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ABSTRACT ON COSMOPARTICLE PHYSICS

MECHANISMS OF DARK MATTER GENESIS

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INTRODUCTION

The generation of DM in the early universe can proceed via thermal or non-thermal production, or both, or it may result from a particle-antiparticle asymmetry. None of the mechanisms proposed below explains the approximate relation

$$\rho_{B,0} = \rho_{DM,0}, \tag{1}$$

where $\rho_{B,0}$ and $\rho_{DM,0}$ energy density of baryons and dark matter particles in the modern Universe. This approximate equality was also fulfilled in the past, at rather late stages of the expansion of the Universe. Several possible mechanisms for the generation of dark matter and baryon asymmetry have been proposed in the literature, leading to the relation 1.1. However, a convincing and natural explanation for the approximate coincidence has not yet been found. Perhaps this is indeed a coincidence.

1. GENESIS OF DARK MATTER

Freeze-out: The process of chemical decoupling from the high-temperature, high-density thermal bath (freeze-out) as a paradigm for particle production in the early universe is both a predictive and a successful one. The possibility that just like light elements, neutrinos, and CMB photons, particle DM also originated from a thermal decoupling process has thus garnered significant attention.

A particle species chemically decouples when the rate for the species' number-changing processes drops below the Hubble rate H. Rough estimates for the abundance of relics can be obtained by calculating the freeze-out (i.e. "decoupling") temperature $T_{f.o.}$, corresponding to $H(T_{f.o.}) \sim (T_{f.o.})$, equating the comoving number density at freeze-out and today, eventually obtaining the physical density of relic particles today. This procedure assumes that entropy is conserved between $T_{f.o.}$ and today, an assumption that could well be violated, especially for heavy relics that decouple early, for instance by entropy injection episodes. Notice also that the freeze-out calculation strongly depends on the assumed background cosmology, and changes e.g. if the early universe is not radiation-dominated around DM decoupling.

The calculation of the freeze-out relic abundance hinges on a Boltzmann equation relating the Liouville operator to the collision operator acting on the phase space density. Under a variety of simplifying assumptions including homogeneity and isotropy, it is possible to reduce the relevant equation for the number density n of a single species pair-annihilating with particles in the thermal bath via 2-to-2 processes to

$$\frac{dn}{dt} - 3Hn = -\langle \sigma v \rangle (n^2 - n_{eq}^2), \tag{1.1}$$

where $\langle \sigma v \rangle$ is the thermally-averaged pair-annihilation cross section times relative velocity, and n_{eq} is the equilibrium number density. Relics for which the freeze-out temperature is much larger than the particle mass (and thus that freeze-out as ultra-reltivistic) are called hot relics; if the opposite is true, the relic is instead considered cold.

A straightforward calculation shows that to leading order the frozen-out density of hot relics is linearly proportional to the relic particle mass. The comoving number density $Y = n/s$, where s is the entropy density, for a hot relic is approximately given by its equilibrium value,

$$Y_{f.o.} \simeq Y_{eq} \simeq 0.278 \frac{g_{eff}}{g_{*s}}, \tag{1.2}$$

where g_{eff} is the relic's effective number of degrees of freedom, and g_{*s} is the number of entropic relativistic degrees of freedom, both calculated at $T_{f.o.}$. The resulting relic abundance, assuming an iso-entropic expansion, is

$$\Omega_{hot} h^2 = \frac{m Y_{f.o.} s_0 h^2}{\rho_0} \simeq \frac{m}{93\, eV}, \tag{1.3}$$

with s_0 the entropy density today, and with the latter equality holding for the case of SM neutrinos, with a freeze-out temperature around 1 MeV (which enters in the final relic abundance through the degrees of freedom dependence on the right-hand-side of Eq. (1.3)).

For cold relics, the leading-order dependence of the relic abundance on the DM particle properties is an inverse proportionality relation to the pair-annihilation cross section,

$$\Omega_{cold} h^2 \simeq 0.1 \left(\frac{x_{f.o.}}{20}\right) \left(\frac{10^{-8} GeV^{-2}}{\sigma_{DM+DM \leftrightarrow DM}}\right), \tag{1.4}$$

where $x \equiv m_{DM}/T$. In turn, the freeze-out temperature is approximately given by the solution to the equation

$$\sqrt{x} \cdot e^{-x} = (m_{DM} \cdot M_P \cdot \sigma_{DM+DM \leftrightarrow DM})^{-1}, \tag{1.5}$$

where $M_P \simeq 2.43510^{18}\, GeV$ is the reduced Planck mass. As a result, $T_{f.o.} \simeq m_{DM}/x_{f.o.}$, with $x_{f.o.}$ a number between 10 and 50, depending on the cross section, with only a logarithmic dependence on the DM mass. Since for electroweak-scale cross sections and masses $\sigma_{DM+DM} \simeq 10^{-8}\, GeV^{-2}$, "weakly-interacting massive particles", or WIMPs have gained exceptional popularity. Notice that Eq. (1.4) bears, however, no connection to the weak scale, despite the relation being known as "WIMP miracle".

Numerous scenarios exist, including notably supersymmetry and models with universal extra dimensions where the relic abundance of the DM is controlled by processes involving a slightly heavier, unstable, co-annihilating species. In this case the calculation of the abundance of the stable species proceeds similarly to what outlined above, with an effective pair-annihilation cross section that captures the effects of co-annihilation replacing the pair-annihilation cross section.

Freeze-in: Collisional processes can lead to the production of out-of-equilibrium particles that progressively accumulate over cosmic time, a process sometimes called freeze-in. The abundance of the frozen-in particles produced at a given redshift depends on the product of the production rate times the Hubble time at that redshift. Freeze-in generally implies that the lightest observablesector particles decay to the DM with relatively long lifetimes, giving peculiar signals at colliders. Gravitinos are an example of DM candidates possibly produced via a freeze-in type scenario, albeit the portal coupling is in that case via a higher dimensional, Planck-suppressed operator.

Cannibalization and other dark-sector number-changing processes: Thermal processes can drive the abundance of the DM beyond simple 2-to-2 number-changing interactions. For instance, DM can "cannibalize" itself if n $\rightarrow$ 2 processes exist. In this case, a critical aspect is whether or not the DM sector is in thermal contact with the Standard Model thermal bath. If it is, n $\rightarrow$ 2 processes can drive the relic abundance, e.g. in the Strongly Interacting Massive Particles (SIMP) scenario. Models exist where the kinetic decoupling (i.e. the decoupling from the thermal equilibrium velocity distribution) of the two sectors drives the abundance of the DM (elastically decoupling relics, or ELDERs). When the two sectors are not in thermal contact, n $\rightarrow$ 2 processes heat the DM sector dramatically, rapidly affecting the temperature ratio between the visible and dark sectors. If the relevant cross sections are large enough, and the DM mass light enough, significant effects can arise in structure formation.

Non-thermal production: DM production can proceed via processes out of thermal equilibrium ("non-thermal" production). These include DM production via the decay of a "mother" particle (or of topological defects[5], moduli etc.) to the DM, or production via gravitational effects.

Asymmetric DM: An enticing alternative possibility for DM production is that of asymmetric DM: the relic DM abundance arises from an asymmetry between anti-DM and DM. This asymmetry may or may not be related to the baryon-antibaryon asymmetry. If it is, then depending on the model and its thermal history, a relation exists between the mass of the DM and the proton mass. A variety of proposals have been put forward where alternately baryogenesis is explained from a DM sector asymmetry, or vice-versa.

Primordial Black Holes production: A qualitatively stand-alone class of DM candidates, primordial black holes (PBHs), arises from entirely different mechanisms from what reviewed above. PBHs are thought to originate from gravitational collapse of large density fluctuations in the early universe.[4] The over-densities could be produced in a variety of ways, such as topological defects like cosmic strings, necklaces or domain walls, curvature fluctuations from a period of ultra-slow-roll, a sound speed "resonance" , an early phase of matter domination, or subhorizon phenomena including a phase transition and preheating.Albeit the calculation depends on the details of gravitational collapse, the formation time is connected to the PBH mass via $M = \gamma M_{PBH} \simeq 2 \cdot 10^5 \gamma M/odot$ with $\gamma \simeq (1/\sqrt{3})^3$ during radiation domination.[1]

2. FREEZE-IN DURING AN EARLY MD ERA

Dark matter, X, may be generated by new physics at the TeV scale during an early matter-dominated (MD) era that ends at temperature $T_R \ll$ TeV. Compared to the conventional radiation-dominated (RD) results, yields from both Freeze-Out and Freeze-In processes are greatly suppressed by dilution from entropy production, making Freeze-Out less plausible while allowing successful Freeze-In with a much larger coupling strength. Freeze-In is typically dominated by the decay of a particle B of the thermal bath, B $\rightarrow$ X

When inflation ended, provided the inflaton decays were not extremely rapid, there was an era of matter domination (MD) which ended at the reheat temperature T_R. It is commonly assumed that T_R is very high, many orders of magnitude above the TeV scale, but observationally the most stringent constraints are from the effective number of neutrino species, $T_R > 4$ MeV. This early MD era could extend for many decades in temperature above T_R, and include the TeV epoch. There are several alternative origins for a long early MD era, including long-lived heavy particles that were once in thermal equilibrium and oscillating fields composed of light bosons. When the MD era results from inflation it has an evolution that is purely Non-Adiabatic in character, giving a MD_{NA} era. The more general MD era splits into two, beginning with Adiabatic evolution, MD_A, and ending with MD_{NA}, as illustrated in Fig. 2.1.

Particle dark matter requires an addition to the Standard Model with a cosmologically stable particle X of mass m_X. How is the abundance of X determined? The Freeze-Out mechanism results if X has sufficient interactions with the known particles that it is in thermal equilibrium at temperatures T of order m_X. As T drops below m_X, X tracks a Boltzmann distribution for a while but, as it becomes more dilute, its annihilation rate drops below the expansion rate and it freezes out of thermal equilibrium. This mechanism has great generality, applying to a wide range of theories where the interactions of X are sufficient to put it in thermal equilibrium at T $\sim m_X$. Furthermore, it

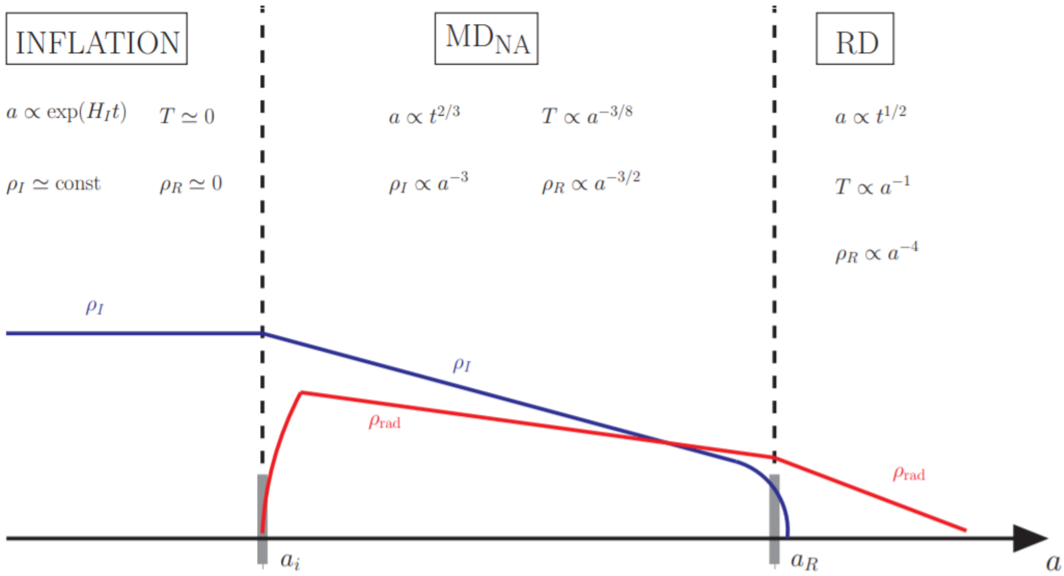

Figure 2.1: A long MD era after inflation

is highly predictive since it is an IR dominated process - it does not depend on any physics at energy scales above its mass. Indeed, Freeze-Out during the RD era suggests m_X is broadly of order a TeV and that X has an interaction rate that can lead to signals at both collider detectors and direct and indirect detection experiments.

Freeze-In provides another general production mechanism and results when the interactions of X are insufficient to bring it into thermal equilibrium. At temperatures T above m_X these feeble interactions allow the production of X from decays or scatterings of some bath particle B of mass m_B at rate $\Gamma(T)$. This produces a yield of X particles at time t(T) of

$$Y_X^{Prod}(T) \sim \Gamma(T)t(T) \tag{2.1}$$

which is IR dominated when the interactions between B and X are of dimension 4 or less. Taking B to be the lightest observable sector particle carrying the stabilizing symmetry, the production rate becomes Boltzmann suppressed below m_B, so that the dominant contribution to Y_X arises at $\sim m_B$. For T $> m_B$, the yield Y(X) grows towards its equilibrium value, but it never reaches equilibrium, and this Freeze-In towards equilibrium stops once T drops below m_B. Decays generally dominate over scatterings, so that in this paper we study the reaction

$$B \to A_{SM} X \tag{2.2}$$

where A_{SM} is one or more Standard Model particles.

It is assumed that Freeze-In (FI) production of dark matter during the early MD era via the decay B $\to A_{SM}X$. B is the lightest observable sector particle that carries the stabilizing symmetry, X is the dark matter and A_{SM} is one or more Standard Model particles. In addition to FI, there is a Freeze-Out (FO) population of B that eventually decays to X.

The number density evolution of X is described by the Boltzmann equation

$$\frac{dn_X}{dt} + 3Hn_x = \Gamma_B n_B^{eq} \frac{K_1[m_B/T]}{K_2[m_B/T]}, \tag{2.3}$$

with Γ_B the width of B and $K_{1,2}[x]$ the first and second modified Bessel functions of the 2nd kind. The equilibrium number density obtained using Maxwell-Boltzmann statistics reads

$$n_B^{eq} = \frac{g_B}{2\pi^2} m_B^2 T K_2[m_B/T]. \tag{2.4}$$

At high temperatures (T $\gg m_B$) we recover the T^3 abundance for a relativistic species, whereas at low temperatures (T $\ll m_B$) the number density has the Maxwell-Boltzmann exponential suppression. For this reason, the FI production of X is dominated at temperatures $T_{FI} \simeq m_B$.

If FI occurs during the RD era, the Boltzmann equation (2.1) can be easily solved giving a final yield for X

$$Y_X = \frac{n_X}{s} = 4.4 \cdot 10^{-12} \left(\frac{g_B}{2}\right) \left(\frac{106.75}{g_*}\right)^{3/2} \left(\frac{300GeV}{m_B}\right) \left(\frac{\Gamma_B/m_B}{1.8 \cdot 10^{-25}}\right) \tag{2.5}$$

where s is the entropy density. Observations fix the DM density but the yield is fixed once the mass of X is known using

$$\xi_{DM} = \frac{\rho_{DM}}{s} = m_{DM} Y_{DM} = 0.44\, eV \tag{2.6}$$

which is close to the usual temperature of matter radiation equality, $T_{eq} \simeq 1\, eV$. We can thus rewrite eq. (2.5) as

$$\xi_X = \xi_{DM} \cdot \left(\frac{g_B}{2}\right)\left(\frac{106.75}{g_*}\right)^{3/2}\left(\frac{m_X}{100GeV}\right)\left(\frac{300GeV}{m_B}\right)\left(\frac{\Gamma_B/m_B}{1.8\cdot 10^{-25}}\right) \tag{2.7}$$

The coupling λ, defined by

$$\Gamma_B = \frac{\lambda^2}{8\pi} m_B, \tag{2.8}$$

must be very small to avoid overclosure. For the reference masses and number of spin states shown in (2.5) the observed DM density results for $\lambda \simeq 2 \cdot 10^{-12}$.

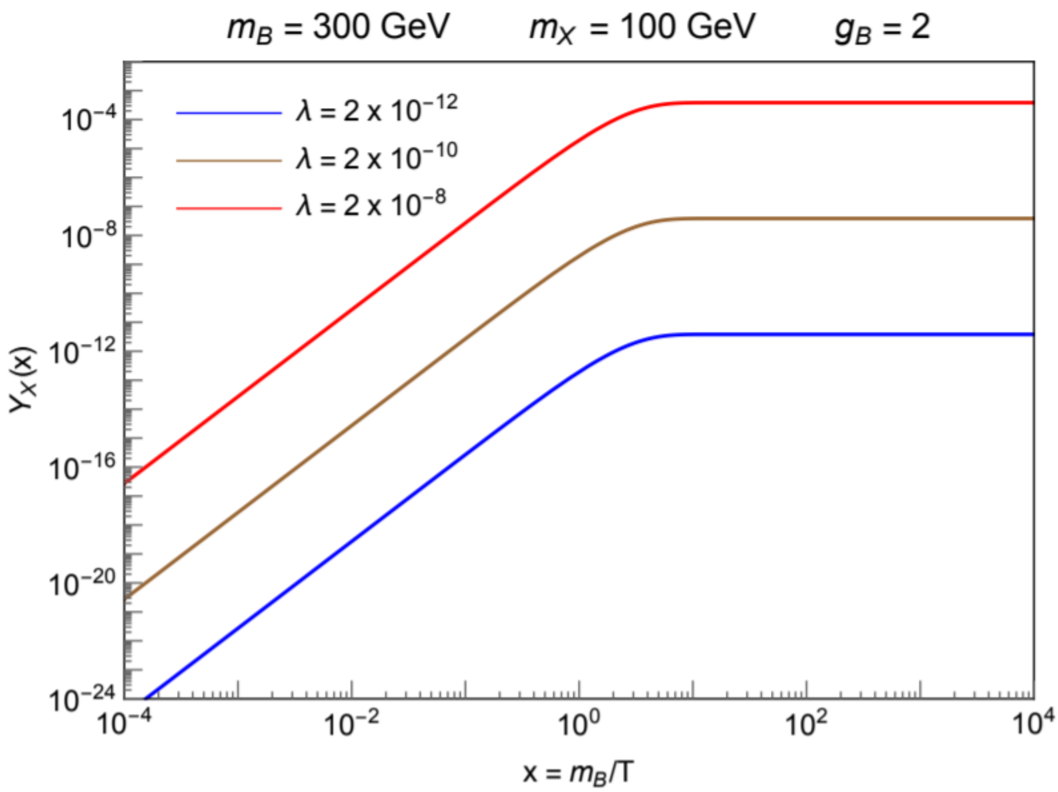

Figure 2.2: Yields $Y_X(x)$ for FI during the RD era for three values of the decay width, as defined in Eq. (2.8), with fixed reference values m_B = 300 GeV, m_X = 100 GeV, and g_B = 2.

In Fig.2.2 the FI solutions for the DM yield $Y_X(x)$ are shown, with time variable $x = m_B/T$. There are three different lines corresponding to three different values of λ (namely three different decay widths as in Eq. (2.8)). In all cases we see that the process is maximally active at temperatures of order

m_B. For the reference parameters of (2.5), the blue line reproduces the observed DM density, whereas the brown and red lines overproduce the amount of dark matter. In particular, the red line reaches an asymptotic value of Y_X which is 10 % the equilibrium value Y_X^{eq}, and above this point the approximation that DM arises from FI through the decay in Eq. (2.2), and without the inverse reaction, breaks down.

2.1. FREEZE-IN DURING NON-ADIABATIC EVOLUTION

Freeze-In during MD_{NA} occurs between a_{NA} and a_R, and the final DM density depends on only one parameter of the background, the reheat temperature T_R. This includes the important example of FI during reheating after inflation, as sketched in Fig. 2.1. In this case inflation could be the conventional one with more than 60 e-foldings with the matter M identified as the usual inflaton, or it could be a later era of inflation with fewer e-foldings.

The FI yields $Y_X(x)$ are shown in Fig.2.3 for three different values of the decay width

$$\lambda = \begin{cases} 2 \cdot 10^{-12} & T_R = 1TeV \\ 2 \cdot 10^{-8} & T_R = 4.6TeV \\ 2 \cdot 10^{-7} & T_R = 2.4TeV \end{cases} \tag{2.9}$$

In each case we choose TR in such a way that we reproduce the observed DM abundance. The values of λ span a wider range than the case of Fig.2.2 since Freeze-In during a MD era allows for larger couplings. The blue line has $T_R > m_B$, so that FI occurs during the RD era. As λ is increased, the FI process becomes more powerful, as illustrated by the brown and red lines, and to obtain the observed DM abundance much lower values of T_R must been taken, so that the X abundance is diluted by entropy production after FI. However, it is not possible to arbitrarily increase λ and still be in the Freeze-In regime. For sufficiently large λ, the peak of the Y_X functions in Fig.2.3 will reach the equilibrium value and FI is no longer the applicable production mechanism. For reference values $m_B = 300$ GeV, $m_X = 100$ GeV, and $g_B = 2$) this gives

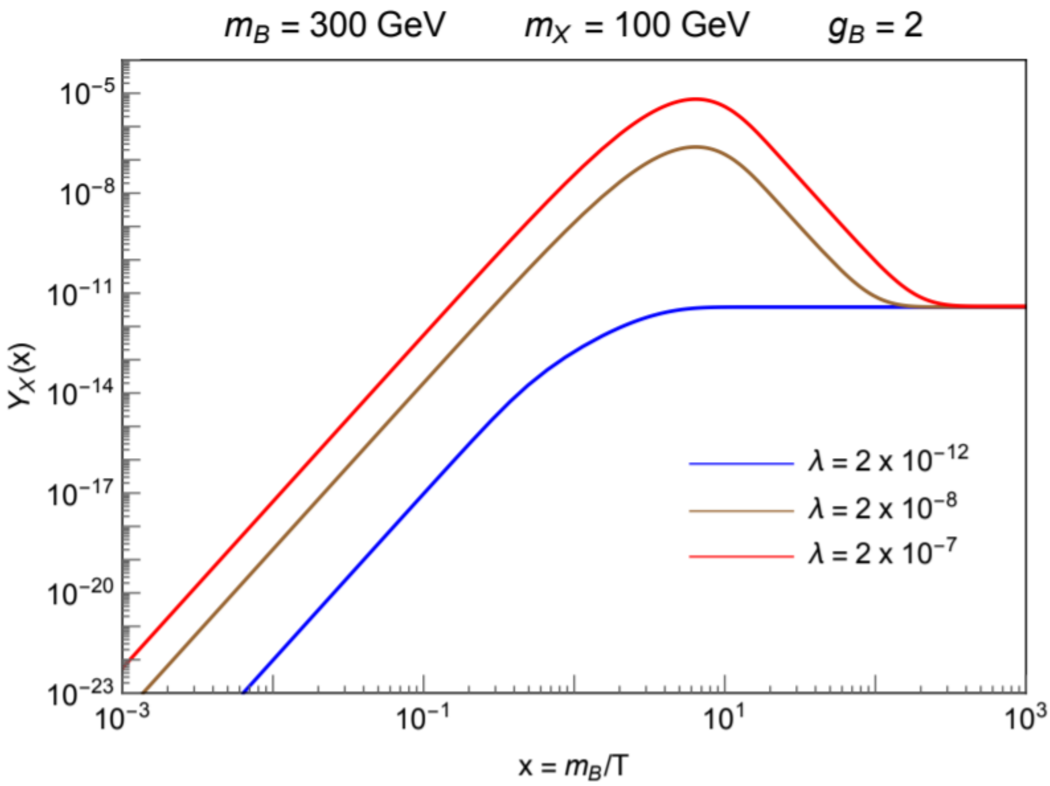

Figure 2.3: Yields $Y_X(x)$ for FI during the MD_{NA} era, for example during reheating after inflation. For each of three λ, T_R is chosen to give the observed DM abundance with fixed reference values $m_B = 300$ GeV, $m_X = 100$ GeV, and $g_B = 2$.

a lower limit $T_R \geq 0.6$ GeV. At the same time we have to keep $T_R \leq m_B$; otherwise FI occurs during the RD era and there would be no dilution.[2][10]

3. NON-THERMAL PRODUCTION

There is strong evidence for the existence of a substantial amount of cold dark matter (CDM). The leading candidates for CDM are weakly interacting massive particles (WIMPs), such as the neutralino. The neutralino is the lightest supersymmetric particle. In models with R parity it is stable, and its mass density in the universe is generally assumed to be a relic of an initially thermal distribution in the hot early universe. Assuminig, in addition, the presence of a small cosmological constant, the CDM scenario is consistent with both the observations of the large scale structure of the universe ($\gg$ 1Mpc) and the fluctuations of the cosmic microwave background. Many experiments searching for neutralino dark matter particles are under way.

The collisionless CDM scenario, however, predicts too much power on small scales, such as a large excess of dwarf galaxies, the over-concentration of dark matter in dwarf galaxies and in large galaxies. Recently Spergel and Steinhardt proposed a new concept of dark matter with strong self interaction. This puts WIMPs as candidates for dark matter in considerable jeopardy.

It is possible to propose a scenario with non-thermal production of WIMPs. These WIMPs could be relativistic when generated. Their comoving free-streaming scales could be as large as of the order 0.1 Mpc or larger. The density fluctuations on scales less than the free-streaming scale would then be severely suppressed. Consequently the discrepancies between the observations of dark matter halos on the sub-galactic scales and the predictions of the standard WIMPs dark matter picture could be resolved.

To begin with, supposed that a general case of non-thermal production of the neutralinos by the decay of topological defects such as cosmic string, by the decay of an unstable heavy particle, or produced non-thermally by the reheating process in a scenario of inflation at low energy scale. The momentum distribution function of the neutralinos is for simplicity assumed to be Gaussian:

$$f(p) = \frac{A}{\sqrt{2\pi}\sigma} exp\left(\frac{(p-p_c)^2}{2\sigma^2}\right), \tag{3.1}$$

where p_c is the central value and σ describes the width of the distribution.

Given a model, the parameters p_C and σ can be determined. For instance, in the supersymmetric version of the $U_{B-L}(1)$ model, the (Higgsino-like) neutralinos arise directly from the decay of the right-handed neutrinos and their superpartners. In this model, p_C is about a half of the mass of the mother particles. For a two-body decay, if the mother particle is at rest, the distribution function $f(p)$ is a δ-function. The value of σ characterizes the average non-vanishing velocity of the mother particles (when $\sigma \to 0$, $f(p)$ approaches a δ-function).

In Eq.(3.1), A is a normalization factor determined by the energy density of the non-thermal component

$$\rho(NT) = 4\pi \int E(p) f(p) p^2 dp, \tag{3.2}$$

where $E(p) = (p^2 + m^2)^{1/2}$ and m is the rest mass of the dark matter particle. Given that the physical momentum $p(t)$ scales as the inverse of the cosmic scale factor $a(t)$, also define $r = a(t)p(t)/m$. During cosmic evolution r is a constant. Throughout this paper we set a $(t_0) = 1$, so r can be understood as the velocity of the particles at the present time (note that the dark matter particles are non-relativistic now even though they are relativistic when generated).

The comoving free-streaming scale R_f for the non-thermal particles can be calculated as follows:

$$\begin{aligned} R_f &= \int_{t_i}^{t_{eq}} \frac{v(t')}{a(t')} dt' \simeq \int_{0}^{t_{eq}} \frac{v(t')}{a(t')} dt' \\ &\simeq 2 r_c t_{eq} (1 + z_{eq})^2 \ln \sqrt{1 + \frac{1}{r_c^2 (1 + z_{eq})^2}} + \frac{1}{r_c (1 + z_{eq})} \end{aligned} \tag{3.3}$$

where $r_c \equiv a(t) p_c(t)/m = p_c(t_0)/m$ and the subscript '$_{EQ}$' denotes radiation-matter equality. Below the free-streaming scale, the power spectrum will be severely damped. To account for the lack of substructure in the Local Group, N-body simulations study show that the free-streaming scale of the dark matter should be ≈ 0.1 Mpc. This corresponds to $r_c \approx 10^{-7}$, which gives rise to a constraint on the parameters of the model.

The model with non-thermal production of WIMPs provides a promising scenario for large-scale structure formation of the universe. However, consistency needs to be checked with observations on small scales, especially on scales of the Lyman-α forest. In Fig. 3.1 shown the power spectrum of this model and the observed power spectrum of the Lyman-α system at z = 2.5 shown as the filled circles with error bars. For comparison the power spectra for the conventional CDM and WDM models also given. In fitting to the observed data, the primordial spectral index n = 0.97 for all models was choosen. The mass for the WDM particles is choosen as 750 eV, and the parameters for the NTDM models are $r_c = (1.3, 1.4, 1.5) \cdot 10^{-7}$, respectively.

For a flat universe with $\Omega_\Lambda = 0.6$ the range of these scales corresponds to $0.4h^{-1}Mpc \leq k \leq 12.8h^{-1}Mpc$. From Fig. 3.1 one can see that the larger the value of r_c is, the lower the small-scale power spectrum becomes. By comparing the values of the power spectra at the upper limit of the above range in k with the power spectrum of the WDM model with $m_W = 750$ eV, an upper limit on r_c of $r_c \leq 1.5 \cdot 10^{-7}$ was obtained.

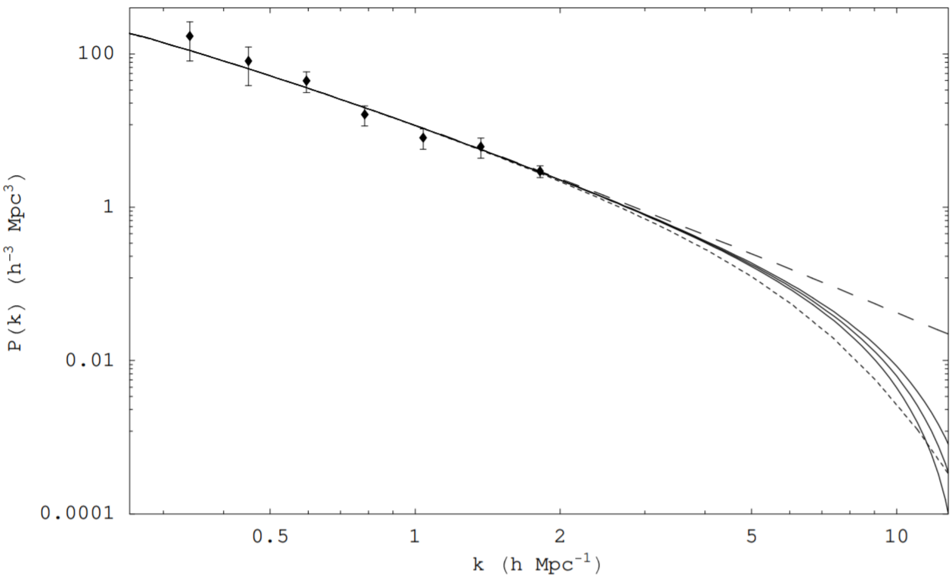

Figure 3.1: The power spectra of the CDM model (long dashed curve), the WDM model with $m_W = 750$ eV (short dashed curve) and the NTDM models with $r_c = (1.3, 1.4, 1.5) \cdot 10^{-7}$ (solid curves, from top down), compared to the observed lyman-α P(k) at z = 2.5 (filled circles with error bars).

There were studies another constraints. The particle phase-space density

Q is defined as $Q \equiv /\langle v^2\rangle^{3/2}$, where ρ is the energy density and $\langle v^2\rangle$ is the mean square value of the particle velocity. The astronomically observable quantity is the mean coarse-grained phase-space density. In the absence of dissipation, the coarse-grained phase space density can only decrease from its primordial value. Thus, one can use the observed maximum phase-space density to set a lower limit on the phase-space density for the dark matter particles. The highest observed phase-space density is obtained from dwarf spheroidal galaxies: $Q_{obs} \approx 10^{-4} M_\odot pc^3 (km/s)^{-3}$. For this models, $\langle v^2\rangle \simeq r_c^2$ at the present time, and therefore

$$Q_0 \simeq \rho_{NT0}/r_c^3. \tag{3.4}$$

Because the primordial phase space density decreases with time when the particles are relativistic and becomes a constant after the particles become non-relativistic, one requires $Q_0 > Q_{obs}$, which can be translated into a constraint on r_c:

$$r_c < (\rho_{NT0}/Q_{obs})^{1/3} \approx 2.5 \cdot 10^{-7}. \tag{3.5}$$

This limit is of the same order of magnitude but slightly weaker than that from the observations of the Lyman-λ forest.

Compared to the conventional CDM model, the power spectrum of this model on sub-galatic scales is severely damped, which will account both for the lack of substructure in the Local Group and for the observed smooth inner halos. In addition, during the structure formation, the non-vanishing velocity of the WIMPs may further reconcile the discrepancies between the theoretical predictions and observations on sub-galactic scales. To solve the sub-galactic-scale problem, this model requires the value $r_c \approx 1.5 \cdot 10^{-7}$. This has implications on the models of non-thermal production of WIMPs.

In summary, the discrepancies between theory and observations on sub-galactic scales disfavores the conventional WIMPs CDM model. In this section wes shown that if the dark matter particles have a non-thermal origin, these discrepancies can be resolved and the WIMPs remain good candidates for the dark matter particle.[7]

4. CANNIBAL DARK MATTER

A thermally decoupled hidden sector of particles, with a mass gap, generically enters a phase of cannibalism in the early Universe. The Standard Model sector becomes exponentially colder than the hidden sector. We propose the Cannibal Dark Matter framework, where dark matter resides in a cannibalizing sector with a relic density set by 2-to-2 annihilations. Observable signals of Cannibal Dark Matter include a boosted rate for indirect detection, new relativistic degrees of freedom, and warm dark matter.

In the non-relativistic dark matter class of models, under the generic requirement that number changing interactions are active, the hidden sector undergoes a phase of cannibalism. The properties of cannibalism are determined by the requirement that the dark sector and SM separately preserve their comoving entropy densities. This leads to different scalings of the temperature versus the scale factor a,

$$T_\gamma \approx 1/a \text{ and } T_d \approx 1/\log(a) \tag{4.1}$$

where $T_\gamma(T_d)$ represents the temperature of the SM (cannibalizing) sector. The hidden sector temperature stays almost constant as the Universe expands as number changing interactions efficiently convert the rest mass of non-relativistic particles into kinetic energy. The different dependence on the scale factor implies that the SM particles are exponentially colder than the cannibalizing sector, as a function of the hidden sector temperature. In first studies and followup studies with cannibalism, it is assumed that DM annihilates through a 3-to-2 (or 4-to-2) process. Previous studies of cannibalism assume that the DM relic density comes from number changing interactions, such as 3-to-2 or 4-to-2 annihilations. In this section proposed a new class of DM models, where DM resides in a cannibalizing sector with relic density determined by 2-to-2 annihilations. In general, cannibal dark matter is realized in hidden sectors that contain both metastable states, ϕ, undergoing number changing interactions such as $\phi\phi\phi \to \phi\phi$, and a stable DM candidate, χ, annihilating into the metastable states through 2-to-2 annihilations, $\chi\chi \to \phi\phi$.

The exponential cooling of the SM, during cannibalism, has dramatic implications for DM phenomenology. DM must have a larger annihilation rate than conventional scenarios, in order to reproduce the observed relic density. Therefore, Cannibal DM predicts boosted rates for indirect detection. The Universe is exponentially older at DM freeze-out, implying less redshifting between DM decoupling and the start of structure formation and therefore the possibility of warm DM. Cannibal DM can also lead to new relativistic degrees of freedom, leaving imprints on the Cosmic Microwave Background (CMB).

Cannibalism requires the following conditions:

- The dark sector is kinetically decoupled from the SM sector.
- The dark sector has a mass gap.
- The dark sector remains in chemical equilibrium, through number changing interactions, at temperatures below the mass of the LDP.

As a simple example, we consider a dark sector with a real scalar LDP (Lightest Dark sector Particle) and generic interactions,

$$V_\phi = \frac{m_\phi^2}{2}\phi^2 + \frac{A}{3!}\phi^3 + \frac{\lambda}{4!}\phi^4, \tag{4.2}$$

where $A = \sqrt{3\lambda}m_\phi$. Cannibalism begins when T_d drops below m_ϕ, and chemical equilibrium is maintained through $\phi\phi\phi \to \phi\phi$ annihilations.

In the following, we assume that ϕ is metastable and eventually decays to either SM states or dark radiation.

In order for cannibalism to occur, ϕ must be out of kinetic equilibrium with its decay products and therefore have lifetime longer than Hubble, $\tau_\phi > H^{-1}$, when $T_d = m_\phi$. For decays to SM, we focus on $\phi \to \gamma\gamma$. For the dark radiation case, we assume that ϕ decays to a light species that is kinetically decoupled from ϕ and begins with zero abundance (for exampl $\phi \to \gamma'\gamma'$, where γ' is a light hidden photon).

Since the dark sector and the SM are kinetically decoupled, they have different temperatures and entropies. Assuming for simplicity that there are no entropy injections, the comoving entropies of the two sectors are separately conserved. This implies that the ratio of SM to dark entropy densities is fixed,

$$\xi \equiv \frac{s_{SM}}{s_d}. \tag{4.3}$$

If the two sectors were in thermal contact in the past, ξ is the ratio of the sum of degrees of freedom within each sector. However, ξ can be much larger if the two sectors reheat to asymmetric temperatures and remain thermally decoupled.

During cannibalism, T_d and s_d are related,

$$s_d \equiv \frac{2\pi^2}{45} g_{*s}^d T_d^3 \simeq \frac{m_\phi^3}{(2\pi)^{3/2} x_\phi^{1/2}} e^{-x\phi} \tag{4.4}$$

where $x_\phi \equiv m_\phi / T_d$ and we assume, for simplicity, that ϕ dominates the hidden sector entropy. Conservation of entropy within each sector implies that the SM becomes exponentially colder than the dark sector,

$$\frac{T_\gamma}{T_d} \sim 0.5 \xi^{1/3} g_*^{-1/3} x_\phi^{5/6} e^{-x\phi/3}, \tag{4.5}$$

where g_* is the effective number of relativistic degrees of freedom of the SM. Since $T_\gamma \approx 1/a$, Eq. (4.5) leads to Eq. (4.1). The energy density in the dark sector is approximately dominated by the LDP and decreases as

$$\rho_d \sim \frac{1}{a^3 \log a}, \tag{4.6}$$

in striking contrast to the energy density of nonrelativistic matter in thermal equilibrium with a relativistic plasma, which decreases as an exponential function of the scale factor. The dark sector starts to drive the expansion of the Universe when its energy density become equal to that of the SM, $T_\gamma^E / T_d^E = 4/(3\xi)$, corresponding to $x_\phi^E \approx -2.8 + \log(g_*^{-1} \xi^4) + 2.5 \log x_\phi^E$.

Cannibalism ends when the dark sector drops out of chemical equilibrium. This can happen for one of two reasons, (1) $\phi\phi\phi \to \phi\phi$ annihilations decouple, or (2) the Universe becomes older than the ϕ lifetime, and ϕ decays away. The second possibility requires that $\Gamma_\phi \sim H$ when T_d is between m_ϕ and the temperature that 3-to-2 processes decouple. This may seem to require a coincidence of scales between the ϕ lifetime and Hubble. However, the scale factor (and therefore the age of the Universe) changes by an exponentially large amount during cannibalism, as a function of the dark sector temperature.

Therefore, option (2) is generic and is realized for an exponentially large range of values for Γ_ϕ.[8]

4.1. FREEZE-OUT OF CANNIBAL DM

We now include DM in the cannibalizing sector. We consider a simple example where DM is a Majorana fermion, χ, that has a Yukawa interaction with ϕ,

$$\mathcal{L} \supset -V_\phi - (\frac{m_{chi}}{2}\chi^2 + \frac{y}{2}\phi\chi^2 + h.c.). \tag{4.7}$$

V_ϕ is defined in Eq. (4.2), $m_\chi > m_\phi$ and y is a complex number. The relic abundance of χ is determined by the freeze-out of 2-to-2 annihilations, $\chi\chi \to \phi\phi$. This cross section is s-wave if either Im yA $\neq 0$ or Im $y^2 \neq 0$. With purely imaginary y and $m_{phi} = 0$ this cross section reads $\langle \sigma \nu \rangle = \frac{A^2|y^2|}{1024\pi m_{chi}^4}$. We assume that DM annihilations decouple during cannibalism. Therefore, DM freezeout occurs after ϕ becomes non-relativistic, but before $\phi\phi\phi \to \phi\phi$ annihilations decouple and before ϕ decays. The cosmological stages of our scenario are depicted in Fig. 4.1.

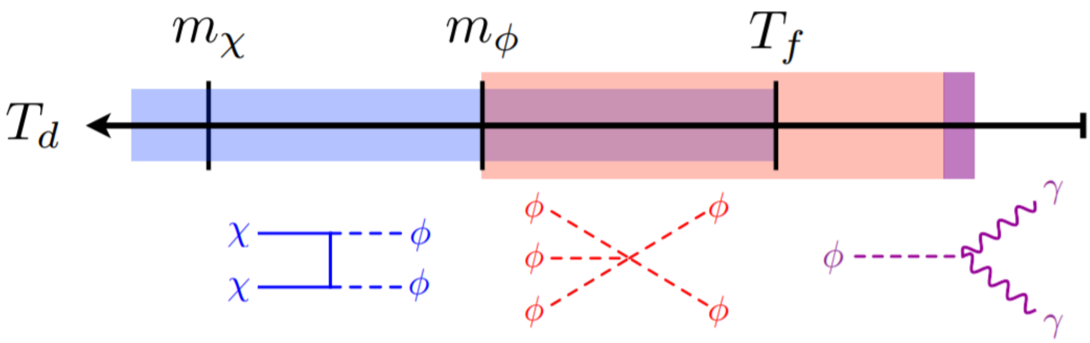

Figure 4.1: Cannibal DM has three stages. (1) DM annihilations, $\chi\chi \to \phi\phi$, are in equilibrium and the LDP, ϕ, is relativistic. (2) Cannibalism begins when ϕ becomes non-relatavistic, and then DM annihilations freeze-out, at temperature T_f, during cannibalism. (3) Cannibalism ends when ϕ decays away or $\phi\phi\phi \to \phi\phi$ annihilations decouple.

Under the assumption that $\chi\phi \to \chi\phi$ and $\phi\phi\phi \to \phi\phi$ are in equilibrium, the evolution of the number density of χ is described by a single Boltzmann equation,

$$\frac{dY_\chi}{d\log x} = -\frac{\kappa(x)s_d\langle\sigma\nu\rangle}{H}(Y_\chi^2 - Y_\chi^{eq2}), \tag{4.8}$$

where $Y_\chi \equiv n_\chi/s_d$, $x \equiv m_\chi/T_d$, and $\kappa(x) \equiv (1 - 1/3d\log g_{*s}^d/d\log x)$. The χ relic density is given by $\Omega_\chi/\Omega_{DM} = m_\chi Y_\chi/(0.4\,eV\,\xi)$, where $\Omega_{DM}h^2 \approx 0.12$ corresponds to the observed DM density. A numerical solution is shown in Fig. 4.2.

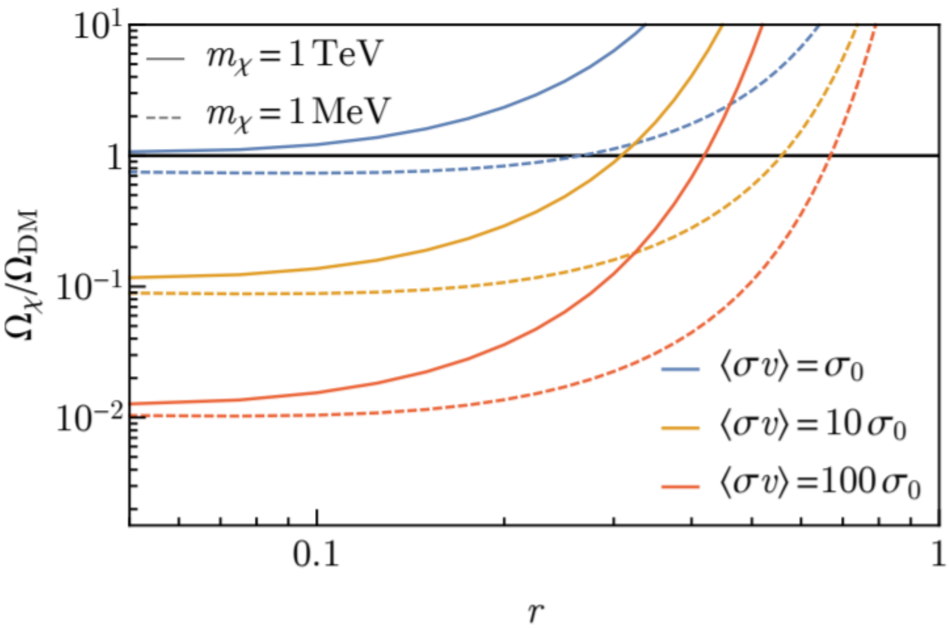

Figure 4.2: Relic density versus $r \equiv m_\phi/m_\chi$ for m_χ = (1 MeV, 1 TeV) and $\langle\sigma\nu\rangle = (1, 10, 100)\sigma_0$, where $\sigma_0 = 3 \cdot 10{-}26\mathrm{cm}^3\mathrm{s}^{-1}$. The relic density grows exponentially with r, as implied by Eqs. (4.5), (4.9), and (4.10).

The Boltzmann equation can be solved analytically in the sudden freeze-out approximation, $n_\chi^{eq}(x_f)\langle\sigma\nu\rangle = H$. There are two regimes depending on whether the SM or ϕ dominate the energy density of the Universe when DM annihilations decouple,

$$\frac{\Omega_\chi}{\Omega_{DM}} \approx \begin{cases} 0.3\frac{x_f}{g_*^{1/2}}\frac{\sigma_0}{\langle\sigma\nu\rangle}\frac{T_d}{T_\gamma}\cdot\frac{1}{D} & \rho_{SM} > \rho_d \\ 0.3\frac{x_f}{g_*^{1/2}}\frac{\sigma_0}{\langle\sigma\nu\rangle}\frac{T_d^{3/2}}{T_{\xi^{1/2}\gamma}^{3/2}}\cdot\frac{1}{D} & \rho_{SM} < \rho_d \end{cases} \tag{4.9}$$

where all quantities are evaluated at $x = x_f$ and $\sigma_0 = 3 \cdot 10^{-26}$ $\mathrm{cm}^3\mathrm{s}^{-1}$. D is a dilution factor that accounts for the entropy which is injected in the SM plasma when ϕ decays. $D = 1$ if ϕ energy density does not dominate at the time it decays; otherwise we evaluate D using a sudden decay approximation, $D = T_\gamma^E/T_{RH}$, where $T_{RH} \approx 0.8 g_*^{-1/4} \Gamma_\phi^{1/2} M_P^{1/2}$ is the temperature that the SM is reheated to after ϕ decays.

According to Eq. (4.5), $T_d/T_\gamma \sim e^{r x_f/3}$, which is naturally very large during cannibalism, requiring a boosted annihilation cross section compared to conventional scenarios. Eq. (4.9) depends on the temperature DM annihilations decouple, $x_f = \frac{m_\chi}{T_f}$, which in the sudden freeze-out approximation is,

$$x_f \approx \delta^{-1} \log[h(r) m_\chi M_P \langle \sigma \nu \rangle], \tag{4.10}$$

where $\delta = 1 - \frac{2}{3}r$ or $1 - \frac{1}{3}r$ for SM or *phi* domination, respectively, and

$$h(r) \approx \begin{cases} 0.3 g_*^{1/6} \xi^{-2/3} r^{-5/3} (1 - \frac{2}{3}r)^{7/6} & \rho_{SM} > \rho_d \\ 0.2 r^{-5/4} (1 - \frac{1}{2}r)^{3/4} & \rho_{SM} < \rho_d \end{cases} \tag{4.11}$$

Eqs. (4.5), (4.9), and (4.10) imply that the final DM abundance is exponentially sensitive to r, as shown in Fig. 3.

4.2. PHENOMENOLOGY

Cannibal DM is not presently testable by direct detection, because the SM and hidden sector are kinetically decoupled, implying a small cross section. However, Cannibal DM predicts rich signals in indirect detection and cosmology, driven by the boosted DM annihilation rate, the (exponentially large) age of the Universe at DM freeze-out, and the decay of relic LDPs.

When ϕ decays, its energy density can modify the number of relativistic degrees of freedom, which is constrained by the CMB: $N_{eff} = 3.15 \pm 0.23 (1\sigma)$. If ϕ decays to photons after neutrino decoupling, the photons are heated relative to the neutrinos, lowering N_{eff}. Alternatively, if ϕ decays to dark radiation the resulting energy density increases N_{eff}.

When ϕ decays to photons, there are indirect detection constraints on the process $\chi\chi \to \phi\phi \to 4\gamma$. Energy injection at the recombination epoch can distort the CMB as measured by Planck. Fermi observations of dwarf galaxies bound DM annihilations at the present epoch.

An important constraint on Cannibal DM follows from the exponentially long time it takes DM to freezeout. While the average velocity of Cannibal DM at freeze-out is of the same order as conventional scenarios, $v_f \sim x_f^{-1/2}$, the SM is exponentially colder at freeze-out, Eq. (4.5). Between kinetic decoupling and matter-radiation equality, DM free-streaming damps density perturbations, resulting in a cut-off in the matter power-spectrum on scales smaller than the free-streaming length, λ_{fs}. This is typically a negligible effect for standard scenarios with DM mass above a few keV since $\lambda_{fs} \propto T_k^{-3/2}$, with T_k the temperature of kinetic decoupling of DM (however heavy and warm DM is possible if DM is produced from late decays of a heavier state). The strongest constraint on the free-streaming length comes from observations of the Lyman-alpha forest: $\lambda_{fs} \leq 0.1 Mpc (2\sigma)$. In order to evaluate the bound we use the approximate expression

$$\lambda_{fs} = \int_{t_k}^{t_{eq}} \frac{v(t)}{a(t)} dt \approx 125\, Mpc\, v_k \frac{\log(1.3 T_k^{eV})}{T_k^{eV}} \tag{4.12}$$

where v_k is the average DM velocity at kinetic decoupling and T_k^{eV} is the relative SM plasma temperature in eV. Eq. 4.12 assumes the universe is radiation dominated between decoupling and equality, which is always true in our case.[3]

5. ASYMMETRIC DARK MATTER

Asymmetric dark matter models are based on the hypothesis that the present-day abundance of dark matter has the same origin as the abundance of ordinary or "visible" matter: an asymmetry in the number densities of particles and antiparticles. They are largely motivated by the observed similarity in the mass densities of dark and visible matter (VM), with the former observed to be about five times the latter.

The motivation comes from the observation that the present-day mass density of DM is about a factor of five higher than the density of VM

$$\Omega_{DM} \simeq \Omega_{VM} \tag{5.1}$$

where Ω as usual denotes the mass density of a given component relative to the critical density. The similarity in these observed densities suggests a common origin, some kind of a unification or strong connection between the physics and cosmological evolution of VM and DM. The present-day density of VM has long been established as due to the baryon asymmetry of the universe: some time during the early universe, a tiny excess of baryons B over antibaryons $\overline{B}$ evidently developed, parameterized by

$$\eta(B) \equiv \frac{\eta_B - \eta_{\overline{B}}}{s} \simeq 10^{-10} \tag{5.2}$$

where number densities are denoted η, and s is entropy density. The baryons in the universe today constitute the excess remaining after all of the antibaryons annihilated with the corresponding number of baryons. The ADM hypothesis simply states that the present-day DM density is similarly due to a DM particle-antiparticle asymmetry, and that these asymmetries are related due to certain processes that occurred rapidly during an early cosmological epoch but later decoupled.

Asymmetric DM may be contrasted with weakly-interacting massive particle or WIMP DM. The latter postulates that the DM is a thermal, non-relativistic relic particle (usually self-conjugate) with mass in the GeVTeV range that decouples when its weak-scale annihilations fall out of equilibrium due to the Boltzmann suppression of the WIMP population. Famously, this cold DM (CDM) scenario “miraculously” provides about the correct DM mass density for generic weakscale annihilation cross-section, with a specific value (weakly dependent on WIMP mass) derived by fitting the abundance exactly. Furthermore, the idea fits in well with independent particle-physics motivations for new weak-regime physics such as supersymmetry. However, in almost all WIMP scenarios the similarity of the DM and VM densities then must be taken to be a coincidence.

Asymmetric DM is one of a number of well-motivated alternatives to the WIMP solution – other examples are keV-scale sterile neutrinos, axions, and Q-balls – which all deserve serious attention. It does not have the WIMP feature of a tight connection between the thermal freeze-out process, indirect detection, direct detection and collider production. Indirect detection through dark particle-antiparticle annihilations is obviously irrelevant, because there are no DM antiparticles left to annihilate with (co-annihilations with, for example, nucleons are possible but not required). Many models of asymmetric DM can be tested through direct detection and collider signatures, but the parameters involved typically include some that are independent of the physics behind the cosmological DM abundance.

The idea that the universe may contain a relic, dark, particle–anti-particle asymmetric component has in one form or another been considered for a few decades. One of the oldest ideas is that of mirror matter, where the dark sector has identical microphysics to the visible sector. Though subsequent observational data have ruled out the specific scenarios, these early papers contemplated a universe with equal amounts of (asymmetric) matter and mirror matter. Another relatively early idea was that DM stability may arise from an analog of baryon-number conservation in a technicolor sector, with the DM state being a neutral techni-baryon. The modern era of ADM research occurred after the observational result that the VM and DM densities were as close together as a factor of five was established.

The VM abundance in the universe today is made from a small number of SM fields: protons and bound neutrons (formed mainly from valance up and down quarks, and gluons), and electrons (with neutrinos and photons comprising the current radiation content of the VM). These are the stable relics of a much larger SM particle content. Since ADM models seek to draw a connection between DM and VM, it is natural to suppose that the DM may also be the stable member(s) of some relatively complicated gauge theory constituting a hidden sector. In general, ADM models have gauge groups that contain the product structure

$$G_V \times G_D \tag{5.3}$$

where the first factor is the SM gauge group or some extension thereof, and the second factor is a dark gauge group. Some models have a gauge force that couples to both sectors, with an extended $U(1)_{B-L}$ being a common example. The dark sector in general may have various fermion and scalar multiplets in representations of GD, and spontaneous gauge-symmetry breaking may occur. Many models have the dark sector as simply just fermions or scalars or a mixture of the two. Taking our cue once again from the visible world, it could well be that the DM is multi-component and that there is dark radiation (bosonic and/or fermionic) as well as dark matter. A relatively complicated dark sector is not mandatory, but it is perfectly consistent with the ADM philosophy. There are two features that are not optional: a conserved or approximately-conserved dark global quantum number so that a dark asymmetry can be defined in the first place, and an interaction that annihilates away the symmetric part of the dark plasma, just as strong and electroweak interactions annihilate the symmetric component of the SM plasma into radiation.

Why certain visible-sector particles are stable (or at least very long lived)? Protons are stable because they are the lightest particles carrying conserved baryon number B. Neutrons are the next-to-lightest baryons, and they are unstable unless they are bound in appropriate nuclei. Electrons are stable because they are the lightest electrically-charged particles, with electric-charge conservation mandated. Note that the simultaneous existence of stable protons and electrons permits the universe to be electrically neutral. The least massive neutrino mass eigenstate has its stability ensured through angular momentum conservation because it is the lightest fermion: its decay products would have to contain a half-integer spin particle. Photons and gluons, being massless, are stable by kinematics (gluons are also stable because they are lightest colored states). The reasons for stability are several: exact (or nearly exact) global symmetry, unbroken gauge symmetry, angular momentum conservation, kinematics, and bound state effects (nuclear physics).

The diversity of causes for stability suggests that the dark sector may similarly contain a number of stable particles, as well as structures analogous to atoms or nuclei from unbroken dark-sector gauge interactions. Rather than getting bogged down in possible complexities, we shall begin thinking about dark sectors by focusing simply on baryon number conservation. In terms of mass, B conservation is the most important of the stabilizing influences in the visible sector: the cosmological B asymmetry (with the proton mass) determines the VM mass density. Since we want to understand Eq. (5.1), it seems sensible to start by postulating a "dark baryon number" and establishing a relationship with "visible" baryon number. Other considerations may be brought in as necessary. Denote the visible and dark baryon numbers by BV and BD, respectively. At low energies and temperatures, both of these must be conserved to ensure the separate stability of VM and DM.

5.1. SYMMETRY STRUCTURE

In the very early universe, there are four generic possibilities for initial asymmetry creation:

1. A non-trivial linear combination of B_V and B_D is exact, but a linearly-independent combination is broken.

2. B_V is broken, while B_D is not.

3. B_D is broken, while B_V is not.

4. Both are broken.

The breaking of at least one baryon-number symmetry is required for creating an asymmetry. Recall that the Sakharov conditions for generating an X-particle number asymmetry are: violation of X, C and CP through processes that occur out of equilibrium.

Once the initial asymmetry creation has happened in cases 2-4, the ADM framework requires some new interactions to then become rapid so that chemical exchange between the two sectors relates the asymmetries. For case 2, these interactions have to reprocess some VM asymmetry into a DM asymmetry, while for case 3 the reverse must happen. For case 4, the initial asymmetries may be quite different, so subsequent interactions should drive the asymmetries towards some kind of equilibration to comply with the standard ADM philosophy. If this does not happen, then while the DM is certainly asymmetric, the lack of relation between the asymmetries means that one is not addressing the primary motivation given by Eq. (5.1).

Case 1 is qualitatively distinct from the others, because correlated asymmetries are created simultaneously in the visible and dark sectors via common interactions, with the universe always being symmetric in some linear combination of the visible and dark baryonic numbers. This very interesting scenario represents the unified generation of both visible and dark matter, and is arguably the most elegant implementation of ADM. Because the universe is symmetric in one linear combination of baryon numbers, this scenario is also said to produce a "baryon-symmetric universe".

Let the conserved and broken linear combinations be, respectively,

$$\begin{aligned} B_{con} &= \alpha B_V + b B_D, \\ B_{bro} &= c B_V + d B_D, \end{aligned} \tag{5.4}$$

where $a, b \neq 0$ to ensure that the conserved quantity is a non-trivial combination, and $ad - bc \neq 0$ to ensure that the two combinations are independent. Abelian charges are of course defined up to a normalization convention, and B_{bro} may be redefined by the addition of a piece proportional to B_{con} because the result is still a broken charge. This permits us to simplify the definitions of the conserved and broken baryon-number charges without loss of generality. By scaling B_{con} we may set $a = 1$ and by scaling B_D we may put $b = -1$. Adding $(d - c)B_{con}/2$ to B_{bro} and then setting $d + c = 2$ completes the simplification process. The result is that Eq. (5.4) is equivalent to

$$B_{con} = B_V - B_D, \tag{5.5}$$

$$B_{con} = B_V + B_D, \tag{5.6}$$

The Sakharov conditions can now be used to engineer dynamical schemes that give rise to an asymmetry in B_{bro} while maintaining $\eta(B_{con}) = 0$ as a constraint, leading to

$$\eta(B_{bro}) \equiv \eta \Rightarrow \eta(B_V) = \eta(B_D) = \frac{\eta}{2}. \tag{5.7}$$

Thus baryon-symmetric models can always be interpreted as implying that the asymmetries in the two sectors are equal, and that the DM is concealing a (generalized) baryon number that cancels the baryon asymmetry of VM (which is what $\eta(B_{con}) = 0$ means).

The fundamental feature of having B_{con} always conserved now requires some additional discussion. With B_V identified as visible-sector baryon number, the B_{con} defined in Eq. (5.5) is anomalous and thus not actually conserved at the quantum level. In the early-universe, an important aspect related to this is the reprocessing of B_V asymmetry by electroweak sphalerons into visible lepton number. In order for B_{con} to also be quantally conserved, we must replace B_V with a suitable, related anomaly-free quantity. The obvious choice is $(B - L)_V$, and indeed this identification is made in many specific baryon-symmetric schemes.

The connection of the visible-sector generator withinB_{con} to proton stability is then more subtle, since $(B - L)_V$ conservation alone cannot enforce it. At the non-perturbative level, the sphaleron or related zero-temperature instanton process conserves a certain $\mathbb{Z}_3$ discrete symmetry that ensures the absolute stability of the proton within the SM, even in the face of the anomalous $(B + L)_V$ violation. At low energies and low temperatures, we must also arrange any model to produce the SM as the effective theory, and thus perforce have a conserved B_V at the perturbative level (in that limit). Depending on how it is constructed, similar issues could arise in the dark sector, and it is worth noting that the absolute stability of DM (if that is what is desired) may be guaranteed by a discrete subgroup of $U(1)_{B_D}$ rather than the full parent group.

With B_{con} now conserved at both the quantal and classical levels, one is free to gauge the associated abelian symmetry. This further deepens the fundamental nature ascribed to this symmetry. Because the gauged U(1) has to be spontaneously broken to ensure that the associated Z' boson is sufficiently massive, but we still want to have the global U(1), the gauged model has to be constructed in a particular way for a full discussion). Essentially, the enlarged symmetry should be the product of the gauged U(1) and the related global

U(1). The scalar field whose vacuum expectation value will spontaneously break the gauged U(1) must have zero charge under the global U(1) to ensure that the latter remains exact. The Z', which has decay channels into dark-sector particles and thus a substantial invisible width, is a generic feature of gauged baryon-symmetric models, and provides one very interesting way of searching for experimental evidence for at least that kind of ADM. In all four cases, depending on the context and the temperature regime of interest, the interactions that lead to a relation between the VM and DM asymmetries are either described by an explicit renormalizable theory, or by effective operators of the form

$$\mathcal{O}_{(B-L)_V}\mathcal{O}_{B_D} \tag{5.8}$$

where $\mathcal{O}_{(B-L)_V}$ is formed from visible-sector fields in a combination that carries nonzero $(B-L)_V$, while $\mathcal{O}_{B_D}$ is a dark-sector analog. The interactions must preserve some linear combination of the $(B-L)_V$ and B_D numbers, otherwise they would wash out both asymmetries. Some of these operators lead to interesting collider signatures if the effective scale is in the TeV regime. In some models, non-perturbative sphaleron processes are used to relate the asymmetries.[6]

5.2. ASYMMETRY GENERATION

A full ADM theory should specify the dynamics of asymmetry generation, although a number of works simply assume an initial asymmetry was created by some means, and focus instead on how the asymmetry gets distributed between the visible sector and the dark sector. The most common asymmetry creation scenarios are out-of-equilibrium decays, Affleck-Dine dynamics, bubble nucleation during a firstorder phase transition, asymmetric freeze-out, asymmetric freeze-in, and spontaneous genesis. We now briefly review the basic idea behind each of these. We shall call the relevant particle number X, which may be B_V, B_D or some linear combination depending on the model.

Out-of-equilibrium decays of heavy particles. This scenario is adapted from Fukugita-Yanagida-style leptogenesis as can occur in the type I seesaw model of neutrino mass generation. The idea is that there is a massive unstable particle that decouples from the thermal plasma, and then decays through interactions that violate X, C and CP. Typically, the decaying particle ψ is self-conjugate and the decay rates for the process $\psi \to x_1 x_2 ...$ and its charge-conjugate $\psi \to x_1^* x_2^* ...$ are unequal due to CP violation, where x_i denotes a particle whose X-charge equals X_i. The unequal decay rates to final states of opposite X create the asymmetry, and because the decays happen after the particle has lost thermal contact with its daughter particles, there is no wash out due to inverse decays. As in standard leptogenesis, the decay amplitude must involve interference between at least two Feynman graphs in order for CP-violating phases to have physical consequences.

Affleck-Dine mechanism. Affleck-Dine (AD) dynamics is a very plausible mechanism in supersymmetric (SUSY) theories, and it is worth noting at the outset that for the purposes of AD asymmetry generation the supersymmetry breaking scale is allowed to be beyond the reach of the LHC. Take a complex scalar field ϕ that carries nonzero X, and compute the Noether charge density

$$J^0 = i(\dot{\phi}^* \phi - \phi^* \dot{\phi}) = R^2 \dot{\theta}, \tag{5.9}$$

where $\phi \equiv (R/\sqrt{2}) \exp(i\theta)$ is an amplitude and phase decomposition. In the AD mechanism, suitable conditions for generating a time-dependent phase exist, thus creating X charge carried by the coherent oscillations of the scalar field (which initially is a spatially-homogeneous condensate). To create the correct amount of charge, the amplitude R needs to be large to compensate for the fact that the violation of X-charge conservation must for phenomenological reasons be small. This is assisted in supersymmetric theories by the generic existence of flat directions for renormalizable scalar potentials. The flat directions are lifted by supersymmetry-breaking soft masses and by effective, non-renormalizable terms that also provide explicit X violation. The coupling of the AD field ϕ to the inflaton helps to set up an initial state of that field during inflation to be at the required high value, and it also implies that some of the parameters in the scalar potential change with time as the universe expands. At those high field values, the effect of the small X-violating terms is amplified, and these terms together with CP-violating and time-dependent parameters kick the AD field in the angular direction and thereby create the X charge or asymmetry. The AD mechanism has been used in several ADM models.

First-order phase transition. This is analogous to electroweak baryogenesis. The idea is that the phase transition from the symmetric phase of a gauge theory to the broken phase proceeds via bubble nucleation seeded by quantum tunneling through a potential barrier. The Higgs field driving the phase transition Yukawa couples to fermions in a CP violating way, and X conservation is violated in the symmetric phase via rapid sphaleron transitions associated with a triangle anomaly between the X current and the gauge fields. The movement of the bubble walls creates departures from equilibrium that partner with the X- and CP-violating interactions to create an X asymmetry carried by the fermions. In electroweak baryogenesis, the observed high value of the electroweak Higgs mass implies that the electroweak phase transition is not first-order (although it may be in the context of an extended Higgs sector). For the ADM application however, the phase transition in some models occurs for a new gauge force, either in the dark sector or in a third sector that mediates between the visible and dark worlds. The resulting parameter freedom makes it trivial to arrange for the phase transition to be as strongly first order as desired.

Asymmetric freeze-out The asymmetric freeze-out mechanism uses the same interactions to generate the asymmetry and to eliminate the symmetric part. Consider the DM particle χ and its antiparticle $\overline{\chi}$. As well as self-annihilations of χ with $\overline{\chi}$ to SM states, χ and $\overline{\chi}$ may also experience X-violating co-annihilations with SM species. Taking χ to be spin-1/2 and assigning $B_D(\chi) = 1$, the crossing-symmetry-related reactions

$$\chi + u_i \to \overline{d}_j + \overline{d}_k, \;\; \chi + d_j \to \overline{u}_i + \overline{d}_k, \tag{5.10}$$

where $i, j, k = 1, 2, 3$ are family indices, are the simplest co-annihilations involving quarks that conserve electric charge, color and $B_V - B_D$, but violate $B_V + B_D$. They are generated through the "neutron portal" type of effective operator schematically written as $udd\chi$. By considering other gauge-invariant combinations of SM fields carrying nonzeroBV, more complicated co-annihilation reactions can be systematically identified. Through C and CP violation, the rates for the χ and $\overline{\chi}$ co-annihilations can be unequal, and if the co-annihilation rates dominate over the self-annihilation rates, then χ and $\overline{\chi}$ will decouple at different temperatures and thus have exponentially different relic number densities. This scenario gives rise to a baryon-symmetric universe when the asymmetry-creating coannihilations preserve a linear combination of B_V and B_D as in the example of Eq. (5.10). In the ADM context, for the cases where the DM mass is in the few-GeV regime, the asymmetric freeze-out co-annihilations must create an asymmetry in visible-baryon number directly, because no sphaleron reprocessing will be possible.

Asymmetric thermal production or asymmetric freeze-in. A particle that is too weakly coupled to ever attain thermal equilibrium with the cosmological bath can nevertheless be slowly produced from processes involving the bath particles. This thermal production process has recently also been called "thermal freeze-in" and applied within the ADM paradigm. Thermal freeze-in has been argued to be the inverse of thermal freeze-out. Freeze-out occurs when a species χ that starts off in thermal equilibrium subsequently decouples from the bath. Freeze-in sees the particle χ being so weakly coupled to the bath that while it is being slowly produced and heading towards thermal equilibrium, its comoving number density becomes a constant – freezes in – before it is actually able to reach equilibrium. The freeze-in happens when the temperature drops below the mass of the heaviest particle involved in the χ production process, thus Boltzmann suppressing that heaviest species. If χ can be produced by the decays or inverse decays of bath particles, then that process will dominate over χ scattering processes. Most of the χ production occurs just before freeze-in, partly because the characteristic Hubble time is longest at that point and often also because the microscopic process is most rapid then. For example, if χ is produced through the decay of a particle of mass m, the decay rate is suppressed by an m/T time dilation factor when $T \gg m$. In this case the largest mass is m, and thus $T = m$ is the approximate freeze-in temperature (precise calculations show that χ is dominantly produced when m/T is in the range 2-5). Asymmetric freeze-in simply means that χ and $\overline{\chi}$ freeze in with unequal co-moving number densities. This is achieved in specific models through the decays of bath particles and antiparticles which violate C and CP and conserve only one linear combination of B_V and B_D. The daughter DM particles may be in thermal equilibrium with other species inhabiting the dark sector provided that the dark-bath temperature is lower than that of the visible bath.

Spontaneous genesis. The Sakharov conditions presuppose CPT symmetry. While this is rigorously a symmetry of all local, relativistic quantum field theory Lagrangians, it is spontaneously broken by the expanding universe solutions used in cosmology. Through appropriate interactions, this can induce effective CPT violation, together with T violation, in the particle physics of the early universe. Once CPT invariance does not (effectively) hold, the Sakharov conditions need not all be obeyed in order to dynamically obtain a particle-number asymmetry. In "spontaneous genesis", the violation of particle number is, of course, still required, but the CP-violation and out-of-equilibrium conditions are not in general obeyed. The basic mechanism requires an effective term of the form $\mathcal{L} \supset \partial_{mu}\phi J_X^\mu/\Lambda$, where J_X^{mu} is the current corresponding to the particle number X that will develop an asymmetry, ϕ is some scalar field, and Λ is a high scale of new physics. For spatially-homogeneous but time-dependent ϕ solutions, this term reduces to $(\dot{\phi}/\Lambda)(n_X - n_{\overline{X}})$. Thus $\dot{\phi}/\Lambda$ is an effective chemical potential for X number, leading to the approximate result $n_X - n_{\overline{X}} \sim T^2\dot{\phi}/\Lambda$. Specific models have to arrange for a suitable background scalar field configuration ϕ to develop. Clearly, they must ensure that $\dot{\phi} \to 0$ at late times. To stop wash out in this limit, the X-violating interactions have to decouple before the effective chemical potential becomes too small.[9]

6. CONCLUSION

In this paper considered various approaches to the issue related to the generation of dark matter. The generation of DM in the early universe can proceed via thermal or non-thermal production, or both, or it may result from a particle-antiparticle asymmetry. Due to the fact that dark matter particles have not yet been experimentally detected, each approach deserves attention. Each of the approaches has cosmological and astrophysical implications and bounds, and touches on direct detection prospects and collider signatures.

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