Asymmetric Dark Matter

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We present notes on asymmetric dark matter (ADM), with a particular focus on elaborating on the points made in [1]. We focus undestanding how to derive cosmological properties of general models of dark matter, as well as particle physics and cosmological considerations that must be taken into account when working with models of ADM. We also summarize some possible signatures of models of ADM and mention some follow up work on the subject.

Component	Ω
Cosmological Constant/Dark Energy	0.742
Dark Matter	0.214
Baryonic Matter	0.0441

TABLE I. Energy distributions of the major components of the universe relative to the critical density $\rho_{\rm crit} \sim 1$.

I. MOTIVATION

Recent results from cosmological observations have led to several new puzzles. It is possible that some of these puzzles are connected to the major puzzles of particle physics at the weak scale. One such puzzle that may be plausibly connected to the weak scale relates to the abundances of particles in the universe. Using data from astrophysical observations from several sources, including the CMB specturm, structure on super-galactic scales, and redshifts of distant luminous sources, the relative energy density distributions of the major components of the present-day universe have been determined accurately[2], as shown in Table I. Note that three components of the universe that we might think are unrelated all have contributions within an order of magnitude of each other. The order 1 contribution of the cosmological constant could plausibly be a coincidence, albeit a very striking one, since its contribution was small in the early universe and only relatively recently on cosmological scales became important.

The dark matter density, on the other hand, tracks the baryon density until a time when they become relativistic or strongly interacting, since both can be treated as non-interacting dust before that. Thus, the baryon and dark matter densities should be generated with comparable sizes. In the standard WIMP with separate baryogenesis paradigm, they are in fact generated in very different ways, so it is striking that they come out to be the same order of magnitude. The WIMP density is generated by thermal freeze-out, while the baryon density is generated by CP violating, out-of-equlibrium processes in the very early universe that lead to a baryon-antibaryon asymmetry. The question we would like to address is whether it is possible to naturally link the two densities and explain why $\rho_B \sim \rho_{\rm DM}$.

The obvious way to do this is to charge the dark matter under $B-L$. Scenarios involving this mechanism are called models of Aymmetric Dark Matter (ADM) [1]. The dark matter-anti dark matter asymmetry is then transferred from the baryon sector via higher dimension operators. Alternatively, it can be generated in the dark matter sector itself and then transferred to the baryons, as mentioned at the end of these notes. Either mechanism fixes $n_{\rm DM} \approx n_B$. The energy densities are then related with an extra factor of mass giving

$$\rho_{\rm DM} \sim \frac{m_{\rm DM}}{m_p} \rho_B. \tag{1}$$

The ratio of the two energy densities is known and fixes the size of the dark matter mass to be around 5 GeV, which happens to be in the exciting region indicated by DAMA and CoGENT.

To summarize, we are motivated to look at models of ADM to accomplish three goals: to explain $\rho_{\rm DM} \sim \rho_B$, to see how light dark matter in the DAMA-CoGENT area could arise, and perhaps to come up with a new mechanism of baryogenesis.

In what follows we take the following philosophy. We view ADM as a framework within which we might explain the above issues. We do not, however, take any given model as the final word on ADM. Therefore, the emphasis in these notes is on how to determine the important features of a given ADM model as opposed giving a survey of all the ADM models that have been proposed so far. There are several techniques of particle thermodynamics that are widely employed in studying models of ADM and we give a review of the most important ones. Section 2 reviews non-equilibrium thermodynamics and the Boltzmann equation. Section 3 reviews equilibrium thermodynamics, charge asymmetries, and chemical potentials. Section 4 gives an example of how to apply these techniques in a somewhat "minimal" context. Section 5 gives an overview of a few other possible models, highlighting the differences in phenomenology and potential issues that one must address in models of ADM. Finally, section 6 discusses some phenomenology for direct detection and collider experiments.

II. COSMOLOGY AND BOLTZMANN EQUATIONS

We review some basic aspects of non-equilibrium thermodynamics in an expanding universe. The starting point of all cosmology is the FRW metric

$$ds^2 = dt^2 - a(t)^2 d\vec{x}^2, \tag{2}$$

where $a(t)$ is the scale factor which describes the scale of the universe. Einstein's equations give the two Friedmann equations for the scale factor a, one of which tells us that

$$H^2 = \frac{8\pi}{3M_{\rm pl}^2}\rho, \tag{3}$$

where ρ is the density of matter in the universe and $H \equiv \dot{a}/a$ is the Hubble rate. We are generally concerned with times when the temperature is above the matter-radiation equality time, which is around 1 eV. Therefore, I will assume from here on that matter is predominantly radiation. Then, the density of photons in the universe is largest and we know that $\rho_\gamma \sim T^4$ from the Stefan-Boltzmann equation. Thus, from (3), we get the following extremely important relation to keep in mind:

$$H \sim \frac{T^2}{M_{\rm pl}}. \tag{4}$$

Why do we care so much about the Hubble rate? We are concerned with when processes are in or out of equilibrium. This is determined by the Boltzmann equation. For a particle with density $n(t)$ and decay rate Γ,

$$\frac{dn}{dt} = -3Hn - \Gamma(n - n_{\rm eq}). \tag{5}$$

Note that the Boltzmann equation as written above is an approximation of a more general formula. It also must be modified when discussing annihilations. The first term tells us that the expansion of the universe tends to dilute the density of matter. The second term tells us that interactions between particle species tend to push the density toward its equilibrium value at a given temperature T.

So long as $\Gamma \gtrsim H$, the particle will be predominantly pushed toward equilibrium: it reacts fast enough on cosmological times to maintain equilibrium. If H becomes larger than Γ, then the first term begins to dominate and the abundance of the particle will decrease out-of-equilibrium. Another way of saying this is that the reverse process that reforms the particle ceases to happen. Thus, our condition for out-of-equilibrium dynamics, using $H \sim T^2/M_{\rm pl}$ is

$$\frac{T^2}{M_{\rm pl}} \gtrsim \Gamma. \tag{6}$$

Next, we estimate the rate in various circumstances, starting with the most simplifying assumptions possible and slowly breaking those assumptions. Consider a model with a cutoff Λ and $T \ll \Lambda$. We want to consider the rate of decay/annihilation of particle X with mass M_X larger than the masses of the particles it decays/annihilates to. Consider the simplest case: $M_X \ll T$ and X unstable. In this case, the rate of decay for X is the usual rate Γ calculated from the underlying QFT multiplied by a relativistic γ factor, given by $\gamma \sim T/M_X$. If the decay is two body, then $\Gamma \sim \alpha/4\pi M_X$. The final point to make here is that both the mother and daughters will be highly relativistic and thus have comparable densities. The final result for this case:

$$\Gamma_D \sim \Gamma_{ID} \sim \frac{\alpha}{4\pi}\frac{M_X^2}{T} \tag{7}$$

For this to go out of equilibrium some time during the relativistic era of X, we need

$$\frac{T^2}{M_{\rm pl}} \gtrsim \frac{\alpha}{4\pi}\frac{M_X^2}{T} \to \alpha \lesssim 4\pi\frac{M_X}{M_{\rm pl}}. \tag{8}$$

In other words, the particle must be extremely weakly coupled. If the particle is relatively heavy, but decays by higher dimension operators, this becomes possible. More typically, the particle remains close to equilibrium throughout its relativistic era.

When $T \lesssim M_X$, things begin to change. The relatively minor change is that the decay rate of the particle itself no longer has a time dilation suppression. The big change is that now the inverse process has a different rate. Detailed balance tells us that near equilibrium $\Gamma_D n_X \approx \Gamma_{ID} n_D$, where D are the decay products of X. This means that

$$\Gamma_{ID} \approx \frac{n_X}{n_D}\Gamma_D. \tag{9}$$

But if $M_D \ll M_X$, then D will be relativistic and X will not be. Lesson 1 for non-relativistic particles: their densities are Boltzmann suppressed. This means that $\Gamma_{ID} \ll \Gamma_D$ exponentially in M_X/T. For temperatures even a decade below the mass of X, this suppression is huge and can push the inverse decay rate below the Hubble rate. Let's see this in a little more detail:

$$\frac{T^2}{M_{\rm pl}} \gtrsim \frac{n_X}{n_D}\frac{\alpha}{4\pi}M_X \approx \left(\frac{M_X}{T}\right)^{3/2} e^{-M_X/T}\frac{\alpha}{4\pi}M_X. \tag{10}$$

This equation has no analytic solution in T, but the right-hand side clearly goes to 0 faster than the left, so there will be a temperature where this is satisfied. What point is that? For masses within a few orders of magnitude of the weak scale and a large range of relatively natural renormalizable couplings, the decouplings happens a temperature order 10 below the mass of the particle: it is very fast, as you might expect given the exponential.

Next, we consider a particle that is (nearly) stable. Suppose, however, that it can annihilate by interacting with an $\overline{X}$ to some much lighter decay products. The Boltzmann equation in this case is

$$\dot{n} + 3Hn = \langle\sigma v\rangle(n^2 - n_{\rm eq}^2). \tag{11}$$

We define the rate of reactions to be $\Gamma = \langle\sigma v\rangle n$

There are two possible scenarios. If the decay proceeds by an onshell light particle, then the cross-section goes like $\sigma \sim \alpha^2/E^2$. If the decay proceeds offshell by a non-renormalizable operator of dimension d, then the cross-section goes like $\sigma \sim$ P.S. $\alpha E^{d-6}/\Lambda^{d-4}$, where P.S. is some factors of phase space. For $T \gg m_X$, the rate of annihilation is given by σn_X, so that

$$\Gamma \sim \begin{cases} \alpha^2 T, & \text{on-shell light propagator} \\ \text{P.S. } \alpha\frac{T^{2d-7}}{\Lambda^{2d-8}}, & \text{off-shell non-renormalizable interaction} \end{cases}. \tag{12}$$

This reaction goes out of equilibrium as in the case of decay when $H \sim T^2/M_{\rm pl} \gtrsim \Gamma$, so that

$$\begin{array}{ll} T \gtrsim \alpha^2 M_{\rm pl}, & \text{on-shell light propagator} \\ T \lesssim \left(\frac{\Lambda^{2d-8}}{\text{P.S.}\alpha M_{\rm pl}}\right)^{1/(2d-9)}, & \text{off-shell non-renormalizable interaction} \end{array}. \tag{13}$$

Unless α is extremely small, the first scenario will not occur within the realm of applicability of our theory. The second case, however, can definitely occur and will be important. For a concrete example, suppose the interaction given by a dimension 6 operator (i.e. four-fermi). Then the condition is

$$T \lesssim \left(\frac{\Lambda^4}{\alpha M_{\rm pl}}\right)^{1/3}. \tag{14}$$

For $\alpha = 1/4\pi$ and $M_X = 1$ TeV, this is satisfied for $T > M_X$ so long as $\Lambda \gtrsim 5 \times 10^6$ GeV, which is entirely plausible.

Lastly, let's suppose that the interactions do not go out of equilibrium before $T \sim M_X$. In this case, as before, detailed balance tells us that the inverse decay will be highly Boltzmann suppressed and will surely go out of equilibrium by the time $T \sim M_X/(\text{few} \times 10)$. Furthermore, the *forward* decay rate $\langle\sigma v\rangle n$ also becomes highly Boltzmann suppressed and goes out of equilibrium. This process is known as thermal freeze-out. Essentially, just as the density of X and $\overline{X}$ starts to get Boltzmann suppressed, the H term in the Boltzmann equation begins to dominate and the density in a proper volume gets frozen at the proper density for the freeze-out temperature $T_{\rm freeze-out} \sim M_X/(\text{few} \times 10)$. The expansion of the universe becomes too fast for the dark matter particles to find each other and annihilate in appreciable amounts.

Note that all of the above formalism applies equally well for other types of scattering processes, with the appropriate reinterpretations. For another example we will care about, there may be B- or L-violating scattering processes. The conditions for these to be in- or out- of equilibrium are exactly the same.

III. EQUILIBRIUM THERMODYNAMICS AND CONSERVATION OF CHARGE

The other aspect of thermodynamics that will be important for us is the consequence of conservation of charge and how to deal with matter-anti matter asymetries in thermodynamics. For those, such as myself, who are not very familiar with the formalism for dealing with this, namely chemical potentials, I give a brief review, following [3] and [4].

In particle physics (and also many condensed matter systems), the number of particles (or excitations) of a given type is not conserved. There can be many interactions that convert one type of particle to another. However, if there is a conserved charge, then the total charge asymmetry must be conserved. More explicitly, if there are two species particles charged + and − under a $U(1)$ symmetry, then N_+ and N_- are not controlled variables for a given system, but $N_+ - N_-$ is. That means that we can set $N_+ - N_-$ however we like, but the actual number of each type of particle is determined by minimizing the free energy.

Now, we define the chemical potential μ_i of a particle species i as $\partial F/\partial N_i$ while holding the controlled parameters fixed. We can talk about the free energy in two ways. Consider once again the situation with two particles species with charges + and −. We can consider both N_+ and N_- as external variables to be fixed by minimizing the free energy. But we know there is a constraint that $N_0 \equiv N_+ - N_-$ is fixed, so we can also treat N_0 as an controlled parameter and only have one of the particle numbers, say N_-, fixed by minimizing the free energy. Of course, the two must match up so that

$$F(T, V, N_0; N_-) = F(T, V; N_-, N_+). \tag{15}$$

The semi-colon seperates variables that define the state of the system ("controlled parameters") from variables that we can determine from thermodynamics ("internal variables"). The state of the system is determined by minimizing the free energy with respect to the internal variables while holding the control parameters fixed, so that either way of writing it, we must have

$$\frac{\partial F}{\partial N_-}|_{T,V,N_0} = \frac{\partial F}{\partial N_-}|_{T,V} + \frac{\partial F}{\partial N_+}|_{T,V}\frac{dN_+}{dN_-} = \mu_- + \mu_+ = 0, \tag{16}$$

where we use $dN_+/dN_- = 1$ since charge is conserved and the definition of chemical potential

$$\mu_i \equiv \frac{\partial F}{\partial N_i}|_{V,T}. \tag{17}$$

We learn something very general: chemical potentials are associated with conserved quantities. More generally, equation (16) holds for any interaction in the system. That is, for an interaction with particles a going in and paricles b going out, then

$$\sum_a \mu_a = \sum_b \mu_b. \tag{18}$$

Ultimately, we would like to find quantities like $n_+ - n_-$, switching now to number densities, which are more natural in cosmology. These are related to the chemical potential. Supposing that the chemical potential is small compared to the temperature (which is reasonable since the baryon asymmetry is small), then by integrating the Fermi-Dirac and Bose-Einstein distributions, we find that the asymmetry between particles with positive and negative charge under a symmetry is

$$n_+ - n_- = \begin{cases} \frac{gT^3}{3}\frac{\mu}{T}, & \text{for bosons} \\ \frac{gT^3}{6}\frac{\mu}{T}, & \text{for fermions} \end{cases}, \tag{19}$$

where g is the spin degeneracy and μ is the total chemical potential of the particles charged under the corresponding symmetry. If we consider temperatures above the electroweak phase transition (EWPT), then our conserved charge is $U(1)_Y$. Below the EWPT, the conserved charge is $U(1)_{EM}$. We are further interested in writing down the baryon asymmetry B and lepton asymmetry L. Below the scale of baryogenesis/leptogenesis, these are conserved as well at some fixed amount. By applying (16), we can see that the total chemical potential contributing to a given asymmetry must be weighted by the charge of the particle under the asymmetry. For example, consider the MSSM in thermal equilibrium above the scale of the EWPT. We can write the asymmetries as follows (keeping in mind that we sum over all the degrees of freedom with the particular charge and making the simplifying assumption that all generations have the same chemical potential):

$$B = 6\mu_q - 3\mu_{u^c} - 3\mu_{d^c} \tag{20}$$
$$L = 6\mu_\ell - 3\mu_{e^c} \tag{21}$$
$$Q_Y = 3\mu_q - 6\mu_{u^c} + 3\mu_{d^c} - 3\mu_\ell + 3\mu_{e^c} + \mu_{H_u} - \mu_{H_d}. \tag{22}$$

The goal is to then use the conservation of chemical potential in all the non-gauge interactions of the MSSM, the conservation of chemical potential in sphaleron processes, and the fact that $U(1)_Y$ is conserved in all processes above the weak scale giving $Q_Y = 0$ to write all of these asymmetries in terms of a single chemical potential, say μ_q. We get the following relations from the MSSM superpotential and the sphaleron transitions:

$$\mu_{u^c} + \mu_q + \mu_{H_u} = 0 \tag{23}$$

$$\mu_{d^c} + \mu_q + \mu_{H_d} = 0 \tag{24}$$

$$\mu_{e^c} + \mu_\ell + \mu_{H_d} = 0 \tag{25}$$

$$\mu_{H_d} + \mu_{H_u} = 0 \tag{26}$$

$$3\mu_q + \mu_\ell = 0 \tag{27}$$

We see that we now have 6 equations (including the fact that $Q_Y = 0$) in 7 variables, so that we can eliminate 6 of the chemical potentials as promised. The remaining chemical potential can only be set by understanding baryogenesis. However, this is enough to relate the various asymmetries to one another. For example, we can determine that

$$\frac{B}{B-L} \approx 0.31, \tag{28}$$

which is a result that we will use later.

IV. THE "REFERENCE" MODEL

Now, we return to the idea of asymmetric dark matter. Let's briefly recall the plan: add a new particle to the spectrum which is charged under B or L and is therefore stable. It's number density is tied to the number density of baryons since the dark matter-anti dark matter asymmetry is communicated to the dark matter by the baryon sector. This gives a scenario where dark matter has mass of a few GeV and a density about a few times the density of baryons. We would like to construct some specific models employing these dynamics. We start with a reference model given by [1] for which we demonstrate in detail the mechanisms at work and then give brief overviews of other possibilities.

The reference model is an extension of the MSSM to include a new pair of gauge singlet chiral fermions X and $\overline{X}$. We charge these fields under baryon number with $L = \pm\frac{1}{2}$. The superpotential is then modified to include the following terms up to dimension 5:

$$\Delta W_{\text{eff}} = m_X \overline{X} X + \frac{1}{M_i}\overline{X}^2 L H_i + \overline{1}\Lambda X^4 + \overline{\overline{1}\Lambda}^4\overline{X}^4. \tag{29}$$

Mass terms of the form X^2 and $\overline{X}^2$ are forbidden even in the presence of L-violating neutrino masses due to an unbroken Z_4 subgroup of $U(1)_L$: the Majorana neutrino mass breaks L with $\Delta L = 2$, while the X's have $L = \pm\frac{1}{2}$. Another way to think about this is that the only L-violating spurion we can write has L charge 2. That allows us to see why the additonal quartic superpotential terms should be allowed. We need these L-violating interactions to drop out of equilibrium before the ones that of the form $\overline{X}^2 L H_i$ do so that the dominant L asymmetry is tied to the SM asymmetry.

The paper [1] discusses several aspects about the implementation of this model, including possible UV completions. However, in the interest of time, I will focus on the cosmology and phenomenology. Suppose at some high temperature in the early universe baryon asymmetry is transferred to the SM via one of the many mechanisms of baryogenesis. The SM speaks to the X sector via the dimension 5 coupling between L and $\overline{X}^2$ in the superpotential. Thus, in equilibrium, there will be some dark matter asymmetry which will be related to the $B-L$ standard model asymmetry by some order 1 number that we will calculate shortly. Once these interactions fall out of equilibrium, the dark matter asymmetry gets frozen in. The value to which the asymmetry freezes in depends on whether it happens before or after the EW phase transition, at which the L and H_u fields become non-relativistic and EWSB sets it. So we would like to figure out when the interactions go out of equilibrium above the EW scale ~ 100 GeV. Applying (13), we find that this occurs when

$$\Lambda \gtrsim \sqrt{\frac{M_{\text{pl}} T}{16\pi}} \gtrsim 10^9 \text{ GeV}. \tag{30}$$

For $\Lambda \ll 10^9$ GeV and $M_X \ll T \ll 100$ GeV, there is a small, but interesting region of parameter space. In this part of parameter space, there are two types of interactions that are relevant. The first are sneutrino decays to dark matter. These decays go out of equilibrium a little below $T \sim m_{\tilde{\nu}}$. There are also scattering processes between light particles obtained by integrating out the heavy weak scale particles. The dominant one is the dimension 9 operator

$$\mathcal{L}_{\rm eff} \supset \frac{v_u^2}{\Lambda^2 m_{\tilde{nu}}^4 m_{\tilde{B}}} (\overline{X}X)^2 \nu\nu. \tag{31}$$

The rate for this interaction, including non-trivial phase space factors and setting $m = m_{\tilde{\nu}} = m_{\tilde{B}}$ at the soft SUSY scale, is

$$\Gamma \sim \frac{1}{16\pi}\left(\frac{1}{8\pi^2}\right)^2 \left(\frac{v_u^2}{M^2 m^5}\right)^2 T^{11}, \tag{32}$$

giving a corresponding temperature at which these interactions go out of equilibrium of

$$T \lesssim 20 \text{ GeV} \left(\frac{M}{\text{TeV}}\right)^{4/9} \left(\frac{m}{100 \text{ GeV}}\right)^{10/9}, \tag{33}$$

which is comparable to the temperature at which sneutrino decays go out of equilibrium. Of course, pushing M so low we begin to worry about contributions to processes like $\mu \to e\gamma$, so we have to worry about the flavor physics of these interactions and tune them accordingly.

Next, we would like to determine X, the charge asymmetry of X and $\overline{X}$. All the machinery is in place from Section 3. We need only add the additional relations that

$$2\mu_{\overline{X}} + \mu_{H_u} + \mu_\ell = 0 \tag{34}$$
$$\mu_X + \mu_{\overline{X}} = 0 \tag{35}$$

Putting all these pieces together and solving a set of linear equations, we get relations (2.9) and (2.10) from [1], which immediately yield $M_X \approx 11$ GeV. (I'm a little confused as to why we don't include the lepton charge of the X particles when writing down the asymmetry relations, but to get reproduce results of Kaplan et. al., you must assume that the lepton asymmetry relation is unmodified by X). This is somewhat larger than 5 GeV because the conserved charge density is proportional to the number of degrees of freedom carrying the charge, so the baryon number density is somewhat larger than the X number density. A similar analysis can be done below the EWPT giving a mass of $M_X \approx 13$ GeV.

The last thing to worry about, ironically, is that we want the X and $\overline{X}$ to annihilate efficiently enough that the residual dark matter density is almost completely determined by the asymmetry. The way our model has been written down so far, there is actually no way for the X and $\overline{X}$ to annihilate since they are neutral under all SM gauge groups. There is a natural solution to this problem that also solves the problem of the light mass of the X. We simply extend our model to the MSSM and have the X mass determined by the VEV $s/\sqrt{2}$ of a scalar S.

We need a to be light, but this can be achieved naturally in the NMSSM. Recall that the superpotential of the Higgs sector of the NMSSM, adding in a term to generate the mass of the X, is given by

$$W = \lambda S H_u H_d + \frac{\kappa}{3} S^3 + \lambda' S X \overline{X}. \tag{36}$$

This clearly has an R-symmetry when all the fields get the same charge 2/3 (W should have R charge 2). Of course, the VEVs of the fields spontaneously break this symmetry. There would be a massless Goldstone boson except that the SUSY-breaking sector also explicitly breaks the R-symmetry. The soft terms for the Higgs sector are

$$V = \sum_{\rm i} m_i^2 |\phi_i|^2 + A_\lambda S H_u H_d + A_\kappa S^3 + A_{\lambda'} S X \overline{X} + \text{h.c.}. \tag{37}$$

The scalars clearly respect the R-symmetry, but the A terms break it. However, we know of many scenarios in which the A terms are small. In that case, the lightest pseudo-scalar of the Higgs sector is a light pseudo Nambu-Goldstone boson. This boson couples to $X\overline{X}$ after spontaneous symmetry breaking via

$$\mathcal{L} \supset m_X \overline{X} X e^{ia/s} + \text{h.c.}. \tag{38}$$

At relativistic energies (but below the weak scale), the annihilation would go out of equilibrium for

$$T \lesssim \frac{16\pi s^4}{M_X^2 M_{\rm pl}} \sim 10^{-12} \text{ GeV}, \tag{39}$$

for $M_X \sim 10$ GeV, $s \sim 100$ GeV. So it will remain in equilibrium throughout the relativistic phase. When $T \lesssim M_X$, this decay should go out of equilibrium quickly, but it will freeze out like a normal WIMP, except with a density suppressed by a factor of about $(s/M_X)^2 \sim 100$. Thus, some of the dark matter density would be due to a symmetric density, but it would be about 100 times less the asymmetric density.

This is not the only possibility, but it is the most immediately satisfying, even though it messes with perturbative unification. The other possibilities mentioned all have their own issues. One possibility is that the singlet S couples only weakly to the Higgs. This does not explain the mass of the Higgs. It also destroys any direct detection or collider signals. Another is to have a light enough UV completion that allows the X to annihilate. However, this possibility tends to run into flavor problems, such as large $\mu \to e\gamma$.

The upshot is that even within this reference model, there is a wide range of phenomenological possibilities, as well as issues. The basic model is extremely simple; there is only one new term in the superpotential. However, it likely requires an extension beyond the MSSM and additional assumptions about the size of the scales in the theory. It would also be nice if this model could explain the series of direct detection signals that have been observed recently [5–7]. Of course, the model is flexible enough to not be excluded if these results go away: it is easy to dial down the cross-section further. That means that it is also possible for such a model to evade obvious collider or astrophysical signatures. We discuss potential signatures of the model shortly.

V. OTHER MODELS

In this short section, we discuss a few other possible models proposed by [1]. One similar possibility is to introduce fields X and $\overline{X}$ as before, but this time charge the fields under baryon number with charge $\pm 1/2$. The operator that transfers the asymmetry is then even mor suppressed:

$$W \supset \frac{1}{\Lambda^2}\overline{X}^2 udd \tag{40}$$

(Is this operator even possible? To write something gauge invariant, don't we need to write say $X^2 u^c d^c d^c$ instead?) It gives the same mass $M_X \approx 11$ GeV for the dark matter particle as the lepton number model. One problem with this scenario is that the operator is more suppressed than in the lepton case. We need to push Λ down as low as or even lower than TeV.

Another possibility is to charge the X and $\overline{X}$ with charge ± 1 under lepton number. In this case, X behaves like a sterile neutrino. In order to avoid getting a large Majorana mass of the form X^2 or $\overline{X}^2$, we need to assume Dirac neutrino masses. Then, the operator transmitting the asymmetry is

$$\mathcal{L} \supset \frac{1}{\Lambda^4}\overline{X}^2 (LH)(LH) + \text{h.c.}. \tag{41}$$

A similar calculation to that performed in the previous section goes through and we find that $M_X \approx 6$ GeV. It is most interesting to consider a UV completion of this scenario. We add an addition scalar doublet H' that does *not* get a VEV. The relevant part of the Lagrangian is then

$$\mathcal{L} \supset y' L_i H' \overline{X} - \frac{\lambda'}{4}[(H^\dagger H')^2 + \text{h.c.}]. \tag{42}$$

The low energy coupling is then given by $y'y'\lambda'/m_{H'}^4$. Note that if the new scalar got a VEV, then $\overline{X}$ would mix with neutrinos and would decay. Then come the slew of constraints based on cosmological observations. To annihilate the symmetric component efficiently, H' must be fairly light, less than around $y' \times 200$ GeV. We must also have $\lambda' \ll 1$ to decouple the transfer of the asymmetry above M_X. Lepton flavor violation is also an issue here and the flavor structure must be non-trivial. Lastly, the model should probably be embedded in supersymmetry to allow for all the light scales involved. The supersymmetric version of this model is rich with phenomenology, including an LSP in addition to X.

As discussed in Section 1, the main idea here was to see what sort of issues we want to think about after writing down a model of ADM. I present an incomplete list:

- Keeping X stable
- Stabilizing the various scales (often involving the complications of an NMSSM);
- Suppressing "bad" flavor violation (i.e. $\mu \to e\gamma$);

- Annihilating the symmetric component of the dark matter density efficiently;
- Preserving some kind of R-symmetry in SUSY versions;
- Ensuring that the $B-L$ asymmetry transfer to the dark sector freezes out before the X goes non-relativistic;
- Keeping the effects on other cosmological observables like BBN and matter-radiation equality within bounds;
- Ideally, ensuring that there are signals for direct detection and/or collider experiments.

VI. SIGNATURES OF ADM

Lastly, we consider the possibility of discovering ADM. There is a wide range of possibilities for phenomenology in these models because there are very weak requirements on their interactions with the SM. The major constraint is that the symmetric component must annihilate away sufficiently that the residual density is determined by the asymmetry and not by the freeze-out relic density. That is, we do *not* want a WIMP! If this annihilation is to SM particles, then it is possible to have both direct detection and collider signatures in the usual WIMP way.

There are two main possibilities for annihilation to the SM. The first, as discussed above, is that the dark matter annihilates to the lightest pseudoscalar Higgs. The dark matter can then interact with nuclei by t channel Higgs exchange. The other, which I only mentioned very briefly, involves the UV completion of the model. The class of UV completions of interest involves the addition of new elecroweak doublets to the model. More concretely, for the reference model we add a vector-like pair of electroweak doublets D, $\overline{D}$ at the weak scale. These would be just like adding another pair of Higgs doublets, except that they do not get a VEV and they have $L = \pm 1/2$. The superpotential is then

$$W \supset M\overline{D}D + \lambda' \overline{X} D H_u + y' L \overline{D}\overline{X}. \tag{43}$$

Then, the dark matter could annihilate by t-channel exchange of a D to leptons or neutrinos. The short summary is that the first scenario generally gives direct detection cross sections orders of magnitude below current limits, while the second scenario gives cross sections closer to current bounds. We now estimate these cross sections more explicitly.

The coefficient for t-channel exchange of a heavy scalar particle between X and a q can be estimated as $a_q \sim g_X g_q / M_H^2$ where M_H is the mass of the heavy particle. This is related to the cross-section for dark matter to scatter off a nucleon (which is what experimentalists generally quote) by some complicated form factors [8]:

$$\sigma(Xn \to Xn) = \frac{4m_r^2}{\pi} \left(\sum_q f_{Tq} a_q \frac{m_p}{m_q} \right)^2, \tag{44}$$

where I make the good approximation that scattering off protons and neutrons is the same, the sum is over light quarks, f_{Tq} are some form factors obtained by integrating over the speeds of incoming dark matter particles, and m_r is the reduced mass of the nucleon and dark matter. This is dominated by coupling to the strange quark so we make a further approximations and say

$$\sigma(Xn \to Xn) \sim \frac{4}{\pi} f_{Ts}^2 \frac{a_s^2 m_p^4}{m_s^2}. \tag{45}$$

Numerically, $f_{Ts}^2 \sim 0.1$. Furthermore, $a_s = g_X g_S / M_H^2$. Thus, we get the final approximation that

$$\sigma(Xn \to Xn) \sim \frac{0.04}{\pi} \frac{g_X^2 g_s^2 m_p^4}{m_s^2 m_H^4}. \tag{46}$$

For the case where the dark matter annihilates to the lightest Higgs scalar, the coupling $g_X \sim \sqrt{2} m_X / s \sin\theta$, where s is the VEV of the NMSSM Higgs S and θ is the mixing angle between this Higgs and the MSSM down Higgs. The coupling $g_s = \sqrt{2} m_s / v_d$ as usual. For $g_X \sim 0.1$ and $g_s \sim 0.01$, this gives a cross-section of

$$\sigma(Xn \to Xn) \sim 10^{-44} \mathrm{cm}^2 \left(\frac{m_{h^0}}{100 \text{ GeV}} \right)^4. \tag{47}$$

A more careful calculation finds a result slightly smaller than this. This is to be compared to current bounds which sit around 10^{-43} cm^2 for a particle of mass around 10 GeV. The cross-sections are predicted are typically a couple of order of magnitudes lower. However, they do depend on several parameters, including the mixing angles in the Higgs sector, the VEV of S, and the mass of the lightest Higgs scalar. There is also a spin-dependent interaction due to pseudo-scalar exchange, but it is well below current bounds.

The second case is trickier. The dominant coupling is through a one-loop magnetic moment and charge radius. Unfortunately, I haven't had time to look through the calculation in detail. However, the result is a cross-section below current bounds, but not by as much.

Now, we move on to discuss some collider signatures. The most striking would be new and significant decay modes for Higgs boson, modified kinematics and lepton counting SUSY cascades, and new states charged under SM groups beyond the superpartners within reach of the LHC.

In the Higgs sector, the Higgs can couple to dark matter through mixing between the MSSM CP-even Higgs and the singlet Higgs S. Bounds from requiring enough X annihilation and placing the model above chargino mass bounds restrict this mixing and place the model in a region where Higgs decays to X are competitive with decays to b quarks. The Higgs can also have competitive decays to the light pseudoscalar, which further decays to b's or τ's giving a signature $h \to 4b$ or $h \to 4\tau$.

The collider phenomenology of SUSY decays is also modified in significant ways. The would-be LSP (say a neutralino) would no longer be exactly stable, since dark matter would be charged under, for example, a discrete Z_4 symmetry. This means that the lightest super-partner could decay to pairs of dark matter particles. The challenge is to then extract the dark matter mass from these decays. The decays are also highly suppressed by higher scales and by the fact that they are generally many-body decays. Therefore, there is the possibility that the vertex of the lightest superpartner (henceforth referred to as the NLSP) is displaced.

In the standard scenario with a neutral NLSP, the results could actually be very boring. The dominant decay to dark matter and neutrinos, which would be completely invisible. If the NLSP is sufficiently heavy, it could also decay to a Higgs and invisible particles. If the NLSP is either a slepton or a squark, it will typically decay to its standard model partner and invisible particles. There can be displaced vertices in this case of order mm to cm. In $B = 1/2$ models, it is possible to get spectacular displaced vertices with multiple jets emerging also at mm to cm. In this case, if there is a slepton NLSP, then the displaced vertices can be on even longer distances, around 1 m or longer.

All of the estimates of collider signatures are highly model and spectrum dependent. The point is to say that they can significantly modify collider signatures from the usual MSSM paradigm and can in some scenarios be observed at the LHC.

Finally, to conclude, we mention an incomplete list of follow up work to [1]. We simply mention the words to entice the reader to learn more and to get an idea of the directions that have been explored:

- Asymmetric freeze-in of FIMPs [9];
- Baryogenesis via the dark sector [10–12];
- Anti-baryonic dark matter [13];
- Explaining the lightness of dark matter [14];
- Determining indirect detection signatures [15];
- Quirky/Bound-state dark matter [16–19];
- Leptophilic dark matter [20];
- Discovering ADM through anti-neutrinos [21].

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