard Model baryon asymmetry is the result of a dark phase transition (“darkogenesis”) was originally proposed in Ref. [303]; other mechanisms have been developed in *e.g.*, [304–307]. However, whereas darkogenesis models typically rely on higher-dimensional operators or a messenger sector in order to transfer the baryon asymmetry to the Standard Model, we use a neutrino portal in a minimal, renormalizable model.

The rest of this Chapter is organized as follows. In Sec. 4.2, we define the model content and investigate the conditions for dark-sector baryogenesis. In Sec. 4.3, we calculate the resulting Standard Model- and dark-sector asymmetries and investigate the features and signatures of two possible dark matter scenarios in depth. Conclusions follow in Sec. 4.4.

4.2 Dark-Sector Baryogenesis

In this section, we introduce the particle content of our model and discuss how the dark and Standard Model baryon asymmetries are generated. The dark sector contains an $SU(3)' \times SU(2)' \times U(1)'$ gauge group, one Standard Model-like matter generation (including a right-handed singlet neutrino), and two Higgs doublets:

$$Q', u'_R, d'_R, L', e'_R, N'_R, \Phi_1, \Phi_2. \quad (4.1)$$

Throughout this Chapter, superscripts $'$ on Standard Model particles refer to their dark-sector counterparts which are charged analogously under the dark gauge group. N'_R refers to the right-handed neutrino singlet, while $\Phi_{\{1,2\}}$ refer to the two dark Higgs doublets. In contrast with the usual Standard Model mass hierarchy, we assume that leptons are heavy while quarks are light. In particular, e' is assumed to have a dark Yukawa coupling similar to that of the Standard Model top quark.

The dark gauge sector is directly analogous to the Standard Model with the exception that the dark $U(1)'_{\text{EM}}$ photon is massive and dark hypercharge kinetically mixes with Standard Model hypercharge. After electroweak symmetry breaking in both the dark and

Standard Model sectors, these features may be parameterized in terms of the dark photon as

$$\mathcal{L} \supset \frac{\epsilon}{2} F_{\mu\nu} F'^{\mu\nu} + \frac{1}{2} m_{\gamma'}^2 A'_\mu A'^\mu. \quad (4.2)$$

The right-handed neutrino singlet is coupled to both dark sector Higgses as well as the Standard Model Higgs

$$\mathcal{L} \supset Y_n^a \bar{L}' \tilde{\Phi}_a N'_R + y_N \bar{L} \tilde{H} N'_R + c.c., \quad (4.3)$$

for $a \in \{1, 2\}$ where $\tilde{H} = i\sigma_2 H^*$ and similarly for $\tilde{\Phi}_a$. Each particle may possess distinct Yukawa couplings Y^a to the two Higgs doublets. It is also possible that the dark neutrino has a Majorana mass term; we will consider this to be small.

Dark-sector baryogenesis proceeds as follows. At high temperatures, the dark sector is in the unbroken electroweak phase. At some temperature T_C , the dark sector undergoes a strongly first-order electroweak phase transition, which provides a departure from thermal equilibrium. At the phase transition, explicit CP-violating terms in the dark Higgs potential induce a changing CP-violating phase in the fermion mass terms across the bubble wall; this in turn leads to a CP-violating force across the bubble wall. Together with the dark electroweak sphaleron, this results in a $B' + L'$ asymmetry that is primarily driven by the dark electron, which we take to have an $\mathcal{O}(1)$ Yukawa coupling.

The precise parameter space over which the two-Higgs doublet mechanism may generate a sufficient baryon asymmetry has been the subject of extensive study; see *e.g.* [111, 277–279, 281–283, 292, 293]. Baryogenesis with two Higgs doublets favors light Higgs masses and large quartic couplings, and in the context of extensions of the Standard Model Higgs sector, this can cause issues such as Landau poles; together with recent electric dipole measurements [113], this leads to significant constraints on the parameter space over which baryogenesis may occur. However, in the present setup, electroweak baryogenesis is easier to realize. Among other things, leptons diffuse further into the symmetric phase and do not suffer from suppression by the strong sphalerons [119]. Besides, EDM constraints do not apply and the parameter space in the dark scalar sector is almost entirely unconstrained. Hence, we expect that it should not be difficult to achieve the required baryon asymmetry and will not perform an in-depth analysis.

Acoustic sound waves and magnetohydrodynamic turbulence generated in the aftermath of a first-order phase transition will generically produce a gravitational wave signal with a characteristic spectrum determined by the rate and energy release of the phase transition. With the advent of gravitational wave astronomy, gravitational wave signals from dark phase transitions have become a subject of considerable interest; see *e.g.* [308–311]. The spectrum of the gravitational waves from this model could likely fall in the detection range of future gravitational wave observatories such as LISA, BBO, and DECIGO. Because the spectrum is generic and our model space is so unconstrained, we will not perform an in-depth investigation of gravitational wave signals here; see [297] for the expected spectrum.

Following dark-sector baryogenesis [297], the dark leptons are in equilibrium with the Standard Model through the neutrino portal via the process $\Phi'\nu' \leftrightarrow H\nu$ and $N'_R \leftrightarrow H\nu$.

Together with the Standard Model sphaleron, this will transfer some of the initial dark lepton asymmetry into a Standard Model baryon and lepton asymmetry. At some temperature, the remaining leptons will decay to the Standard Model through the processes $e' \rightarrow \nu' \bar{u}' d'$ and $\nu' \rightarrow \nu H$, leaving only quarks and photons in the dark sector. Following hadronization, the symmetric component of the dark baryons will annihilate into dark photons (through *e.g.*, $\pi'^+ \pi'^- \rightarrow \gamma' \gamma'$ and $\pi'^+ \pi'^- \rightarrow \pi'^0 \pi'^0$, $\pi'^0 \rightarrow \gamma' \gamma'$) which in turn decay into the Standard Model.

The remaining baryonic asymmetry forms asymmetric dark matter with the dark matter mass set by the relative Standard Model and dark matter abundances. The precise behavior of the dark hadronic content and the nature of dark matter depend on the parameters of the model and is discussed in depth in Sec. 4.3.

We now recall the values for the Standard Model lepton and baryon asymmetries [297]. We assume there are no new particles in the Standard Model and hence the Standard Model electroweak phase transition is a crossover. If the dark neutrino decays after the Standard Model sphaleron has decoupled, we find

$$B = \frac{36}{133} B', \quad L = -\frac{97}{133} B'. \quad (4.4)$$

If on the other hand the dark neutrino is heavy and decays before the Standard Model sphaleron has decoupled, we obtain asymmetries

$$B = \frac{12}{37} B', \quad L = -\frac{25}{37} B'. \quad (4.5)$$

The asymmetries will be different if the Standard Model electroweak phase transition is strongly first-order instead of a crossover and may be found in Ref. [297]; we do not discuss these cases further as they require additional extensions of the Standard Model sector.

Although the model presented in this section is in some sense the most minimal, it generalizes quite straightforwardly to a fully mirrored model in which the dark sector has three Standard Model-like generations with a similar mass hierarchy. A full mirror sector with three families of quarks would also motivate some GUT-scale equivalence of the $SU(3)$ and $SU(3)'$ gauge couplings, which could in turn follow similar RG flows to the IR. This would explain the coincidence of the dark and Standard Model matter densities and is an advantage of mirror world models generally.

One might also want to consider a model with three generations but only one Higgs doublet in which the dark baryon asymmetry is generated according to the mechanism originally suggested by Farrar and Shaposhnikov for minimal-Standard Model baryogenesis [284, 289–291]. In this mechanism, the CKM or PMNS matrix would lead to CP-violating coefficients for fermionic reflections off the bubble wall of a strongly first-order electroweak phase transition. However, our non-perturbative analysis following [291] found that even with the largest possible degree of CP-violation and $\mathcal{O}(1)$ Yukawa couplings, it was impossible to generate a sufficiently large dark baryon asymmetry even in the optimistic limit of a thin wall, a fast sphaleron, and no diffusion.

4.3 Dark Matter and Experimental Signatures

In this section, we discuss the fate of the dark sector following dark-sector baryogenesis. The remaining asymmetric hadronic content will be dark matter and is overall neutral due to individual conservation of both dark and Standard Model $U(1)_{\text{EM}}$ charges. Below the confinement scale of $SU(3)'$, $\Lambda_{SU(3)'}$, the only remaining dark sector particles are hadrons and photons. Since the dark quark masses do not affect the baryogenesis mechanism (as long as their Yukawas are sufficiently smaller than the much heavier dark leptons), there are different viable dark matter scenarios. We will enumerate the few simplest cases below, but first discuss the phenomenology that is common to all of them.

Since the dark leptons are heavy, much of the dark hadronic symmetric entropy density is transferred into π^0 . To ensure the dark sector does not overclose the universe, we require π^0 to decay. This is easily achieved through the dominant decay mode to two dark photons as long as $m_{\gamma'} \leq m_{\pi^0}/2$. In order for the dark photon to decay into the Standard Model, we require $m_{\gamma'} \geq 2m_e$. The decay rate of the dark photon to a pair of Standard Model leptons is

$$\Gamma_{\gamma' \rightarrow \bar{l}l} = \frac{\alpha \epsilon^2 (m_{\gamma'}^2 + 2m_l^2)}{3m_{\gamma'}} \sqrt{1 - \frac{4m_l^2}{m_{\gamma'}^2}}. \quad (4.6)$$

For dark photon masses below a GeV, the decay rate into hadronic channels is non-perturbative. We infer the decay rates from the branching ratios derived from measured ratios of hadronic final-state cross sections to those of muons in e^+e^- collisions [312]. We require the resulting total decay rate of the dark photon to be faster than Hubble before Standard Model neutrinos decouple around $T \sim 3$ MeV [313], which is true for all dark photon masses we consider as long as $\epsilon \gtrsim \mathcal{O}(10^{-10})$.

With these common considerations outlined, we discuss two distinct limits: one in which all of dark matter is the dark neutron, n' , and the other in which dark matter is comprised of equal numbers of dark protons, p' , and π'^- , assuming $|m_{n'} - m_{p'}| \gtrsim 100$ MeV. These scenarios predict different dark matter masses and direct detection constraints and prospects. If n' and p' masses are closer, the situation is between these two limits.

Dark Neutron Dark Matter

In this scenario, the lightest dark baryon is the neutron with $m_{p'} - m_{n'} \approx m_{u'} - m_{d'} \gtrsim 100$ MeV, while both quarks are light, $m_{u'}, m_{d'} < \Lambda_{SU(3)'}$. After annihilations $p' \pi'^- \rightarrow n' \gamma'$, $\pi'^+ \pi'^- \rightarrow \pi^0 \pi^0$ and decays $\pi^0 \rightarrow \gamma' \gamma'$, $\gamma' \rightarrow \text{SM}$, the entire dark baryon asymmetry is in n' which forms all of dark matter. While there is a subcomponent of p' , the strong interactions and large dark neutron-proton mass splitting allow us to safely assume n' comprises the vast majority of dark matter. Since the relative dark and Standard Model baryon asymmetries are set above, the dark matter mass is precisely determined by the relative

baryon and dark matter abundances [11],

$$\frac{\Omega_c}{\Omega_b} = \frac{B'}{B} \frac{m_{n'}}{m_p} = 5.238. \quad (4.7)$$

In the case Eq. (4.4) that N'_R is light, we predict a dark matter mass¹

$$m_{n'} = 1.33 \text{ GeV}. \quad (4.8)$$

In the case Eq. (4.5) that N'_R is heavier than the scale of the Standard Model electroweak crossover, we predict

$$m_{n'} = 1.59 \text{ GeV}. \quad (4.9)$$

Although the n' is neutral, it should possess a magnetic moment similar to that of the Standard Model neutron. This, combined with the γ' - γ kinetic mixing, allows n' to scatter off protons in nuclei with a cross section²

$$\sigma_{n'p} \approx \epsilon^2 e^2 e'^2 F_2^{n'2} v^4 \frac{m_p^4 m_{n'}^2 (3m_p^2 + 2m_p m_{n'} + 5m_{n'}^2)}{6\pi m_{\gamma'}^4 (m_p + m_{n'})^6}, \quad (4.10)$$

where v is the incoming dark matter velocity and $F_2^n \approx -1.913$ for the Standard Model neutron. The most stringent spin-independent, per-nucleon cross section constraint on dark matter with masses $m_\chi \sim 1 \text{ GeV}$ comes from XENON1T [314]. In particular, for $m_{n'} = 1.33 \text{ GeV}$, the bound on the dark matter-nucleon cross section is $\sigma^{\text{SI}} < 7.6 \times 10^{-40} \text{ cm}^2$. This bound assumes equal couplings of the dark matter to neutrons and protons, but n' only scatters off protons, so the upper limit for $n'p$ scattering is slightly larger:

$$\sigma_{n'p} < \left(\frac{A}{Z}\right)^2 7.6 \times 10^{-40} \text{ cm}^2 \approx 4.4 \times 10^{-39} \text{ cm}^2. \quad (4.11)$$

In addition to constraints from direct detection, there are also limits on the self-interaction among dark matter particles from galaxy clusters $\sigma \lesssim 0.2 \text{ cm}^2/\text{g}$ [315–317]. The neutrons in the Standard Model have an astoundingly large cross section at low energies, $\sigma \approx 4.5 \times 10^{-23} \text{ cm}^2$, much larger than the geometric cross section $\approx 10^{-25} \text{ cm}^2$. This is regarded as a consequence of accidental (and unnatural) cancellations in the effective field theory (see, *e.g.*, [318–320]) and is not generic. According to recent lattice QCD calculations from the HAL QCD collaboration [321], the self-interaction among n' is below the limit for rather heavy dark pions, $m_{\pi'} \gtrsim 0.4 m_{n'}$.³ Therefore, this scenario prefers $m_{u'} \gtrsim 100 \text{ MeV}$,

¹These predictions are subject to calculable α_s/π corrections in chemical equilibrium at the percent level.

²Both the dark neutron charge radius and the possible Higgs portal give subdominant contributions to this scattering.

³This is still subject to uncertainties given disagreements with the NPLQCD collaboration [322, 323] (with a possible resolution [324]), and the calculations are in the flavor $SU(3)$ limit. It is also possible that much smaller $m_{\pi'}$ leads to small self-interaction, but it is currently beyond what can be studied on lattice.

which in turn allows for larger dark photon masses since $m_{\gamma'} < m_{\pi'_0}/2$. However, it is difficult to have a larger cross section at lower velocities to address the small-scale structure problems as shown with the effective range theory framework [325].

The viable dark photon parameter space for the neutron dark matter scenario is shown in Fig. 4.1 (**Left**) with current constraints from experiments [326, 327], supernovae [270, 271], and BBN [328], as well as the projected sensitivities of upcoming experiments including LHCb [329, 330], Belle-II [326, 331], AWAKE [332] (10^{16} electrons of 50 GeV), HPS [333], SeaQuest [207], LDMX [334] (HL-LDMX with $E_{\text{beam}} = 16$ GeV), FASER [335] (LHC Run 3 with 150 fb^{-1}), NA62 [336], and SHiP [206]. Additionally, the NA64 bounds should improve soon [337]. Note also that even spectroscopy of resonance states is possible at e^+e^- colliders [172, 177]. The figure assumes the scenario in which $m_{n'} = 1.33$ GeV (cf. Eq. (4.8)) and u' prefers $m_u \gtrsim 100$ MeV so that $m_{\pi'} \sim 0.5m_{n'}$. In addition to making the n' self-interactions consistent with constraints, this allows dark photons as heavy as $m_{\gamma'} \sim 0.25m_{n'} = 0.3$ GeV.⁴

Interestingly, while there is currently decades of viable parameter space in which the dark photon mass and kinetic mixing can achieve the asymmetric dark neutron dark matter, much of this will be probed by future experiments. Since the cross section in Eq. (4.10) is velocity-suppressed, current and future direct detection experiments are far from probing the viable dark photon parameter space. To illustrate this, we naïvely assume the XENON1T bound in Eq. (4.11) scales linearly with exposure and project the constraint for XENON1T with 100 times its current exposure (as in DARWIN [339]) as a thin dashed line in the upper left of Fig. 4.1. Additionally, we incorrectly assume that all incoming dark matter have the largest possible velocity $v_{\text{max}} = v_{\text{esc}} + v_{\text{E}} \sim (550 + 240) \text{ km/s}$ (the sum of the escape and Earth velocities in the galactic frame). Clearly, such neutral dark matter seems well outside the current direct detection bounds and future dark photon searches will better probe the viable parameter space.

Dark Proton & Pion Dark Matter

Next, we consider the case $m_{u'} < m_{d'} < \Lambda_{SU(3')}$ so that the proton is the lightest dark baryon. Similar to the dark neutron case, we assume $m_{d'} - m_{u'} \gtrsim 100$ MeV to guarantee that the n' abundance is negligible. Conservation of $U(1)'_{\text{EM}}$ charge implies an equal number of π'^- comprising a subcomponent of dark matter. Even though the relative dark and Standard Model baryon asymmetries are set above, the additional pion dark matter component gives a range of possible dark matter subcomponent masses. To reproduce the observed relic abundance, they satisfy

$$\frac{B' m_{p'} + m_{\pi'^-}}{B m_p} = 5.238. \quad (4.12)$$

It is interesting to note that p' and π'^- may scatter resonantly in the p -wave through Δ'^0 . In this case, it becomes an ideal resonant self-interacting dark matter to address the

⁴This upper bound on the dark photon mass will relax at higher values of kinetic mixing because $\pi'^0 \rightarrow \gamma' \gamma'^* \rightarrow \gamma' e^+ e^-$ would be possible, but we do not consider this further.

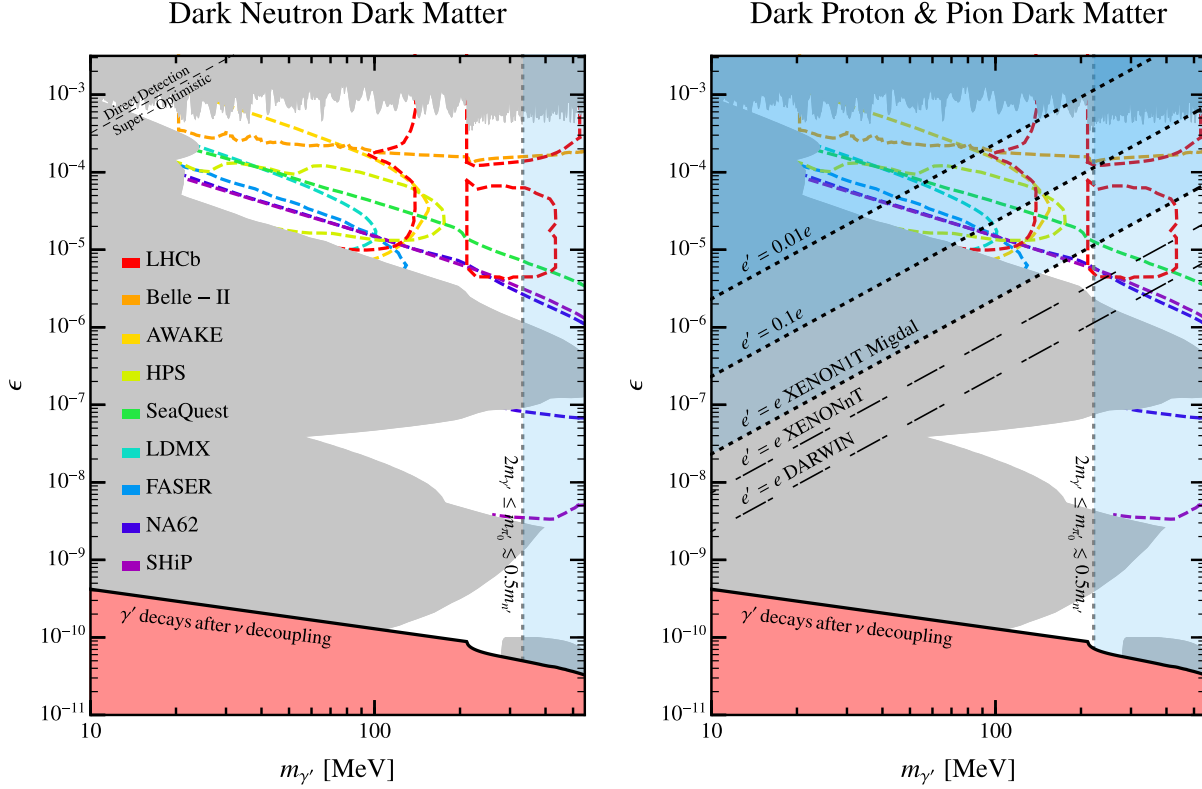


Figure 4.1: Viable dark photon parameter space for asymmetric dark hadron dark matter. Existing constraints on dark photons from experiments [326, 327], supernovae [270, 271], and BBN [328] are dark gray. The constraints specific to our models, namely that $m_{\gamma'} \leq m_{\pi^0}/2$ and that the dark photon decays before Standard Model neutrinos decouple around $T \sim 3$ MeV, are in light blue and red, respectively. Also shown in rainbow colors are projections from future experiments [206, 207, 326, 329–336]. **Left:** Viable parameter space for the dark neutron dark matter case with the predicted $m_{n'} = 1.33$ GeV and assuming $m_{\pi'} \sim 0.5m_{n'}$. **Right:** Viable parameter space for the dark proton and pion dark matter case assuming $m_{p'} = 2m_{\pi'} = 0.887$ GeV. The direct detection constraint from XENON1T [314] for various e'/e is shown in blue, as are naïve projections for XENONnT [338] and DARWIN [339].

small-scale structure problems if the resonant velocity is $v_R \sim 100$ km/s with a constant s -wave cross section with $\sigma/m \sim 0.1$ cm²/g [340]. The possibility of this threshold resonance prefers $m_{\pi'}/m_{p'} \sim 0.4$, which is also desirable to permit larger dark photon masses. The “Coulomb” potential barrier may also lead to a $p'p'$ resonance in the s -wave.

For illustrative purposes, we pick the dark proton and pion masses to be $m_{p'} = 2m_{\pi'} = 0.887$ GeV to yield the observed relic abundance in the case that N'_R is light. While any masses satisfying Eq. (4.12) with the baryon asymmetry ratio given by Eq. (4.4) are possible, a larger dark pion mass leads to a wider viable dark photon parameter space. Likewise, in

the case Eq. (4.5) that N'_R is heavier than the scale of the Standard Model electroweak crossover, we set the masses to be $m_{p'} = 2m_{\pi'} = 1.06$ GeV.

For the available parameter space, the would-be Bohr radius $\alpha'/m_{\pi'}$ of the $p'\pi'^-$ “atom” is longer than the range of the dark-photon exchange force, $1/m_{\gamma'}$, so we do not expect these atoms to form. Therefore, direct detection experiments can probe p' - p scattering with

$$\sigma_{p'p} \approx \epsilon^2 e^2 e'^2 \frac{m_p^2 m_{p'}^2}{\pi(m_p + m_{p'})^2 m_{\gamma'}^4}, \quad (4.13)$$

and similarly for π' - p scattering. The XENON1T [314] bound is quite weaker in this heavy-pion case since for $m_{p'} = 0.887$ GeV, the bound on the dark matter-nucleon cross section is $\sigma^{\text{SI}} < 2.0 \times 10^{-39}$ cm². Additionally, the upper limit for $p'p$ (or $\pi'p$) scattering is larger by $(A/Z)^2$ due to the lack of coupling to neutrons.

The viable dark photon parameter space for the dark proton and pion dark matter scenario is shown in Fig. 4.1 (**Right**) for $m_{p'} = 2m_{\pi'} = 0.887$ GeV to give the widest parameter space. The direct detection limit also weakens if e' is much smaller than e . To demonstrate this, we show the XENON1T constraint assuming $e'/e = \{1, 0.1, 0.01\}$ as dashed black contours. There is still a large viable parameter space and future improvements in the limits appear promising. We naïvely assume that XENONnT [338] with its larger exposure will increase the current bound by an order of magnitude, though the exact improvement in sensitivity from this Migdal effect analysis is not so obvious [314]. We also show what DARWIN [339] may probe with its possible additional order of magnitude improvement. Interestingly, it appears that future direct detection experiments may be competitive with and even exceed the sensitivity of dark photon experiments.

Other Dark Hadron Dark Matter

Yet another possibility is that there is only one light dark quark, let's say u' . Then, dark matter is partially comprised of a dark $\Delta'^{++}(u'u'u')$ baryon whose abundance comes from the dark baryon asymmetry. There are also twice the number of dark “pions” $\pi'^-(\bar{u}'d')$, now heavier than $\Lambda_{SU(3)'}$. To produce the observed dark matter relic abundance, we require

$$\frac{B'}{B} \frac{m_{\Delta'^{++}} + 2m_{\pi'^-}}{m_p} = 5.238, \quad (4.14)$$

where $m_{\Delta'^{++}} \approx \Lambda_{SU(3)'}$. Besides the possible difference in masses, the direct detection would be similar to the p' and π'^- dark matter case above. The symmetric component annihilates into $\eta(\bar{u}'u') \rightarrow \gamma'\gamma' \rightarrow 2(e^+e^-)$. We do not discuss this and other variants further.

4.4 Conclusion

We have introduced a minimal renormalizable model of asymmetric matters from a dark first-order phase transition which leads to a detectable gravitational wave signature. Electroweak-like baryogenesis in the dark sector generates a dark-sector asymmetry which is then ferried

to the Standard Model via a neutrino portal. Kinetic mixing between the dark and Standard Model photons allows for the symmetric dark-sector entropy to safely transfer to the Standard Model, while also providing a means for the remaining asymmetric dark matter to scatter in direct detection experiments. In the case of dark neutron dark matter, we find decades of viable dark photon parameter space which will be explored in the near future by many upcoming experiments. If instead the asymmetric dark matter is comprised of dark protons and pions, the parameter space is currently being tested by direct detection experiments. In the latter case, self interactions among the dark matter may also ameliorate small-scale problems such as the diversity problem. If the energy scale of the dark first-order phase transition temperature is below 1000 TeV, we expect a gravitational wave signal at future observatories.

Chapter 5

Directly Detecting Signals from Absorption of Fermionic Dark Matter

5.1 Introduction

As discussed in Chapter 1, the null results of WIMP dark matter searches have sparked a renaissance in dark matter model building in search of alternative thermal histories which predict lighter dark matter [126–135]. Indeed, we have expounded such light dark matter models with alternate thermal histories in the last three chapters. For masses below the GeV scale, dark matter which scatters off a target will typically deposit energy below the threshold of the largest direct detection experiments ($\mathcal{O}(\text{keV})$), significantly relaxing the direct constraints.

To discover these lighter dark matter candidates, the direct detection program is moving toward detecting smaller energy deposits with novel scattering targets and lower-threshold detectors [139–148]. Current technology is already sensitive to energy deposits of $\mathcal{O}(\text{eV})$ [341] and new proposals could detect energy deposits of $\mathcal{O}(\text{meV})$ [149–155, 158]. As the direct detection program pushes the low-mass frontier, it can also broaden its searches for different signals to increase its impact with little additional cost.

Particle dark matter detection strategies can be grouped into two classes: scattering and absorption. Searches for scattering look for a dark matter particle depositing its kinetic energy onto a target within the detector, typically a nucleus or an electron. In contrast, searches for absorption look for signals in which a dark matter particle deposits its mass energy. Absorption signals have primarily been considered for bosonic dark matter candidates with studies of fermionic absorption signals limited to induced proton-to-neutron conversion in Super-Kamiokande [159] and sterile neutrino dark matter [161, 342–348] (see also exothermic dark matter [349] and self-destructing dark matter [350] for related signals).

In this Chapter, we systematically study direct detection signals from the absorption of fermionic dark matter. We describe novel signals and their corresponding lowest-dimension operators; project the sensitivities of ongoing and proposed dark matter direct detection

and neutrino experiments to these signals; and demonstrate the consistency of these signals with the issues of dark matter stability and abundance. Our approach in this Chapter is to consider these new signals broadly, trying to emphasize the UV-independent physics. By contrast, we will begin the next chapter with a careful look at two specific UV completions and their model-specific constraints.

5.2 Signals and Operators

For simplicity, we take dark matter (χ) to be a Dirac fermion charged under lepton number, and impose only Lorentz, $SU(3)_C \times U(1)_{EM}$, CP, lepton and baryon number symmetries. Baryon number conservation is necessary to avoid proton decays while lepton number allows the (Dirac) neutrino to remain light. We enumerate operators in the effective theory with the fields $\{\chi, n, p, e, \nu, F_{\mu\nu}\}$, where $F_{\mu\nu}$ is the EM field strength tensor. We do not include other QCD resonances as they have no bearing on direct detection.

Consider first dimension-6 operators of the form, $[\bar{\chi}\Gamma_i\nu][\bar{\psi}\Gamma_j\psi]$, where $\psi \supset \{n, p, e, \nu\}$ and $\Gamma_i = \{\mathbf{1}, \gamma_5, \gamma_\mu, \gamma_\mu\gamma_5, \sigma_{\mu\nu}\}$ denotes the different possible Lorentz structures of the bilinear. These “neutral current” operators generate the first class of new signals we consider; $\bar{\chi} + T \rightarrow \bar{\nu} + T$, where T is a target nucleus or electron which absorbs a fraction of the dark matter mass energy. We will focus on nuclear absorption, where the rates may be coherently enhanced, and postpone the study of electron absorption.

Next, consider dimension-6 operators of the form, $[\bar{\chi}\Gamma_i e][\bar{n}\Gamma_j p]$. These generate a class of “charged current” signals; $\bar{\chi} + \frac{A}{Z}X \rightarrow e^\pm + \frac{A}{Z\mp 1}X_*^\mp$, in which dark matter can induce $\beta^\pm$ decay in nuclei which are stable in vacuum. This process potentially has multiple correlated signals: a detectable $e^\pm$, a nuclear recoil, a prompt γ decay from the excited final nucleus, and further nuclear decays if the final nucleus is unstable. Induced β^+ decays have significantly smaller rates relative to β^- due to the Coulomb repulsion between the emitted e^+ and the nucleus, so we focus on dark matter-induced β^- decays and leave the β^+ decays for the next chapter¹. The same charged current operators can also shift the endpoint of the $\beta^\pm$ distribution for nuclei which already undergo $\beta^\pm$ decays in vacuum. While this might be detectable at PTOLEMY [351, 352], these kinds of experiments have small exposures and large backgrounds as we discuss in detail in the next chapter.

Finally, dark matter candidates which have fermionic absorption signals decay. At dimension-5, the operator $\bar{\chi}\sigma^{\mu\nu}\nu F_{\mu\nu}$ induces decays of χ as do the dimension-6 operators, $\bar{\chi}\gamma^\nu\Gamma_{(5)}\partial^\mu\nu F_{\mu\nu}$, where $\Gamma_{(5)} \equiv \{\mathbf{1}, \gamma_5\}$. At higher dimensions, there exist operators allowing multiphoton decays. The single photon channel can be detected with the usual line search, while the multiphoton channels are constrained by diffuse photon emission. Detectable fermionic absorption signals, consistent with indirect detection bounds, typically require lighter dark matter as the decay rates scale with a large power of m_χ . We include a discussion of decays below for each signal and operator we consider.

¹ β^+ decays induced by dark matter with $m_\chi \gg \text{MeV}$ were proposed for Hydrogen targets in Super-Kamiokande [159].

5.3 Neutral Current Signals: Nuclear Recoils

We first study the process $\chi + N \rightarrow \nu + N$, where N is a target nucleus. We will focus on two operators:

$$\mathcal{O}_{\text{NC}} = \frac{1}{\Lambda^2} \bar{\chi} \gamma_\mu P_R \nu (\bar{n} \gamma^\mu n + \bar{p} \gamma^\mu p) + \text{h.c.} . \quad (5.1)$$

We choose this Lorentz structure for concreteness. However, the neutral current signal is not highly dependent on it as long as there is some amount of vector coupling to the nucleons. These can arise from a theory of a heavy Z' coupled to quarks and χ with some mixing between the right handed components of χ and ν . The incoming χ is non-relativistic, so its mass dominates its energy resulting in a momentum transfer (q) and nuclear recoil energy (E_R):

$$q \simeq m_\chi, \quad E_R \simeq \frac{m_\chi^2}{2M}, \quad (5.2)$$

where M is the mass of the nucleus. For contrast, elastic scattering off a nucleus yields at most $E_R = 2v^2\mu^2/M$, where μ is the reduced mass and v is the dark matter velocity (see [122] for a recent review). This $1/v^2$ increase in E_R relative to WIMP scattering allows searches for lighter dark matter with both direct detection experiments and higher-threshold, neutrino experiments.

The differential rate of neutral current nuclear recoils from absorbing fermionic dark matter is;

$$\frac{dR}{dE_R} = N_T \frac{\rho_\chi}{m_\chi} \frac{|\overline{\mathcal{M}_N}|^2}{16\pi M^2} \delta(E_R - E_R^0) \Theta(E_R^0 - E_{\text{th}}), \quad (5.3)$$

where N_T is the number of target nuclei, $\rho_\chi \simeq 0.4 \text{ GeV}/\text{cm}^3$ is the local dark matter energy density, $E_R^0 \equiv m_\chi^2/2M$, E_{th} is the experiment's threshold, and $|\overline{\mathcal{M}_N}|^2$ is the matrix element squared (at q) averaged over initial spins and summed over final spins. In elastic scattering, the spread in incoming dark matter velocities causes a spread in recoil energies, but in fermionic absorption, the rate is sharply peaked at $E_R = E_R^0$. Every isotope in an experiment has a distinct peak with width (ΔE_R) determined by higher order corrections to Eq. (5.2), corresponding to $\Delta E_R/E_R \sim 10^{-3}$. There are no modulation signals or rate uncertainties arising from the dark matter velocity distribution.

The total rate for absorption by multiple nuclei is

$$R = \frac{\rho_\chi}{m_\chi} \sigma_{\text{NC}} \sum_j N_{T,j} A_j^2 F_j^2 \Theta(E_{R,j}^0 - E_{\text{th}}), \quad (5.4)$$

where $N_{T,j}$, A_j , $E_{R,j}^0$, and F_j , are the number, mass number, recoil energy, and Helm form factor [370] (evaluated at $q = m_\chi$ and normalized to 1) of target isotope j . The cross section per-nucleon is $\sigma_{\text{NC}} = m_\chi^2/(4\pi\Lambda^4)$. Absorption has the unique signature of correlated, peaked counts in dR/dE_R bins containing $E_{R,j}^0 = m_\chi^2/(2M_j)$ for the different target isotopes with

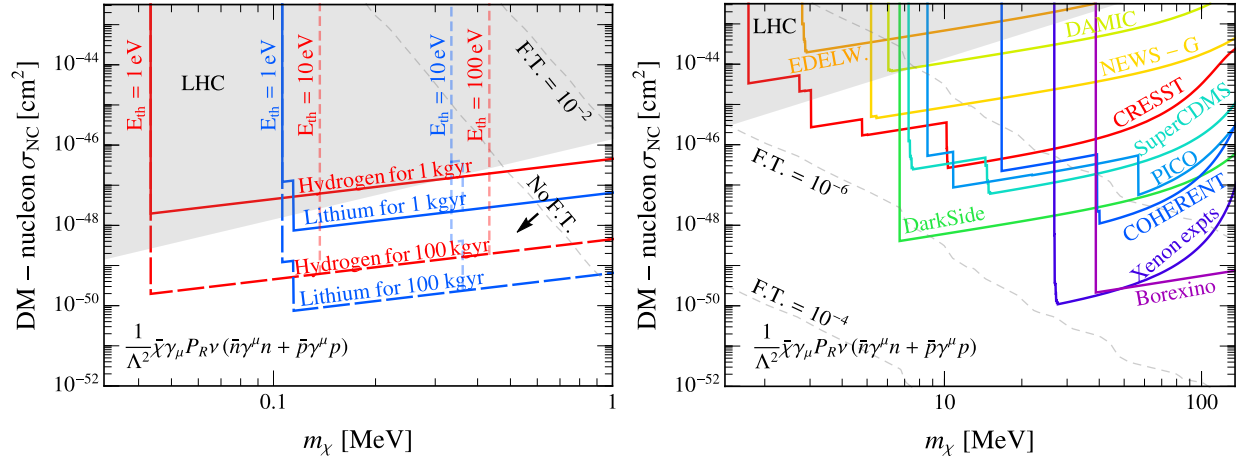


Figure 5.1: **Left:** Projected sensitivities of future experiments to σ_{NC} . We show two exposures (1/100 kg-yr) of two different target materials (Hydrogen in red and Lithium in blue) with three possible nuclear-recoil energy thresholds (1, 10, and 100 eV). **Right:** Projected sensitivities of current experiments to σ_{NC} , including CRESST III [353] and CRESST II [354] (“CRESST” in red); EDELWEISS-SURF [355] (orange); NEWS-G [356] (yellow); DAMIC [357] (lime); DarkSide-50 [358, 359] (green); CDMSliteR2 [360] and SuperCDMS [361] (“SuperCDMS” in aqua); PICO-60 run with C₃F₈ [362] and PICO-60 run with CF₃I [363] (“PICO” in sky blue); COHERENT [364, 365] (blue); LUX [123], PandaX-II [366], and XENON1T [338] (“Xenon expts” in navy blue); and Borexino [367] (purple; see [368] to extract nuclear recoil threshold). Both panels include LHC bounds [369] and the indirect detection constraints from χ decay [27] which require different levels of fine-tuning between the UV and IR contributions to kinetic mixing between the photon and Z' for the Z' model as described in the text.

masses M_j . This can be a powerful discriminator from backgrounds since the relative heights and spacing of the peaks is completely determined. Whether an experiment can resolve these distinct peaks depends on its energy resolution and the mass splitting between the target isotopes.

For $m_\chi \lesssim \text{MeV}$, future experiments are necessary to probe the small nuclear recoil energy. Detailed projections are challenging due to the breadth of proposals and possible absorption by collective modes of nuclei. So, we roughly estimate the sensitivity of such future detectors in Fig. 6.4 (**Left**) where, for simplicity, we require at least 10 events to set our projections, independent of mass or experiment. The cross sections are smaller than those in typical WIMP searches due to the larger number densities of lighter dark matter. We consider Hydrogen and Lithium targets with energy thresholds of eV – 100 eV for 1 kg-year and 100 kg-year exposures (see [146] for one possible realization). In the Lithium target, the detection of two correlated signals from both isotopes is possible. For $m_\chi \gtrsim \text{MeV}$, we project the sensitivity of current experiments in Fig. 6.4 (**Right**). Interestingly, which experiments

best probe the neutral current absorption signal are not always the same as those which best probe WIMPs (e.g., Borexino). Each experiment can only detect the absorption signal off a target isotope when its distinct nuclear recoil energy is larger than the threshold energy, hence the edges in Fig. 6.4.

We now address the stability of χ . For concreteness, we consider a model where a heavy Z' couples in an isospin-invariant way to quarks, with gauge coupling $g_{Z'}$, and to χ and $P_R\nu$ in an off-diagonal way². Quark loops induce a kinetic mixing, ϵ , between the Z' and the photon of order $\epsilon \sim g_{Z'}e/16\pi^2$ allowing the decay $\chi \rightarrow \nu e^+e^-$. Without additional Z or Z' mass suppressions, the decay $\chi \rightarrow \nu\gamma$ is forbidden by gauge invariance while $\chi \rightarrow \nu\gamma\gamma$ is forbidden as a consequence of charge conjugation (also known as Furry's theorem). For $m_\chi \lesssim \text{MeV}$, the electron channel is kinematically forbidden and the dominant decay is $\chi \rightarrow \nu\gamma\gamma\gamma$, whose primary contribution proceeds through a kinetic mixing and the Euler-Heisenberg Lagrangian, yielding the approximate rate

$$\Gamma_{\chi \rightarrow \nu\gamma\gamma\gamma} \simeq (g_{Z'}\epsilon)^2 10^{-19} \frac{m_\chi^{13}}{m_e^8 m_{Z'}^4}. \quad (5.5)$$

Estimating the dark matter decay rates in this simple UV completion, we find future experiments can quickly probe new parameter space while cross-sections accessible to current experiment are ruled out by indirect detection bounds [27].

However, it is possible to suppress dark matter decays by fine-tuning the UV contribution to the kinetic mixing against the IR piece estimated here. Concretely, we define this fine-tuning as $\text{F.T.} \equiv |\epsilon_{\text{UV}} - \epsilon|/\epsilon$ and we show the fine-tuning necessary to evade indirect detection constraints with dashed gray lines labeled “F.T.” in Fig. 6.4. We note that the projected direct detection sensitivities in Fig. 6.4 are insensitive to the details of the UV completion. We will study ways to reduce fine-tuning by incorporating flavor-dependent couplings to suppress ϵ in future work.

Also shown in Fig. 6.4 are direct constraints from LHC mono-jet searches on the Z' model, which bound new neutral currents below the TeV scale [369]. Dijet constraints are model dependent: in the Z' model, dijet bounds are suppressed in the limit where the quark coupling to the Z' is much smaller than the χ coupling to Z' . The excess neutrino flux emanating from the Sun and Earth due to Eq. (6.29) is not quite large enough to be seen in neutrino observatories in the near future. Cosmological bounds depend on initial conditions (e.g., the reheat temperature) and the UV completion. While we postpone a detailed study of the dark matter relic abundance, we comment that a simple way to populate such light dark matter is through its thermal production and relativistic decoupling followed by its dilution from the decay of another heavy particle—a mechanism considered for sterile neutrinos [372]. If the dark matter χ decouples while it is relativistic, which with only the quark coupling, will occur at the latest around the QCD phase transition, it will be overproduced. After relativistic decoupling of χ another state becomes non-relativistic leading it to quickly dominate the

²This suppresses the $\chi \rightarrow \nu\nu\nu$ decay; this off-diagonal coupling appears in inelastic dark matter models [371].

energy density of the universe. This mechanism can produce and sufficiently dilute dark matter to achieve its observed relic abundance over the entire dark matter mass range shown in Fig. 6.4. While other production mechanisms are possible, one must carefully avoid spoiling BBN for dark matter lighter than $\mathcal{O}(1 \text{ MeV})$ [373], *e.g.* via freeze-in. Neutron absorption of χ through $n + \chi \rightarrow p + e$ could also spoil BBN, but the number density of χ is substantially less than that of Standard Model neutrinos and the rate of scattering is suppressed relative to that of neutrinos by $(m_W/\Lambda)^4$ making this negligible.

5.4 Charged Current Signals: Dark-Matter-Induced β^- Decays

Next, consider signals from $\chi + {}^A_Z\text{X} \rightarrow e^- + {}^A_{Z+1}\text{X}_*^+$ (or at the nucleon level, $\chi + n \rightarrow p + e^-$), which we refer to as an induced β^- decay. This process can cause stable elements to become unstable in the presence of dark matter if m_χ is large enough to overcome the kinematic barrier. Such a signal may proceed through the dimension-6 charged current vector operator,

$$\mathcal{O}_{\text{CC}} = \frac{1}{\Lambda^2} [\bar{\chi}\gamma^\mu e] [\bar{n}\gamma_\mu p] + \text{h.c.} . \quad (5.6)$$

This can be generated by a W' which can appear if the electroweak gauge group is embedded in a larger gauge group which subsequently breaks into the Standard Model.

We consider the vector operator in Eq. (5.6) to leverage known results from the neutrino and nuclear physics literature. The vector-vector interaction primarily induces Fermi transitions which are characterized by their conservation of spin (J) and parity (P) of the nucleus [374], also known as $J^P \rightarrow J^P$ transitions. However, we emphasize that the dark matter induced β^- decay signal is more general, with different vertex structures allowing different transitions. We leave a study of additional interactions to future work.

Denoting the mass of a nucleus of mass number A and atomic number, Z , by $M_{A,Z}$, we focus on isotopes which satisfy $M_{A,Z} < M_{A,Z+1}^{(*)} + m_e$, such that the nucleus is stable against β^- decay in a vacuum (the $(*)$ is included to emphasize the daughter nucleus may be in an excited state, typically $200 \text{ keV} - 1 \text{ MeV}$ heavier in mass). Then dark matter induced β^- decay is kinematically allowed if

$$m_\chi > m_{\text{th}}^\beta \equiv M_{A,Z+1}^{(*)} + m_e - M_{A,Z} . \quad (5.7)$$

In these induced decays, χ is absorbed by the target nuclei and transfers the majority of its rest mass to the outgoing electron. In the limit where $m_\chi - m_{\text{th}}^\beta \gg m_e$, the electron and nuclear recoil energies are analogues to the neutral current case with $m_\chi \rightarrow m_\chi - m_{\text{th}}^\beta$, and are given by;

$$E_R \simeq \begin{cases} m_\chi - m_{\text{th}}^\beta & (\text{electron}) \\ (m_\chi - m_{\text{th}}^\beta)^2 / 2M_{A,Z+1}^{(*)} & (\text{nucleus}) \end{cases} . \quad (5.8)$$

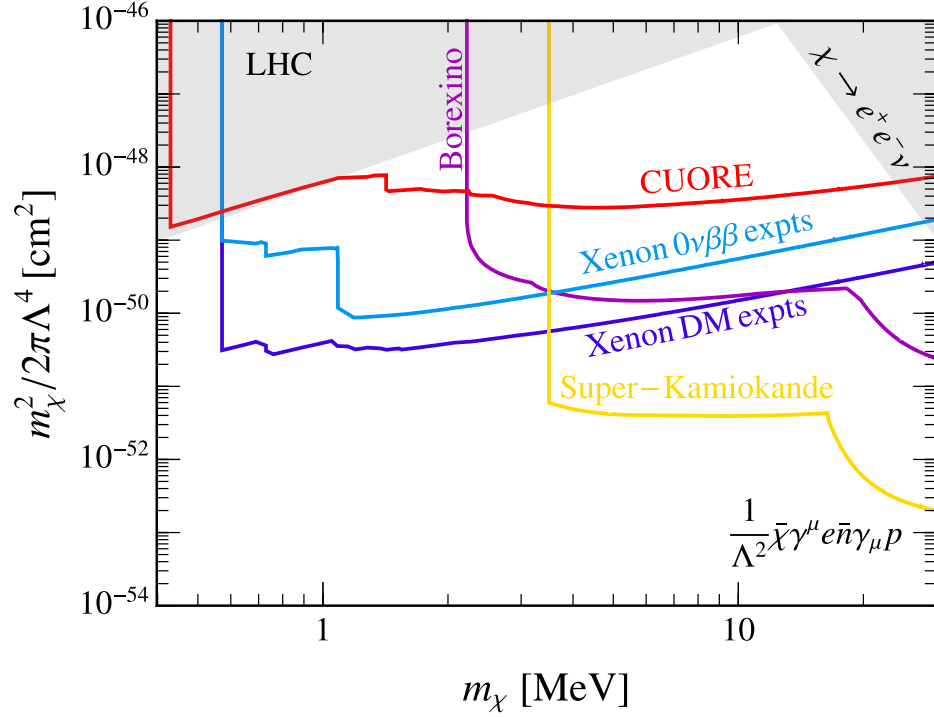


Figure 5.2: Projected sensitivities to $m_\chi^2/2\pi\Lambda^4$ from a dedicated search for the charged current induced β^- transition at Cuore [375] (red); LUX [123], PandaX-II [366], and XENON1T [338] (“Xenon DM expts” in navy blue); EXO-200 [376] and KamLAND-Zen [377] (“Xenon $0\nu\beta\beta$ expts” in sky blue); SuperKamiokande [378] (yellow); CDMS-II [379] (aqua); DarkSide-50 [358, 359] (green); and Borexino [367] (purple). Also shown are LHC bounds [380] and indirect constraints from χ decays in our simple UV model [27].

Therefore, the energetic outgoing electron will shower in the detector, and can be searched for. The nuclear recoil energy, as with the neutral current case, is independent of dark matter velocity, and can be searched for as well. Additional correlated signals result from the possible de-excitation of the daughter nucleus and its subsequent decay (typically many days later). These multiple signals make possible correlated searches to reduce backgrounds. The specific signals depend on the experiment, the particular isotope, and the dark matter mass.

The rate for dark matter-induced β^- decays is;

$$R = \frac{\rho_\chi}{2m_\chi} \sum_j N_{T,j} (A_j - Z_j) \langle \sigma v \rangle_j, \quad (5.9)$$

where we sum over all isotopes in a given target material, $N_{T,j}$ is the number of target

isotope j , and $\langle\sigma v\rangle_j$ is an isotope's velocity averaged cross-section

$$\langle\sigma v\rangle_j = \frac{|\vec{p}_e|_j}{16\pi m_\chi M_{A_j, Z_j}^2} \overline{|\mathcal{M}_{N_j}|^2}, \quad (5.10)$$

where $|\vec{p}_e|_j^2 = (m_{\text{th}, j}^\beta - m_\chi)(m_{\text{th}, j}^\beta - m_\chi - 2m_e)$ is the electron's outgoing 3-momentum in the center of mass frame (which is approximately the lab frame), in the limit that $m_e, m_\chi, m_{\text{th}, j}^\beta \ll M_{A_j, Z_j}$. The amplitude $\mathcal{M}_N$ is for absorption by the whole nuclei (the momentum transfer is not enough to resolve individual nucleons), which can be related to the nucleon level amplitude $\mathcal{M}$ (with the spinors normalized to $p_\mu p^\mu = M_{A_j, Z_j}^2$) through the Fermi function, $\mathcal{F}(Z, E_e)$ and a form factor, $F_V(q^2)$:

$$\mathcal{M}_N = \sqrt{\mathcal{F}(Z+1, E_e) F_V(q^2)} \mathcal{M}. \quad (5.11)$$

The Fermi function accounts for the Coulomb attraction of the ejected electron and can enhance the cross-section by several orders of magnitude for heavier elements. The form factor is equal to 1 for small momentum transfer relative to the nucleon mass, $q^2 \ll m_n^2$, while for larger q^2 the dependence can be extracted from the neutrino literature [381]. In principle, (6.54) must contain a sum over all possible nuclear spin states. The assumption made here is that this sum will be dominated by $\Delta J^P = 0$ transitions as is the case of a vector coupling [374]. Excitation of additional final states is possible if $q \gtrsim r_N^{-1}$, where $r_N \simeq 1.2A^{1/3}\text{fm}$ is the nuclear radius [382], however for simplicity we focus on lighter masses such that these do not contribute significantly to the rate for any isotope considered here.

The total rate is found by summing over the contributions from each isotope. Evaluating (6.54) in the limit where $m_e, m_\chi, m_{\text{th}}^\beta \ll M_{A, Z}$, the total rate is;

$$R = \frac{\rho_\chi}{2m_\chi} \sum_j N_{T, j} (A_j - Z_j) \frac{|\vec{p}_e|_j^3 \mathcal{F}(Z_j + 1, E_e)}{2\pi\Lambda^4 (m_\chi - m_{\text{th}, j}^\beta)}, \quad (5.12)$$

where we have integrated over all energies with the assumption that such a signal could be detected by most experiments under consideration here given the multitude of correlated high energy signals.

We project the sensitivity of current experiments to the charged current signal in Fig. 5.4 where we again require at least 10 events to set our projections, independent of isotope mass or experiment. Sensitivities are displayed in terms of the theoretically interesting quantity $m_\chi^2/2\pi\Lambda^4$ (to which Eq. (6.54) reduces in the limit of large M_{A_j, Z_j} and $m_\chi \gg m_{\text{th}, j}^\beta$, modulo the Fermi function). As with the neutral current case, limits depend on the different isotopes in a given experiment. In particular, the kinks in Fig. 5.4 occur at $m_\chi \sim m_{\text{th}, j}^\beta$ for every relevant isotope in a given experiment.

To estimate the dark matter decay constraints from a typical UV completion, we consider a model with a W' coupled vectorially to up and down quarks without any direct couplings to leptons. When kinematically allowed, the dominant decay is $\chi \rightarrow e^+ e^- \nu$ which arises

from a kinetic mixing between W' and the Standard Model W boson of order $\sim g_{W'e}/16\pi^2$. The decay $\chi \rightarrow \nu\gamma$ is subdominant since it is at two-loop order, making it roughly $(4\pi)^2$ smaller. We estimate the decay rate and show the resulting indirect constraints [27] in gray in Fig. 5.4. The decay bounds are much weaker than in the neutral current case as they are suppressed by both the weak scale and the W' mass.

In addition to decays, there are direct bounds from LHC searches for $pp \rightarrow \ell\nu$. A search was done by CMS at 8 TeV looking for helicity-non-conserving contact interaction models which have contact operators with vertex structure different than that of the Standard Model [380] which sets a powerful constraint on the charged current operators. For the W' model, this constraint corresponds to a scale in Eq. (5.6) of $\Lambda \gtrsim 3.2$ TeV. In the W' model, there is also a Z' which could lead to direct bounds. However, direct searches for Z' are not as stringent as those for the W' as the Z' can be somewhat heavier than the W' and elastic scattering constraints are negligible for the masses and Λ of interest here. We also consider low energy searches for modifications to the $V - A$ gauge structure of the Standard Model [383, 384] and for light fermions in charged pion decays: $\pi^\pm \rightarrow e^\pm\chi$ [385], but find they are subdominant to the CMS constraint. Cosmological constraints require detailed assumptions about the initial conditions and the full set of interactions. As for $\mathcal{O}_{\text{NC}}$, a simple production mechanism with $\mathcal{O}_{\text{CC}}$ has the dark matter decouple from the Standard Model bath while relativistic, followed by late decays of a dominating particle [386–388].

5.5 Discussion

In this Chapter, we have introduced a novel class of signals from fermionic dark matter absorption in direct detection and neutrino experiments. We have studied the sensitivities of future and current experiments to neutral current signals from the process $\chi + \text{N} \rightarrow \nu + \text{N}$, as shown in Fig. 6.4. This neutral current causes target isotopes to recoil with distinct energies and correlated rates, enabling significant background reduction in searches. We have also studied the sensitivities of current experiments to induced β^- decays from the process $\chi + {}^A_Z\text{X} \rightarrow e^- + {}^A_{Z+1}\text{X}_*^+$, as shown in Fig. 5.4. This charged current enjoys multiple signatures from a sequence of events starting with a nuclear recoil and ejected e^- , followed by a likely γ decay and often a final β decay or electron capture event several days later. For both signals, ongoing experiments can probe orders of magnitude of unexplored parameter space by performing dedicated searches.

Without yet knowing the true nature of dark matter, it is impossible to know how it will appear in an experiment. Perhaps, it has been a fermion, depositing its mass energy into unsuspecting targets all along.

Chapter 6

Absorption of Fermionic Dark Matter by Nuclear Targets

6.1 Introduction

As discussed in Chapter 1 and the previous chapter, null WIMP searches have forced dark matter theories and experiments to evolve. Sub-GeV dark matter models with novel thermal histories have come in vogue [126–138], as partially evidenced by Chapters 2, 3, and 4. Simultaneously, existing experiments have been getting bigger and lowering their energy thresholds while a variety of light dark matter experiments have been proposed [139–158].

Another strategy for progress is to explore novel dark matter direct detection signals (see e.g., [159–161]). With this direction in mind, in the previous chapter we considered the absorption of a fermionic dark matter particle in a detector — a scenario in which the dark matter mass energy is available to the target (in contrast to the well studied elastic scenario where only its velocity-suppressed kinetic energy may be imparted). In this Chapter, we build on this foundation by exploring and cataloging all such absorption signals off nuclear targets, leaving the consideration of electron targets to future work.¹ We organize these signals by the corresponding types of higher dimensional operators which lead to fermionic absorption.

We begin by considering signals from absorption of fermionic dark matter from “neutral current” processes of the form

$$\bar{\chi} + {}^A_Z\text{X} \rightarrow \bar{\nu} + {}^A_Z\text{X}, \quad (6.1)$$

where ${}^A_Z\text{X}$ is the nuclear target with atomic number Z and atomic mass number A , χ is the dark matter, and ν is a Standard Model neutrino. This signal is generated by dimension-6 neutral current operators of the form $[\bar{\chi}\Gamma_i\nu][\bar{n}\Gamma_j n]$ and $[\bar{\chi}\Gamma_i\nu][\bar{p}\Gamma_j p]$, where $\Gamma_i = \{\mathbf{1}, \gamma_5, \gamma_\mu, \gamma_\mu\gamma_5, \sigma_{\mu\nu}\}$ contains all possible Lorentz structures. Since dark matter must be lighter than the nucleons to avoid rapid decays, energy momentum conservation ensures

¹See [389] for recent work on MeV sterile neutrino dark matter elastically scattering off electrons.

that the outgoing neutrino carries away most of the dark matter (mass) energy. Nevertheless, a fraction of the χ mass is still converted into kinetic energy for the recoiling nucleus, resulting in a distinct signal. Like spin-independent WIMP scattering, the absorption rates can enjoy a coherent enhancement for larger nuclei. In this Chapter, we study the neutral current in detail by surveying current experiments and discussing the types of future experiments best suited to detect these processes.

Another set of fermionic absorption signals are induced β decays

$$\bar{\chi} + {}^A_Z\text{X} \rightarrow e^\pm + {}^A_{Z\mp 1}\text{X}^{(*)}, \quad (6.2)$$

where $(*)$ denotes a possible excited state of the nucleus (which range from below an MeV to 10s of MeV above the ground state depending on the isotope). Such “charged current” processes are generated by dimension-6 operators of the form $[\bar{\chi}\Gamma_i e][\bar{n}\Gamma_j p]$. The induced decay can occur in isotopes that are stable or unstable in the vacuum. Stable (or meta-stable) isotopes exist in macroscopic quantities in current experiments and so these can be employed to look for multiple correlated signals: the energetic ejected $e^\pm$, the recoil of the daughter nucleus, a γ from the decay of the excited daughter nucleus, and another β decay of the final nucleus, if it is unstable. Due to these multiple signals and large $e^\pm$ (and potentially photon) energy, dedicated searches for induced β decays do not rely on the nuclear recoils being above a given experimental threshold. However, these processes themselves have kinematic thresholds allowing them to only probe dark matter masses larger than ~ 400 keV. This kind of signal has been considered in the context of sterile neutrino dark matter detection [161], where it was concluded that immense quantities of Dysprosium (which is rare but has an anomalously small β decay energy threshold of ~ 2.5 keV) is needed to probe the parameter space consistent with indirect detection bounds from sterile neutrino decay. In this Chapter, we study alternative dark matter candidates, and find that current experiments can easily observe signals consistent with other constraints.

In principle, one can look for induced β^- or β^+ decays for all isotopes within a detector. Indeed, we find many induced β^- decay targets in current dark matter direct detection and neutrino experiments. However, induced β^+ decay rates suffer relative to those of β^- due to the Coulomb repulsion of the e^+ by the nucleus as well as Pauli blocking effects of the outgoing neutron. As such, we will primarily be interested in signals from induced β^+ decays off of Hydrogen targets in neutrino experiments, as considered in [159] for Super-Kamiokande. Nonetheless, induced β^+ decays are worth consideration since they allow complementary isotope targets in experiments to probe the same operators and might be necessary to search for the asymmetric dark matter scenario, where only χ or $\bar{\chi}$ may be present today. In this Chapter, we survey the current experiments which can be used to look for induced β decays of stable isotopes and their projected reach.

For unstable isotopes, it is more challenging to accumulate macroscopic targets in detectors. Nevertheless, since they have no induced β decay thresholds, they may be used to detect arbitrarily light dark matter. To find these signals, one can look for outgoing β with energies beyond the kinematic endpoint of the target isotope’s β decay spectra. Despite

these practical challenges, there are proposals with other primary physics goals which rely on β^- decaying isotopes, such as PTOLEMY [351, 352]. Previous studies have focused on sterile neutrinos where it's challenging to compete with current decay bounds [345, 347]. We propose to test light dark matter using this signal and overview the types of experiments necessary to probe parameter space consistent with other constraints.

Importantly, dark matter candidates which allow either the neutral or charged current fermion absorption signals are inevitably unstable and their decays can be searched for using telescope observations. Since these indirect detection bounds are inherently model-dependent, we treat all of the dominant decays in concrete UV completions of the above dimension-6 fermionic absorption operators. For the neutral current, we present a model of gauged baryon-number with additional coupling to χ and introduce a mixing of χ with the (Dirac) neutrino. For the charged current, we present a modification of left-right symmetric models where χ is put into a right handed doublet with the electron instead of the neutrino. In all cases, the decays depend on large powers of m_χ and so, requiring dark matter to be sufficiently long-lived leads us to consider masses well below the GeV scale.

This Chapter is organized as follows. For the neutral current and charged current dimension-6 operators which yield these unique signals, we present simple UV completions in Section 6.2. Additionally, Section 6.2 contains a detailed discussion of the UV model-dependent dark matter decay modes that constrain our parameter space. We briefly summarize all relevant details of current experiments for which we make projections in Section 6.3. In Sections 6.4 through 6.6, we comprehensively consider all possible signals from absorption by nuclear targets at current and future direct detection, neutrino, and neutrino-less double beta decay experiments. We conclude in Section 6.7.

6.2 UV Completions

In this section, we present two UV completions that realize the neutral and charged current signals presented in this Chapter. After presenting the models, we discuss in detail the various cosmological and collider constraints, with a particular emphasis on implications for dark matter stability. In general, a variety of thermal (and non-thermal) production mechanisms can accommodate the observed dark matter relic abundance. Therefore, accommodating $\Omega_\chi h^2 \sim 0.1$ does not place any restrictions on the model parameter space that is of interest to fermionic absorption signals. As such, we omit a detailed discussion of production mechanisms.

Neutral Current

Simple models that generate the neutral current operator can be built through the introduction of additional $U(1)$ symmetries broken above the weak scale, and a mass-mixing between the dark matter candidate and a neutrino. For simplicity, consider a scenario where only χ and the Standard Model quarks are charged under a new $U(1)'$, with all of the quarks

charged equally (i.e., gauging baryon number):

$$\mathcal{L} \supset g_\chi \left(\frac{1}{3} \sum_q \bar{q} \gamma_\mu q + Q_\chi \bar{\chi} \gamma_\mu \chi \right) Z'^\mu + \frac{\epsilon}{2} Z'_{\mu\nu} F^{\mu\nu} + \frac{m_{Z'}^2}{2} Z'_\mu Z'^\mu, \quad (6.3)$$

where g_χ is the $U(1)'$ gauge coupling, we have taken the dark matter to have charge Q_χ under the $U(1)'$, and we have set the quark charge to unity without loss of generality. We also include a kinetic mixing, ϵ , which has a natural value $\epsilon \sim eg_\chi/16\pi^2$ arising from the running of quarks within the loop. Integrating out the Z' yields the following dimension-6 operator:

$$\mathcal{L} \supset \frac{g_\chi^2}{m_{Z'}^2} \frac{1}{3} Q_\chi \sum_q \bar{q} \gamma^\mu q \bar{\chi} \gamma_\mu \chi. \quad (6.4)$$

Note that Eq. (6.4) is an operator typically considered in elastic scattering. Now suppose that χ mixes with the Standard Model neutrinos through a Yukawa interaction of a scalar, ϕ (with charge Q_χ under the $U(1)'$) which gains a vacuum expectation value (giving the Z' a mass contribution). For simplicity, we consider a model with lepton number charged dark matter and approximately massless Dirac neutrinos such that the $U(1)'$ invariant mass term is given by:

$$\mathcal{L}_{\text{mass}} \supset m_\chi \bar{\chi} \chi + (y \phi \bar{\chi} P_R \nu + \text{h.c.}) = \begin{pmatrix} \bar{\nu} & \bar{\chi} \end{pmatrix} \begin{pmatrix} 0 & 0 \\ y \langle \phi \rangle & m_\chi \end{pmatrix} P_R \begin{pmatrix} \nu \\ \chi \end{pmatrix} + \text{h.c.} + \dots \quad (6.5)$$

After diagonalization, there is one massless state (identified with the Standard Model neutrino) and one massive state with mass $\sqrt{m_\chi^2 + y^2 \langle \phi \rangle^2}$. Furthermore, a mixing is induced between $\chi_R \equiv P_R \chi$ and ν_R with a mixing angle, θ_R given by:

$$s_{\theta_R} = \frac{y \langle \phi \rangle}{\sqrt{y^2 \langle \phi \rangle^2 + m_\chi^2}}. \quad (6.6)$$

Since the mixing is only between the right handed fields, the W -induced $\chi \rightarrow \nu \gamma$ decay rate is heavily suppressed while maintaining a large direct detection signal (in contrast to the case of sterile neutrinos).

We now discuss the phenomenology of this model. The direct detection signal is primarily governed by the effective operator:

$$\mathcal{L} \supset \frac{Q_\chi g_\chi^2 s_{\theta_R} c_{\theta_R}}{m_{Z'}^2} (\bar{n} \gamma_\mu n + \bar{p} \gamma_\mu p) \bar{\chi} \gamma_\mu P_R \nu + \text{h.c.} \quad (6.7)$$

There will also be an elastic scattering mode but it is challenging to see for the masses of interest here, as it produces a smaller energy deposit. Consequently, searches looking for

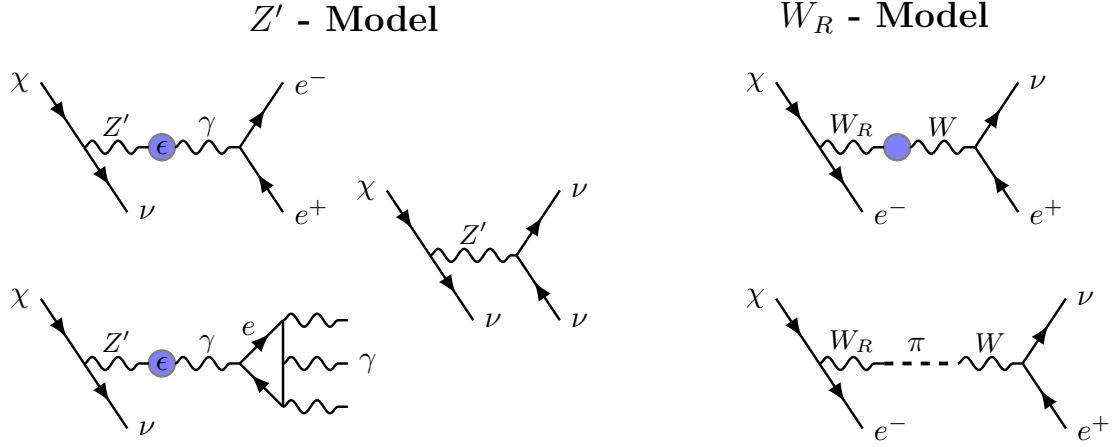


Figure 6.1: Most constraining decays for χ in the neutral current model (**left**) and charged current model (**right**).

elastic scattering will generally be weaker than a dedicated fermion absorption search in this model.

As alluded to above, dark matter is unstable as the $\chi - \nu$ mixing can lead to various decays of χ depicted in Fig. 6.1 (**left**). Consider first those decays in Fig. 6.1 (**left**) induced by 1-loop kinetic mixing without additional insertions of Z' or Z propagators. A curious feature of this Z' model is it *does not* induce 1 photon or 2 photon decay channels up to these additional insertions — the single photon channel through kinetic mixing is forbidden by gauge invariance (this is equivalent to the usual statement that particles charged under a new $U(1)'$ do not couple to the Standard Model photon after diagonalization), while the 2 photon channel is forbidden by charge conjugation (also known as Furry's theorem). Considering higher orders, we find the decay of dark matter to 1, or 2 photons up to neutrino mass insertions:

$$\Gamma_{\chi \rightarrow \nu \gamma} = 0 + \mathcal{O}\left(\frac{m_\chi^{13}}{(4\pi)^{13} m_{Z'}^{12}}\right), \quad \Gamma_{\chi \rightarrow \nu \gamma \gamma} = 0 + \mathcal{O}\left(\frac{m_\chi^9}{(4\pi)^{11} m_{Z'}^8}\right). \quad (6.8)$$

For the 1 photon channel, the dominant decay is through a 3-loop diagram with 3 Z' 's. The leading contribution to the 2 photon channel comes from a 2-loop diagram with 2 Z' 's. All of these contributions are negligible for the dark matter masses of interest to us here ($m_\chi \lesssim m_\pi$). Including neutrino mass insertions induces decays through a W loop analogous to those of sterile neutrinos but suppressed by an additional mixing angle and dependent on the flavor structure between the right and left handed neutrinos. Since the neutrino that enters the effective operator in Eq. (6.7) via mixing with χ can be massless, these potential decay channels can be made arbitrarily small. We assume this here for simplicity.

For $m_\chi \gtrsim 2m_e$ the dominantly constraining decay mode is $\chi \rightarrow \nu e^+ e^-$ induced by kinetic

mixing with decay rate given by

$$\Gamma_{\chi \rightarrow \nu e^+ e^-} = \left(16 \log 2 - \frac{31}{3}\right) \frac{m_\chi^5}{512\pi^3} \left(\frac{\epsilon e Q_\chi g_\chi s_{\theta_R} c_{\theta_R}}{m_{Z'}^2}\right)^2, \quad (6.9)$$

For lower dark matter masses the dominant visible decay is $\chi \rightarrow \nu \gamma \gamma$ through kinetic mixing in conjunction with the Euler-Heisenberg Lagrangian [390] (one can also circumvent kinetic mixing by attaching external photons to a loop of quarks however this diagram involves parametric suppressions by meson masses and we estimate it to be subdominant). The decay rate is estimated as (computing the phase space factor numerically with the aid of *MadGraph* [267]):

$$\Gamma_{\chi \rightarrow \nu \gamma \gamma} \simeq 10^{-7} m_\chi^{13} \left(\frac{(8 Q_\chi g_\chi \epsilon) s_{\theta_R} c_{\theta_R} \alpha^2}{360 m_e^4 m_{Z'}^2}\right)^2. \quad (6.10)$$

In addition, χ can decay invisibly to neutrinos, $\chi \rightarrow 3\nu$, which proceeds through a large power of the $\nu_R - \chi_R$ mixing angle:

$$\Gamma_{\chi \rightarrow \nu \nu \nu} = (16 \log 2 - 11) \frac{m_\chi^5}{128\pi^3} \left(\frac{Q_\chi^2 g_\chi^2 s_{\theta_R}^3 c_{\theta_R}}{m_{Z'}^2}\right)^2. \quad (6.11)$$

All the decays arise from irrelevant operators, and as such the rates are proportional to large powers of m_χ . Therefore, ensuring a stable dark matter candidate leads us to consider lighter dark matter candidates. The limits on dark matter decay rates depends sensitively on the dark matter mass and particular decay channel. For $\chi \rightarrow \nu e^+ e^-$ and $\chi \rightarrow \nu \gamma \gamma$ decays we recast constraints from [27] while for $\chi \rightarrow \nu \nu \nu$ we use bounds from the non-observation of an anomalous change in the equation of state of the Universe from the era of the Cosmic Microwave Background until present day [391].

The particular decay rates computed here clearly depend sensitively on the particular model chosen. As a striking example, note that in the case of a scalar mediator its possible to completely eliminate the $\chi \rightarrow 3\nu$ decay mode by choosing a scalar which does not carry a coupling to two neutrinos. In an effort to not let the specifics of the model overshadow the signal regions observable in experiments, we allow for the possibility of fine-tuning away decays by introducing a UV kinetic mixing parameter and a UV contribution to the 3ν operator that can cancel these decays modes to some level. As we show, this will be necessary in all the detectable parameter space of neutral current absorption for $m_\chi \gtrsim \text{MeV}$.

Another possible tension could be that the production of dark matter results in too great an energy density in the right-handed neutrinos. Assuming that dark matter is produced via UV freeze-in, we find that for the lightest dark matter masses we consider for the NC operators, the energy density in ν_R relative to that in a Standard Model ν_L is always less than $\sim 10^{-3}$. Thus, though the energy density in ν_R depends on the initial dark matter production mechanism, in general, it does not have to be in tension with measurements of the early radiation energy density.

In addition to constraints from indirect detection, there are bounds on this UV completion that do not depend on χ being dark matter. For $m_{Z'}$ well above the weak scale, the dominant constraints arise from mono-jet searches (see [369] for a recent summary). For lighter Z' , the dominant constraints arise from flavor changing meson and Z decays induced by a Wess-Zumino-Witten term present in theories which gauge an anomalous combination of Standard Model charges [392, 393]. Constraints also come from looking for heavy anomaly-canceling fermions directly in colliders [394]. Since the scale of the effective operators we consider here are above the weak scale, we do not expect significant constraints from star cooling, beam dump, or supernovae which are often crucial when discussing light dark matter.²

Charged Current

UV completions which result in a charged current signal typically require new states charged under electromagnetism. Such a situation is a prediction of an extended electroweak sector, one example of which we explore here. A simple extended breaking pattern is [395]³

$$SU(2)_L \times SU(2)_R \times U(1)_X \xrightarrow{\langle \Phi \rangle} SU(2)_L \times U(1)_Y \xrightarrow{\langle H \rangle} U(1)_{\text{EM}}. \quad (6.12)$$

The initial breaking can be accomplished when an $SU(2)_R$ doublet scalar, Φ charged as $(\mathbf{1}, \mathbf{2}, 1/2)$, gets a vev;

$$\Phi = \begin{pmatrix} \phi^+ \\ \phi^0 \end{pmatrix}, \quad \langle \Phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ u \end{pmatrix}. \quad (6.13)$$

In this stage of breaking $U(1)_Y$ is formed out of a linear combination of $SU(2)_R \times U(1)_X$ charges:

$$Y = X + T_R^3. \quad (6.14)$$

The second stage of breaking can be accomplished with a H charged as $(\mathbf{2}, \bar{\mathbf{2}}, 0)$. This corresponds to:

$$H = \begin{pmatrix} h_1^0 & h_1^+ \\ h_2^- & h_2^0 \end{pmatrix}, \quad \langle H \rangle = \frac{v}{\sqrt{2}} \begin{pmatrix} c_\beta & 0 \\ 0 & s_\beta \end{pmatrix}, \quad (6.15)$$

where the EM charge is given by,

$$Q = X + T_L^3 + T_R^3. \quad (6.16)$$

The lepton number carrying dark matter χ in this set-up is identified with the right handed component $\chi_R \equiv P_R \chi$ charged as $(\mathbf{1}, \mathbf{2}, 0)$ and is assumed to complete the lepton

²In principle, this conclusion may be too hasty since, while the fermion absorption operator scale we consider will always be above the weak scale, the $\bar{q}\gamma^\mu q \bar{\chi}\gamma_\mu \chi$ operator could have a scale a little below the weak scale. Nevertheless, since we work in a regime where it is at most comparable to the weak scale and Z' only couples to baryons we estimate there are no additional strong constraints.

³More generally, the fermions may be charged under $SU(2)_{L/R}$ in alternative structures (see e.g. [396] for a review).

right handed doublets. We do not need to introduce gauge singlet right-handed neutrino partners for the Standard Model neutrinos, but may do so to realize a standard seesaw mechanism. Many known mechanisms may be used to generate Standard Model neutrino masses and we do not prefer a particular one as they do not affect the fermionic absorption phenomenology. Additionally, there will be an inert left handed component χ_L *i.e.*, a singlet under all gauge symmetries. As is typical for left right symmetric models we place the Standard Model right-handed fermions (we will consider only one generation here) into right handed doublets ($\equiv R$) under $SU(2)_R$ and left handed fermions in doublets ($\equiv L$) under $SU(2)_L$.

The $SU(2)_R$ gauge boson masses primarily arise in the usual way, from the kinetic term once Φ develops a vev $\langle \Phi \rangle$, and are given by

$$M_{W_R} = \frac{1}{2}g_R u, \quad M_{Z_R} = \frac{1}{2}(g_R^2 + g_\chi^2)^{1/2}u. \quad (6.17)$$

In addition, there is a mass mixing between the W and W_R at tree level given by,

$$\mathcal{L} \supset -\frac{1}{4}g_L g_R v^2 s_{2\beta} W_\mu W_R^\mu. \quad (6.18)$$

For a generic scalar potential $s_{2\beta}$ is $\mathcal{O}(1)$ and hence there is a mixing angle between W and W_R of $\mathcal{O}(m_W^2/m_{W_R}^2)$, which leads to χ decay. To minimize this mixing we work in limit that $s_{2\beta} \rightarrow 0$, which can be achieved if H contributes negligibly to the breaking of $SU(2)_R$ such that $c_\beta = 1$ and $s_\beta = 0$ as in the inert doublet model [397].

Recall that we place the quarks into right handed multiplets, while the dark matter χ_R completes the lepton right handed doublets. This leads to the following term allowed by all the symmetries:

$$\mathcal{L} \supset \frac{g_R}{\sqrt{2}} W_{R\mu} (\bar{\chi} \gamma^\mu P_R e + \bar{u} \gamma^\mu P_R d) + \text{h.c.} \quad (6.19)$$

Additionally, fermion masses are generated from Yukawa interactions with H and $\tilde{H} \equiv \sigma_2 H^* \sigma_2$ of the form $\bar{L} H R$ and $\bar{L} \tilde{H} R$ as follows:

$$\mathcal{L} \supset \frac{y_u v}{\sqrt{2}} \bar{u} u + \frac{y_d v}{\sqrt{2}} \bar{d} d + \frac{y_\ell v}{\sqrt{2}} \bar{\ell} \ell + \frac{y_\nu v}{\sqrt{2}} (\bar{\nu} P_R \chi + \text{h.c.}) \quad (6.20)$$

Unlike standard studies, instead of considering right handed neutrinos in the lepton doublets we have introduced χ_R states as well as the additional inert χ_L . In this sense this model explicitly breaks the true left-right symmetric nature of the setup. Since χ_L is a singlet it forms a Yukawa coupling with the Φ and the right handed doublet,

$$\mathcal{L} \supset y_\chi \Phi \bar{\chi} P_R \begin{pmatrix} e \\ \chi \end{pmatrix} + \text{h.c.} \supset \frac{y_\chi u}{\sqrt{2}} \bar{\chi} \chi, \quad (6.21)$$

preventing ν_L and χ_R from forming a Dirac fermion. We assume $y_\chi u \gg y_\nu v$ such that the Standard Model neutrino is effectively massless and m_χ is a free parameter.⁴ After integrating out the W_R boson we get the quark level interactions:

$$\mathcal{L} \supset \frac{g_R^2}{M_{W_R}^2} \frac{1}{2} [\bar{u} \gamma_\mu P_R d] [\bar{e} \gamma^\mu P_R \chi] + \text{h.c.} . \quad (6.22)$$

In terms of the nucleons this gives the interaction:

$$\mathcal{L} \supset \frac{g_R^2}{4M_{W_R}^2} [\bar{p} \gamma_\mu (1 + \lambda \gamma_5) n] [\bar{e} \gamma^\mu P_R \chi] + \text{h.c.} , \quad (6.23)$$

where $\lambda \simeq 1.2694 \pm 0.0028$ is the axial to vector coupling from data [381].

We now consider possible decays as shown in Fig. 6.1 (**right**). The safest possibility is to only charge the first generation under the new $SU(2)_R$ to minimize the mixing between the W and W_R , and so we focus on this case⁵. One loop radiative corrections induce a log-divergent mixing between W_R and the Standard Model W boson which vanishes at u , and at low energies, is approximately

$$\mathcal{L} \supset \left(\frac{g_L^2 g_R^2 m_u m_d}{(4\pi)^2 M_{W_R}^2 M_W^2} \log \frac{u}{\Lambda_{\text{QCD}}} \right) g_{\mu\nu} j_L^\mu j_R^\nu \quad (6.24)$$

where $j_{L,R}^\mu$ are the left and right gauge currents (defined without the couplings). Below the QCD scale there is an additional contribution from the running which we estimate at leading order using chiral perturbation theory. Starting with the chiral Lagrangian ($\Sigma \equiv e^{i\pi^a \sigma^a / f_\pi}$) we can extract the mixing with the pions:

$$\frac{f_\pi^2}{4} \text{Tr} [(D_\mu \Sigma)^\dagger D^\mu \Sigma] \supset \frac{f_\pi}{2} [g_L W_\mu^+ + g_R W_{R\mu}^+] \partial^\mu \pi^- + \text{h.c.} \quad (6.25)$$

Integrating out the pions induces a coupling between the left and right handed currents:

$$\mathcal{L} \supset \left(\frac{g_L^2 g_R^2 f_\pi^2}{8m_\pi^2} \frac{1}{M_{W_R}^2 M_W^2} \right) \partial^\mu j_{R,\mu} \partial^\nu j_{L,\nu} + \text{h.c.} , \quad (6.26)$$

⁴ y_ν is not needed, but is included since no symmetry forbids it. However, we do assume that it is sufficiently small to prevent significant decays to 3 left-handed neutrinos through the Standard Model 4-neutrino coupling.

⁵In addition to making the W - W_R mixing worse, charging more Standard Model generations would introduce new, accompanying dark-sector states for the heavier Standard Model leptons. If they comprise a fraction of dark matter, they would not produce charged current signals at experiments as processes with a dark matter absorbed and a heavy lepton emitted would be kinematically forbidden. Thus, the charged current signals would only come from the fraction of “first generation” dark matter.

which of the two terms dominates will depend on the mass of χ . The decay rates (ignoring the interference terms) are:

$$\Gamma_{\chi \rightarrow e^+ e^- \nu}^{(1\text{-loop})} = (16 \log 2 - 11) \frac{m_\chi^5}{512\pi^3} \left(\frac{g_L^2 g_R^2 m_u m_d}{(4\pi)^2 M_{W_R}^2 M_W^2} \log \frac{m_u}{\Lambda_{\text{QCD}}} \right)^2, \quad (6.27)$$

$$\Gamma_{\chi \rightarrow e^+ e^- \nu}^{(\pi)} = \left(\log 4 - \frac{31}{24} \right) \frac{m_\chi^7 m_e^2}{256\pi^3} \left(\frac{g_L^2 g_R^2 f_\pi^2}{8m_\pi^2} \frac{1}{M_{W_R}^2 M_W^2} \right)^2. \quad (6.28)$$

Note that Eq. (6.27) will in general be less constraining than decays in the neutral current model as Eq. (6.27) contains additional factors of inverse mediator when compared to Eq.(6.9)- (6.11).

As in the case of the neutral current, one can look directly for the operators we consider here without requiring the presence of χ as dark matter. The most powerful direct search arise from collider physics from searches for heavy charged states. For simplicity we focus on the limit where W_R is heavy such that it is never produced on-shell. In this case, there are limits using the energy-enhanced nature of the $u\bar{d} \rightarrow e\chi$ process however we note that this process does not interfere with any Standard Model rate resulting in most collider searches being inapplicable (as they rely on a final state neutrino). Nevertheless, there are searches at 8 TeV which look for helicity-non-conserving contact interactions which should be roughly applicable here [380]. We estimate that these restrict $g_R^2/4M_{W_R}^2 \lesssim (4.5 \text{ TeV})^{-2}$.

In addition to direct collider searches, one can look for deviations from the Standard Model in known β decays. For any such decay, the Standard Model prediction is hard to evaluate rendering it challenging to use these process for precision searches for new physics in the limit that the scale of the higher dimensional operator is well above the weak scale. Nevertheless, it was suggested to use super-allowed (Fermi) transitions in between isotopes with vanishing spin and unit parity ($I^P = 0^+ \rightarrow 0^+$) [383] (see also [398–403] for earlier work). Such transitions are insensitive to the axial part of the operator and the vector contribution does not get renormalized under QCD [404] (what became known as the *conserved vector current* hypothesis), which makes it possible to compute the rates to the sub-percent level. In [383], constraints are put on operators of the form, $\Lambda_\nu^{-2} [\bar{p}\Gamma_i n] [\bar{e}\Gamma_j \nu]$, which can interfere with the Standard Model amplitudes resulting in a limit on the operator cutoff scale: $\Lambda_\nu^2 \gtrsim 10^3 G_F$ ($G_F \simeq 10^{-5} \text{ GeV}^{-2}$ is the Fermi constant). Computing these constraints for the operators of interest here is an involved task and beyond the scope of this Chapter. Instead, we roughly estimate the sensitivity β decay experiments can have assuming a similar analysis can be done for operators involving χ . The constraint on Λ_ν is sensitive to the interference term between the Standard Model and new operator term which is $\mathcal{O}(G_F \Lambda_\nu^{-2})$, while for χ operators there is no interference. Defining the higher dimensional operator for charged current fermion absorption with a scale Λ , the leading term in the β experiments is $\mathcal{O}(\Lambda^{-4})$. Equating the observed limit to this operator we find that if such a search were carried out we would expect a sensitivity of order, $\Lambda \gtrsim 1.5 \text{ TeV}$, which is weaker than present collider bounds. We also emphasize that such constraints depend critically on the mass of χ — when $m_\chi \gtrsim \mathcal{O}(\text{MeV})$ different β decay channels become kinematically unavailable, quickly

weakening the constraints. Other possible ways to handle the nuclear uncertainties are using the neutron lifetime and angular distributions in nuclear decays, however the constraints using these techniques are weaker than the ones estimated above [384].

Outside of nuclear decays its possible to use charged pions decay searches looking for $\pi^\pm \rightarrow e^\pm \chi$ [405], however these searches are not able to extend to χ masses below 60 MeV due to backgrounds from muon decays, which will be outside our range of interest for the charged current operator. Lastly, we comment that, as for the neutral current operator UV completion, we do not expect significant constraints from star cooling, beam dump, or supernovae since the scale of the effective operators we consider here are above the weak scale.

6.3 Relevant Current Experiments

In this Section, we summarize the relevant details of all current experiments for which we project sensitivities in Sections 6.4 through 6.6 (Figs. 6.4, 6.6, and 6.7). A few additional comments are in order for some of these experiments. We give projections based on Run 2 of CDMSlite [360] and not Run 3 [406] since Run 2 had a larger exposure and roughly the same threshold. We conservatively underestimate Borexino's exposure by assuming that the 3218 days which had at least an 8-hr exposure only had an 8-hr exposure. Borexino's electron equivalent energy threshold is close to 70 keV [407], which corresponds to a proton recoil threshold of 500 keV [408]. It's also worth noting that the Carbon recoil threshold is too high thanks to its poor relative light yield. EXO-200 is 80% $^{136}_{54}\text{Xe}$, while KamLAND-Zen is 91%.

6.4 Neutral Current Nuclear Recoils

We first study the nuclear recoils from the dimension-6 neutral current operator generated by the UV model discussed in Sec. 6.2 with the identification $1/\Lambda^2 \equiv Q_\chi g_\chi^2 s_{\theta_R} c_{\theta_R} / m_{Z'}^2$,

$$\frac{1}{\Lambda^2} (\bar{n} \gamma^\mu n + \bar{p} \gamma^\mu p) \bar{\chi} \gamma_\mu P_R \nu + \text{h.c.} \quad (6.29)$$

This operator leads to the the nuclear recoil process;

$$\chi(m_\chi \vec{v}) + N(\vec{0}) \rightarrow \nu(\vec{p}_\nu) + N(\vec{q}), \quad (6.30)$$

in which an incoming dark matter with velocity $\vec{v}$ is absorbed by a target nucleus N at rest which then recoils with momentum $\vec{q}$ against the light ν of momentum $\vec{p}_\nu$. While this vector structure is inspired by the Z' model, we emphasize that these signals can arise from operators with a more general Lorentz structure for which the formalism that follows can also be applied.

Experiment	Goal	Exposure	Target	$E_{\text{NR}}^{\text{th}}$	Refs
CRESSTII	DM	52 kg day	CaWO ₄ crystals	307 eV	[354]
CRESSTIII	DM	2.39 kg day	CaWO ₄ crystals	100 eV	[353]
DAMIC	DM	0.6 kg day	Si CCDs	0.7 keV	[357]
DarkSide-50	DM	6786 kg day	Liquid Ar	0.6 keV	[358, 359]
EDELWEISS	DM	.0334 kg day	Ge	0.06 keV	[355]
LUX	DM	91.8 kg yr	Liquid Xe	4 keV	[123]
NEWS-G	DM	9.7 kg day	Neon	720 eV	[356]
PandaX-II	DM/ $0\nu 2\beta$	150 kg yr	Liquid Xe	3 keV	[366]
PICO-60	DM	3420 kg day	Superheated CF ₃ I	13.6 keV	[363]
PICO-60	DM	1167 kg day	Superheated C ₃ F ₈	3.3 keV	[362]
SuperCDMS	DM	577 kg day	Ge crystals	1.6 keV	[361]
CDMSlite	DM	70 kg day	Ge crystals	0.4 keV	[360]
XENON1T	DM	1.0 t yr	Liquid Xe	3 keV	[338]
CUORE	$0\nu 2\beta$	86.3 kg yr	TeO ₂ crystals	100 keV	[375]
EXO-200	$0\nu 2\beta$	233 kg yr	Liquid $^{136}_{54}\text{Xe}$	—	[376]
KamLAND-Zen	$0\nu 2\beta$	504 kg yr	$^{136}_{54}\text{Xe}$ in LS	—	[377]
Borexino	solar ν	817 t yr	C ₆ H ₃ (CH ₃) ₃	500 keV	[367]
COHERENT	CE ν NS	6726 kg day	CsI[Na]	6.5 keV	[364, 365]
Super-Kamiokande	ν	171,000 t yr	H ₂ O	—	[378]

Table 6.1: Experiments which can probe fermionic dark matter absorption signals for which we show projected sensitivities in Figs. 6.4, 6.6, and 6.7. Experiments without an explicit nuclear recoil threshold, $E_{\text{NR}}^{\text{th}}$, are not used for neutral current projections.

The relevant experimental observable is the differential scattering rate per nuclear recoil energy. We begin with the usual differential cross section,

$$\begin{aligned} d\sigma &= \frac{|\overline{\mathcal{M}_N}|^2}{4E_\chi E_N v} \prod_j \frac{d^3 p_j}{2E_j (2\pi)^3} (2\pi)^4 \delta^4(p_\chi^\mu + p_i^\mu - p_\nu^\mu - p_f^\mu) \\ &= \sqrt{\frac{E_R}{2M}} \frac{|\overline{\mathcal{M}_N}|^2 dE_R d(\cos \theta_{qv})}{16\pi v m_\chi p_\nu} \delta(E_R + p_\nu - m_\chi(1 + v^2/2)), \end{aligned} \quad (6.31)$$

where $|\overline{\mathcal{M}_N}|^2$ is the matrix element squared averaged over initial and summed over final spins, $p_{i(f)}^\mu$ is the initial (final) four-momentum of the nucleus, $E_R = q^2/2M$ is the energy of the recoiling nucleus, M is the mass of the nucleus, θ_{qv} is the angle between $\vec{v}$ and $\vec{q}$, and $p_\nu = \sqrt{m_\chi^2 v^2 + q^2 - 2m_\chi v q \cos \theta_{qv}}$. The incoming dark matter is non-relativistic, so its energy is roughly equal to its mass. Dropping $\mathcal{O}(v)$ terms, the energy-conserving δ function simplifies to⁶

$$\delta(E_R + p_\nu - m_\chi(1 + v^2/2)) \simeq \frac{m_\chi}{M} \delta(E_R - E_R^0), \quad (6.32)$$

where $E_R^0 = m_\chi^2/2M$ since $m_\chi \ll M$. Thus, the differential cross section reduces to

$$\frac{d\sigma}{dE_R} = \frac{|\overline{\mathcal{M}_N}|^2}{16\pi v M^2} \delta(E_R - E_R^0). \quad (6.33)$$

The differential scattering rate per nuclear recoil energy in an experiment is related to this differential cross section by

$$\frac{dR}{dE_R} = N_T n_\chi \left\langle \frac{d\sigma}{dE_R} v \right\rangle \Theta(E_R^0 - E_{\text{th}}), \quad (6.34)$$

where N_T is the number of nuclear targets in the experiment, n_χ is the local number density of dark matter, the average is performed over the incoming dark matter's velocity distribution, and $\Theta(E_R^0 - E_{\text{th}})$ approximates the nuclear recoil energy threshold of the experiment with a step function (see Section 6.3 for a summary of E_{th} for the experiments considered here). The average over the dark matter velocity distribution is trivial and yields

$$\frac{dR}{dE_R} = N_T \frac{\rho_\chi}{m_\chi} \sigma_{\text{NC}} A^2 F(q)^2 \delta(E_R - E_R^0) \Theta(E_R^0 - E_{\text{th}}), \quad (6.35)$$

where $\rho_\chi \simeq 0.4 \text{ GeV/cm}^3$ is the local dark matter energy density, $\sigma_{\text{NC}} = m_\chi^2/(4\pi\Lambda^4)$ is the absorption cross section per nucleon, A is the atomic mass number of the target nucleus, and $F(q)$ is the Helm form factor [370] of the target nucleus (normalized to 1).

⁶ ν_R could in fact be any neutral, light fermion in the dark sector. As long as it is much lighter than dark matter, $q \sim m_\chi \gg m_{\nu_R}$ gives rise to the distinctive fermionic absorption nuclear recoil spectrum.

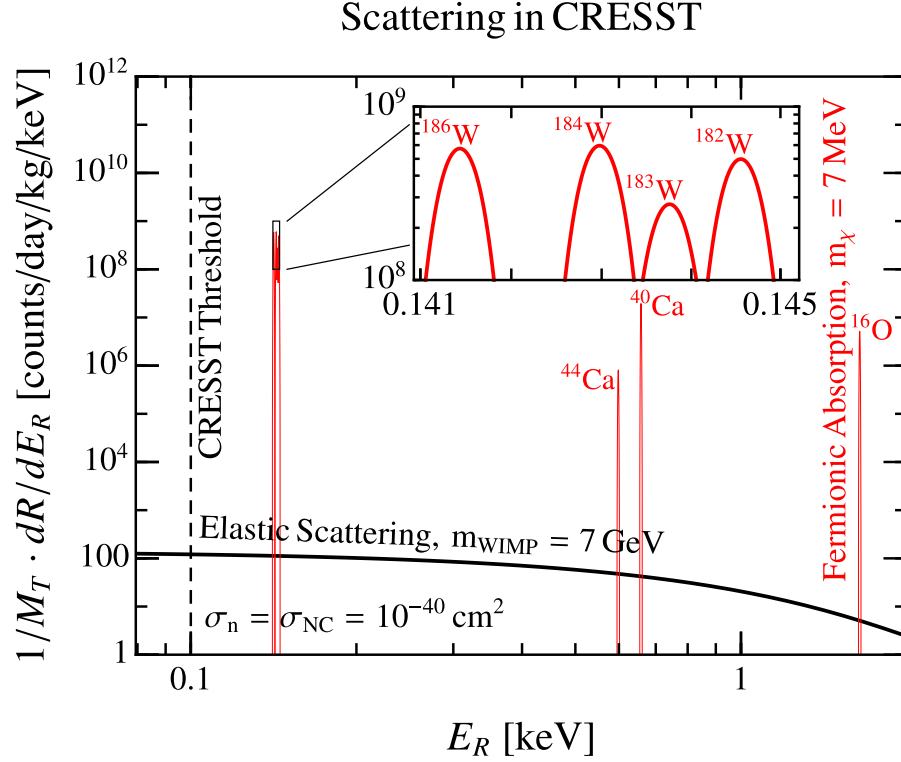


Figure 6.2: Differential scattering rate per recoil energy per detector mass at CRESST [353] from fermionic absorption of a $m_\chi = 7$ MeV dark matter with $\sigma_{\text{NC}} = 10^{-40} \text{ cm}^2$. Also shown is the elastic scattering rate for a WIMP with mass $m_{\text{WIMP}} = 7$ GeV and spin-independent cross section per nucleon $\sigma_n = 10^{-40} \text{ cm}^2$, along with the CRESSTIII nuclear recoil threshold at 100 eV. The figure inset zooms in on the bunched peaks corresponding to the four isotopes of Tungsten.

This scattering rate is different from the usual elastic scattering rate for dark matter (see [122] for a recent review) since the typical recoil energy from an elastic scatter is of the order $v^2 \mu_{\chi N}^2 / M$ (where $\mu_{\chi N}$ is the χ - N reduced mass), while the fermionic absorption recoil is peaked at $m_\chi^2 / 2M$. Therefore, for a fixed dark matter mass, the nuclear recoil energy for fermionic absorption is $1/v^2 \sim 10^6$ times larger than that of usual elastic recoil. This allows direct detection experiments to probe dark matter candidates roughly $1/v \sim 10^3$ times lighter than normal, in addition to allowing neutrino detectors with larger exposures but higher thresholds to make competitive searches. In order to highlight the differences between fermionic absorption and elastic scattering rates, we compare the differential scattering rates per recoil energy per detector mass (M_T) at one particular experiment, CRESST [353], in Fig. 6.2. To make an illustrative comparison, we set the spin-independent WIMP cross-section equal to the absorption cross section per nucleon, which we set as $\sigma_{\text{NC}} = 10^{-40} \text{ cm}^2$, and show the elastic rate for a heavier WIMP, $m_{\text{WIMP}} = 7$ GeV, while taking 7 MeV for

the fermion absorption signal. To obtain the finite heights and widths of the fermionic absorption peaks which are not given by the δ function in Eq. (6.35), we calculate the differential scattering rate after expanding the energy-conserving δ function to first order in v (see below for details). CRESST illustrates the differences in scattering rates well because it contains multiple target isotopes in its CaWO_4 crystals which give rise to four peaks from absorbing fermionic dark matter which are distinguishable if the energy resolution is less than 50 eV [353]. The figure demonstrates the relative ease with which experiments looking for fermionic absorption nuclear recoils can see the signal above the background by correlating the locations and heights of scattering rates off multiple target isotopes. Even in the absence of multiple distinguishable peaks, detectors can still use the peaked nature of the fermionic absorption differential scattering rate to differentiate the signal from the noise.

To produce Fig. 6.2, we need to evaluate the energy-conserving δ function in Eq. (6.31) at $\mathcal{O}(v^1)$ since evaluating it at $\mathcal{O}(v^0)$ yields a differential scattering rate proportional to a delta function in E_R (see Eq. (6.35)). At $\mathcal{O}(v^1)$, we find

$$\delta \left(E_R + p_\nu - m_\chi \left(1 + \frac{v^2}{2} \right) \right) \simeq \frac{\delta (\cos \theta_{qv} - \cos \theta_{qv}^0)}{m_\chi v}, \quad (6.36)$$

where

$$\cos \theta_{qv}^0 = \frac{E_R + \sqrt{2ME_R} - m_\chi}{m_\chi v}. \quad (6.37)$$

The superscript indicates that this is the value for $\cos \theta_{qv}$ at which the energy-conserving δ function's argument vanishes. This cosine's allowed range places a minimum condition on v :

$$v_{\min} = \frac{|E_R + \sqrt{2ME_R} - m_\chi|}{m_\chi}. \quad (6.38)$$

The differential scattering rate at $\mathcal{O}(v)$ is then

$$\frac{dR}{dE_R} = N_T \frac{\rho_\chi}{m_\chi} \sqrt{\frac{E_R}{2M}} \frac{|\mathcal{M}_N|^2}{16\pi m_\chi^2 p_\nu} \int d^3v \frac{f(v)}{v} \theta(v - v_{\min}) \quad (6.39)$$

where $p_\nu = \sqrt{2\sqrt{2}m_\chi\sqrt{E_RM} - 2E_R(M + \sqrt{2ME_R})}$. We approximate the dark matter velocity distribution with a capped Maxwell distribution (see [122] for a review)

$$f(\vec{v}) = \frac{1}{N} \exp \left[-\frac{(\vec{v} + \vec{v}_e)^2}{v_0^2} \right] \theta(v_{\text{esc}} - |\vec{v} + \vec{v}_e|), \quad (6.40)$$

where $N = \pi^{3/2} v_0^3 \left(\text{erf}[v_{\text{esc}}/v_0] - \frac{2v_{\text{esc}}}{\sqrt{\pi}v_0} \exp[-v_{\text{esc}}^2/v_0^2] \right)$ normalizes the velocity distribution to unity, $v_e \simeq 240$ km/s is the Earth's approximate galactic velocity (dominated by the Sun's),

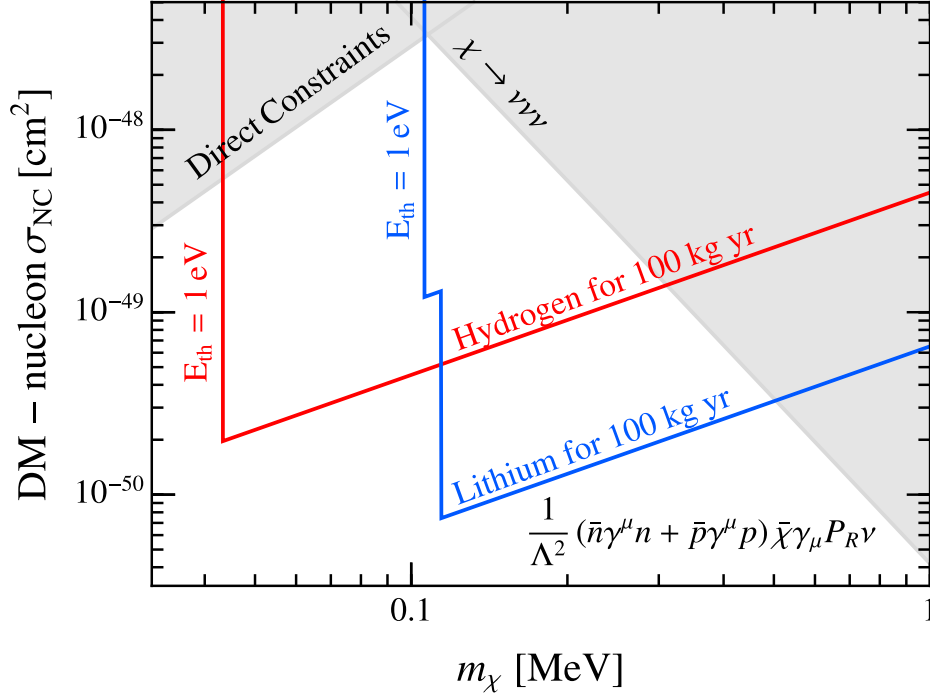


Figure 6.3: Projected upper bound on σ_{NC} as a function of m_χ at future detectors with Hydrogen or Lithium targets. Bounds for both potential targets are shown assuming $E_{\text{th}} = 1$ eV and $M_T T = 100$ kg yr. Also shown in gray are the constraints from direct searches for Z 's and decays of χ for the benchmark UV completion with $m_{Z'} = 18$ GeV, $s_{\theta_R} = 10^{-1.5}$, and $Q_\chi = 0.1$, as described in the text.

$v_0 \simeq 220$ km/s, and $v_{\text{esc}} \simeq 550$ km/s is the galactic escape velocity. Thus, the differential scattering rate on a single isotope j per target mass is

$$\frac{1}{M_T} \frac{dR_j}{dE_R} = \frac{\rho_\chi}{m_\chi} \sigma_{\text{NC}} \frac{N_j M_j \sqrt{2E_R M_j}}{2M_T m_\chi^2 p_\nu} A_j^2 F_j^2 \left\langle \frac{1}{v} \right\rangle_{v > v_{\text{min}}}, \quad (6.41)$$

where v_{min} is given by Eq. (6.38). With this, we produce the differential scattering rates from fermionic absorption off the few target isotopes in CRESST in Fig. 6.2.

Having discussed the novel signature of neutral current nuclear recoils from fermionic absorption, we now project the sensitivities of future and current experiments to this signal. Integrating the differential scattering rate over all recoil energies and summing over all isotopes j present in an experiment, we find the total event rate is

$$R = \frac{\rho_\chi}{m_\chi} \sigma_{\text{NC}} \sum_j N_{T,j} A_j^2 F_j(q)^2 \Theta(E_{R,j}^0 - E_{\text{th}}). \quad (6.42)$$

For simplicity, we project bounds on σ_{NC} by requiring < 10 events occur in a given experiment.

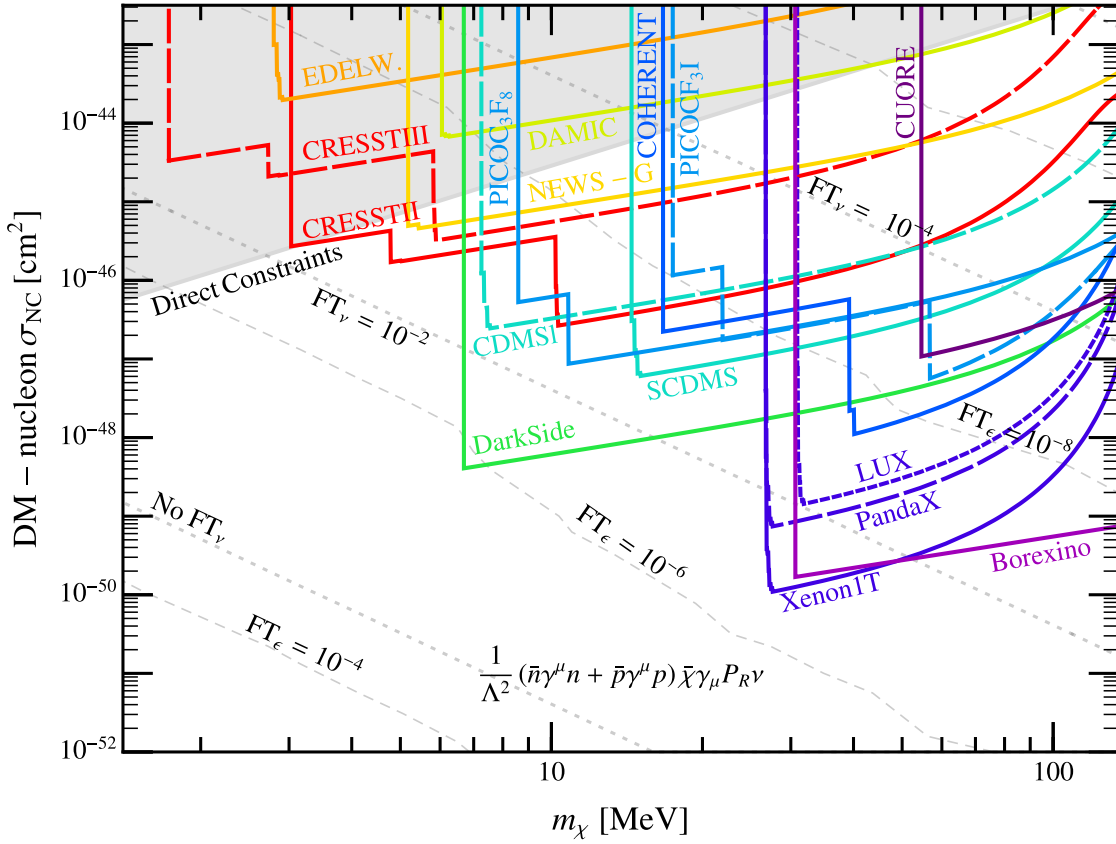


Figure 6.4: Projected upper bound on σ_{NC} as a function of m_χ at current experiments, including CUORE [375] (dark purple), Borexino [367] (purple), LUX [123] (dotted navy blue), PandaX-II [366] (dashed navy blue), XENON1T [338] (solid navy blue), COHERENT [364, 365] (blue), PICO-60 run with CF_3I [363] (dashed sky blue), PICO-60 run with C_3F_8 [362] (solid sky blue), SuperCDMS [361] (solid aqua), CDMSlite Run 2 [360] (dashed aqua), DarkSide-50 [358, 359] (green), DAMIC [357] (lime), NEWS-G [356] (yellow), EDELWEISS-SURF [355] (orange), CRESST II [354] (solid red), and CRESST III [353] (dashed red). Here we have taken $m_{Z'} = 18$ GeV, $s_{\theta_R} = 10^{-2}$, and $Q_\chi = 0.1$. Also shown are the constraints from direct searches for Z' s and decays of χ for the benchmark point chosen, as described in the text. The dashed grey contours show the level of fine-tuning necessary in our UV completion to avoid rapid $\chi \rightarrow \nu e^+ e^-$ and $\chi \rightarrow \nu \nu \nu$ decays in tension with indirect detection bounds [27, 391].

We start by considering the regime $m_\chi \lesssim \text{MeV}$. Since $E_R \propto 1/M$, lighter isotopes are particularly useful in probing such light candidates. Even still, reaching such light masses requires nuclear recoil thresholds lower than those of current experiments. We do not make detailed projections for any specific future experiment since they are diverse and would possibly involve absorption by collective modes rather than individual nuclei. Instead, we

simply make projections for scattering off individual nuclei of Hydrogen or Lithium in future experiments with $E_{\text{th}} = 1$ eV and $M_T T = 100$ kg yr in Fig. 6.3 (see [146, 156] for proposals). The kink in the Lithium line is due to the two naturally occurring isotopes, ${}^6\text{Li}$ and ${}^7\text{Li}$. Since $E_R \propto 1/M$ and we approximate the energy threshold with a step function, the kink occurs when the dark matter is too light to cause ${}^7\text{Li}$ to recoil with an energy above the threshold, but is still heavy enough to push ${}^6\text{Li}$ above the threshold.

While the projected sensitivities only rely on the defining neutral current absorption operator in Eq. (6.29), any UV completion will have other relevant constraints. In particular for our UV completion, setting $m_{Z'} = 18$ GeV, $s_{\theta_R} = 10^{-1.5}$, and $Q_\chi = 0.1$, the $\chi \rightarrow \nu\nu\nu$ decay is the most constraining indirect bound [391]. Additionally, searches for such a Z' [369] and bounds on Standard Model four-fermion interactions [393] are the most stringent direct constraints. Both direct and indirect bounds are shown as gray regions in the figure. We see that moderate, achievable energy thresholds and exposures of light isotopes can quickly probe viable, unexplored parameter space for $m_\chi \lesssim$ MeV.

Next, we consider heavier dark matter with $m_\chi \gtrsim$ MeV. This dark matter is sufficiently heavy to cause nuclear recoils above detector thresholds at existing experiments, as shown in Fig. 6.4. In fact, the recoil energies can be so large that they allow non-dark matter specialized experiments, such as Borexino [367], COHERENT [364, 365], and CUORE [375], to probe viable parameter space. We summarize the relevant details of each current experiment in Section 6.3. Similar to our discussion of the kinked Lithium line, every kink in this figure corresponds to a particular isotope's recoil energy dropping below its corresponding experiment's recoil threshold.

As discussed before, any particular UV completion of the neutral current absorption operator will have additional constraints. We set $m_{Z'} = 18$ GeV, $s_{\theta_R} = 10^{-2}$, and $Q_\chi = 0.1$ for our simple UV completion in the figure. Fine-tuning of ϵ in Eq. (6.9) is needed to avoid indirect detection bounds on $\chi \rightarrow \nu e^+ e^-$ decays [27], denoted with dashed grey contours labeled FT_ϵ in the figure. Additionally, fine-tuning against the IR contribution to the $\chi \rightarrow \nu\nu\nu$ decay in Eq. (6.11) is necessary to varying degrees [391], denoted with dashed grey contours labeled FT_ν . These fine-tuning contours further motivate future iterations of current experiments with larger exposures and lower thresholds, such as Argo [409], DARWIN [339], PICO-500, and SuperCDMS SNOLAB [410]. However, the requirement of fine-tuning to evade indirect detection bounds is highly model-dependent. For example, introducing flavor-dependent couplings might greatly reduce the need for any fine-tuning. Regardless, the projected bounds on σ_{NC} are model-independent and encourage both searches for these neutral current fermionic dark matter absorption signals at current experiments and the study of UV completions of these operators which more naturally suppress decays bounded by indirect detection.

6.5 Charged Current: Induced β Decays

We now study the signals from the dimension-6 charged current operator generated by the UV model discussed in Sec. 6.2:

$$\frac{1}{\Lambda^2} [\bar{p} \gamma_\mu (1 + \lambda \gamma_5) n] [\bar{e} \gamma^\mu P_R \chi] + \text{h.c.}, \quad (6.43)$$

where $\lambda \simeq 1.2694 \pm 0.0028$, with we identify $1/\Lambda^2 \equiv g_R^2/4M_{WR}^2$. For sufficiently massive dark matter, scattering on a nucleus can result in the conversion of a neutron/proton within the nucleus into a proton/neutron, accompanied by the emission of an energetic $e^\pm$ in the final state — analogous to the familiar induced $\beta^\mp$ processes in neutrino physics. This processes can lead to a variety of possible correlated signals depending on the target nucleus and dark matter mass. We now consider the scenario in which dark matter induces β transitions in isotopes that are stable against β decay in a vacuum. Signals arising from the decays of stable isotopes are particularly appealing as such nuclei exist in large abundances within the target material of current direct detection and neutrino experiments. We reserve study of unstable isotopes and the effect of dark matter transitions on the kinematic endpoint of their β decay spectrum for Sec. 6.6.

Induced β transitions will occur if the dark matter mass is above the kinematic threshold given by:

$$m_\chi > m_{\text{th}}^{\beta^\mp} \equiv M_{A,Z\pm 1}^{(*)} + m_e - M_{A,Z}. \quad (6.44)$$

Throughout this Chapter, we take $M_{A,Z}$ to be the mass of the *nucleus* of the isotope ${}^A_Z\text{X}$. Note that if $M_{A,Z-1} < M_{A,Z} - m_e$, then the nucleus ${}^A_Z\text{X}$ can undergo electron capture. These isotopes are generally not long-lived, so the scenario of induced β^+ decay is most interesting for isotopes where electron capture is kinematically forbidden. Thus, for dark matter induced β^+ decay, we limit ourselves to isotopes with $M_{A,Z-1} > M_{A,Z} - m_e$. An additional complication occurs for β^+ decays in heavy isotopes where in general the number of neutrons is far greater than the number of protons, and Pauli Blocking effects would make β^+ transitions into the ground state or lowest lying excited states of the daughter nucleus disfavored. This motivates us to focus on signals of induced β^+ decays in experiments containing Hydrogen as a target.

If the energy imparted upon the proton/neutron by the dark matter, *i.e.* m_χ , is significantly less than the binding energy of the nucleus $\lesssim 10$ MeV, the proton/neutron will not have enough energy to escape and the dominant process will be from the outbound nucleon remaining bound to the nucleus, leading to two possible processes:

$$\begin{aligned} \beta^-: \quad \chi + n &\rightarrow p + e^- \Rightarrow \chi + {}^A_Z\text{X} \rightarrow {}^A_{Z+1}\text{X}^{(*)} + e^-, \\ \beta^+: \quad \bar{\chi} + p &\rightarrow n + e^+ \Rightarrow \bar{\chi} + {}^A_Z\text{X} \rightarrow {}^A_{Z-1}\text{X}^{(*)} + e^+. \end{aligned} \quad (6.45)$$

Note that as a result of angular momentum conservation considerations, the daughter nucleus will generically be produced in an excited state. For dark matter masses greater than 10

MeV (where the incoming dark matter also begins to resolve the individual nucleons upon scattering), other signals are in general possible. In particular, with enough energy the incoming dark matter particle can break apart the nucleus and the ejected proton/neutron will then hadronize and shower for energies above the QCD scale, leading to an array of possible new signals. The inclusive cross section in this case may be computed using standard techniques [381]. However, dark matter decay rates scale as a large power of m_χ (see Eq. (6.27) for the expressions for our charged current UV completion), and so will tend to be in conflict with astrophysical constraints at larger m_χ for detectable cross-sections. We leave a detailed study of this regime to future work.

The kinematics of the induced β process in the limit that $m_\chi \ll M_{A,Z}$ are straightforward to compute. The energy of the outgoing $e^\pm$ and daughter nucleus are:

$$E_R \simeq \begin{cases} m_\chi - m_{\text{th}}^\beta & (\text{electron}) \\ (m_\chi - m_{\text{th}}^\beta)^2 / 2M_{A,Z-1}^{(*)} & (\text{nucleus}) \end{cases} \quad (6.46)$$

Note that the recoil energies parallel the neutral current case with the replacement $m_\chi \rightarrow m_\chi - m_{\text{th}}^\beta$ for $m_\chi > m_e$. As can be seen from Eq. 6.46 the recoiling nucleus signal will be, as with the neutral current signal, velocity independent to leading order.

We now compute the inclusive rate for dark matter induced β decays and make projections for current experimental sensitivities. We defer discussion of specific signals to the end of this section but note that these charged current signals are typically well above experimental thresholds and have several correlated signals. Hence for most experiments of interest, the events are striking enough that they should easily pass experimental cuts. The signal rate for an experiment carrying a set of target isotopes parameterized by j is given by:

$$R = \frac{\rho_\chi}{2m_\chi} \sum_j N_{T,j} n_j \langle \sigma v \rangle_j, \quad (6.47)$$

where ρ_χ is the local dark matter density, and $N_{T,j}$ is the number of targets of a given isotope. Here we have assumed that any $e^\pm$ energy could potentially be detected, and integrate over all energies and angles of $e^\pm$ emission). An additional factor n_j accounts for the total number of parton level targets for the β decay:

$$n_j \simeq \begin{cases} A_j - Z_j & (\beta^-) \\ Z_j & (\beta^+) \end{cases} \quad (6.48)$$

We emphasize that the rate scales with target volume and experimental exposure. Therefore, computing the rate as in Eq. (6.47) requires experimental input along with the computation of the scattering cross section for dark matter off nucleons.

For m_χ below the binding energy of nucleons ($m_\chi \lesssim 10$ MeV), the absorption process cannot resolve the constituents of the nucleons and we consider only scattering off entire nuclei. In this regime, we may write the differential cross section (in the center of mass

frame) as⁷

$$\frac{d\sigma}{d\Omega} = \frac{1}{64\pi^2 E_{\text{cm}}^2} \frac{|\vec{p}_e|}{|\vec{p}_\chi|} \sum_{\text{transitions}} |\overline{\mathcal{M}_N}|^2, \quad (6.49)$$

where Ω is the solid angle θ is the angle between the incoming dark matter and the emitted electron/positron, $\mathcal{M}_N$ the amplitude is for scattering off of nucleons, $\vec{p}_e$ ($\vec{p}_\chi$) is the momentum of the electron (dark matter), and here we sum over possible nuclear spin states which manifests as a sum over the allowed transitions. Note that Eq. (6.49) holds for both induced β^- and β^+ decays. However, the nuclear rate will depend on which of the two processes is under-consideration, and we discuss this further in what follows.

We emphasize that any large exposure experiments (both neutrino and designated dark matter direct detection) can be re-purposed to search for induced beta decays from fermionic absorption provided the kinematic threshold, Eq. (6.44), is low enough. A list of select stable (or meta-stable) isotopes in which light dark matter could induce β transitions is summarized in Table 6.2. The isotopes are grouped by their stability against β^- or β^+ decay, and for each possible transition we quote the value of the threshold given by Eq. (6.44) for the $\Delta I = 0$ and $\Delta I = \pm 1$ transition with the lowest threshold — (though other transitions are also generally accessible — see Fig. 6.5 for a summary of various transitions corresponding to the experimental targets we consider here).

The plethora of different isotopes that undergo induced β^- transitions leads to a wide range of signals at various experiments utilizing different target materials. For a list of experiments and corresponding target materials and exposures, see Section 6.3. From Table 6.2 we see that the lowest possible dark matter mass that the induced beta decay signal can probe is about $m_\chi \sim 355$ keV from absorption by a $^{131}_{54}\text{Xe}$ nucleus at a Xenon based experiment. Additionally note that $^{125}_{52}\text{Te}$ has a particularly low threshold of 374 keV, making CUORE, which utilizes TeO_2 crystals and was designed to search for neutrino-less double β decay, particularly suited to probe sub-MeV dark matter masses. For the charged current process, we focus entirely on current experiments. Interestingly, there is one transition in the Standard Model with an anomalously small threshold of 2.5 keV that can employ $\Delta I = \pm 1$ transitions, $\chi + {}^{163}_{66}\text{Dy} \rightarrow {}^{163}_{67}\text{Ho} + e^-$. This was considered in [161] as a way to probe sterile neutrinos. Unfortunately, the tremendous expense of building large volume experiments filled with Dysprosium make it a challenging direction to observe fermionic absorption.

Induced β^- Decays

We now focus specifically on signals from induced β^- decays

$$\chi(m_\chi \vec{v}) + {}^A_Z\text{X}(\vec{0}) \rightarrow e^-(\vec{p}_e) + {}^A_{Z+1}\text{X}(\vec{q}), \quad (6.50)$$

⁷Note that to $\mathcal{O}(v^0)$, the center of mass frame is approximately the lab frame, and we drop any indices to this affect.

	Process	Isotope (Threshold $\Delta I = 0$, $\Delta I = \pm 1$)
β^- :	${}^A_6\text{C} \rightarrow {}^A_7\text{N}$	${}^{12}_6\text{C}(18.3 \text{ MeV}, 17.3 \text{ MeV}), {}^{13}_6\text{C}(2.22 \text{ MeV}, 5.7 \text{ MeV})$
	${}^A_8\text{O} \rightarrow {}^A_9\text{F}$	${}^{16}_8\text{O}(16.4 \text{ MeV}, 16.4 \text{ MeV}), {}^{17}_8\text{O}(2.76 \text{ MeV}, 3.75 \text{ MeV}),$ ${}^{18}_8\text{O}(2.70 \text{ MeV}, 1.65 \text{ MeV})$
	${}^A_{52}\text{Te} \rightarrow {}^A_{53}\text{I}$	${}^{130}_{52}\text{Te}(1.42 \text{ MeV}, 6.62 \text{ MeV}), {}^{128}_{52}\text{Te}(2.25 \text{ MeV}, 1.25 \text{ MeV}),$ ${}^{126}_{52}\text{Te}(3.1 \text{ MeV}, 1.41 \text{ MeV}), {}^{125}_{52}\text{Te}(430 \text{ keV}, 374 \text{ keV})$
	${}^A_{54}\text{Xe} \rightarrow {}^A_{55}\text{Cs}$	${}^{129}_{54}\text{Xe}(1.19 \text{ MeV}, 1.33 \text{ MeV}), {}^{131}_{54}\text{Xe}(570 \text{ keV}, 355 \text{ keV}),$ ${}^{134}_{54}\text{Xe}(2.23 \text{ MeV}, 490 \text{ keV}), {}^{136}_{54}\text{Xe}(1.09 \text{ MeV}, 1.06 \text{ MeV})$
	β^+ : ${}^1_1\text{H} \rightarrow n$	${}^1_1\text{H}(1.8 \text{ MeV})$

Table 6.2: Here we summarize notable isotopes that could undergo induced β transitions, with a focus on the most abundant isotopes of materials interesting for the experiments we consider here (see Section 6.3). The right hand column displaces the threshold Eq. (6.44) for the transition with the smallest splitting between the ground state of the parent nucleus and the excited state of the daughter allowed by the selection rule in question ($\Delta I = 0, \Delta I = \pm 1$). Note that for Gamow Teller $\Delta I = \pm 1$ ground state to ground state transitions are possible leading to the sub-keV thresholds. For induced β^- transitions many processes exist with low threshold (see Fig. 6.5 for additional possible higher threshold transitions), however for β^+ transitions Pauli blocking effects in heavy isotopes will generally disfavor transitions into the lowest lying excited states of the daughter nucleus. This leads us, in the present chapter, to consider only induced β^+ processes involving target materials with Hydrogen.

and compute projected experimental limits by computing the expected rate Eq. (6.47) at specific experiments given current exposures. The nucleon level amplitude, required to compute the differential scattering cross section in Eq. (6.49), may then be written as follows:

$$\mathcal{M}_N = \sqrt{\mathcal{F}(Z+1, E_e)} \mathcal{M}. \quad (6.51)$$

Here $\mathcal{M}$ is the parton level scattering amplitude for β^- transitions $\chi + n \rightarrow p + e^-$ generated by the operator in (6.43), with the p/n momentum normalized to nuclei mass *i.e.* $p^\mu p_\mu = M_{A_j, Z_j}^2$ ⁸. The factor $\mathcal{F}(Z+1, E)$ in Eq. (6.51) is the usual Fermi function accounting for

⁸Note that the vector and axial vector form factors are implicitly defined in (6.43) for low momentum transfer $q^2 = (k_\chi - k_e)^2$, and agree with results from the neutrino literature namely $f_V(q^2 \sim 0) = 1$ and $f_A(q^2 \sim 0) = -1.2694$ [381]. Note that while the technology exists to compute form factors for scattering between different nuclei generated by general operators, we are unaware of such a computation carried out in the literature.

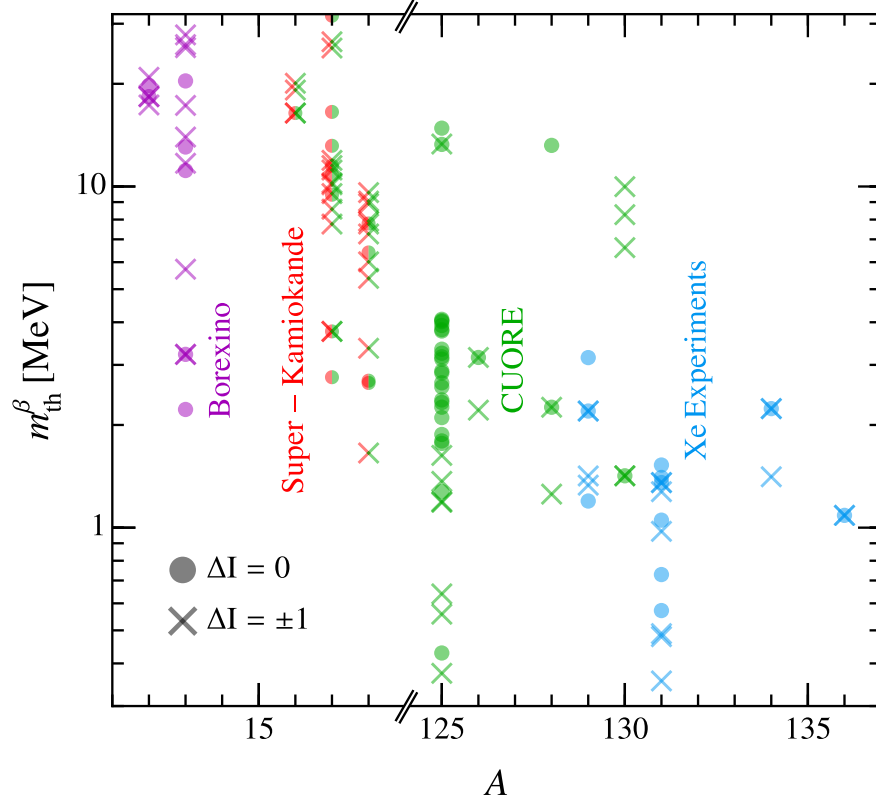


Figure 6.5: A comprehensive summary of every induced β^- decay we consider, plotted in the A versus m_{th}^β . Circles represent nuclear transitions in which the nuclear spin does not change, while crosses represent those which change by 1. Note that Super-Kamiokande and CUORE both contain oxygen, hence their overlap from $A = 16 - 18$.

Coloumb interactions between nucleons [381] for β transitions, and is given by:

$$\mathcal{F}(Z, E_e) = 2(1 + S) \frac{|\Gamma(S + i\eta)|^2}{\Gamma(1 + S)^2} (2r_N |\vec{p}_e|)^{2S-2} e^{\pi\eta}, \quad (6.52)$$

where $\eta = \alpha Z E_e / |\vec{p}_e| = \alpha Z E_e / \sqrt{E_e^2 - m_e^2}$ and $S = \sqrt{1 - \alpha^2 Z^2}$. Here the nuclear radius is given by $r_N = 1.2 \text{ fm } A^{1/3}$. For large $E_e \gg m_e$, the Fermi function asymptotes to a larger for isotopes with larger Z — therefore experiments utilizing light target nuclei (such Super-Kamiokande and Borexino) are particularly sensitive at larger dark matter masses (as they result in higher energy electrons).

As discussed above, Eq. (6.49) includes a sum over all possible nuclear spin states. The Lorentz structure of a given charged current operator dictates what kind of angular momentum selection rules are in effect, and therefore which nuclear transitions are allowed. For the model considered here the operator in Eq. (6.49) contains both vector and axial vector

couplings. In the case of a pure vector operator and light dark matter, Fermi transitions will dominate (other transitions will be suppressed by factors of $e^{-m_\chi r_N}$). Fermi transitions are transitions in which the spin of the dark matter and the electron are parallel so that the spin angular momentum of the initial and final nucleus is unchanged $\Delta I_F = 0$. Meanwhile, for axial and axial-vector couplings, Gamow-Teller transitions become possible and contribute to the sum. In these transitions the dark matter and electron have anti-parallel spins so that the nucleus spin must change to conserve angular momentum $\Delta I_{GT} = 1$. Both Fermi and Gamow-Teller transitions preserve parity (π). Note that here we have focused entirely on the low energy limit. For $m_\chi \gg r_N^{-1}$ all angular momentum transitions are accessible. A summary of all possible Fermi and Gamow-Teller transitions for the experimental target materials considered here is presented in Fig. 6.5, where we also show the kinematic threshold for β^- decays to occur given each excited state mass $M_{Z_j+1, A_j}^{(*)}$ as dictated by Eq. (6.44) (the smallest thresholds possible for Fermi and Gamow-Teller transitions is also quoted in Table. 6.2)

The parton spin averaged nucleon level matrix element arising from the charged current operator Eq. (6.43) for induced β^- decay is given by:

$$\begin{aligned} |\overline{\mathcal{M}}|^2 = & \frac{4m_n m_\chi}{\Lambda^4} \left[E_e (2m_n - m_p + 2m_\chi - E_e) - m_e^2 \right. \\ & \left. + 2\lambda (E_e^2 - m_e^2) + \lambda^2 (E_e (2m_n + m_p + 2m_\chi - E_e) - m_e^2) \right], \end{aligned} \quad (6.53)$$

Note that Eq. (6.53) contains both vector ($\propto \lambda^0$), axial vector ($\propto \lambda^2$), and interference terms ($\propto \lambda$). All terms contribute to $\Delta I = 0$ and $\Delta\pi = 0$ transitions, while the axial term additionally allows for Gamow-Teller transitions $\Delta I = \pm 1$ and $\Delta\pi = 0$. Therefore, the daughter nucleus is often necessarily be formed in an excited state — one that obeyed the Fermi or Gamow-Teller spin angular momentum selection rule.

With Eq. (6.53) we can compute thermally averaged cross section from Eqs. (6.49) - (6.51) for induced β^- decays, which we find to be:

$$\langle \sigma v \rangle_j = \frac{|\vec{p}_e|_j}{16\pi m_\chi M_{A_j, Z_j}^2} |\overline{\mathcal{M}}_{N_j}|^2, \quad (6.54)$$

where as the rate is independent of solid angle to leading order we able to trivially carry out the angular integral. Here, $|\vec{p}_e|_j^2 = (m_{\text{th}, j}^\beta - m_\chi)(m_{\text{th}, j}^\beta - m_\chi - 2m_e)$ is the electron's outgoing 3-momentum in the center of mass frame (which is approximately the lab frame), in the limit that $m_e, m_\chi, m_{\text{th}, j}^\beta \ll M_{A_j, Z_j}$. The j index refers to a specific transition (Fermi or Gamow-Teller) of a given isotope of an experimental target material. The total event rate for induced β^- may now be computed by plugging in Eq. (6.54) into Eq. (6.47) and summing over the contributions from each isotope and each possible transition a given isotope could undergo under the angular momentum selection rules:

$$R = \frac{\rho_\chi}{2m_\chi} \sum_j N_{T,j} n_j \frac{|\vec{p}_e|_j}{16\pi m_\chi M_{A_j, Z_j}^2} \mathcal{F}(Z+1, E_e) |\overline{\mathcal{M}}|^2. \quad (6.55)$$

In contrast to the neutral current signal rate there are no implicit experimental cuts imposed on the induced β^- event rate at this level *i.e.* we assume all events are above the experimental energy threshold (with the notable exception being Super-Kamiokande which has a higher energy detection threshold of 3.5 MeV [378]).

We now comment on the possible signals due to a charged current event. From the kinematics in Eq. (6.46) we see that there is significant velocity-independent nuclear recoil energy, analogous to the situation in neutral current processes. One can then look for correlated peaks between different isotopes in the target material as discussed in Sec. 6.4. Even more striking, charged current processes result in an emitted energetic $e^\pm$ which could be observed at experiments. The emitted energetic electron can shower in the detector, and one may then search for this electron in parallel with the nuclear recoil. Additionally, the daughter nucleus will typically be produced in an excited state (as determined by angular momentum selection rules) which will then decay (typically with a known lifetime). This secondary decay will emit a photon which may be searched for at experiments sensitive to photon emission. Furthermore, the produced nucleus itself may be unstable and decay on time scales of interest to the experiment. In summary the possible signals are summarized as follows:

- Emitted high energy electron
- Recoiling nucleus at recoil energies peaked around a single bin (analogous to the neutral current case discussed above).
- Photon from nucleus being produced in an excited state
- Decay of unstable nucleus

The plethora of correlated signals makes it possible to trigger on several different signals. The optimal search strategy will depend on the specific target being considered, experimental capabilities, and the dark matter mass.

As an example consider β^- induced decays within a Xenon based detector. Stable (or meta-stable) isotopes can undergo the charged current absorption process:

$$\chi + {}^A_{54}\text{Xe} \rightarrow e^- + {}^A_{55}\text{Cs}^{(*)}, \quad (6.56)$$

where $A = \{126 (28.4\%), 131 (21.2\%), 134 (10.4\%), 136 (8.8\%)\}$ are the dominant isotopes given in their natural abundance (note that neutrino-less double beta decay experiments use enriched ${}^{136}_{54}\text{Xe}$). Depending on the Xenon isotope and the transition that takes place, the resultant Cesium nucleus may or may not be in an excited state. As an explicit example, consider ${}^{131}_{54}\text{Xe}$. The ground state of this isotope is a $I^P = \frac{3}{2}^+$ state while that of ${}^{131}_{55}\text{Cs}$ is in a $\frac{5}{2}^+$ state. The Fermi $\Delta I_F = 0$ transition $\frac{3}{2}^+ \rightarrow \frac{3}{2}^+$ produces an excited state of Cesium 60 keV above the Xenon ground state. Adding the electron mass results in a threshold of $m_{\text{th}}^\beta = 571$ keV for the process to occur.

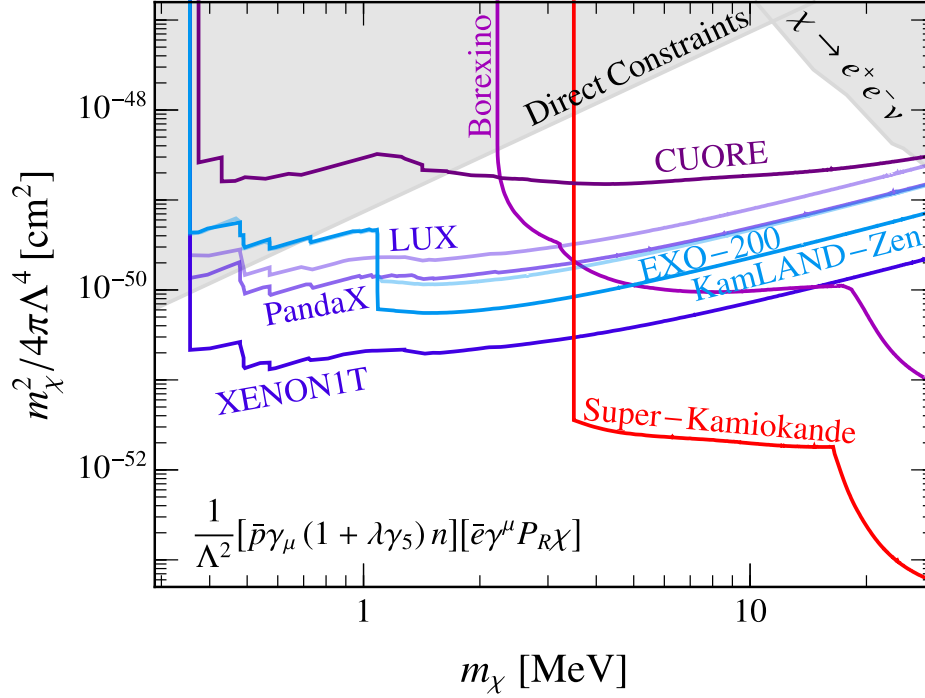


Figure 6.6: The projected constraints from a dedicated search for induced β^- signals at XENON1T, LUX, Panda-XII, EXO, and KamLAND-Zen for Xenon conversion. Super-Kamiokande and Borexino for absorption by Hydrogen, and by TeO_2 crystals in CUORE. The various experiments, their exposures and physics goals are summarized in Section 6.3.

For an axial vector coupling the Gamow-Teller, $\Delta I_{GT} = 0, \pm 1$ transitions are also possible and one must take into account the contribution of all such possible transitions in the amplitude. In particular, the $\frac{3}{2}^+ \rightarrow \frac{5}{2}^+$ transition from the ground state of $^{131}_{54}\text{Xe}$ to the ground state of $^{131}_{54}\text{Cs}$ contributes with threshold 360 keV, additionally transitions to the third excited state of $^{131}_{54}\text{Cs}$ also occur at relatively low threshold 490 keV. For setting limits, we sum over all relevant contributions from the various allowed transitions. Fig. 6.5 shows the values of the various $\Delta I = 0$ and Gamow-Teller only $\Delta I = \pm 1$ transitions for a given isotope atomic number that correspond to relatively low thresholds and significant abundances.

To project the sensitivity of current experiments to charged current signal we require at least 10 events with the results shown in Fig. 6.6 for various experiments summarized in Section 6.3, where note that again (with the exception of Super-Kamiokande) these projections assume no experimental cuts which is motivated due to the large number of correlated possible signals; nuclear recoil, energetic e^- (searches for high energy electrons could in principle even be done using an existing analysis, for instance the S2 XENON1T data set [411], but we leave a detailed analysis to future work), emitted γ , and decay of an unstable daughter nucleus. For some isotopes, the excited states have not yet been fully mapped out.

In general, however for heavier elements an excited state with matching angular momentum for each transition should exist within $\sim \text{MeV}$ of the ground state. For practical purposes when the data on the excited state corresponding to the $\Delta I = 0$ transition is not available we take the splitting to be 1 MeV.

As with the neutral current case, limits are sensitive to the different isotopes in a given experiment, where m_{th}^β for all the possible transitions in the experiments considered here are summarized in Fig. 6.5. In particular, the discontinuities in the limits of Fig. 6.6 occur at $m_\chi \sim m_{\text{th}}^\beta$, below which transitions become inaccessible for a given isotope. Projected scale linearly with exposure of a given experiment as is expected from (6.47). In particular, consider the projections for searches for induced β^- signals at XENON1T, LUX, Panda-XII, EXO, and KamLAND-Zen for Xenon conversion. Note that KamLAND-Zen and EXO are dedicated neutrino-less double beta decay experiments which utilize enriched $^{136}_{54}\text{Xe}$ target material, while Xenon based dark matter detectors contain Xenon isotopes in their natural abundance within. As expected the projected limits scale with exposure with XENON1T being more sensitive than LUX or Panda-XII. At high energies EXO and KamLAND-Zen do better than LUX and Panda-XII due to their larger exposures. However once energies fall below the threshold for induced beta decays by absorption off $^{136}_{54}\text{Xe}$ at about an 1 MeV, absorption occurs through the significantly sub-dominate isotopes, weakening the projected limits. CUORE, an experiment employing TeO_2 crystals looking for neutrino-less double beta decay, is also shown and could search for absorption of dark matter off both the Tellurium and Oxygen nuclei. Note that Super-Kamiokande and Borexino do particularly well simply due to their enormous Hydrogen detector volume and in particular enjoy an enhancement at large energies due to the Fermi function relative to detectors with larger Z target materials. The capabilities of Super-Kamiokande become more pronounced for $m_\chi \gtrsim 16 \text{ MeV}$ when it becomes kinematically possible to induce β^- decays off $^{16}_8\text{O}$ (which comprises about 99% of the oxygen in its natural abundance).

In Sec. 6.2 we discussed constraints from colliders as well as indirect detection constraints due to dark matter stability for the case of the charged current UV model at hand. Regions excluded due to these constraints are shaded out, and as expected stability considerations favor dark matter masses lighter than about 10 MeV and collider constraints favor mediator scales heavier than about a TeV.

Induced β^+ Decay

We now briefly consider signals from induced β^+ decays. In this process, an incoming dark matter converts a proton into a neutron and a positron. For heavy isotopes whose nuclei contain a significantly larger fraction of neutrons over protons, Pauli blocking of the outgoing neutron disfavors the production of a neutron in the lowest-lying energy states, thereby resulting in larger thresholds. Additionally, while such transitions have been studied in the context of supernova neutrinos, a detailed analysis of the favored transitions is not available in the literature and we leave this to future work. For the present, we simply focus on the case of induced β^+ transitions in experiments employing a target material consisting

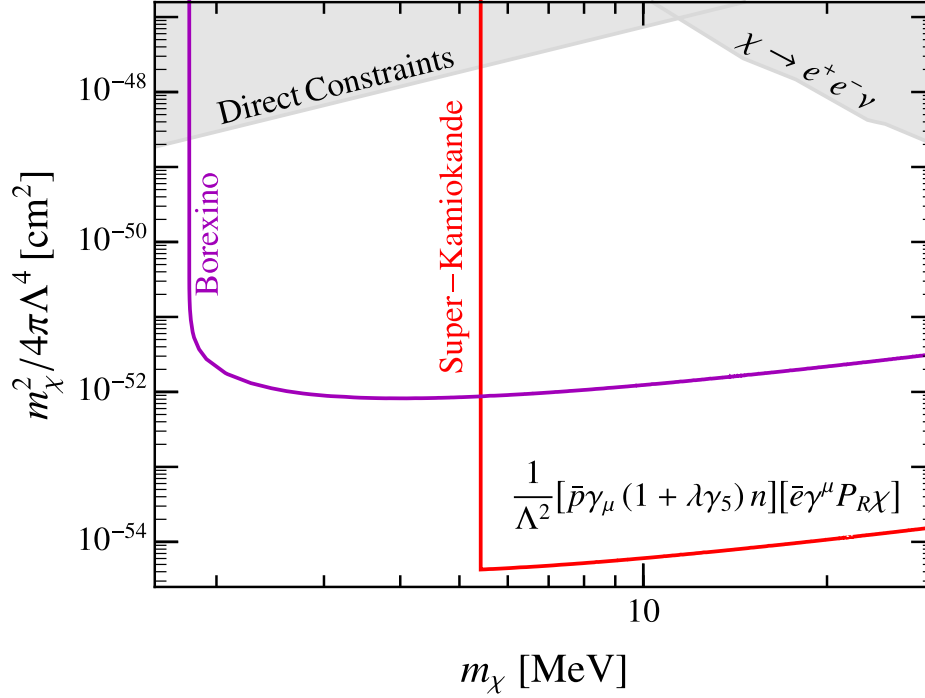


Figure 6.7: The projected constraints from a dedicated search for induced β^+ signals from Hydrogen at Super K and Borexino.

of Hydrogen, and consider the process;

$$\chi + {}^1_1\text{H} \rightarrow n + e^+, \quad (6.57)$$

with threshold of 1.8 MeV.

The kinematics, matrix element and rate can be calculated from the discussion above. In Fig. 6.6 we show projected limits for induced β^+ decay in the liquid H_2O of Super-Kamiokande, and the $\text{C}_6\text{H}_3(\text{CH}_3)_3$ target material of Borexino. In this way, two large volume experiments can now probe low dark matter masses — down to the kinematic threshold of 1.9 MeV in Borexino (which has a detection threshold of 70 keV), and the 5.3 MeV in Super-Kamiokande (which has a detection threshold of 3.5 MeV). Projected limits for these to experiments are shown in Fig. 6.7, once again, to set the limit, we assume that a given experiment can resolve and measure the energy of the produced $e^\pm$. Relative to β^- , β^+ processes tend to have larger thresholds leading to complementarity between the two types of searches. Lastly we note that asymmetric dark matter models may result in only β^- or β^+ , further motivating carrying out both types of searches.

6.6 Charged Current: β Endpoint Shifts

Having already considered β transitions induced by charged current operators, we now consider the possible signals in isotopes which β decay without the presence of dark matter. In these isotopes, light dark matter absorption with an unstable parent nucleus causes a shift in the kinematic endpoint of the β spectrum. Since the decay is allowed in vacuum, this is a threshold-less process and can occur for arbitrarily light dark matter. In particular, kinematic endpoint shifts can probe $m_\chi \lesssim \text{MeV}$, where induced β transitions are kinematically forbidden by Eq. (6.44). Isotopes which β^+ decay have smaller scattering rates than those which β^- decay due to the Fermi function, while those which electron-capture decay still have kinematic thresholds (albeit smaller ones). Thus, we ignore targets which naturally β^+ or electron-capture decay and just focus on those which β^- decay⁹.

Detecting rare β^- spectra endpoint shifts is difficult due to the lack of experiments with large exposures of unstable targets. There are only a few well motivated physics goals which employ β^- decaying isotopes: measuring the Standard Model neutrino masses [414–416], detecting the Cosmic Neutrino Background (C ν B) [417], and producing light sterile neutrinos [346, 418]. In this section we consider the possibility of using one of these existing or proposed experiments to look for fermionic absorption. Needless to say, the experimental requirements for detecting dark matter in this way are often quite relaxed relative to those needed for other experimental physics goals. For example, C ν B detection experiments must resolve β^- energies at the possible scale of Standard Model neutrinos’ masses, $m_\nu \sim 0.1 \text{ eV}$, thus requiring a target with a low Q value: tritium. By contrast, the fermionic dark matter we consider is heavier: $m_\chi \gtrsim 190 \text{ eV}$ [419, 420], so this target selection criterion is irrelevant. This motivates us to consider different β^- decaying isotopes more generally than the present proposals.

Toward this end, we categorize every isotope which dominantly β^- decays and has a half-life between 10^8 and 10^{13} seconds in Table 6.3¹⁰. In addition to dominantly decaying via Fermi or Gamow-Teller transitions, some of these isotopes predominantly undergo first forbidden transitions ($\Delta\pi = 1$, $\Delta I = 0, \pm 1, \pm 2$) or second forbidden transitions ($\Delta\pi = 0$, $\Delta I = \pm 2, \pm 3$), where $\Delta\pi$ is the parity change and ΔI the spin change. The Fermi and Gamow-Teller transition isotopes in Table 6.3 are of particular interest to us since they are the most long lived β^- -decaying isotopes for which the dark matter capture rates would not be momentum suppressed. We systematically checked the most recent β^- -decay experiments to verify that none had significant exposures of these interesting targets¹¹. This is expected since the exposures required to determine β^- spectra and half-lives are generally small.

⁹One could also consider re-purposing sterile-neutrino search experiments which used β^+ [412] and electron-capture [413] decaying isotopes. Unfortunately, the smallest mixing angle they constrain is 4×10^{-3} [412] which corresponds to an already-constrained Λ in the charged current operator in Eq. (6.43) from direct searches.

¹⁰We do not consider shorter half-lives since tritium’s half-life is roughly 4×10^8 seconds and would therefore be a better target than any shorter lived isotope. Isotopes with longer half-lives than 10^{13} seconds decay via even higher order forbidden transitions.

¹¹References for each isotope are ^{14}C [422], ^{32}Si [423], ^{60}Co [424], ^{63}Ni [425], ^{154}Eu [426], and ^{228}Ra [427].

Fermi	${}^3\text{H}$
Gamow-Teller	${}^{14}\text{C}$ ${}^{32}\text{Si}$ ${}^{60}\text{Co}$ ${}^{63}\text{Ni}$ ${}^{154}\text{Eu}$ ${}^{228}\text{Ra}$
1st forbidden	${}^{39}\text{Ar}$ ${}^{42}\text{Ar}$ ${}^{79}\text{Se}$ ${}^{85}\text{Kr}$ ${}^{90}\text{Sr}$ ${}^{137}\text{Cs}$ ${}^{151}\text{Sm}$ ${}^{194}\text{Os}$ ${}^{204}\text{Tl}$ ${}^{210}\text{Pb}$ ${}^{227}\text{Ac}$ ${}^{241}\text{Pu}$
2nd forbidden	${}^{36}\text{Cl}$ ${}^{94}\text{Nb}$ ${}^{99}\text{Tc}$ ${}^{126}\text{Sn}$

Table 6.3: Every isotope which dominantly β^- decays and has a half-life between 10^8 and 10^{13} seconds, categorized by their main transition [421]. Only Fermi and Gamow-Teller transitions are not momentum-suppressed for our charged current operators and therefore of interest to us.

The largest experimental proposals with targets which can undergo a Fermi or Gamow-Teller transition are made of tritium. The proposal with the largest tritium exposure is PTOLEMY which hopes to be the first experiment to measure the $\text{C}\nu\text{B}$ [417]¹². There has also been recent interest in using a comparable exposure of tritium in an experiment to measure coherent neutrino-atom scattering [428]. Regardless, we will do a proposal-independent analysis below when projecting sensitivities to the charged current signal which only depends on the exposure of tritium.

Tritium β^- decays to Helium via

$${}^3\text{H} \rightarrow {}^3\text{He} + e^- + \bar{\nu}_e, \quad (6.58)$$

where ${}^3\text{H}$ and ${}^3\text{He}$ refer to the nuclei (and not the atoms) of those isotopes. The Q -value for this decay is

$$Q = m_{{}^3\text{H}} - m_{{}^3\text{He}} - m_e - m_\nu \simeq 18.6 \text{ keV}, \quad (6.59)$$

where $m_{{}^3\text{H}} \simeq 2808.921 \text{ MeV}$ and $m_{{}^3\text{He}} \simeq 2808.391 \text{ MeV}$ are the nuclear masses. The charged current operators in Eq. (6.43) allow tritium to capture incoming dark matter via

$$\chi + {}^3\text{H} \rightarrow {}^3\text{He} + e^-. \quad (6.60)$$

Since the tritium nuclear transition is from $1/2^+ \rightarrow 1/2^+$, both the vector and axial-vector operators contribute to the fermionic absorption rate.

¹²While KATRIN [414] and Project 8 [416] also use tritium, they need far less than PTOLEMY's proposed amount for their measurements of the electron antineutrino's mass and will not probe charged current operator parameter space which is not already ruled out by LHC constraints.

Using Eqs. (6.47) and (6.54) from above, we find the rate for tritium to absorb fermionic dark matter is

$$R = \frac{\rho_\chi}{m_\chi} N_{3\text{H}} \frac{|\vec{p}_e|}{4\pi m_{3\text{H}} \Lambda^4} F(Z+1, E_e) \left[E_e \left(2m_{3\text{H}} - m_{3\text{He}} + 2m_\chi - E_e \right) - m_e^2 \right. \\ \left. + 2\lambda (E_e^2 - m_e^2) + \lambda^2 \left(E_e \left(2m_{3\text{H}} + m_{3\text{He}} + 2m_\chi - E_e \right) - m_e^2 \right) \right]. \quad (6.61)$$

Unlike in heavy elements where corrections due to the Fermi-Dirac distribution of nucleons in a nucleus are difficult to compute, for tritium, these have been well studied. To account for such corrections, we map $\lambda \rightarrow \sqrt{2.788/3}\lambda$ [429]. Then, in the light dark matter limit, we reproduce the standard neutrino capture cross section on tritium [347].

To project the sensitivity of future tritium-based experiments, we again require the number of absorption events to be less than 10. To connect the projected sensitivities of this low- m_χ region to those already considered above from the induced β signals, we show projected bounds on $\sigma \equiv m_\chi^2 / (4\pi \Lambda^4)$ in Fig. 6.8 for a tritium exposures of 100 g yr. PTOLEMY [417] expects to have an exposure of at least 100 g yr. We also show the direct searches bound [380] on the UV completion from Sec. 6.2. It is interesting that with less than 1 kg yr of tritium, a future experiment could start probing the lightest possible fermionic dark matter [419, 420] with these charged current interactions. This provides further motivation to pursue proposals such as PTOLEMY.

6.7 Discussion

In this Chapter, we comprehensively consider signals from the absorption of fermionic dark matter by nuclear targets at direct detection and neutrino experiments. These signals arise from a set of dimension-6 operators which do not conserve dark matter number and can be broadly classified into “neutral current” and “charged current” varieties. We present simple UV completions which lead to these operators and consider bounds from indirect searches for dark matter decays, as well as bounds coming from searches at collider experiments.

The neutral current operators induce dark matter velocity-independent nuclear recoils at distinct energies with relative spacing and peaks which result in a distinguishable signal. We present the general expressions for the rates as well as study the kinematics. We find that future (lower threshold) dark matter experiments employing lighter targets can achieve sensitivity to $m_\chi \lesssim \text{MeV}$ while remaining consistent with bounds from collider searches and indirect detection. Due to decay rates scaling with large powers of m_χ , above an MeV, the bounds from indirect detection become stringent. However, these constraints depend on the UV completion, while our projected sensitivities do not, so indirect detection constraints can in principle be fine-tuned away. Regardless, current dark matter and neutrino experiments can similarly probe a large unexplored parameter space.

In the presence of a dark matter background, the charged current operators can induce β decays in otherwise stable isotopes. This yields multiple possible signals: the ejected

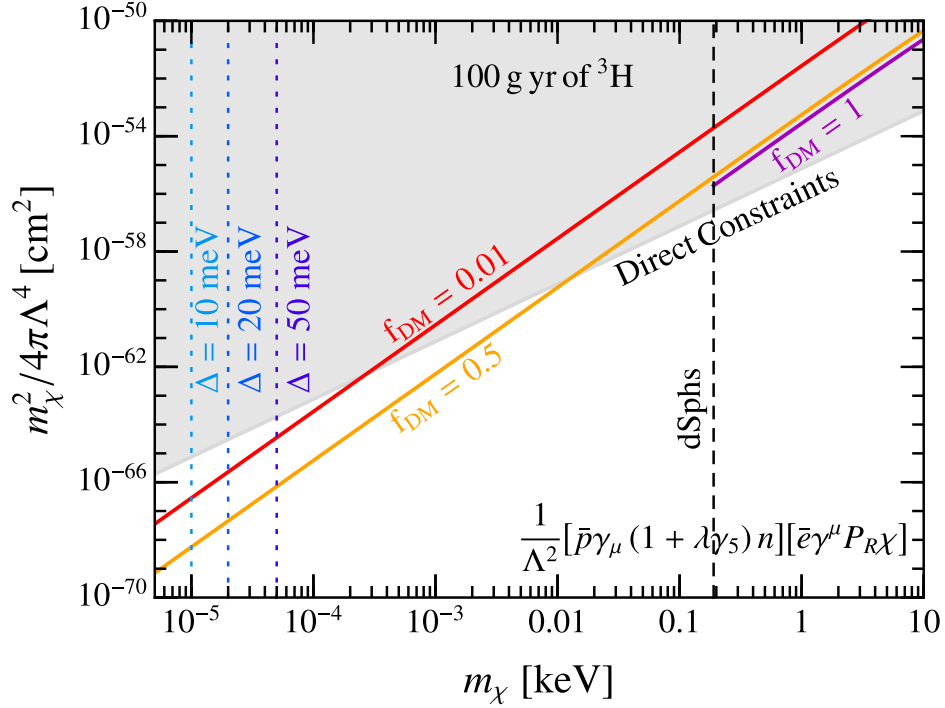


Figure 6.8: Bounds from dark matter capture events on tritium inside an experiment with a 100 g yr exposure assuming the dark matter captured is a (sub)component of dark matter which makes up 100%, 50%, and 1% of dark matter’s energy density (labeled $f_{\text{DM}} = 1$, $f_{\text{DM}} = 0.5$, and $f_{\text{DM}} = 0.01$ respectively). For reference, PTOLEMY has a proposed exposure of at least 100 g yr [417]. The Δ ’s show the lowest detectable dark matter mass for an experiment with that corresponding energy resolution. Also shown is the LHC bound [380] on our UV completion and the lightest possible fermionic dark matter $m_\chi \sim 190$ eV consistent with dwarf spheroidal galaxies assuming $f_{\text{DM}} = 1$ [419, 420].

energetic $e^\pm$, the nuclear recoil of the daughter nucleus, a prompt γ from the decay of the excited daughter nucleus, and further decay if its unstable. These correlated signals (unique for every isotope in an experiment) could be searched for simultaneously to reduce possible backgrounds. While one may consider both β^- and β^+ decays for any element, large suppressions in β^+ rates in heavier isotopes makes β^- the more promising candidate for every element, other than Hydrogen. We project sensitivity for a variety of current dark matter and neutrino experiments, finding powerful sensitivities, easily surpassing current direct and indirect constraints for $300 \text{ keV} \lesssim m_\chi \lesssim 30 \text{ MeV}$. In addition, we make projections for induced β^+ decays in Hydrogen for Borexino and Super-Kamiokande. Due to their sheer size, we find these experiments can probe deep into unexplored parameter space, having the largest potential impact for heavier masses.

While induced β decays are prominent signals that can be seen in almost any dark matter

or neutrino experiment, they inevitably require dark matter that is sufficiently heavy to induce such a transition putting the rough lower bound on the sensitivities of $m_\chi \gtrsim 500$ keV. To probe lower masses one can instead look for shifts in the kinematic endpoint of β spectra in isotopes that are already unstable in a vacuum. Due to the lack of existing or future experiments with such targets, we focused on the projected sensitivity of a tritium-based experiment, such as PTOLEMY. We find that such experiments could probe the lightest possible fermionic dark matter consistent with phase space packing bounds which interacts with the Standard Model through these charged current operators.

There is a host of current experiments which could discover dark matter from dedicated analyses for signals from the absorption of fermionic dark matter on nuclear targets. As such, different experiments could probe complementary unexplored regions of parameter space. The possibility of dark matter which interacts with the Standard Model through either neutral or charged current operators further motivates many proposed future experiments which have other concrete physics goals. As the quest for dark matter leads us away from the WIMP paradigm and into the ocean of light dark matter scenarios, fermionic absorption represents an exciting new class of signals that could, in the near future, discover the nature of dark matter.

Chapter 7

Summary

Some of the greatest problems in physics have lingered for decades: the electroweak hierarchy, the origin of the baryon asymmetry, and the identity of dark matter. In this dissertation, we have focused on the last of these and proposed mutual solutions which address the first two as well. Each of our dark matter candidates also has large self interactions which could address CDM small-scale structure issues. We have expounded sub-GeV dark matter models with $3 \rightarrow 2$ dark matter relic-setting processes and detectable ALP mediators; frozen-in twin (anti)electron dark matter in a natural, minimal, and cosmologically viable MTH framework; and asymmetric dark matter whose origin is unified with the baryon asymmetry's via a simple and predictive dark-sector baryogenesis. We have also seen that the stubborn persistence of the dark matter conundrum is partially due to null search results at large, sensitive detectors which are often optimized for specific dark matter candidates such as WIMPs. Taking this as an opportunity, the latter half of this dissertation has focused on novel dark matter models and signals that could be discovered in these current experiments, but may have been overlooked. Specifically, we have considered signals from the absorption of sub-GeV fermionic dark matter which include unique nuclear recoils, induced β decays, and β endpoint shifts.

Not knowing has always been a driving force of discovery. We find ourselves in a uniquely prescient age of modern science which is full of not knowing and must therefore precede great discovery; what a time for physics! With this optimistic view, we await Nature's next grand reveal while dreaming of all the things dark matter could be and how to find it.

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