

# Effectively Stable Dark Matter

Clifford Cheung and David Sanford  
*Walter Burke Institute for Theoretical Physics*  
*California Institute of Technology, Pasadena, CA 91125\**  
 (Dated: July 6, 2015)

We study dark matter (DM) which is cosmologically long-lived because of standard model (SM) symmetries. In these models an approximate stabilizing symmetry emerges accidentally, in analogy with baryon and lepton number in the renormalizable SM. Adopting an effective theory approach, we classify DM models according to representations of  $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_B \times U(1)_L$ , allowing for all operators permitted by symmetry, with weak scale DM and a cutoff at or below the Planck scale. We identify representations containing a neutral long-lived state, thus excluding dimension four and five operators that mediate dangerously prompt DM decay into SM particles. The DM relic abundance is obtained via thermal freeze-out or, since effectively stable DM often carries baryon or lepton number, asymmetry sharing through the very operators that induce eventual DM decay. We also incorporate baryon and lepton number violation with a spurion that parameterizes hard breaking by arbitrary units. However, since proton stability precludes certain spurions, a residual symmetry persists, maintaining the cosmological stability of certain DM representations. Finally, we survey the phenomenology of effectively stable DM as manifested in probes of direct detection, indirect detection, and proton decay.

## I. INTRODUCTION

Dark matter (DM) is elegantly accounted for by a neutral, cosmologically long-lived particle beyond the standard model (SM). In most circumstances, however, DM stability is ensured by a symmetry that is simply imposed by fiat. While as much can be expected from DM effective theories, similarly ad hoc choices are often needed for their ultraviolet completions, for instance in theories of supersymmetry or extra dimensions where  $R$ -parity or  $KK$ -parity are assumed.

In contrast, the SM implements stability with less contrivance: charge stabilizes the electron, while angular momentum stabilizes the neutrino. The proton is not guaranteed to be stable, but like DM it is cosmologically long-lived, with current limits bounding its lifetime to be greater than  $\sim 10^{33}$  years. Famously, the SM gauge symmetry explicitly forbids baryon and lepton number violation at the renormalizable level, suggesting an argument for proton stability from effective theory. This mechanism, whereby an approximate symmetry arises as the byproduct of existing symmetries, is sometimes referred to as an accidental symmetry.

In this paper we argue that DM, like the proton, can be cosmologically stable as an accident of SM symmetries. For our analysis we consider the symmetry group

$$SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_B \times U(1)_L, \quad (1)$$

which is exactly preserved in the renormalizable SM to all orders in perturbation theory. Beyond the renormalizable level,  $B$  and  $L$  may be approximate or exact, depending on whether they are gauged in the ultraviolet.

While SM quark and lepton flavor are approximate symmetries, we will not consider them here.

In our analysis, we enumerate models according to the quantum numbers of DM under the SM symmetry group. We take the stance of effective theory throughout: all interactions, renormalizable and non-renormalizable, are to be included with order one coefficients. We assume a DM mass  $M$  near the weak scale and an effective theory cutoff  $\Lambda$  at or below the Planck scale.

To start, we discard all representations without a component neutral under color and electromagnetism. We then discard all representations that permit DM decay into SM particles via dimension four or five operators. In such cases even a Planck scale cutoff is insufficient to prevent cosmologically prompt DM decay. On the other hand, decays induced by dimension six operators hover right at the boundary of current bounds from indirect detection, assuming a cutoff near the scale of grand unification (GUT) [1].

Applying these criteria, we enumerate models of effectively stable DM, focusing on all possible fermionic and scalar DM candidates whose leading decays occur at dimension six or seven. These results are summarized in Tabs. I and II for dimension six and Tabs. III and IV for dimension seven decays, respectively.

Our approach resonates with that of minimal DM [2, 3] but with the crucial difference that we incorporate both  $B$  and  $L$  and faithfully assume the presence of all interactions not forbidden by symmetry. Without  $B$  and  $L$ , representations smaller than the quintuplet of  $SU(2)_L$ , DM will promptly decay to SM particles via renormalizable interactions [2–4]. This is consistent with our own finding that effectively stable DM in small representations of  $SU(2)_L$  must carry  $B$  or  $L$ . For exactly this reason many of these models may be generated through the mechanism of asymmetric DM [5, 6]. Previous authors have also considered accidental stabilization of DM

\* clifford.cheung@caltech.edu, dsanford@caltech.edu

by flavor symmetries [7] or new gauge symmetries [8].

Finally, we consider the possibility that  $B$  and  $L$  are merely approximate. For this analysis we introduce a spurion parameterizing the hard breaking of baryon and lepton number by arbitrary units  $\Delta B$  and  $\Delta L$ , respectively. In many models of effectively stable DM, this induces prompt decays. However, not all values of  $(\Delta B, \Delta L)$  are phenomenologically safe: the non-observation of proton decay suggests that dimension six operators of the form  $Q^3 L$  should be forbidden, thus precluding hard breaking by units of  $(\Delta B, \Delta L) = (1, 1)$ . Remarkably, given arbitrary breaking by any  $(\Delta B, \Delta L) \neq (1, 1)$ , we still maintain a handful of viable candidates for effectively stable DM. In these models, DM is long-lived because of the SM gauge symmetry together with the stability of the proton.

The remainder of this paper is as follows. In Sec. II we enumerate viable models of effectively stable DM using the criteria described above. In Sec. III we study experimental constraints on these theories from direct detection, indirect detection, and proton decay. Finally, we summarize our results and discuss future directions in Sec. IV.

## II. CLASSIFICATION OF MODELS

In this section we enumerate representations of the SM symmetry group with a DM component that is neutral, cosmologically long-lived, and generated with the observed relic abundance. We adopt an effective theory perspective in which any operator allowed by symmetry is present with a strength set by a sub-Planckian cutoff.

### A. Neutrality

The quantum numbers of DM are parameterized by a discrete choice of representations for  $SU(3)_C$  and  $SU(2)_L$  together with a continuous choice of charges for  $U(1)_Y$ ,  $U(1)_B$ , and  $U(1)_L$ . For simplicity, we focus on pure gauge eigenstates, assuming that DM is either a complex scalar or a Dirac fermion.

To begin, we restrict to representations that have a neutral component under the unbroken SM gauge symmetry. Thus, DM is an  $SU(3)_C$  singlet. For DM that is a  $k$ -plet under  $SU(2)_L$ , we require that  $k \geq 2|Y| + 1$ , where the hypercharge  $Y$  is quantized to an integer or half-integer value if  $k$  is odd or even, respectively.

Charge neutrality places no constraints on the  $B$  or  $L$  charges of DM. While irrational values of  $B$  and  $L$  are allowed a priori, this is literally equivalent to enforcing DM number as an exact symmetry of the Lagrangian. Furthermore, exact global symmetries are known to conflict with black hole no-hair theorems [9]. For these reasons we restrict to rational values  $B$  and  $L$ .

While  $B$  and  $L$  are symmetries of the renormalizable SM, they may of course be spontaneously or explicitly broken in the full theory. To parameterize these effects conservatively, we introduce effective hard breaking of  $B$  and  $L$  into the low energy theory with a dimensionless spurion for symmetry breaking by units of  $(\Delta B, \Delta L)$ . For remainder of this section we assume that  $(\Delta B, \Delta L) = (0, 0)$ , but return to the issue of explicit breaking later on in Sec. III.

### B. Stability

We now determine the leading operators that are allowed by symmetry and mediate DM decay. An operator that induces DM decay takes the form

$$\mathcal{O}_{\text{DM}} = X \mathcal{O}_{\text{SM}}, \quad (2)$$

where  $X$  is the fermion or scalar field that contains DM and  $\mathcal{O}_{\text{SM}}$  is an operator composed entirely of SM fields. For later convenience, we define

$$N = [\mathcal{O}_{\text{DM}}], \quad (3)$$

to be the dimension of the DM decay operator. The quantum numbers of  $X$  are equal and opposite to those of  $\mathcal{O}_{\text{SM}}$ , so to enumerate all decay operators it suffices to determine all operators  $\mathcal{O}_{\text{SM}}$  of a given charge and operator dimension.

We define a fiducial decay rate into SM particles,

$$\Gamma(X \rightarrow \text{SM}) \sim \frac{M}{4\pi} \left( \frac{M}{\Lambda} \right)^{2(N-4)}, \quad (4)$$

corresponding to two-body decay via  $\mathcal{O}_{\text{DM}}$ . Depending on the precise operator, this decay may be three-body or higher. Moreover, decays into SM fermions will involve flavor structures that may further suppress the width. In any case, the fiducial decay rate in Eq. (4) should be taken as an overestimate.

By definition, a cosmologically stable DM particle  $X$  has a lifetime of order the age of the universe,

$$\tau(X \rightarrow \text{SM}) \gtrsim 10^{18} \text{ sec} \quad (\text{age of universe}). \quad (5)$$

However, this bound is weaker than experimental constraints on cosmic ray production from DM decay into positrons [10], gamma rays [11], antiprotons [12], and neutrinos [13]. These limits all place similar constraints on the DM lifetime, of order

$$\tau(X \rightarrow \text{SM}) \gtrsim 10^{26} \text{ sec} \quad (\text{indirect detection}), \quad (6)$$

for  $M$  of order the weak scale. While CMB bounds are also stronger than the one from the age of the universe, they are still weaker than indirect search bounds by  $\sim 2 - 3$  orders of magnitude [14].

Comparing Eq. (6) to Eq. (4), we obtain upper bounds on the cutoff of higher dimension operators. For dimension five, six, and seven operators, this implies a schematic lower bound on the cutoff,

$$\Lambda \gtrsim \begin{cases} \left(\frac{M}{1 \text{ TeV}}\right)^{\frac{26}{3}} 10^{29} \text{ GeV}, & N = 5 \\ \left(\frac{M}{1 \text{ TeV}}\right)^{\frac{4}{3}} 10^{16} \text{ GeV}, & N = 6 \\ \left(\frac{M}{1 \text{ TeV}}\right)^{\frac{2}{3}} 10^{12} \text{ GeV}, & N = 7 \end{cases} \quad (7)$$

Since Eq. (4) is an overestimate, this bound on the cutoff is conservative.

For dimension five operators, a sub-Planckian cutoff  $\Lambda$  would require  $M \lesssim 10 \text{ keV}$  for DM, so hereafter we consider only DM decay via dimension six and seven operators. Notably, dimension six operators are a particularly intriguing portal for weak scale DM because GUT-suppressed dimension six operators induce decays that lie just at the boundary of current experimental limits. For dimension seven operators, the bound on  $\Lambda$  is even smaller. Fig. 1 shows the lower bound on  $\Lambda$  as a function of  $M$  for dimension six and seven decay operators. The bounds for dimension five operators are not shown because they require a cutoff above the Planck scale. Also shown are the natural scale for DM mass  $M$  near the weak scale and cutoff scale  $\Lambda$  at the GUT scale, as well as regions excluded by experimental limits from LEP and direct detection, to be discussed later.

To summarize, a sub-Planckian cutoff implies that cosmologically stable, weak scale DM forbids all dimension four or five decay operators. This is a stringent constraint on the DM representation. For example, since a pure singlet scalar  $X$  can couple to any dimension four operator  $\mathcal{O}_{\text{SM}}$  in the SM, the associated DM will promptly decay at dimension five. At the opposite extreme is scalar or fermionic DM with extremely large charges. In this case lower dimensional operators are forbidden simply because so many SM fields have to be included in the decay operator just to preserve the symmetry. As noted in [2], this occurs for  $k$ -plets of  $SU(2)_L$  with large values of  $k$ .

All of our fermionic DM candidates and most of our scalar DM candidates carry non-zero  $B$  or  $L$ . The reason is that most representations with zero  $B$  or  $L$  typically have very low dimension decay operators. For example whenever  $\mathcal{O}_{\text{SM}}$  is a SM fermion bilinear that is  $B$  and  $L$  neutral, it is always possible to construct a lower dimension operator via replacements of the form  $QU \rightarrow H^\dagger$ ,  $LE \rightarrow H$ ,  $\bar{Q}\bar{\sigma}^\mu Q \rightarrow B^\mu, W^\mu$ , etc. As a result, viable DM tends to carry  $B$  or  $L$ , or reside in a large representation of  $SU(2)_L$  so that gauge invariance requires many Higgs fields in the operator.

Large hypercharge can serve the same role in terms of stability, but as noted earlier, a neutral DM component requires that  $k \geq 2|Y| + 1$  for a  $k$ -plet under  $SU(2)_L$ .

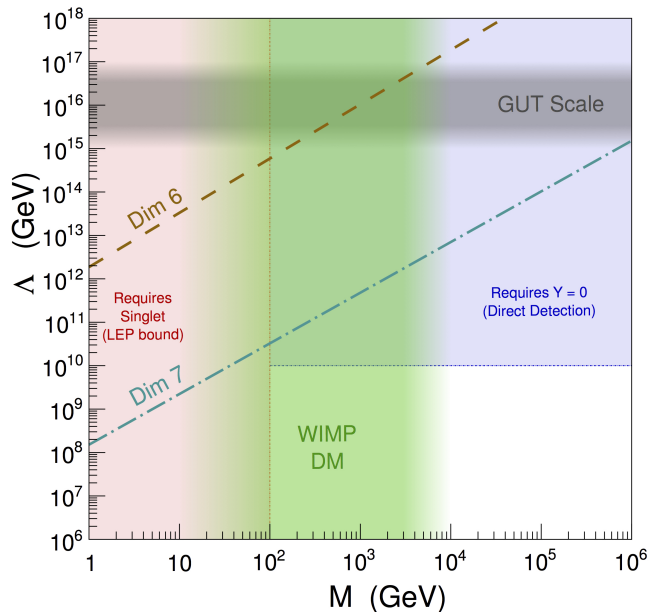


Figure 1. Lower bound on the effective theory cutoff  $\Lambda$  as a function of the DM mass  $M$  in order to avoid cosmologically prompt decay via dimension six (gold dashed) and seven (teal dot-dashed) operators. LEP (red shaded) excludes light, weakly interacting or hypercharged DM, while direct detection (blue shaded) excludes hypercharged DM that cannot accommodate an inelastic splitting to evade bounds.

This in turn bounds the net  $SU(2)_L \times U(1)_Y$  charges of the SM fields that couple to  $X$ . In particular, it largely excludes decay operators involving  $E$ , since this SM field has a sizable hypercharge and no  $SU(2)_L$  charge.

Tabs. I and II list all fermionic and scalar DM representations, respectively, whose leading decay is mediated by a dimension six operator. Every representation is color neutral and every representation carries  $B$  or  $L$  except the scalar  $SU(2)_L$  sextet. For the fermionic DM we forbid both members of the Dirac pair from decaying via dimension five or lower operators. Also, shown are various attributes of the model regarding direct detection and explicit  $B$  and  $L$  breaking, to be discussed later. Tabs. III and IV are the same as Tabs. I and II except they apply to DM representations whose leading decay is mediated by a dimension seven operator.

### C. Relic Abundance

Let us comment briefly on the origin of the DM relic abundance. DM with non-trivial SM gauge charges will be equilibrated with the thermal plasma in the early universe, in which case DM annihilations are mediated by gauge interactions whose strength is fixed by the charges of  $X$ . In general there may be additional interactions between the DM and SM fields, either through the Higgs boson or through direct couplings to quarks and leptons.

Fermionic DM ( $N = 6$ )

$SU(2)_L$	$U(1)_Y$	$U(1)_B$	$U(1)_L$	$\mathbb{B}, \mathbb{L}$	$\sigma_{\text{SI}}$	$\mathcal{O}_{\text{DM}}$
1	0	1	0	✓	✓	$UD^2X$
2	-1/2	-1	0	✗	✗	$Q^3X$
2	1/2	-1	0	✗	✗	$Q\bar{D}^2X$
3	-1	0	-1	✗	✗	$H^3LX$
3	0	1	0	✗	✓	$\bar{Q}^2DX$
3	1	1	0	✓	✗	$\bar{Q}^2UX$
4	-1/2	-1	0	✗	✗	$Q^3X$
4	3/2	0	-3	✗	✗	$L^3X$
5	-1	0	-1	✓	✗	$H^3LX$
5	0	0	-1	✓	✓	$H^\dagger H^2LX$
5	1	0	-1	✓	✗	$H^{\dagger 2}HLX$
5	2	0	-1	✓	✗	$H^{\dagger 3}LX$

Table I. Classification of fermionic DM which is stable up to dimension six decays. Columns 1-4 list DM charges under the SM symmetry group. Column 5 has a ✓ for DM that remains stable up to dimension six decays even when  $B$  and  $L$  are explicitly broken while preserving proton stability. Column 6 has a ✓ for DM that is safe from spin-independent DM-nucleon scattering through the  $Z$  boson. Column 7 lists an example operator mediating DM decay.

However, the DM charges are chosen by construction to eliminate renormalizable decays, so the latter are highly suppressed. On the other hand, the former are allowed for scalar DM, and from an effective theory perspective should be included if they are permitted by the symmetries of the theory.

If gauge interactions dominate DM annihilations, then the relic abundance is fixed by thermal freeze-out [2, 3]. Since the strength of the  $SU(2)_L \times U(1)_Y$  gauge interactions are known, there is only one free parameter available to tune the relic abundance: the DM mass. As discussed in [2, 3], to obtain DM of the observed relic abundance today requires  $M$  of order the weak scale or slightly above, depending on the DM spin and quantum numbers. Of course, direct couplings to the Higgs [15] can strongly affect the relic abundance. However, annihilations mediated by the Higgs generally add to the DM annihilation cross-section, depleting the DM abundance and thus requiring a larger DM mass.

Another approach to DM generation is asymmetric DM, which fits quite naturally into the effective theory picture presented here. In models of asymmetric DM there is an approximate DM number  $X$  that is conserved along with  $B$  or  $L$ . In the early universe, some set of interactions, for example higher dimension operators or sphalerons, break some combination of  $X$ ,  $B$ , and  $L$ , thus sharing particle asymmetries among the quarks, leptons, and DM. As we have shown, the majority of effectively stable DM candidates carry non-zero  $B$  or  $L$ , which implies that the very same operators that mediate DM decay can also share the asymmetries in the early universe. A similar observation was noted about asymmetric DM models in [6]. Alternatively, these kinds of

Scalar DM ( $N = 6$ )

$SU(2)_L$	$U(1)_Y$	$U(1)_B$	$U(1)_L$	$\mathbb{B}, \mathbb{L}$	$\sigma_{\text{SI}}$	$\mathcal{O}_{\text{DM}}$
1	0	0	-2	✗	✓	$H^2L^2X$
3	0	0	-2	✗	✓	$H^2L^2X$
5	0	0	-2	✗	✓	$H^2L^2X$
5	1	0	-2	✗	✗	$H^\dagger HL^2X$
5	2	0	-2	✗	✗	$H^{\dagger 2}L^2X$
6	1/2	0	0	✓	✗	$H^{\dagger 3}H^2X$
6	3/2	0	0	✓	✗	$H^{\dagger 4}HX$
6	5/2	0	0	✓	✗	$H^{\dagger 5}X$

Table II. Classification of scalar DM which is stable up to dimension six decays. Notation from Tab. I.

higher dimension operators can also be applied in the reverse direction to produce the baryon asymmetry [16].

If the decay operator is in chemical equilibrium in the early universe, then the DM and  $B$  or  $L$  asymmetries will be shared according to their charges [5]. As is well-known, for efficient sharing this typically requires light asymmetric DM of order the  $\sim$  GeV scale rather than the weak scale. For DM that is a gauge singlet, such as the first entry in Tab. I, this offers a viable model of asymmetric DM, provided additional annihilation modes to deplete the symmetric abundance.

On the other hand, models with DM carrying SM gauge charges are excluded by LEP [17, 18] for  $\sim$  GeV scale masses. In this case, asymmetric DM is possible only if asymmetry sharing through the decay operator is inefficient, so the abundance of DM is less than the amount prescribed by chemical equilibrium, thus requiring a larger DM mass.

### III. EXPERIMENTAL CONSTRAINTS

In this section we survey the bounds on effectively stable DM from laboratory and telescope experiments. As we will see, bounds from direct detection and proton decay have an interesting connection to explicit breaking of  $B$  and  $L$ .

#### A. Direction Detection

Experimental bounds on spin-independent DM-nucleon scattering are extremely stringent. In particular, Dirac or complex scalar DM with a spin-independent coupling to the  $Z$  boson is excluded by many orders of magnitude. Such an interaction arises when  $Y \neq 0$ , so direct detection constraints are trivially evaded by DM with vanishing hypercharge.

In Tabs. I and II and Tabs. III and IV, the column denoted  $\sigma_{\text{SI}}$  carries a ✓ for models which have zero hypercharge and are thus safe from direct detection, and ✗ otherwise. For the latter, models of hypercharged DM

Fermionic DM ( $N = 7$ )

$SU(2)_L$	$U(1)_Y$	$U(1)_B$	$U(1)_L$	$\mathcal{B}, \mathcal{L}$	$\sigma_{SI}$	$\mathcal{O}_{DM}$
3	1	-1	0	✗	✗	$H^\dagger Q D^2 X$
3	1	0	-3	✗	✗	$HL^3 X$
4	-3/2	0	-1	✓*	✗	$H^4 LX$
4	3/2	1	0	✓*	✗	$H^\dagger \bar{Q}^2 UX$
4	-1/2	1	0	✓*	✗	$H \bar{Q}^2 DX$
5	-1	-1	0	✓*	✗	$H Q^3 X$
5	0	-1	0	✓*	✓	$H^\dagger Q^3 X$
5	1	0	-3	✓*	✗	$HL^3 X$
5	2	0	-3	✓*	✗	$H^\dagger L^3 X$
6	-3/2	0	-1	✓	✗	$H^4 LX$
6	-1/2	0	-1	✓	✗	$H^\dagger H^3 LX$
6	1/2	0	-1	✓	✗	$H^{\dagger 2} H^2 LX$
6	3/2	0	-1	✓	✗	$H^{\dagger 3} HLX$
6	5/2	0	-1	✓	✗	$H^{\dagger 4} LX$

Table III. Classification of fermionic DM which is stable up to dimension seven decays. Notation from Tab. I, but with a ✓\* to indicate models which are cosmologically stable at dimension five or less but may still decay at dimension six when  $B$  and  $L$  are violated consistent with proton stability.

can evade these direct detection if we allow for additional structures which we now discuss.

In particular, limits on spin-independent scattering via the  $Z$  boson are null if there is an even tiny mass splitting  $\delta$  between the components of the Dirac fermion or complex scalar [19, 20]. In this case, scattering through the  $Z$  boson is inelastic, requiring additional energy to excite the incoming DM particle into a neighboring mass eigenvalue. In particular, scattering is kinematically forbidden provided the mass splitting exceeds

$$\delta \geq \frac{\beta^2 \mu}{2}, \quad (8)$$

where  $\beta \sim 200$  km/sec is the DM velocity and  $\mu$  is the reduced mass of the DM-nucleus system. For the typical atomic weights of targets such in experiments like CDMS [21], XENON100 [22], and LUX [23], for  $M \sim 100$  GeV this translates into a bound  $\delta \gtrsim 10$  keV. For very small  $M$ , the reduced mass will decrease and so too will the lower bound on  $\delta$ , but these regions of light DM are excluded by LEP for  $Y \neq 0$  [17, 18].

However, inducing the requisite mass splitting  $\delta$  requires new operators that explicitly break the DM particle number associated with Dirac fermions or complex scalars. For DM candidates that carry  $B$  or  $L$ , this implies explicit breaking of  $B$  or  $L$ . However, the spurion responsible for explicit breaking,  $(\Delta B, \Delta L)$ , enters with twice the  $B$  and  $L$  charge of the DM. Consequently, even with these splitting operators, there is still an unbroken  $Z_2$  subgroup of  $B$  or  $L$  that maintains DM stability.

For example, consider a fermionic DM particle that is a doublet of hypercharge  $Y = 1/2$ . The leading operator that can split its components is  $H^{\dagger 2} X^2 / \Lambda$ . A mass

Scalar DM ( $N = 7$ )

$SU(2)_L$	$U(1)_Y$	$U(1)_B$	$U(1)_L$	$\mathcal{B}, \mathcal{L}$	$\sigma_{SI}$	$\mathcal{O}_{DM}$
1	0	-1	-1	✓*	✓	$Q^3 LX$
2	-1/2	1	-1	✗	✗	$D^3 LX$
2	-1/2	0	-2	✗	✗	$H^3 L^2 X$
2	1/2	1	-1	✗	✗	$UD^2 LX$
3	-1	-1	-1	✗	✗	$Q\bar{U}^2 LX$
3	-1	1	1	✗	✗	$\bar{Q}D^2 LX$
3	0	-1	-1	✓*	✓	$Q^3 LX$
4	-3/2	-1	1	✗	✗	$Q^3 EX$
4	-1/2	-1	1	✗	✗	$Q^2 \bar{U} LX$
4	-1/2	0	-2	✗	✗	$H^3 L^2 X$
5	0	-1	-1	✓*	✓	$Q^3 LX$
5	2	0	-4	✗	✗	$L^4 X$
6	-1/2	0	-2	✓*	✗	$H^3 L^2 X$
6	1/2	0	-2	✓*	✗	$H^\dagger H^2 L^2 X$
6	3/2	0	-2	✓*	✗	$H^{\dagger 2} H L^2 X$
6	5/2	0	-2	✓*	✗	$H^{\dagger 3} L^2 X$
7	0	0	0	✓	✓	$H^{\dagger 3} H^3 X$
7	1	0	0	✓	✗	$H^{\dagger 4} H^2 X$
7	2	0	0	✓	✗	$H^{\dagger 5} H X$
7	3	0	0	✓	✗	$H^{\dagger 6} X$

Table IV. Classification of scalar DM which is stable up to dimension seven decays. Notation from Tab. I.

splitting sufficient to evade direct detection requires

$$\Lambda \lesssim \frac{v^2}{\delta} \sim 10^9 \text{ GeV}, \quad (9)$$

which is a low cutoff in the context of DM stability. For even larger hypercharges, the level splitting operator involves more Higgs fields, and the requisite cutoff is even lower, dropping to  $\sim 30$  TeV for  $Y = 1$ .

According to the philosophy of effective theory, the bound in Eq. (9) defines a cutoff scale at which all higher dimension operators allowed by symmetry should be present—including those which can mediate DM decay. For dimension five, six, and seven operators, such a low cutoff is inconsistent with cosmological limits on weak scale DM. This is depicted in Fig. 1, where the blue shaded region shows that below above  $\Lambda \gtrsim 10^9$  GeV, the cutoff is too high to induce a sufficient mass splitting to evade direct detection, so  $Y = 0$  is forbidden. Conversely, effectively stable DM with non-zero hypercharge can only evade direct detection if there is a low cutoff, in which case cosmological stability requires the DM decay operator be dimension eight or higher.

In general, if the mass splitting requires a higher dimension operator then direct detection is inconsistent with the criterion of cosmologically stable DM. On the other hand, there is no issue if the splitting operator is renormalizable. For example, this is possible for complex scalar DM with  $Y = 1/2$ , which permits the renormalizable operator  $\lambda H^{\dagger 2} X^2$ . For  $M = 100$  GeV, a splitting of  $\delta = 10$  keV is achieved for  $\lambda \gtrsim 10^{-9}$ , which is easily satisfied for order one couplings. If  $X$  carries  $B$  or

$L$ , this mass splitting operator explicitly breaks  $B$  or  $L$  down to a discrete baryon or lepton parity, however still maintaining DM stability. Alternatively, this interaction is symmetry preserving if  $X$  does not carry  $B$  or  $L$ , as is the case for  $k$ -plets of  $SU(2)_L$  with large values of  $k$  which couple only to Higgs bosons in the decay operator.

Even for models which evade bounds on  $Z$ -mediated scattering, direct detection may still impose constraints on Higgs boson exchange. For example, scalar DM-nucleon scattering can occur via the Higgs portal interaction,  $H^\dagger H X^\dagger X$  [24]. For fermionic DM, the analogous interactions are higher dimension. Non-singlet fermionic and scalar candidates can also scatter with nucleons at loop level via multiple gauge boson exchange, which may be observable in the next generation of direct detection experiments [25].

## B. Indirect Detection

Since these DM candidates eventually decay on cosmological time scales, they are naturally probed by cosmic ray telescopes. Conveniently, the authors of [6] studied indirect detection constraints on DM decay via high dimension operators of the very type considered in this paper, albeit with underlying motivation of asymmetric DM. In particular, they considered bounds from FERMI, PAMELA, AMS-02, and HESS on high energy gamma rays and charge particle cosmic rays from electrons, protons, and anti-protons, obtaining a limit on the DM lifetime which we have taken as a loose input for Eq. (6). We refer the reader to [6] for precise numerical bounds, but we summarize the salient takeaways below.

In general, DM carrying  $B$  will decay to quarks, yielding anti-protons, while DM carrying  $L$  will decay to leptons, yielding positrons and neutrinos. All of the DM decays considered here will produce high energy gamma rays from charged particle bremsstrahlung, hadronic decays, and inverse Compton scattering of CMB photons and starlight. However, due to the large  $B$  and  $L$  charges of most our DM candidates, gamma ray lines are not typically expected among the theories considered here. An exception occurs for DM particles which decay through operators involving only the Higgs boson, in which case mixing together with a loop of SM particles will induce two-body decays of DM to photons. Another possibility occurs for DM with unit  $L$  number, which can decay to photon plus neutrino.

## C. Proton Decay

The non-observation of proton decay offers a strong motivation for at least approximate  $B$  and  $L$  conservation. Current limits on  $p^+ \rightarrow e^+ \pi^0$  and related decay modes from the Super-Kamiokande experiment require a lifetime of at least  $\sim 10^{33}$  years [26–29], which

already places significant constraints on the simplest GUTs. From an effective theory viewpoint, proton decay is mediated by dimension six operators of the form

$$\frac{Q^3 L}{\Lambda^2}, \quad (10)$$

which breaks  $B$  and  $L$  but preserves  $B - L$ . Current limits on proton decay imply a lower bound on the cutoff of approximately  $\Lambda \gtrsim 10^{15}$  GeV [30], so  $B$  and  $L$  are very well-preserved symmetries.

As noted earlier, we can parameterize  $B$  and  $L$  violation with a dimensionless spurion characterizing hard breaking by some number of units of  $\Delta B$  and  $\Delta L$ . However,  $(\Delta B, \Delta L) = (1, 1)$  is of particular note because it permits proton decay via the operator in Eq. (10). As a result, if we wish to forbid this then we should avoid breaking by units of  $(\Delta B, \Delta L) = (1, 1)$ . Remarkably, even with  $(\Delta B, \Delta L) \neq (1, 1)$  there are still viable models of effectively stable DM. In these models there are low dimension DM decay operators allowed by the SM gauge symmetry, but these operators require explicit breaking by  $(\Delta B, \Delta L) = (1, 1)$  which would also decay the proton.

For example, this happens in the model described by the first row of Tab. I. Since  $X$  is a gauge singlet, the gauge invariant operator  $HLX$ , if present, would induce catastrophically prompt DM decay. However, the existence of both  $UD^2X$  together with  $HLX$  would require symmetry breaking by  $(\Delta B, \Delta L) = (1, 1)$  which would in turn induce dimension six proton decay. Conversely, if  $B$  and  $L$  are explicitly broken while forbidding proton decay, an accidental symmetry remains which effectively stabilizes DM.

Note that proton decay is also mediated by operators beyond dimension six requiring  $B$  and  $L$  violation by one unit and an odd number of units, respectively. For example,  $H^\dagger D^3 L$  induces proton decay with  $(\Delta B, \Delta L) = (-1, 1)$  breaking. Forbidding this larger class of spurions would allow for even more viable candidates for cosmologically stable DM, but we do not consider this possibility here.

In Tabs. I and II and Tabs. III and IV, the column labelled  $\mathcal{B}, \mathcal{L}$  carries a  $\checkmark$  if DM is still effectively stable—that is, cannot decay at dimension five or less—even after including any hard breaking spurion with  $(\Delta B, \Delta L) \neq (1, 1)$ , and  $\times$  otherwise. The  $\checkmark^*$  entries in Tabs. III and IV indicate models with DM that, while stable at dimension five or less, still decays at dimension six. For these models, DM is still cosmologically stable if the cutoff is higher than was required for dimension seven decays.

## IV. SUMMARY AND OUTLOOK

DM phenomenology often hinges on the assumption of a stabilizing symmetry. Naturally, this leads one to wonder about the underlying reason for cosmological stability. In this paper we present an alternative hypothesis

whereby DM is long-lived as an accident of the SM symmetry group. We have classified all models in which DM decay is cosmologically slow and induced by dimension six or seven operators. In such cases a sub-Planckian cutoff is sufficient to prevent cosmologically prompt DM decays which are excluded by indirect detection. All candidates for effectively stable DM either carry  $B$  or  $L$ , or reside in a high representation of  $SU(2)_L$ . We have identified those models which are consistent with stringent bounds on spin-independent DM-nucleon scattering. Finally, we have accounted for the possibility explicit breaking of  $B$  and  $L$  by arbitrary units. As long as the symmetry breaking spurion still forbids proton decay, models of effectively stable DM persist.

Our analysis leaves a number of avenues for future work. First, there is the question of DM stabilization

by SM quark and lepton flavor symmetries, which are well-preserved for the light generations. Second, there is the question of building these models in a supersymmetric context. Lastly, it would be interesting to see if any of the models presented here arise explicitly in GUT constructions.

## ACKNOWLEDGMENTS

CC is supported by a DOE Early Career Award DE-SC0010255 and a Sloan Research Fellowship. DS is supported in part by U.S. Department of Energy grant DE-FG02-92ER40701 and by the Gordon and Betty Moore Foundation through Grant No. 776 to the Caltech Moore Center for Theoretical Cosmology and Physics.

- 
- [1] A. Arvanitaki, S. Dimopoulos, S. Dubovsky, P. W. Graham, R. Harnik, *et al.*, Phys.Rev. **D79**, 105022 (2009), arXiv:0812.2075 [hep-ph].
  - [2] M. Cirelli, N. Fornengo, and A. Strumia, Nucl.Phys. **B753**, 178 (2006), arXiv:hep-ph/0512090 [hep-ph].
  - [3] M. Cirelli, A. Strumia, and M. Tamburini, Nucl.Phys. **B787**, 152 (2007), arXiv:0706.4071 [hep-ph].
  - [4] L. Di Luzio, R. Grober, J. F. Kamenik, and M. Nardecchia, (2015), arXiv:1504.00359 [hep-ph].
  - [5] D. E. Kaplan, M. A. Luty, and K. M. Zurek, Phys.Rev. **D79**, 115016 (2009), arXiv:0901.4117 [hep-ph].
  - [6] Y. Zhao and K. M. Zurek, JHEP **1407**, 017 (2014), arXiv:1401.7664 [hep-ph].
  - [7] L. Lavoura, S. Morisi, and J. Valle, JHEP **1302**, 118 (2013), arXiv:1205.3442 [hep-ph].
  - [8] O. Antipin, M. Redi, A. Strumia, and E. Vigiani, (2015), arXiv:1503.08749 [hep-ph].
  - [9] T. Banks and N. Seiberg, Phys.Rev. **D83**, 084019 (2011), arXiv:1011.5120 [hep-th].
  - [10] A. Ibarra, A. S. Lamperstorfer, and J. Silk, Phys.Rev. **D89**, 063539 (2014), arXiv:1309.2570 [hep-ph].
  - [11] S. Ando and K. Ishiwata, (2015), arXiv:1502.02007 [astro-ph.CO].
  - [12] G. Giesen, M. Boudaud, Y. Genolini, V. Poulin, M. Cirelli, *et al.*, (2015), arXiv:1504.04276 [astro-ph.HE].
  - [13] C. Rott, K. Kohri, and S. C. Park, (2014), arXiv:1408.4575 [hep-ph].
  - [14] J. M. Cline and P. Scott, JCAP **1303**, 044 (2013), arXiv:1301.5908 [astro-ph.CO].
  - [15] C. Cheung and D. Sanford, JCAP **1402**, 011 (2014), arXiv:1311.5896 [hep-ph].
  - [16] C. Cheung and K. Ishiwata, Phys.Rev. **D88**, 017901 (2013), arXiv:1304.0468 [hep-ph].
  - [17] A. Heister *et al.* (ALEPH), Phys.Lett. **B533**, 223 (2002), arXiv:hep-ex/0203020 [hep-ex].
  - [18] J. Abdallah *et al.* (DELPHI), Eur.Phys.J. **C31**, 421 (2003), arXiv:hep-ex/0311019 [hep-ex].
  - [19] L. J. Hall, T. Moroi, and H. Murayama, Phys.Lett. **B424**, 305 (1998), arXiv:hep-ph/9712515 [hep-ph].
  - [20] D. Tucker-Smith and N. Weiner, Phys.Rev. **D64**, 043502 (2001), arXiv:hep-ph/0101138 [hep-ph].
  - [21] R. Agnese *et al.* (CDMS), Phys.Rev.Lett. **111**, 251301 (2013), arXiv:1304.4279 [hep-ex].
  - [22] E. Aprile *et al.* (XENON100), Phys.Rev.Lett. **109**, 181301 (2012), arXiv:1207.5988 [astro-ph.CO].
  - [23] D. Akerib *et al.* (LUX), Phys.Rev.Lett. **112**, 091303 (2014), arXiv:1310.8214 [astro-ph.CO].
  - [24] R. Schabinger and J. D. Wells, Phys.Rev. **D72**, 093007 (2005), arXiv:hep-ph/0509209 [hep-ph].
  - [25] J. M. Cline, K. Kainulainen, P. Scott, and C. Weniger, Phys.Rev. **D88**, 055025 (2013), arXiv:1306.4710 [hep-ph].
  - [26] K. Abe *et al.* (Super-Kamiokande), Phys.Rev. **D90**, 072005 (2014), arXiv:1408.1195 [hep-ex].
  - [27] V. Takhistov *et al.* (Super-Kamiokande), Phys.Rev.Lett. **113**, 101801 (2014), arXiv:1409.1947 [hep-ex].
  - [28] K. Abe *et al.* (Super-Kamiokande), Phys.Rev.Lett. **113**, 121802 (2014), arXiv:1305.4391 [hep-ex].
  - [29] C. Regis *et al.* (Super-Kamiokande), Phys.Rev. **D86**, 012006 (2012), arXiv:1205.6538 [hep-ex].
  - [30] P. Nath and P. Fileviez Perez, Phys.Rept. **441**, 191 (2007), arXiv:hep-ph/0601023 [hep-ph].