Discovering Asymmetric Dark Matter with Anti-Neutrinos

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Abstract

We discuss possible signatures of Asymmetric Dark Matter (ADM) through dark matter decays to neutrinos. We specifically focus on scenarios in which the Standard Model (SM) baryon asymmetry is transferred to the dark sector (DS) through higher dimensional operators in chemical equilibrium. In such cases, the dark matter (DM) carries lepton and/or baryon number, and we point out that for a wide range of quantum number assignments, by far the strongest constraints on dark matter decays come from decays to neutrinos through the “neutrino portal” operator $H L$. Together with the facts that ADM favors lighter DM masses $\sim$ a few GeV and that the decays would lead only to anti-neutrinos and no neutrinos (or vice versa), the detection of such decays at neutrino telescopes would provide compelling evidence for ADM. We discuss current and future bounds on models where the DM decays to neutrinos through operators of dimension $\leq 6$. For dimension 6 operators, the scale suppressing the decay is bounded to be $\gtrsim 10^{12} - 10^{13}$ GeV.

1 Introduction

It is by now well established that about a quarter of the mass of the universe is in the form of some kind of non-luminous, non-baryonic matter. The standard cosmic history leading to thermal relic dark matter (DM), where DM annihilations freeze out at temperatures near its mass, is certainly an attractive possibility, but it is not the only one. Indeed, the observed baryon and lepton number densities did not arise due to thermal freeze-out of annihilations; they are due to a primordial asymmetry between the number densities of particles and anti-particles. A scenario of this type, “Asymmetric Dark Matter” (ADM), is also perfectly viable to explain the origin of the relic density of dark matter particles.

One compelling idea is to have the dark matter asymmetry arise due to a direct connection to the asymmetry in the baryons and leptons (see e.g. [1, 2, 3, 4, 5, 6, 7, 8, 9, 10, 11]). In particular, one can imagine that the dark matter particle itself carries baryon number and/or lepton number. In that case the dark matter will share in the asymmetry of regular matter in the early universe. This general scenario was proposed in [3], and it is the scenario that we will be assuming throughout this paper. One of its most appealing qualities is that it leads generically to a prediction for the mass of the dark matter particle; since the asymmetry in the dark sector ends up being roughly comparable to that in the baryons, the ratio of
dark matter to baryonic energy densities is given approximately by the ratio of the dark matter mass $m_{\text{DM}}$ to the mass of the proton. This leads to a prediction for $m_{\text{DM}}$ of about $\mathcal{O}(1 - 10) \text{ GeV}$ depending on the details of the model. Here we are simply assuming that all interactions which transfer the asymmetry between the dark and visible sectors have frozen out of equilibrium by the time the temperature has dropped below $m_{\text{DM}}$. It is an amusing coincidence that the prediction here for $m_{\text{DM}}$ is fairly close to the weak scale, and thus could arise in models of electroweak symmetry breaking, much as in the case of the standard weakly interacting massive particle (WIMP).

The fundamental test of the WIMP hypothesis will eventually be an issue of measuring the couplings of a given WIMP candidate, in order to check whether the associated annihilation cross section would indeed lead to an appropriate relic density. Such a confirmation would yield convincing evidence that the dark matter relic density did indeed have a thermal origin. Similarly, obtaining convincing evidence that the ADM mechanism is at play may also be possible. There would be two crucial ingredients which would be important to confirm; first, one would like to check that the dark matter particles do indeed carry a particular baryon number and/or lepton number, and second, one would like to check that $m_{\text{DM}}$ is of an appropriate size.

A difficulty here is that, since we are requiring that the operators responsible for carrying the baryon/lepton asymmetry between the dark and visible sectors froze out at temperatures above $m_{\text{DM}}$, these operators must be fairly suppressed. On the other hand, it is precisely these operators which would reveal the baryon/lepton number carrying properties of the dark sector; the requirement of early freeze-out implies that signatures of these operators will in general be difficult to detect.

In this paper, we will discuss one reasonably generic signature which the asymmetry transferring operators might have. If the dark matter carries baryon/lepton number, then decays or annihilations of this particle into the standard model leptons and/or baryons may in general be made possible. Since electric charge is conserved and the DM particle must be electrically neutral, the excess standard model baryon/lepton number will necessarily show up either in the neutrino sector, or in equal numbers of electrons and protons (or their anti-particles). Due to the requirements on $m_{\text{DM}}$ in this scenario, neutrinos from DM decays would typically have energies of order a few GeV, and could lead to a distinctive bump in cosmic neutrino data in this energy range. Moreover, due to the origin of the signal, this bump would be associated with anti-neutrinos rather than neutrinos (or vice versa), and this is a property which is potentially discernable at the Super Kamiokande detector [12], as well as MINOS [13] and possible future detectors. This feature would make an observation of this type particularly compelling evidence in favor of the ADM mechanism.

Unfortunately, in the case of operators leading to dark matter annihilations, early freeze-out requirements constrain the event rates to be considerably below the reach of both current

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1If the asymmetry transfer freezes out below the DM mass, the densities of DM and baryons are not related to each other in such a straightforward way, and depend on the size of the DM annihilation cross-section, much as in the standard thermal WIMP case.
and upcoming neutrino data. We relegate the details to appendix B. As a result, we shall instead focus on operators leading to dark matter decays.

As implied above, our most important limits will come from the Super Kamiokande (Super-K) neutrino detector. The MINOS detector has about 20% the fiducial volume of Super-K, though this is partly compensated for by its ability to directly distinguish between neutrinos and anti-neutrinos, thereby cutting down on the atmospheric neutrino background in the ADM scenario. Although the IceCube [14] experiment now typically sets more stringent limits on neutrino signals than Super-K, its low energy threshold of \( \sim 100 \text{ GeV} \) is too high for our purposes. We find that operators of dimension 6 leading to DM decays must be suppressed by at least \( \sim 10^{12} - 10^{13} \text{ GeV} \), due to existing neutrino constraints. In many models of the early universe, the reheating temperature is required to be less than about the GUT scale, and thus, upcoming improvements in the constraints will be probing into the remaining window where the operator responsible for decay is also capable of transferring the asymmetry between the SM and the DS.

The outline of this paper is as follows. In section 2, we review the asymmetric dark matter (ADM) paradigm, and how one obtains quantitative predictions for the DM mass, focussing on a specific choice for the operator responsible for DM decay to neutrinos. In section 3, we discuss details of the production and observation of neutrinos from DM decays, and the resulting constraints on the decay spectrum. We also discuss possibilities for distinguishing neutrinos from anti-neutrinos at water-Cherenkov detectors and other possible future detectors. In section 4, we consider more general operators which may lead to the decay of the DM particle. For each operator, our analysis gives a bound on the size of the scale suppressing the interaction. In particular, we consider all possible interactions of dimension \( \leq 6 \) coupling the DM sector to a SM \((B - L)\)-carrying gauge-singlet operator. Finally, in section 5, we speculate on future directions and model-building issues.

2 Asymmetric Dark Matter

Asymmetric dark matter replaces the standard thermal history of dark matter with one more closely analogous to that of baryons. The relic density, rather than being fixed by the freeze-out of DM self-annihilation, is set by a small asymmetry between dark matter particles and anti-particles, with all anti-particles eventually annihilating. Various mechanisms for linking the dark matter asymmetry to the baryon asymmetry of the universe have been proposed in the literature [1, 2, 3, 4], usually involving equilibrium processes that transfer any particle asymmetry from the Standard Model (SM) to the dark sector (DS) and vice versa. We will focus on a specific class of models, where the asymmetry is generated somehow (the details of which will not concern us here) and transferred between the SM and DS by nonrenormalizable operators of the form

\[
\Delta \mathcal{L} = \frac{\mathcal{O}_{\text{DS}} \mathcal{O}_{\text{SM}}}{\Lambda^{d-4}},
\]

(1)
where $\mathcal{O}_{\text{SM}}, \mathcal{O}_{\text{DS}}$ are gauge-invariant operators composed solely of SM or DS fields respectively, and carrying equal and opposite nonzero baryon-minus-lepton ($B - L$) number.

For the most part, we will leave model-building details aside and work only with the operators (1) below the scale at which they are generated. The spectrum and ($B - L$)-numbers of the light DS states then determine which such operators can be generated. We will focus on the lowest-dimension ($d \leq 6$) operators of the form (1). It is of course not difficult for higher-dimension operators to be selected by appropriate ($B - L$) numbers, e.g. if ($B - L$) of the DM particle is $\ll 1$. Also, note that we will consider some instances where some mechanism suppresses the asymmetry-transfering operators beyond just powers of the scale where they are generated, and dimensional analysis is somewhat incomplete in this case. Still, classifying operators by dimension is useful anyway, since it makes it easy to read off the lowest-dimension allowed operators once the quantum numbers of the fields and any spurions are specified. In any case, the phenomenological signature of the decay is more sharply peaked in energy when there are fewer decay products, making the signature more distinctive; this in itself is a reason to focus on operators of relatively low dimension.

Part of what makes the mechanism we are considering interesting for a signature is that the gauge-invariant operator $HL$ is renormalizable and appears in the majority of low-dimension operators of the form (1). The operator $HL$ has been called the “neutrino portal” [15] since it can contract with DS operators to allow dark matter decays to neutrinos. Neutrinos from dark matter decays have the advantage over other decay products that they are not electrically charged, and thus their galactic propagation does not depend on difficult-to-determine astrophysics. While neutrinos are often produced in models as the result of dark matter decays, they are rarely the dominant decay mode or the leading discovery channel. The only exception of which we are aware is the strongly coupled model in [15], where the DS is a strongly coupled QCD analogue that can decay to neutrinos and dark glueball states. While this model is very interesting, the ones that we are considering differ in several aspects. First of all, there is the additional motivation for these models based on the ADM mechanism. Second, the DS in our case is perturbative and perhaps simpler for determining the predictions of specific models. Third, in the strongly coupled models, dark matter decays producing charged baryons are significant enough to be a competitive or leading signature compared to neutrinos in most of the parameter space depending on one’s assumptions about galactic propagation; essentially, the neutrino constraints are stronger partly because they have the advantage of not suffering as much from astrophysical uncertainties. We will see however that in a range of cases for ADM, charged lepton and baryon production is greatly suppressed. This follows from the fact that ADM favors the dark matter to be much lighter than in [15]. Finally, the assumption of ADM will lead to some additional changes in the predictions of the model. In particular, ADM predicts that only anti-neutrinos and not neutrinos (or in some models, vice versa) ought to be produced in dark matter decays.

\footnote{In fact, one could reasonably take the dark matter mass to be lighter in their model as well in order to suppress the non-neutrino decays. The reason [15] focused on $1 \text{ TeV} < m_{\text{DM}} < 5 \text{ TeV}$ was to get the correct thermal relic density, but this could likely be modified with some additional model-building.}
We are interested in the \((B - L)\) “portals” of the SM, i.e. gauge-invariant operators that carry \((B - L)\) number. For this section and the next, we will focus on a simple illustrative example - that of the operator \((LH)^2\), and an associated interaction \(O_{2\nu} \equiv \frac{1}{2}X(LH)^2\), where \(X\) is a complex scalar DM particle. Note that for our purposes, the flavor indices in this interaction are irrelevant, due to neutrino oscillations; all flavors will be equally represented in the signal. In later sections, we will consider more general operators for DM decay.

As usual, chemical equilibrium dictates relations among the chemical potentials of each particle species in the early universe (see e.g. [16, 17]), which we briefly review here. The DM asymmetry then follows from the DM chemical potential, and determines the DM particle mass. In general, the number density asymmetry in a given species is \(n - \bar{n} = 2gT^2\mu_6\) for bosons and \(gT^2\mu_6\) for fermions, where \(g\) counts the number of internal degrees of freedom (e.g., \(g = 1\) for a Weyl fermion). We will calculate the constraints slightly differently from usual, in a manner that emphasizes the dependence on the conserved charges of the model. This will allow us to discuss the implications for the dark matter mass in terms of assumptions concerning the symmetries of the dark sector couplings, rather than having to write down specific dark sector interactions and analyze the resulting constraints. This will be helpful later when we discuss more general models.

The point is that each interaction in chemical equilibrium imposes a linear relation on the chemical potentials, of the form \(\sum_i \mu_i = 0\), where the sum is over all the particles in the interaction. In the absence of any conserved quantities, there are at least as many interaction terms as there are fields with chemical potentials, and so the system is overdetermined and the chemical potentials vanish. However, for each \(U(1)\) conserved charge, it will also be true that \(\sum_i q_i = 0\), where \(q_i\) is the charge of the \(i\)-th particle in the interaction. Therefore, the choice of \(\mu_i = cq_i\) leads to a non-vanishing solution for the chemical potentials, where \(c\) is arbitrary. Since the constraint equations are linear, the general solution will be a sum over the solutions for each abelian symmetry.

In the Standard Model, the only two flavor-universal, linearly independent, anomaly-free \(U(1)\)s are hypercharge and \((B - L)\). Labeling the two corresponding coefficients \(c_Y\) and \(c_{B-L}\), we have

\[
\begin{align*}
\mu_H &= \frac{1}{2}c_Y, \\
\mu_l &= -\frac{1}{2}c_Y - c_{B-L}, \\
\mu_e &= -c_Y - c_{B-L}, \\
\mu_q &= \frac{1}{6}c_Y + \frac{1}{3}c_{B-L}, \\
\mu_u &= \frac{2}{3}c_Y + \frac{1}{3}c_{B-L}, \\
\mu_d &= -\frac{1}{3}c_Y + \frac{1}{3}c_{B-L}.
\end{align*}
\]

(2)

One may easily check that \(c_Y\) and \(c_{B-L}\) cancel out of the constraint imposed by any Standard Model interaction; for instance, \(QUC_H\) imposes \(\mu_q - \mu_u + \mu_H = 0\), which is automatically satisfied by (2). In general, there may be accidental symmetries that can introduce additional parameters. However, the accidental symmetries \(B\) and \(L\) separately are efficiently violated above the electroweak scale by sphaleron processes [18].

\footnote{For simplicity, we assume that the asymmetry is generated to be flavor-universal. Alternatively, if the DM couples to all lepton flavors, then these interactions will enforce flavor-universal asymmetries in the SM regardless of the flavor-independence of the original asymmetry. At any rate, this does not qualitatively affect our results.}
Additionally, the total hypercharge must vanish,

\[ N_f (\mu_q + 2\mu_u - \mu_d - \mu_l - \mu_e) + 2\mu_H = 11c_Y + 8c_{B-L} = 0, \quad (3) \]

leading to a single remaining free parameter required to fix all of the chemical potentials, which must be set by initial conditions. It is convenient to parameterize this in terms of the total baryon asymmetry \( B = N_f (2\mu_q + \mu_u + \mu_d), \) where \( n_B - \bar{n}_B \equiv \frac{T^2 B}{6}, \) leading to

\[ c_{B-L} = \frac{11}{28} B, \quad c_Y = -\frac{2}{7} B . \quad (4) \]

For \( O_{2\nu}, \) the conserved quantum numbers of the dark matter particle are completely determined by the interaction, and one obtains a definite prediction for the dark matter particle chemical potential and mass: \( \mu_X = 2c_{B-L} = \frac{11}{14} B \) and \( m_X = \frac{\Omega_{DM}}{\Omega_b} \frac{B}{2\mu_X} m_p \approx 3 \text{ GeV} \). This is representative of the typical size of mass favored by ADM. The decay spectrum into a particular neutrino flavor is in this case effectively \( \frac{dN}{dE} = \frac{2}{3} \delta(E_n - m_{DM}/2). \) This is independent of the relative branching ratios of decays to different flavors, due to neutrino oscillations.

\section{Neutrino Flux}

\subsection{Halo Flux}

The primary source of dark matter decays is from the galactic halo. The flux of neutrinos of a given flavor detected at angle \( \psi \) is

\[
\frac{d\Phi}{dE} = \frac{\Gamma_{DM}}{4\pi m_{DM}} \frac{dN}{dE} R_{sc} \rho_{sc} \Delta \Omega \mathcal{J}_{\Delta \Omega}(\cos \psi),
\]

where \( dN/dE \) is the spectrum of decays into a specific flavor, \( \Gamma_{DM} \) is the decay width to all neutrino flavors, and \( m_{DM} \) is the dark matter mass. Here we have defined the average flux integral \( \mathcal{J}_{\Delta \Omega} \) as follows. First, we take the conventional definition for the line-of-sight (LOS) flux integral \( \mathcal{J} \): if the direction the telescope is pointed in makes an angle \( \psi \) with the direction toward the galactic center (GC), then

\[
\mathcal{J}(\cos \psi) = \int_0^{l_{max}} \rho(\sqrt{R_{sc}^2 - 2xR_{sc} \cos \psi + x^2})dx,
\]

where the upper limit \( l_{max} = \sqrt{R_{mw}^2 - \sin^2 \psi R_{sc}^2 + R_{sc} \cos \psi} \) depends on the size of the Milky Way halo, which we take to be \( R_{mw} = 34 \text{kpc}, \) and \( R_{sc} = 8.5 \text{kpc}, \rho_{sc} = 0.3 \text{GeV/cm}^3 \) are the galactic radius and density at the solar circle in the NFW halo profile, respectively. For an effective detector resolution of \( \theta, \) we then take the average of \( \mathcal{J}(\cos \psi) \) over a cone of half-angle \( \theta \) (i.e. solid angle \( \Delta \Omega = 2\pi(1 - \cos \theta)) \):

\[
\mathcal{J}_{\Delta \Omega}(\cos \psi) = \frac{1}{\Delta \Omega} \int_0^\theta \sin \theta' d\theta' \int_0^{2\pi} d\phi J(\hat{r} \cdot \hat{r}_{GC}),
\]

\footnote{This prediction for the mass is modified by \( O(20\%) \) if the asymmetry-transferring processes freeze out below the electroweak scale \cite{3}.}
\[ J_{\theta=10^\circ}(1) = 12.6 \quad J_{\theta=30^\circ}(1) = 6.7 \quad J_{\theta=180^\circ}(1) = 1.8 \quad r_s (\text{kpc}) = 20 \quad \rho_0 (\text{GeV/cm}^3) = 0.06 \quad \alpha = 0.17 \quad \beta = n/a \quad \gamma = n/a \]

Table 1: \( J_{\Delta \Omega}(1) \) for various halo models and angular averages (\( \Delta \Omega = 2\pi(1 - \cos \theta) \)).

where \( \hat{r} \cdot \hat{r}_{GC} = \sin \psi \sin \theta' \cos \phi + \cos \psi \cos \theta' \). \( J_{\Delta \Omega} \) is at a maximum in the direction of the GC, at \( \psi = 0 \).

The size of the signal from dark matter decays to neutrinos depends somewhat on the galaxy halo profile. To give a sense of this dependence, we will compare the size of the signal for four different halos: the Einasto profile

\[ \rho_{\text{Ein}} = \rho_0 \exp \left[ -2 \left( \frac{r/r_s}{\alpha} - 1 \right) / \alpha \right] \]

and the NFW, Moore, and Kravstov profiles, which have the parameterization

\[ \rho = \rho_0 \left( \frac{r}{r_s} \right)^{-\gamma} (1 + (r/r_s)^\alpha)^{(\gamma-\beta)/\alpha}. \]

The values of the parameters for these models are given in Table 1. The largest variation in the halo profiles is nearest the galactic center.

For dark matter decay, unlike the case for dark matter annihilation, the size of the signal near the galactic center is not strongly dependent on the halo profile. In Table 1, we show \( J_{\Delta \Omega}(1) \) (i.e., pointed directed at the galactic center) for cones of half-angle 10°, 30°, and 180°. As one can see, the variation is fairly modest and less than a factor of 2 in even the most extreme example.

3.2 Cosmic Flux

The contribution to the neutrino signal from cosmic dark matter decays is small compared to the halo contribution, but for completeness, we will review this here (see e.g. [19] for more details). The cosmic flux signal is given by

\[ \frac{d\Phi}{dE} = \frac{\Gamma_{\text{DM}}}{4\pi m_{\text{DM}}} \frac{\Omega_{\text{DM}} \rho_c}{H_0} \int_0^\infty \frac{dN(E')}{dE'} dz \cdot \frac{dz}{h(z)}, \]

where \( \rho_c \) is the universe critical density, \( H_0 \) is the Hubble constant today, \( \Omega_{\text{DM}} = 0.22 \), \( E' \equiv (1 + z)E \) is the redshifted energy, and \( h(z) = [(1 + z)^3 \Omega_m + \Omega_{\Lambda}]^{1/2} \). For comparison with the halo flux, we may define an effective \( J_{4\pi} \) for the cosmic flux:

\[ J_{4\pi, \text{Eff}} = \frac{\Omega_{\text{DM}} \rho_c}{H_0} \frac{\int_0^\infty \frac{dz}{(1+z)h(z)}}{4\pi R_{sc} \rho_{sc}} = 0.16. \]

(The additional \((1+z)\) in the denominator is from the integration over energy.) This is negligible even compared to the full \( 4\pi \) average \( J \) in any of the halo models, and we will ignore it from now on.
3.3 Neutrino Constraint

To derive our constraints, we will follow the analysis of [19]. To be conservative, we demand that the predicted muon-neutrino flux at Super Kamiokande from dark matter decays be less than the background from atmospheric ($\nu_\mu + \bar{\nu}_\mu$) neutrinos. For two-body decays, neutrinos from the halo decays will be nearly monochromatic. Thus, following [19], we choose an energy bin of width $\Delta \log_{10} E = 0.3$, which is about the width of the bins Super-K uses in their own analyses [12], centered on $E = m_{DM}/2$, and we demand that the total predicted signal in this bin be less than the flux of atmospheric muon-neutrinos. Note that for multi-body decays, this would contain only some fraction of the total decay spectrum, making the bound on the decay rate correspondingly weaker. Since the uncertainty in the background is only $\sim 10$ percent, one expects a dedicated analysis of Super-K data to improve the $2\sigma$ bound on the decay width by roughly a factor of five; such an analysis is currently underway [20]. For the background rate, we use the flux calculated by Monte Carlo in [21], and take the average over the full sky. The resulting constraint is shown in Fig. 1. We have explicitly factored out the model-dependent number $n$ of neutrinos produced in the decay, and the angular integral factor $\mathcal{J}_{\Delta \Omega}$. The above bound translates into a limit on the scale $\Lambda$ appearing in $\mathcal{O}_{2\nu} = \frac{1}{2} \frac{X(HL)^2}{\Lambda^4}$ of $\Lambda > 6 \times 10^{13}$ GeV, taking for example $\mathcal{J}_{\theta=30^\circ}$. This operator freezes out at temperatures of order $T_{fr} \sim (\Lambda^4/M_{pl})^{1/3} \gtrsim 10^{12}$ GeV. Assuming that this operator is responsible for the asymmetry transfer and is still present in the effective theory at high temperatures $T_{rh}$ where reheating occurs, $\Lambda$ must be less than $(T_{rh}^3 M_{pl})^{1/4}$ in order for the asymmetry transfer to occur at all. If $T_{rh}$ is $M_{GUT}$ or less, which is true of most models, then the remaining window for $\Lambda$ is $6 \times 10^{13}$ GeV $\lesssim \Lambda \lesssim 5 \times 10^{16}$ GeV. Thus, future improvement in the neutrino constraint will probe interesting regions of the parameter space, based not only on the closeness of $\Lambda$ to the GUT scale but also due to expected limitations on the freeze-out temperature.

3.4 Distinguishing $\nu$s from $\bar{\nu}$s

While seeing a feature in the neutrino spectrum at $O(1 - 10$ GeV) might give suggestive evidence in favor of the ADM picture, a verification of the lepton number violating origin of the signal would be crucial to making this evidence truly compelling. As we will discuss further in section 4 the value of the lepton number carried by the dark matter particle is model dependant, and may be either positive or negative. With the relatively low dimension operators we consider, this implies that the neutrino signal will be composed entirely of either neutrinos or anti-neutrinos. An important question is then whether or not this feature is distinguishable in present detectors, or at future neutrino experiments.

If a signal were to be seen at MINOS despite its smaller fiducial volume than Super-K (4 kton vs. 20 kton), it would trivially be able to test the neutrino vs. anti-neutrino composition of the excess due its magnetic field (through the charge of the same-family lepton produced in interactions). In fact, since the detected background of atmospheric neutrinos is about 2 times larger than for anti-neutrinos, MINOS would have the advantage that it could cut down on its background relative to an ADM anti-neutrino signal. At water Cherenkov detectors such
Figure 1: Bound on total decay width to all flavors of neutrinos for a monochromatic spectrum, from demanding that the signal that would be observed at Super-K not be greater than the atmospheric muon-neutrino ($\nu_\mu + \bar{\nu}_\mu$) background. The bound depends on the solid angle $\Delta\Omega$ of observed sky around the galactic center, and on the number $n$ of neutrinos produced in each decay. ADM typically favors masses in the lower range of this plot.

as Super-K, there are several possible avenues for separating neutrinos from anti-neutrinos on a statistical basis, even though there is no direct way of measuring the signs of particle charges. For dark matter masses below about 5 GeV, the most promising one involves looking for neutrino interactions in the detector whose final states contain one extra muon, beyond the number expected from the flavor of the original neutrino. Such muons are produced as decay products of charged pions, and, though they are typically produced with little kinetic energy, they are easily identified by the appearance of a decay electron one muon lifetime after the initial event, i.e., a “muon-decay electron”. The dominant neutrino interactions (for, e.g., electron neutrinos) leading to charged pion production are

$$\nu_e + P \rightarrow P + e^- + \pi^+, \quad \nu_e + N \rightarrow N + e^- + \pi^+, \quad \bar{\nu}_e + P \rightarrow P + e^+ + \pi^-, \quad \bar{\nu}_e + N \rightarrow N + e^+ + \pi^-.$$ 

We thus see that neutrinos tend to lead to the production of positively charged $\pi$’s, while anti-neutrinos tend to lead to the production of negatively charged $\pi$’s. The key point now is that, before having a chance to decay, a $\pi^-$ traveling through the detector with around a GeV of energy tends to become absorbed by a proton. Coulomb repulsion of $\pi^+$’s, on the other hand, suppresses inelastic scattering of these particles, so that, after stopping in the
detector, they typically are successfully able to decay. For this reason, looking for events with an extra muon-decay electron in the final state at energies around a GeV tends to efficiently select about 10 times as many neutrinos as anti-neutrino events [20]. In general, since the total cross section for detection of neutrinos is about a factor 2 larger than for anti-neutrinos to begin with, the result is a net increased sensitivity to anti-neutrinos by about a factor of 5. Unfortunately, as the pion energy becomes much higher than a GeV (corresponding to DM masses much higher than roughly 5-10 GeV), both π+’s and π−’s tend to scatter elastically before decaying in the detector, and the method is no longer successful.

A few other methods may be used to distinguish neutrinos from anti-neutrinos at water Cherenkov detectors, but it is unclear if Super Kamiokande would have sufficient statistics to make use of them if they were to find a bump in the neutrino spectrum at the energies suggested by ADM. These might thus only be useful at future detectors with larger volumes, such as Hyper Kamiokande, or one at DUSEL at the Homestake mine [22, 23]. In [24], differences in angular distributions for final state leptons in neutrino vs anti-neutrino interactions were used to look for anti-neutrinos coming from the sun. In the ADM case, one could similarly make use of the increased flux that a signal would lead to coming from the galactic center. Finally, inelastic scattering of neutrinos via the interactions

\[
\begin{align*}
\nu_e + N &\rightarrow P + e^-, \\
\nu_\mu + N &\rightarrow P + \mu^-,
\end{align*}
\]

allows for very efficient selection of \( \nu \) over \( \bar{\nu} \) events, but only if one is able to identify the proton in the final state. This has been shown to be possible for protons with around a GeV of energy by making use of the short length of their Cherenkov tracks, but the required cuts are currently rather inefficient [25].

Finally, we note that, if INO (India Neutrino Observatory) were to be built, its magnetic field would allow for trivial separation of \( \nu \) from \( \bar{\nu} \) events [26]. Similarly a possible future liquid argon detector such as GLACIER [27] or one at DUSEL [28] would be able to efficiently identify protons and would also lead to efficient \( \nu/\bar{\nu} \) tagging.

4 Other Models

The low dimension (\( \leq 5 \)) (\( B-L \)) portals in the Standard Model are \( LH, (LH)^2, LLE^c, U^cU^cD^c \), and \( LQD^c \). The portal is opened, so to speak, by operators such as [1]. Denoting bosonic dark sector states by \( X \)’s and fermionic dark sector states by \( \psi \)’s, we may list all such operators from lowest dimension to highest dimension up to dimension \( d = 6 \) (suppressing flavor
indices)

\[
\begin{align*}
    d = 4 & : \quad \mathcal{O}_1 = \psi H L \\
    d = 5 & : \quad \mathcal{O}_2 = X \psi LH \\
    d = 6 & : \quad \mathcal{O}_3 = \psi L L E^c \\
    \mathcal{O}_4 & = \psi L Q D^c \\
    \mathcal{O}_5 & = \psi U^c D^c D^c \\
    \mathcal{O}_6 & = X_1 X_2 \psi LH \\
    \mathcal{O}_{2\nu} & = \frac{1}{2} X (L H)^2
\end{align*}
\] (12)

In order for \( \mathcal{O}_1 \) to lead to a decay of \( \psi \), an \( H \) must be produced in the interaction. Since the dark matter in the cases we consider will always be lighter than the Higgs, we may as well just integrate it out. Thus, for the sake of considering signatures, \( \mathcal{O}_1 \) is just a special case of \( \mathcal{O}_3, \mathcal{O}_4 \). At any rate, \( \mathcal{O}_1, \mathcal{O}_3, \mathcal{O}_4, \) and \( \mathcal{O}_5 \) are all models with \( (B - L)[\psi] = 1 \), which we will not consider further since neutrinos are not the dominant decay signature. However, there is nothing in principle wrong with \( (B - L)[\psi] = 1 \), provided the marginal operator \( \mathcal{O}_1 \) is sufficiently suppressed to evade particle physics constraints. \( \mathcal{O}_5 \) has been studied in the context of baryon number violation in SUSY theories (see [29] and references therein) and is predicted to result in production of a slow-moving (momentum \( \sim \) GeV) baryon from DM decays.

It might seem based on the ratio of \( \Omega_b/\Omega_{DM} \) that the mass is restricted to \( \sim 5 \) GeV. However, a simple example illustrates that \( m_{DM} = \mathcal{O}(1 - 50) \) GeV may also be obtained fairly easily. The reason is that for many of the operators in (12), the prediction for the mass depends on the model-dependent quantum numbers and spectrum of the dark sector particles. Our example is \( \mathcal{O}_2 \) with \( X \) being the DM particle, and supposing that there are no other conserved abelian symmetries under which \( X \) and \( \psi \) are charged. From our analysis in section 2 of chemical equilibrium, this leads to a prediction for the dark matter mass of \( m_X \approx 6.4(-l_X)^{-1} \) GeV, where \( L(X) = l_X, L(\psi) = -1 - l_X \).\footnote{Our convention for \( H \) is that \( v = 175 \) GeV.} Note that this argument applies generally to bosonic DM with lepton number \( l_X \); for fermion DM, the mass prediction is \( m_X \approx 13(-l_X)^{-1} \) GeV. Such DM masses can quite reasonably be between say \( \sim 1 \) and 50 GeV.

The spectrum of observed neutrinos will be somewhat model-dependent. For illustration, we may continue to consider \( \mathcal{O}_2 \) and note two different kinematic possibilities. First, if \( \psi \) is

\footnote{Obviously, if the SM/DS interactions violate baryon number \textit{and} lepton number, one must take care that any operators leading to proton decay are sufficiently suppressed. None of the operators we will be interested in carry baryon number, so this is a question for the UV completion, which is easily addressed. For instance, a sufficient but not necessary condition is that baryon and lepton number are accidental symmetries in the DS and SM below the GUT scale.}

\footnote{If \( l_X \) is positive instead of negative, then there are more \( X^* \)s than \( X \)s at late times, and \( X \) particles are anti-dark matter.
The former possibility is a two-body decay and thus leads to a monochromatic spectrum, but the second is multi-body and the spectrum is somewhat smeared out in energy, and the bound on $\Gamma_{\text{DM}}$ becomes slightly weaker.\(^8\)

The decay spectra for several operators are listed in Table 2. Using them, the bound on $\Gamma_{\text{DM}}/m_{\text{DM}}$ in Fig. 1 is then easily translated into a bound on the coefficients of the ADM operators, once the DM mass is specified. In the table is indicated the operator that gives rise to the spectrum in the scenario where the asymmetry is transferred by the operator $O$, the third and fourth rows give the decay spectra for the operator $O_6$, in the case of $X_1$ (i.e. bosonic) dark matter and in the case of $\psi$ (i.e. fermionic) dark matter, respectively. The final row presents the spectrum in the scenario where the asymmetry is transferred by the operator $O_2$ but the dark matter particle $\psi$ is lighter than $X$, which is produced off-shell and decays to two Weyl fermions $\chi, \chi^c$ through the interaction $gX^*\chi\chi^c$, generating $\mathcal{L} \supset g/\Lambda m_X^2 (\psi HL)(\chi\chi^c)$. In all but the last example above, the spectrum still has a fairly sharp feature at $E_\nu = m_{\text{DM}}/2$, and so is nearly as easily detected as a two-body decay. In fact, the fraction of events that fall in a bin of size $\Delta \log_{10} E = 0.3$ is 87% for the $O_6$ spectrum and 71% for the last example. We

\(^8\)In this case $\psi$ must have an appropriate mass so that it does not contribute significantly to the late-time density of the universe. Alternatively, $\psi$ may contribute to the dark matter density today, but then the prediction for the mass of $X$ must be correspondingly diminished.

\(^9\) In general, we ought to have a reason why the $X$ decay to light DS states is sufficiently slow. A simple reason is to use $(B - L)$ conservation and have $|l_X|$ smaller than the lepton number of any lighter DS states. For example, if $l_X = -1/3$ and $\psi$ is the only leptonic dark sector state lighter than $X$, then $X$ can decay only through suppressed higher dimensional operators like $O_2$ that involve the Standard Model. This is essentially a dark analogue of invoking baryon triality to forbid proton decay.

\(^10\) We take $\psi$ to be a Dirac fermion when calculating decay rates. If $\psi$ is the DM, then this is necessary anyway in order to give $\psi$ a DM-number conserving mass term.
have included this in the calculation of $\Lambda_{\text{min}}$ in Table 2, and in all cases we have taken $m_{\text{DM}} = 3$ GeV for the sake of concreteness.

While Table 2 shows that there is clearly a lot of model-dependence in the bound on the UV scale $\Lambda$, we will note two points. The first is that for the dimension 5 operator $O_2$, the scale is forced to be much larger than the Planck scale, and so is ruled out in the absence of additional structure in the DS. We will discuss this further in the next section. The second point is that for dimension 6 operators, the bound on $\Lambda_{\text{min}}$ is about $10^{12} - 10^{13}$ GeV, and thus the remaining window of open parameter space under the GUT scale is only about 3-4 orders of magnitude in $\Lambda$.

There will also be some small branching ratio for dark matter to decay through an off-shell Higgs to charged Standard Model particles; however, this will be very suppressed since the dark matter is quite light. Details of the comparison between a signal from neutrinos vs. from charged particles are given in appendix A, with the conclusion that decays to charged particles are negligible for the operators we are considering.

5 Discussion

Asymmetric dark matter is a compelling and simple framework, alternative to the standard thermal WIMP paradigm. In particular, ADM favors lighter DM masses around $O(1 - 10)$ GeV, which have received a recent boost in interest due to hints of a signal coming from the CoGeNT direct detection experiment [32], as well as possible explanations for the DAMA anomaly [33, 34, 35, 36, 37, 38, 39, 40, 41]. The phenomenological implications can differ in significant ways from those of a standard thermal WIMP, and we have noted one such possibility here, in which the DM particle carries lepton number and decays dominantly to anti-neutrinos. Such a scenario is motivated as a reasonably generic consequence of a mechanism for transferring the SM baryon asymmetry to the DS by virtue of the lepton number of the DM particle [3].

There are various model-building issues worth exploring in ADM. One would like to specify the mechanism that generates the asymmetry in the first place, and how the symmetric component is to be removed. Furthermore, one needs to generate a mass for the dark matter particle with the appropriate size to give the correct relic density. The fact that the typical masses required are $O(1 - 10$ GeV) suggests a common origin with the electroweak scale, suppressed by additional loop factors or smaller couplings. If the dark matter is a scalar, the mass also needs to be protected from radiative corrections. This is clearly related to the issue of forbidding or suppressing the marginal coupling $|X|^2|H|^2$ to the Higgs. Also, since conservation of lepton number is protecting the lifetime of the dark matter particle, it may not always be straightforward to take advantage of the see-saw mechanism for Standard Model neutrino masses.

We will comment briefly on how these issues may be addressed in one example using mechanisms already suggested in the literature. As noted in [3], the operator $O_2$ is attractive
since it may easily be UV-completed to a model with Standard Model Majorana neutrino masses arising from the see-saw mechanism. For example, we may add two $SU(2)$ doublet fermions $d, d^c$ with $L = 1 + l_X, -1 - l_X$ and hypercharge $Y = \frac{1}{2}, -\frac{1}{2}$ respectively, and take

$$\Delta L = \lambda \psi dH + \lambda' Ld^cX + m_ddd^c. \quad (13)$$

The right-handed neutrino masses arise from a scalar vev $\phi$ with $L = +2$. Thus no renormalizable couplings of $\phi$ to the fields $X, \psi, d, d^c$ are allowed as long as $-\frac{2}{3} < l_X < 0$, and lepton number can then easily be an accidental symmetry in the dark sector. Note, however, that in this particular example, the operator $O_2$ is parametrically suppressed only by the scale $m_d$, which would be constrained by observations to be above the Planck scale according to Table 2 if the decay of the dark matter particle $X \rightarrow \psi \nu$ were kinematically accessible. On the other hand, simple extensions may suppress $O_2$ without requiring such high scales. For example, consider a model with an extra gauged $U(1)$ broken by a scalar vev $\langle \Phi \rangle$, and two copies $d_i, d^c_i$ of the doublet fields, where the non-zero $U(1)$ charges of the fields are as follows:

|         | $U(1)$ |
|---------|---------|
| $X, d_2$ | 1       |
| $\Phi, d^c_2$ | -1      |

and 0 for the remaining fields. Then symmetry forces the lagrangian to be of the form

$$\Delta L = \lambda \psi d_1H + \lambda' Ld^c_1X + m_{d_1}d_1d^{c_1} + c\Phi d_2d^{c_2}, \quad (14)$$

which upon integrating out the $d_i$ fields leads to $O_2$ having a coefficient whose parametric suppression is $\frac{\lambda \lambda' \langle \Phi \rangle}{m_d^2}$.

So far, in our discussions, the DM mass has been put in by hand. However, if the baryon-to-dark matter density ratio is truly to be explained, then $m_{DM}$ must be dynamically generated at the correct scale. As noted in [42], a mass of $O(1 - 10 \text{ GeV})$ arises quite generically in models of spontaneously broken supersymmetry from gravity-mediated effects, provided that the MSSM masses are generated by gauge mediation with a messenger scale at $M \sim 10^{13} \text{ GeV}$. It would be more satisfying, though, to have a model where the mechanism that generates the DM mass is more predictive, and does not depend on a mass parameter that is free to vary over several orders of magnitude. Supersymmetry has the advantage that in theories with large $\tan \beta \sim 20$, allowing a superpotentiel term of the form $W \supset XHdS$, where $S$ is a doublet, generates a mass for $X$ at the appropriate size. Furthermore, additional contributions to the dangerous marginal term $|X|^2|H|^2$ can come only from superpotential terms, which may be easily controlled.

There are many ways the asymmetry can be generated; for an incomplete list see [43]. It is perhaps worth noting that the asymmetry does not have to originate in the SM and get transferred to the DS but could instead originate due to new sources of CP violation in the DS and then be transferred in the other direction.

Finally, there are other possible ways that neutrino signatures of ADM could appear. One potentially interesting direction currently under investigation involves neutrinos coming from
DM annihilation in the sun. ADM is very interesting in that its annihilation cross-section is not a priori related to the cross-section for capture in the sun. Indeed, because of conservation of lepton number, the scattering process will never contribute radiatively to the annihilation process, and thus there is no theoretical barrier to taking the two cross-sections to be quite different.\footnote{It is conceivable that an additional potentially significant boost in the solar signal could come from exponential growth of the DM occupation in the sun due to WIMP-WIMP scattering [46].}

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A Dominance of the Neutrino Signal

In this appendix, we will demonstrate that the neutrino signal is indeed the most important signal in the models we have considered. The main constraint comes from the requirement that the dark matter decays must have a sufficiently small branching ratio to produce positrons. Even though this branching ratio is very small, the positron background that the signal must overcome is smaller than the neutrino background by a factor of order 1000, and moreover, the slow diffusion of positrons through the galaxy results in a greater number of them remaining around to be detected. The dominance of the neutrino signal is thus something which must be checked. By making use of the “leaky box” model for cosmic ray propagation, in which positrons are considered to undergo a random walk process due to the galactic magnetic field, we will make a rough estimate for the required constraint on the positron branching fraction. This constraint will be seen to be satisfied by many orders of magnitude in the models we consider, and for this reason, a more precise calculation will not be necessary.

Let \( Q_{e^+}(E) \) be the production rate per unit volume per unit energy for positrons due to dark matter decays in the Milky Way. Depletion of cosmic ray positrons of energy \( E \sim O(1-10 \text{ GeV}) \) occurs due to a variety of mechanisms; they may lose energy- due to synchrotron radiation, inverse Compton scattering and bremsstrahlung- or they may actually manage to diffuse out of the galaxy. An equilibrium is established whereby the production and depletion effects reach a balance, yielding the relation \[44\]

\[
Q_{e^+}(E) \approx \frac{d}{dE} \left( \frac{dE}{dt} n_{e^+}(E) \right) + \frac{n_{e^+}(E)}{T}. \tag{15}
\]

Here \( n_{e^+}(E) \) is the number density per unit energy of signal positrons, and \( T \) is the average time needed for a positron to diffuse out of the galaxy. Energy losses due to synchrotron
radiation and inverse Compton scattering have $\frac{dE}{dt} \propto -E^2$. Bremsstrahlung, which begins to dominate at energies below a few GeV, has $\frac{dE}{dt} \propto -E$. The net energy loss rate may then be written as

$$ \frac{dE}{dt} = -bE^2 - cE, \quad (16) $$

with $b \sim \frac{1}{2 \times 10^5 \text{years} \times \text{TeV}}$ and $c \sim \frac{1}{6 \times 10^7 \text{years}}$. It follows then that positrons with energies of order several GeV will have a lifetime for energy loss of order $10^7$ years.

The random walk process for the diffusing positrons has them moving a typical distance

$$ r = \sqrt{D \times t} \quad (17) $$

in a time $t$, where $D \sim 3 \times 10^{28} \text{cm}^2/\text{s}$. The timescale for a positron to diffuse out of the galaxy, and thus move a distance of order the galactic height $10^3$ light years, is then given by

$$ T \sim 10^6 \text{years}. \quad (18) $$

We thus see that, at energies of order several GeV, the dominant depletion term on the right hand side of equation (15) is that due to diffusion. It then follows that $n_{e^+}(E)$ is given by

$$ n_{e^+}(E) \sim Q_{e^+}(E)T \quad (19) $$

Let us assume for simplicity that the source $Q(E)$ is a delta function at an energy $E_0$. In reality decays involving positrons will involve many body final states; this will only serve to spread out the spectrum and dilute the signal somewhat. We thus take

$$ Q_{e^+}(E) = \frac{\rho_{\text{DM}}\Gamma_{\text{DM}}}{m_{\text{DM}}} \gamma_{e^+} \delta(E - E_0), $$

where $\gamma_{e^+}$ is the branching ratio for dark matter decays to positrons. Putting everything together, we obtain

$$ n_{e^+}(E) \sim \gamma_{e^+} \frac{\rho_{\text{DM}}\Gamma_{\text{DM}}}{m_{\text{DM}}} T \delta(E - E_0) \quad (20) $$

For comparison, the dominant decays to neutrinos result in a spectrum (again assuming a delta function source for simplicity, centered on $E_0$)

$$ n_\nu(E) \sim \frac{\rho_{\text{DM}}\Gamma_{\text{DM}}}{m_{\text{DM}}} R_G \delta(E - E_0), \quad (21) $$

where $R_G$ is roughly the galactic scale of order $10^5$ light years. Integrating over energy bins of size of order $E_0$, the ratio of positron to neutrino signals is therefore approximately

$$ \frac{\int n_{e^+}(E) \, dE}{\int n_\nu(E) \, dE} \sim \frac{T}{R_G} \sim 10^{\gamma_{e^+}}. \quad (22) $$

In our models we estimate the positron branching ratio (due to decays through an off-shell W-boson) to be approximately $\gamma_{e^+} \sim \frac{1}{(2\pi)^2} \frac{m_{\text{DM}}^6}{m_{\text{W}}^4 m_{\text{e^+}}^2}$, which is $\sim 5 \times 10^{-10}$ at $m_{\text{DM}} = 10$ GeV. Given equation (22), this easily allows the neutrino signal to dominate, even given the $\sim 1000$ times larger background flux in atmospheric neutrinos compared to cosmic ray positrons.
B Annihilations

The framework for ADM we consider here cannot lead to a neutrino signal from dark matter annihilations in the galaxy at an observable level in the foreseeable future. The reason is that a sufficiently large coupling of dark matter to neutrinos would cause the lepton-asymmetry-transfering interactions to freeze out below the dark matter mass. To see this, recall that the condition of freezing out above the DM mass can be written

$$\frac{\Gamma_0}{m_{DM}} \lesssim \frac{\sigma_{m} m_{DM}^3}{H(m_{DM})},$$

(23)

where $\sigma_T T^3 \sim \langle n \sigma v \rangle_T$ is the annihilation rate at temperature $T$. We therefore must have

$$\frac{\Gamma_0}{m_{DM}} = \frac{\sigma_0}{\sigma_m} \frac{\rho_{DM}}{m_{DM}^2} < \frac{\sigma_0}{\sigma_m} \sqrt{g_*(m_{DM})} 10^{-62} \left( \frac{1 \text{GeV}}{m_{DM}} \right)^3,$$

(24)

where $\Gamma_0, \sigma_0$ are the annihilation rate and cross-section $\langle \sigma v \rangle$ today and $g_*(T)$ is the effective number of relativistic degrees of freedom in the universe at temperature $T$. One can see from Figure 1 that absurd enhancements of the cross-section over the value at momenta near the DM mass would be required in order to give an observable rate of DM annihilations in the galaxy.

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