1
Introduction

Cold collisionless fluid that explains many observations in astrophysics and cosmology. It's existence shouldn't be questioned, only its microphysical constitution.

1.1 Evidence for Dark Matter

Today, the amount of evidence in support of dark matter's existence is overwhelming. This evidence comes from astrophysical and cosmological observations that are inconsistent with a universe composed entirely of visible matter. This section serves as a review of this evidence.

1.1.1 Astrophysical Observations

Galaxy Clusters

Some of the first hints for the existence of dark matter came from observations of galaxy clusters. Perhaps the most famous analysis was performed by Fritz Zwicky [1], who was puzzled by the high rotational velocities of galaxies within the Coma Cluster. By applying the virial theorem, equating the cluster's kinetic and gravitational potential energies, he found that the cluster would need to contain a much more significant amount of *dunkle materie* (dark matter) than visible matter to accommodate these high velocities.

Rotation Curves of Spiral Galaxies

The anomalous rotational velocities observed in galaxy clusters can also be observed at the galactic scale. The rotation curves of spiral galaxies, which relate the rotational velocities of stars to their distance from the galactic centre, were observed to be flat at large distances. From the observed distribution of visible matter, Newtonian mechanics predicts that the orbital velocity of a star a distance r from the galactic centre, $v_\star(r)$, is related to the mass contained within a radius r, $M(r)$, through

$$v_\star(r) = \sqrt{\frac{GM(r)}{r}}, \tag{1.1}$$

indicating that the velocity should fall off as $1/\sqrt{r}$ at the outer regions of the galaxy where $M(r)$ is constant. Instead, observations of many spiral galaxies indicate that this velocity remains constant out to the galaxy's edge.

A simple way to produce such a rotation curve is to introduce a spherically symmetric distribution of dark matter surrounding the galaxy,

$$\rho_{\mathrm{DM}}(r) = \frac{v_0^2}{4\pi G r^2}, \tag{1.2}$$

which results in a constant rotational velocity of v_0 out to the galaxy edge. Detailed simulations of structure formation in a Cold Dark Matter (CDM) Universe indicate that the true distribution is better represented by distribution functions such as the Navaro-Frenk-White (NFW) profile [2, 3] or Einasto [4] profiles, which are commonly used in the literature.

An example rotation curve for galaxy NGC 6503 is presented in Fig. 1.1, with the contributions from each of the matter components to the rotational velocity shown [5, 6]. As can be seen, the visible matter constituting disk and gas components does not explain the observed rotational velocity.

Gravitational Lensing

As described by General Relativity, the curvature of space-time around massive entities causes light to travel along curved paths. As such, the mass of astrophysical structures can be deduced from the extent to which they distort the images of objects in the background. The extent of the distortions depends on how massive the foreground object is, ranging from the shearing of the background image (weak lensing), to multiple copies of the background object appearing (strong lensing) [7]. The disparity between the mass obtained from gravitational lensing and the mass of visible matter in the system is further evidence of dark matter's existence [8, 9].

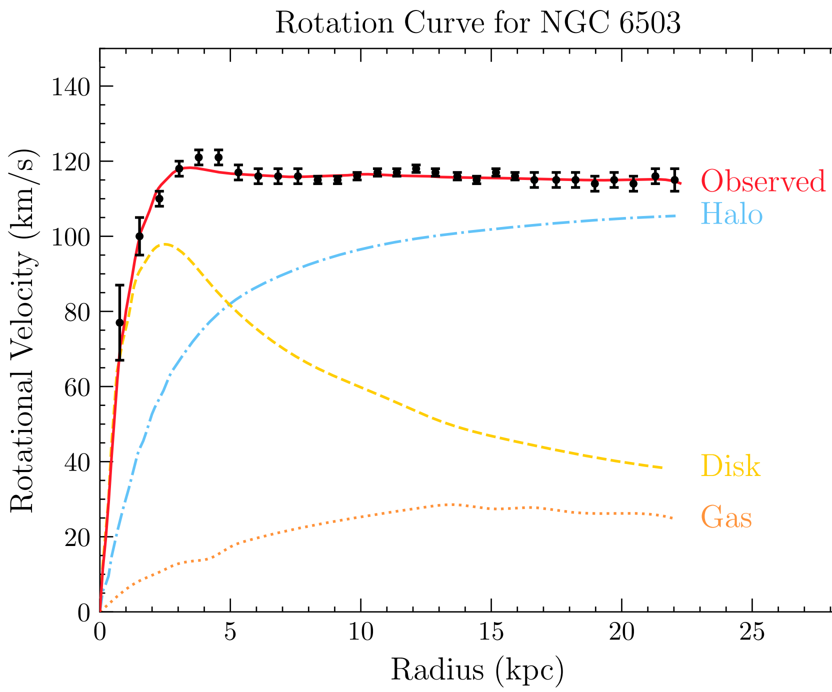

Figure 1.1: Galaxy rotation curve for NGC 6503, showing the contributions to the total velocity (red) from the DM halo (blue), disk (yellow), and gas components. Data used in making this plot was obtained from [5, 6].

The Bullet Cluster

The galaxy cluster 1E 0657-56, commonly referred to as the "bullet-cluster", was formed by the collision of two separate galaxy clusters. The baryonic matter in these clusters is mostly composed of a strongly interacting gas, and as expected produced a significant amount of X-rays during the collision. These X-rays were imaged by the Chandra X-Ray telescope [10], providing information on the resulting distribution of the visible matter. This is shown by the red regions of Fig. 1.2, where it can be seen how the visible matter has been smeared due to the collision. However, when the gravitational potential was mapped using gravitational lensing, it was clear that the majority of the mass was displaced relative to the visible matter. This mass is attributed to the dark matter components of the original clusters. As indicated by the purple regions in Fig. 1.2, the dark matter halos seem to have passed through each other mostly unperturbed. This tells us that not only is the majority of the mass comprised of dark matter, but that the dark matter has extremely weak interactions with both the visible matter and itself.

Figure 1.2: Image of the Bullet Cluster with contours of the gravitational potential superposed. The red regions indicate the baryonic matter after the collision, while the purple regions are the expected DM components deduced from gravitational lensing. [10, 11]

1.1.2 Cosmological Evidence

The current best cosmological model is the Λ-Cold Dark Matter model (ΛCDM), in which Λ refers to the cosmological constant associated with dark energy, and as the name suggests, cold (i.e. non-relativistic) dark matter plays a prominent role. The key components of this model are the aforementioned dark energy and dark matter, along with baryonic matter, and assumes that gravity is described by Einstein's General Relativity. The total energy density of the universe, $\rho_{\text{Univ.}}$, is broken down into three components based on how their density redshifts with the expansion of the Universe. In the ΛCDM model, these components are matter, radiation, and the vacuum energy Λ. The cosmological abundances of each component, (Ω_{m}, Ω_{r}, Ω_Λ respectively), are expressed as a fraction of the critical density, ρ_{crit},

$$\rho_{\text{crit}} = \frac{H^2}{8\pi G_N}, \tag{1.3}$$

$$\Omega_i = \frac{\rho_i}{\rho_{\text{crit}}}, \tag{1.4}$$

where H is the Hubble parameter, such that the total energy density of the Universe satisfies

$$\Omega_{\mathrm{m}} + \Omega_{\mathrm{r}} + \Omega_{\Lambda} = \frac{\rho_{\mathrm{Univ.}}}{\rho_{\mathrm{crit}}}. \tag{1.5}$$

The ratio $\rho_{\mathrm{Univ.}}/\rho_{\mathrm{crit}}$ determines the curvature of the universe, with values greater than 1 corresponding to a closed universe, less than 1 to an open universe, and equal to 1 to a spatially flat universe. Current observations are consistent with a spatially flat universe, and so we have $\sum_i \Omega_i = 1$.

The ΛCDM model has seen huge success as it provides explanations for observed the power spectrum of the Cosmic Microwave Background (CMB), the large-scale structure of the Universe, the abundances of light elements (hydrogen, deuterium, helium, and lithium), and the accelerated expansion rate of the Universe. These observations constrain the parameters of the model, and hence provide a complimentary probe of the properties of dark matter to the astronomical observations discussed above.

The Cosmic Microwave Background

One of the strongest probes of cosmological models is the Cosmic Microwave Background (CMB), relic photons from the time epoch of last scattering. This occurred after recombination, at a temperature of around $\sim 3000\,\mathrm{K}$, once the photons had decoupled from the baryonic matter and could freely propagate through the universe. The photons observed today have been redshifted by the expansion of the Universe, and are well described by a blackbody spectrum with a temperature of $T_{\mathrm{CMB}} = 2.73 \pm 0.0006\,\mathrm{K}$. Observations of the CMB temperature reveal that it is not exactly isotropic, with anisotropies at the level of $\delta T_{\mathrm{CMB}}/T_{\mathrm{CMB}} \sim 10^{-5} - 10^{-6}$ seen on a range of angular scales in the sky. These anisotropies were seeded by the primordial density perturbations that arise during inflation. These perturbations evolve due to the acoustic oscillations of the photon-baryon plasma driven by the interplay between the pressure from the photons and the gravitational attraction of the matter. The oscillations cease once the photons decouple, freezing in their temporal phases that are observed as peaks in the angular power spectrum of the temperature anisotropies.

Measurements of the CMB power spectrum provide information on many of the cosmological parameters. The locations of the acoustic peaks depend on the spatial geometry of the Universe and hence constrains Ω_{tot}. The total matter density, Ω_{m}, affects how the CMB spectrum is gravitationally lensed. The relative amplitudes of the peaks probe the baryon-to-photon ratio and hence the baryon density, Ω_{b}. Finally, the density of dark matter, Ω_{DM}, is obtained by fitting the cosmological parameters to the exact shape of the spectrum [5, 12].

The Planck collaboration most recently performed a precise measurement of the CMB power spectrum in 2018, obtaining best-fit parameters [12, 13]

$$\Omega_{\mathrm{m}} = 0.311 \pm 0.006, \quad \Omega_{\Lambda} = 0.689 \pm 0.006, \tag{1.6}$$

for the matter and dark energy densities. They obtained a total energy density of $\Omega_{\mathrm{tot}} = 1.011 \pm 0.006$ at 68% confidence level, providing strong evidence for a spatially flat Universe. The breakdown of the matter density into the dark and baryonic components is determined by combining the CMB results with constraints from Big Bang Nucleosynthesis (BBN)[1] giving

$$\Omega_{\mathrm{DM}} h^2 = 0.1193 \pm 0.0009, \quad \Omega_{\mathrm{b}} h^2 = 0.02242 \pm 0.00014, \tag{1.7}$$

where h is the dimensionless Hubble constant such that the Hubble parameter today is $H_0 = 100\,h\,\mathrm{km\,s^{-1}\,Mpc}$.

Large Scale Structure

After recombination, the pressure on the baryonic matter from photons began to decrease, eventually allowing the small density perturbations to grow. This led to the growth of stars, galaxies, and the large-scale structure we observe today [14]. N-body simulations of the Universe's evolution require a cold dark matter component for this structure to form. While a small component of the dark matter can be warm, hot dark matter would wash out small-scale structures [15].

1.2 Potential Models of Dark Matter

Given that baryonic matter is composed of particles described by the Standard Model (SM) of particle physics, it is a fair assumption that dark matter will also have a particle nature. Therefore, models of particle dark matter are built by extending the SM in a way consistent with its symmetries. Such models may be as simple as introducing a single new field into the SM, or there may be a more extensive hidden sector with a complicated symmetry structure. Additionally, there are compelling theories in which dark matter is not a fundamental particle, such as primordial black holes (PBHs) formed in the early universe. Given the few details we know about dark matter, there exists an enormous library of viable dark matter candidates. However, there are generic properties a good dark matter candidate must satisfy, namely:

[1] BBN is the process that produced the light elements (D, ^{3}He, ^{4}He, and ^{7}Li) were produced in the early universe. This process is highly sensitive to the physical condition of the universe at that time, allowing for strong constraints to be placed physics beyond the Standard Model.

- **Stable on Cosmological Timescales:** Dark matter must either be stable or have a lifetime significantly longer than the age of the Universe to be present in its current abundance.

- **Neutral or milli-charged under Electromagnetism:** Dark matter, as its name suggests, does not significantly interact with light. Requiring that dark matter be completely decoupled from the Standard Model plasma by the time of recombination yields an upper bound on the electric charge of dark matter [16]

$$q_{\mathrm{DM}}/e < \begin{cases} 3.5 \times 10^{-7} \left(\frac{m_{\mathrm{DM}}}{1\,\mathrm{GeV}}\right)^{0.58}, & m_{\mathrm{DM}} > 1\,\mathrm{GeV} \\ 4.0 \times 10^{-7} \left(\frac{m_{\mathrm{DM}}}{1\,\mathrm{GeV}}\right)^{0.35}, & m_{\mathrm{DM}} < 1\,\mathrm{GeV} \end{cases} \tag{1.8}$$

- **Small Self-Interactions:** The standard ΛCDM cosmology assumes that the dark matter is collisionless. However, small dark matter self-interactions can help resolve existing small-scale structure issues [17, 18]. Current limits on the self-interaction cross-section are $\sigma_{\mathrm{DM-DM}}/m_{\mathrm{DM}} < 0.48\,\mathrm{cm}^2/\mathrm{g}$ come from merging galaxy clusters [11] and the ellipticity of galaxies obtained from X-ray observations [19].

- **Cold:** Dark matter is required to be non-relativistic at the time of structure formation. At most, a small component of the dark matter can be warm (semi-relativistic).

A selection of the more prominent dark matter candidates is shown in Fig. 1.3. The key features of a few of these models are discussed below.

WIMPs

Weakly Interacting Massive Particles (WIMPs) are a class of dark matter candidates that generically have masses and interaction strengths around the weak scale. Many extensions of the SM naturally predict the existence of such a particle, with famous examples being the lightest supersymmetric particle in supersymmetric theories [21], or the lightest stable Kaluza-Klein mode in theories with extra dimensions [22].

Nowadays, WIMP dark matter is used almost synonymously to mean thermal relic, referring to a species whose relic abundance is produced thermally in the early universe through the freeze-out mechanism [23]. In this paradigm, the WIMP is initially in thermal equilibrium with the Standard Model bath. This equilibrium is maintained as long as the interaction rates of the WIMP with the bath, denoted Γ, remain faster than the Hubble expansion of the universe, H. As the universe continues to expand, the temperature of the bath drops slowing down the interaction

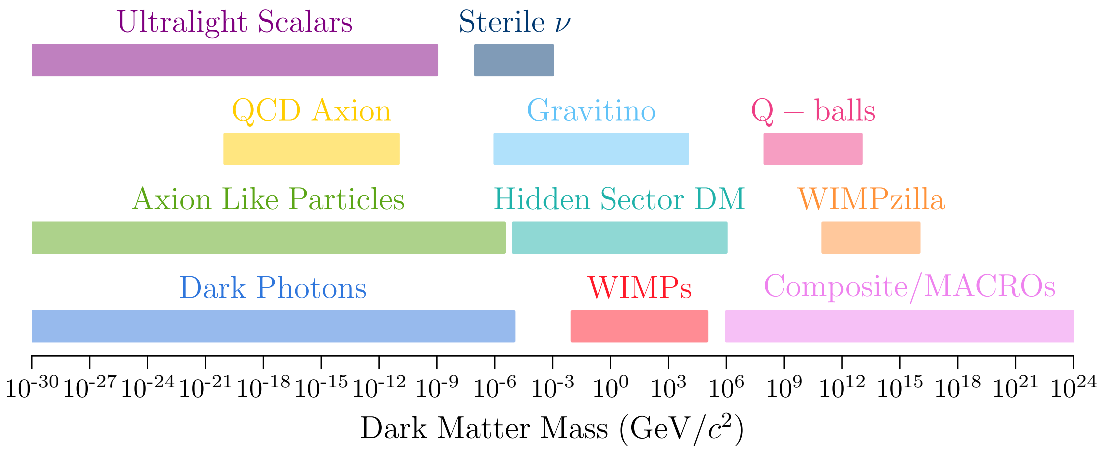

Figure 1.3: Illustrative landscape of dark matter models and the mass range for which they predict a valid candidate. Details can be found in the Dark Matter chapter of the PDG [20].

rates. Eventually, the expansion rate overtakes the interaction rates, $\Gamma/H \lesssim 1$, and the interactions "freeze-out" causing the WIMP to fall out of equilibrium with the bath. At this point, the WIMPs can no longer efficiently annihilate, and their abundance gets "frozen-in" to the value it had at freeze-out, leading to the abundance observed today.

A cold thermal relic, such as dark matter, will freeze-out after it has become non-relativistic. In this scenario, the interaction rates become Boltzmann suppressed[2], and the species rapidly freezes-out. The relic density is therefore sensitive to the annihilation cross-section of the species, $\langle \sigma_{\text{ann}} v \rangle$. More efficient annihilations correspond to larger cross-sections, resulting in the species remaining in equilibrium for longer times. This allows the number density to continue following the exponentially decreasing Boltzmann distribution and yield a smaller relic abundance. The evolution of the abundance of a Majorana fermion WIMP of mass $m_{\text{WIMP}} = 100\,\text{GeV}$ is shown in Fig. 1.4 for three values of the annihilation cross-section. A simple expression relating the annihilation cross-section and the

[2]The number density of a non-relativistic species in thermal equilibrium with the bath will follow $n \propto (mT_{\text{bath}})^{3/2} \exp\left(-m/T_{\text{bath}}\right)$. Once the temperature falls below the mass of the species, the number density becomes exponentially suppressed. This is what is known as "Boltzmann suppression".

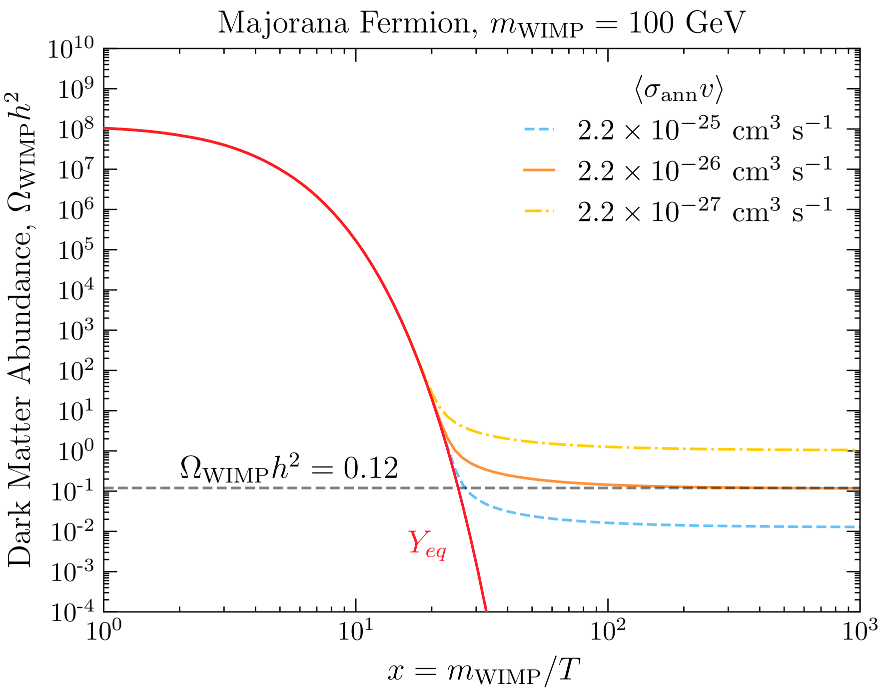

Figure 1.4: Evolution of the DM abundance as a function of $x = m_{\rm DM}/T$. The red line tracks the abundance if the WIMP remained in equilibrium with the bath. The relic abundance for three different annihilation cross-sections is shown in blue, orange, and yellow for $\langle\sigma_{\rm ann}v\rangle = 2.2\times10^{-25}$, 2.2×10^{-26}, and 2.2×10^{-27} cm^{3}s^{-1} respectively.

abundance that is correct to $\sim 5\%$ can be obtained [24]

$$\Omega_{\rm DM}h^2 = \frac{10^{-27}\,{\rm cm}^3\,{\rm s}^{-1}}{\langle\sigma_{\rm ann}v_\chi\rangle}\frac{x_*}{g_*^{1/2}} \tag{1.9}$$

$$\sim 0.12\left(\frac{2.2\times10^{-26}\,{\rm cm}^3\,{\rm s}^{-1}}{\langle\sigma v\rangle}\right), \quad m_\chi \gtrsim 10\,{\rm GeV}, \tag{1.10}$$

where $x_* = m_{\rm WIMP}/T_*$ is evaluated as an intermediate temperature between equilibrium and freeze-out, with g_* the effective number of relativistic degrees of freedom present at this time.

The allowed mass range for a thermal WIMP is between $10\,{\rm MeV} \lesssim m_{\rm WIMP} \lesssim 100\,{\rm TeV}$. Lighter WIMPs will have non-negligible contributions to the effective number of neutrino species at the time of Big Bang Nucleosynthesis, $N_{eff}^{\rm BBN} = 3.044$ [25]. The CMB probes $N_{\rm eff}$ at the time of recombination and can be combined with the BBN result leading to the value $N_{eff} = 2.99 \pm 0.17$ [12]. At the high end of this range, masses larger than $\sim 100\,{\rm TeV}$ are excluded from partial wave unitarity [26].

Axions

The axion originally arose from the Pecci-Quinn solution to the Strong CP problem. This refers to the lack of observed CP-violation in the QCD sector of the Standard Model that arises from the topological term in the Lagrangian

$$\mathcal{L}_{\theta_{\mathrm{QCD}}} = \frac{g_s^2}{32\pi}\theta_{\mathrm{QCD}} G_{\mu\nu}\tilde{G}^{\mu\nu}, \tag{1.11}$$

where g_s is the QCD coupling constant, $G_{\mu\nu}$ is the gluon field strength tensor and $\tilde{G}^{\mu\nu}$ is its dual. This term generates an electric dipole moment for the neutron (nEDM) that has yet to be observed experimentally. The current upper bound on the nEDM is $|d_n| < 0.18 \times 10^{-26}\, e\,\mathrm{cm}$ [27] and can be translated to an upper bound on the CP-violating QCD θ-parameter such that $|\theta_{QCD}| \lesssim 10^{-10}$, raising questions as to why this value seems to be fine-tuned to such a small value.

The Peccei-Quinn solution to this problem introduces a new, anomalous, global $U(1)_{\mathrm{PQ}}$ symmetry and promotes θ_{QCD} to be a dynamical field. Wilczck [28] and Weinberg [29] showed that the axion emerges as the pseudo-Goldstone boson associated with the breaking of $U(1)_{\mathrm{PQ}}$. Though the original axion was quickly out, many modern extensions of the SM predict the existence of a QCD axion. Two of the most prominent UV completions of the axion are the KSVZ [30, 31] and DFSZ [32, 33] models. In these models, the axion produced in the early Universe can serve the role of cold dark matter today. This makes it a very compelling dark matter candidate, as it solves two of the biggest mysteries of physics in one neat package.

However, solving the Strong CP problem can be rather restrictive on the model parameters. For example, the QCD axion's coupling to the photon, $g_{a\gamma\gamma}$, is not a free parameter and depends on the scale at which the PQ symmetry is broken. Many models introduce a light pseudoscalar particle, say a, that couples to the photon in the same way as the QCD axion,

$$\mathcal{L}_{a\gamma\gamma} = -\frac{g_{a\gamma\gamma}}{4} a F_{\mu\nu}\tilde{F}^{\mu\nu}, \tag{1.12}$$

but is not associated with a solution to the Strong CP problem. Such pseudoscalars are known as "Axion Like Particles" (ALPs) and can make a good dark matter candidate.

Primordial Black Holes

Primordial black holes (PBHs) are formed during the early Universe through various mechanisms. The simplest mechanism predicts that PBHs are produced from the gravitational collapse of superhorizon density fluctuations seeded during inflation [34–36]. Unlike black holes that originate from stellar collapse, which have masses $\gtrsim 3M_\odot$, the mass of a PBH can be arbitrary. PBHs can also make a good dark matter candidate, satisfying all the criteria points outlined above. In fact, PBHs with a mass between $\sim (10^{-17} - 10^{-12})\ M_\odot$, dubbed “asteroid mass PBHs”, can account for 100% of the dark matter content in the Universe [37]. Outside this range, PBHs can still make up a small fraction of dark matter [38].

1.2.1 Dark Matter in an Effective Fields Theory Framework

For a dark matter candidate to be truly compelling, it should be able to be embedded into an ultraviolet (UV) complete theory. Such theories are renormalisable[3] and gauge invariant under the SM gauge group $SU(3)_{\text{colour}} \otimes SU(2)_L \otimes U(1)_Y$. This allows them to be predictive up to arbitrarily high energies. These theories are typically quite complex, requiring the introduction of multiple new fields and many more free parameters. As an example, consider the phenomenological Minimal Supersymmetric Standard Model (pMSSM) [38] in which the lightest neutralino[4] can be a thermal WIMP dark matter candidate. In this theory, 19 free parameters are introduced on top of the free parameters in the SM, requiring 38 independent experimental observations to fully constrain the model. Given that at this time, all good dark matter candidates are equally likely to be the correct one, a model-independent approach to interpreting experimental results is desirable. This is achieved by describing the interactions of dark matter with the SM through an effective field theory (EFT).

1.2.2 Overview of Effective Field Theory

Modern physics can be thought of as a ladder of theories that are designed to describe the physics present at a given energy (or length) scale. For example, Newtonian mechanics is a sufficient description of the physics experienced in our everyday lives. However, in situations where the energy is comparable to the mass of the system, Newtonian mechanics breaks down, and Special Relativity must be used to describe the physics. In particle physics, the Standard Model provides an excellent description of particle interactions up to the energies reached at LHC, 13.6 TeV, and perhaps even further beyond. However, even it is expected to break

[3] In a renormalisable theory, the infinities that arise from UV divergences can be absorbed by fixing a finite number of parameters to experimentally observed values.

[4] The neutralinos are the mass eigenstates of the supersymmetric partners of the neutral gauge bosons and the higgsino.

down at higher energy scales, in particular at the Planck scale where a quantum theory of gravity is required. Hence, both Newtonian mechanics and the SM are low-energy, effective descriptions of a more complete theory.

This philosophy based on only describing the physics relevant below some energy scale, Λ, is the core principle of EFTs. The Lagrangian for the effective theory only contains the degrees of freedom that can be produced below the scale Λ, i.e. fields that have masses less than this scale. This low-energy regime described by the EFT is often called the infrared (IR) regime.

In general, the EFT Lagrangian will contain renormalizable terms, $\mathcal{L}_{\text{renorm.}}$, built out of operators that have mass dimension ≤ 4, as well as operators with mass dimension $n > 4$, $\mathcal{O}_i^{(n)}$, that encapsulate the contributions from the UV physics. Each of these higher dimensional operators will be suppressed by the scale of new physics, Λ^{4-n}. The effective Lagrangian can then be written as

$$\mathcal{L}_{\text{EFT}} = \mathcal{L}_{\text{renorm.}} + \sum_{n>4} \sum_{i=1}^{j_n} \frac{C_i^{(n)}(\tilde{\mu})}{\Lambda^{n-4}} \mathcal{O}_i^{(n)}, \tag{1.13}$$

where we sum over all j_n operators present at mass dimension n. The expansion coefficients, $C_i^{(n)}$, are the Wilson coefficients that contain the effects of the UV physics. In general, the Wilson coefficients run with the energy scale they are evaluated at, $\tilde{\mu}$, described by the renormalisation group equations (RGEs). The sum over the mass dimension of the operators is typically terminated at some sensible value, as higher dimensional operators get increasingly suppressed by the cutoff scale Λ.

The series of operators in Eq. 1.13 can be constructed in two different ways. First, we assume some prior knowledge of the underlying UV theory. Then, for a given energy scale Λ, the heavy degrees of freedom are known, and can be explicitly integrated out. There are various methods for performing this process, the simplest being expanding the propagator of the heavy fields in powers of the momenta over the heavy mass, $(p/M)^2$. For the simple case of a heavy scalar, this corresponds to

$$\frac{i}{p^2 - M^2} = \frac{-i}{M^2} \left(\frac{1}{1 - (p/M)^2} \right) \approx \frac{-i}{M^2} \left(1 + \left(\frac{p}{M} \right)^2 + \mathcal{O}\left(\left(\frac{p}{M} \right)^4 \right) \right). \tag{1.14}$$

An alternate, more robust method, is to replace the heavy fields in the Lagrangian with their classical equations of motion. The resulting effective Lagrangian will contain all the operators generated by the UV theory at tree-level. Constructing an EFT in this way is called the *top-down* method.

The second method of constructing an EFT is to be agnostic to the UV physics and write down all possible operators that can be constructed from the IR degrees of freedom. These operators must obey the symmetries of the IR theory, as well as any other constraints one wishes to impose[5]. This is the *bottom-up* approach, offering a more model-independent approach than the top-down method. The Wilson coefficients in this case will be arbitrary functions of the energy scale, determined by solving the RGEs.

In general, the parameter space of the EFT will be lower dimensional than those of the corresponding UV models. This allows for an easier comparison with experimental results, as there are fewer parameters to fit to the data. Once the Wilson coefficients have been constrained at the low energy scale of the experiments, they can be matched to the coefficients generated by some UV theory at another scale, thereby constraining the UV parameters.

1.2.3 Dimension 6 EFT Operators for Dirac Fermion Dark Matter

This work will focus on dimension 6 EFT operators that describe the interactions of Dirac fermion dark matter with standard model fermions. These operators will have a structure

$$\mathcal{L}^{(6)}_{\mathrm{EFT}} \sim \frac{1}{\Lambda^2}(\bar{\chi}\Gamma_{\mathrm{DM}}\chi)(\bar{f}\Gamma_{\mathrm{SM}}f), \tag{1.15}$$

where the Γ_i determines the Lorentz structure of the interaction by taking appropriate combinations from the set

$$\Gamma_i \in \{1, i\gamma_5, \gamma^\mu, i\gamma^\mu\gamma^5, \sigma^{\mu\nu}, i\sigma^{\mu\nu}\gamma^5\}. \tag{1.16}$$

For example, the case of $\Gamma_\chi = \Gamma_{\mathrm{SM}} = 1$ yields scalar currents for both the DM and SM fermions and would correspond to integrating out a heavy scalar mediator in the UV theory. Under the assumption of minimal flavour violation (MFV)[6] are ten such operators at dimension six that form a linearly independent basis. These are given in Table 1.1, along with spin-averaged squared matrix element for dark matter scattering with a fermion. The operators are classified based on the Lorentz nature of the SM fermion bilinear; D1-2: Scalar (S), D3-4: Pseudoscalar (P), D5-6: Vector (V), D7-8: Axial-vector (A), and D9-10: tensor (T). The coupling constant, g_f, for operators that involve the S and P fermion bilinears (operators D1-4) are normalised by the corresponding Yukawa couplings. This is because, in a UV complete theory, these bilinears would couple to the new scalar/pseudoscalar field that mediates the interactions with the dark matter. In many models, this new

[5]For example, one might require that no flavour changing processes are present at dimension 5, despite such operator being allowed by the symmetries.

[6]MFV is the assumption that the only source of flavour violation in the quark sector comes from the SM Yukawa matrices and not from any new physics introduced at a higher scale.

Name	Operator	g_f	$\lvert\overline{M}(s,t,m_i)\rvert^2$
D1	$\bar{\chi}\chi\,\bar{f}f$	$\frac{y_f}{\Lambda_f^2}$	$g_f^2\frac{\left(4m_\chi^2-t\right)\left(4m_\chi^2-\mu^2t\right)}{\mu^2}$
D2	$\bar{\chi}\gamma^5\chi\,\bar{f}f$	$i\frac{y_f}{\Lambda_f^2}$	$g_f^2\frac{t\left(\mu^2t-4m_\chi^2\right)}{\mu^2}$
D3	$\bar{\chi}\chi\,\bar{f}\gamma^5f$	$i\frac{y_f}{\Lambda_f^2}$	$g_f^2t\left(t-4m_\chi^2\right)$
D4	$\bar{\chi}\gamma^5\chi\,\bar{f}\gamma^5f$	$\frac{y_f}{\Lambda_f^2}$	$g_f^2t^2$
D5	$\bar{\chi}\gamma_\mu\chi\,\bar{f}\gamma^\mu f$	$\frac{1}{\Lambda_f^2}$	$2g_f^2\frac{2\left(\mu^2+1\right)^2m_\chi^4-4\left(\mu^2+1\right)\mu^2sm_\chi^2+\mu^4\left(2s^2+2st+t^2\right)}{\mu^4}$
D6	$\bar{\chi}\gamma_\mu\gamma^5\chi\,\bar{f}\gamma^\mu f$	$\frac{1}{\Lambda_f^2}$	$2g_f^2\frac{2\left(\mu^2-1\right)^2m_\chi^4-4\mu^2m_\chi^2\left(\mu^2s+s+\mu^2t\right)+\mu^4\left(2s^2+2st+t^2\right)}{\mu^4}$
D7	$\bar{\chi}\gamma_\mu\chi\,\bar{f}\gamma^\mu\gamma^5f$	$\frac{1}{\Lambda_f^2}$	$2g_f^2\frac{2\left(\mu^2-1\right)^2m_\chi^4-4\mu^2m_\chi^2\left(\mu^2s+s+t\right)+\mu^4\left(2s^2+2st+t^2\right)}{\mu^4}$
D8	$\bar{\chi}\gamma_\mu\gamma^5\chi\,\bar{f}\gamma^\mu\gamma^5f$	$\frac{1}{\Lambda_f^2}$	$2g_f^2\frac{2\left(\mu^4+10\mu^2+1\right)m_\chi^4-4\left(\mu^2+1\right)\mu^2m_\chi^2(s+t)+\mu^4\left(2s^2+2st+t^2\right)}{\mu^4}$
D9	$\bar{\chi}\sigma_{\mu\nu}\chi\,\bar{f}\sigma^{\mu\nu}f$	$\frac{1}{\Lambda_f^2}$	$8g_f^2\frac{4\left(\mu^4+4\mu^2+1\right)m_\chi^4-2\left(\mu^2+1\right)\mu^2m_\chi^2(4s+t)+\mu^4(2s+t)^2}{\mu^4}$
D10	$\bar{\chi}\sigma_{\mu\nu}\gamma^5\chi\,\bar{f}\sigma^{\mu\nu}f$	$\frac{i}{\Lambda_f^2}$	$8g_f^2\frac{4\left(\mu^2-1\right)^2m_\chi^4-2\left(\mu^2+1\right)\mu^2m_\chi^2(4s+t)+\mu^4(2s+t)^2}{\mu^4}$

Table 1.1: Dimension 6 EFT operators [39] for the coupling of Dirac DM to fermions (column 2), together with the squared matrix elements DM-fermion scattering (column 5), where s and t are Mandelstam variables, $\mu = m_\chi/m_T$, and m_T is the target mass.

field will mix with the SM Higgs field, leading to couplings that depend on the fermion masses. The remaining bilinears have coupling constants that depend only on the cutoff scale, Λ_f.

1.2.4 From DM-Quark to DM-Nucleon Interactions

The operators in Table 1.1 describe dark matter interactions at the quark level, as these are the degrees of freedom most models are formulated with. However, we will primarily be interested in dark matter scattering with baryons, which requires taking the matrix element of the quark operators between baryon states, i.e. $\langle\mathcal{B}|\bar{q}\,\Gamma_q q|\mathcal{B}\rangle$. These matrix elements can be calculated through the application of Chiral Perturbation Theory (ChPT), giving a baryon level EFT. The operators of this EFT will have the same form as those in Table 1.1, with the obvious replacement of $f \to \mathcal{B}$, as well as additional form factors that take into account the structure of the baryons.

The required form factors for each operator have been calculated at zero momentum transfer in Ref. [40] and are given by

$$c_{\mathcal{B}}^{S}(0) = \frac{2m_{\mathcal{B}}^2}{v^2}\left[\sum_{q=u,d,s} f_{T_q}^{(\mathcal{B})} + \frac{2}{9} f_{TG}^{(\mathcal{B})}\right]^2, \tag{1.17}$$

$$c_{\mathcal{B}}^{P}(0) = \frac{2m_{\mathcal{B}}^2}{v^2}\left[\sum_{q=u,d,s} \left(1 - 3\frac{\overline{m}}{m_q}\right)\Delta_q^{(\mathcal{B})}\right]^2, \tag{1.18}$$

$$c_{\mathcal{B}}^{V}(0) = 9, \tag{1.19}$$

$$c_{\mathcal{B}}^{A}(0) = \left[\sum_{q=u,d,s} \Delta_q^{(\mathcal{B})}\right]^2, \tag{1.20}$$

$$c_{\mathcal{B}}^{T}(0) = \left[\sum_{q=u,d,s} \delta_q^{(\mathcal{B})}\right]^2, \tag{1.21}$$

where $v = 246$ GeV is the vacuum expectation value of the SM Higgs field, $\mathcal{B}$ is the baryonic species, $\overline{m} \equiv (1/m_u + 1/m_d + 1/m_s)^{-1}$ and $f_{T_q}^{(\mathcal{B})}$, $f_{TG}^{(\mathcal{B})} = 1 - \sum_{q=u,d,s} f_{T_q}^{(\mathcal{B})}$, $\Delta_q^{(\mathcal{B})}$ and $\delta_q^{(\mathcal{B})}$ are the hadronic matrix elements, determined either experimentally or by lattice QCD simulations[7]. The specific values of these matrix elements for various baryons are provided in Appendix **ADD APPENDIX**.

The assumption of zero-momentum transfer is valid when considering interactions with momentum transfers $\lesssim 1\,\mathrm{GeV}$, such as in direct detection experiments. Once the momentum transfer exceeds this, the internal structure of the baryon begins to be resolved, and an additional momentum-dependent form factor is required to account for this [41],

$$F_{\mathcal{B}}(t) = \frac{1}{(1 - t/Q_0)^2}, \tag{1.22}$$

where t is the Mandelstam variable, and Q_0 is an energy scale that depends on the hadronic form factor. For simplicity, we will conservatively take $Q_0 = 1\,\mathrm{GeV}$ for all operators. Putting everything together, the squared coupling constants for dark matter-baryon interactions are obtained by making the replacement

$$g_f^2 \to \frac{c_{\mathcal{B}}^I(t)}{\Lambda_q^4} \equiv \frac{1}{\Lambda_q^4} c_{\mathcal{B}}^I(0) F_{\mathcal{B}}^2(t), \qquad I \in S, P, V, A, T, \tag{1.23}$$

in the matrix elements in the final column of Table 1.1.

[7]The superscript letters S, P, V, A and T stand for Scalar, Pseudoscalar, Vector, Axial-vector and Tensor interactions respectively. The corresponding operators are: D1-2 for S; D3-4 for P; D5-6 for V, D7-8 for A; and D9-10 for T.

1.3 Current Status of Dark Matter Constraints

In broad terms, there are three main ways that we can search for evidence of dark matter, often termed “make it, shake it or break it”. “Make it” refers to dark matter being produced at colliders; “break it” to searching for dark matter annihilation signals; and “shake it” to direct detection of dark matter scattering. An illustrative way of depicting these processes is shown in Fig. add usual diagram. This section discusses the current status of these detection methods.

1.3.1 Collider Bounds

If dark matter is produced in a collider, it will simply leave the detector without depositing any energy. In order to determine if such an invisible particle was produced, conservation of energy-momentum is used to determine if there are any events that are missing energy. In practice, what is searched for is missing momentum that is transverse to the beamline.

Currently, dark matter has not been observed to be produced in particle colliders. This non-observation has instead been used to constrain the dark matter mass and production cross sections or couplings of various models. These limits are typically interpreted in a model-dependent manner, as different dark matter - Standard model couplings can significantly alter the production rates. As mentioned above, EFTs can be used to explore a variety of interactions in a somewhat model-independent way. However, many applications of this nature did not hold up to scrutiny, as the EFTs were being applied at energies outside their regions of validity [42–45], and so care is needed when applying such methods.

The ATLAS and CMS experiments at the LHC have performed analyses on various dark matter production mechanisms, including the exchange of a Z/Z' or Higgs, EFTs and heavy mediators, and mono-jet searches [46][8]. Collider searches also offer complimentary probes of the dark matter-nucleon scattering cross-section [47].

It is important to note that an observation of an invisible massive particle at a collider is not enough to infer that it is dark matter. Such an observation only tells us that such a particle exists but nothing about its abundance, meaning it could just be a sub-component of a larger dark sector. In order to identify whether or not this was a dark matter detection, complimentary observations from direct or indirect detectors would be required.

[8] These searches refer to a single jet being produced alongside a pair of dark matter particles. This jet could be of Standard Model or dark sector origin, with the latter commonly referred to as “mono-X” searches.

1.3.2 Direct Detection Searches

Direct detection experiments vary wildly depending on the dark matter mass range they are trying to probe. For ALP dark matter that is wavelike, haloscope experiments such as ADMX [48] and MADMAX [49] attempt to convert ALPs to photons via the Primakoff effect. Searches for WIMP dark matter look for the dark matter scattering with some detector material, causing it to recoil and release some energy. Given our focus on WIMP dark matter, this section will review the experimental status of these detectors.

The differential rate at which the incoming flux of dark matter will scatter within a detector with N_T targets, as a function of the recoil energy, E_R, is given by

$$\frac{dR(E_R,t)}{dE_R} = N_T \frac{\rho_{\rm DM}}{m_{\rm DM}} \int_{v>v_{\rm min}}^{v_{\rm esc}} v f(\vec{v}+\vec{v}_E) \frac{d\sigma}{dE_R}\, d^3v, \tag{1.24}$$

and depends on the quantities:

- $v_{\rm min}$ is the minimum dark matter velocity required by kinematics for a scattering event to occur;
- $v_{\rm esc} = 528\,{\rm km\,s^{-1}}$ is the Milky Way escape velocity;
- $\vec{v}_E$ is the velocity of the Earth through the dark matter halo[9];
- $f(\vec{v}-\vec{v}_E)$ is the dark matter velocity distribution in the Earth's frame;
- $d\sigma/dE_R$ is the differential scattering cross-section.

Given the low interaction rate of dark matter, the expected event rate in detectors is very low, around one event per day, per kilogram of target material, per kiloelectronvolt deposited. Having such a low event rate requires the detector to be situated in an extremely low background environment, such as underground laboratories.

Direct detection experiments aim to probe two main types of dark matter interactions: Spin-dependent (SD) and spin-independent (SI) scattering. SD interactions couple to the overall spin of the target, while SI interactions are agnostic to this. Therefore, experiments searching for SI interactions benefit from using nuclei with a large atomic number, A, as the interaction cross-section will involve a coherent sum over all nucleons. This leads to an A^2 enhancement of SI interactions compared to the SD counterpart.

[9]This accounts for the orbit of the Earth around the Sun, which induces an annual modulation in the flux of DM.

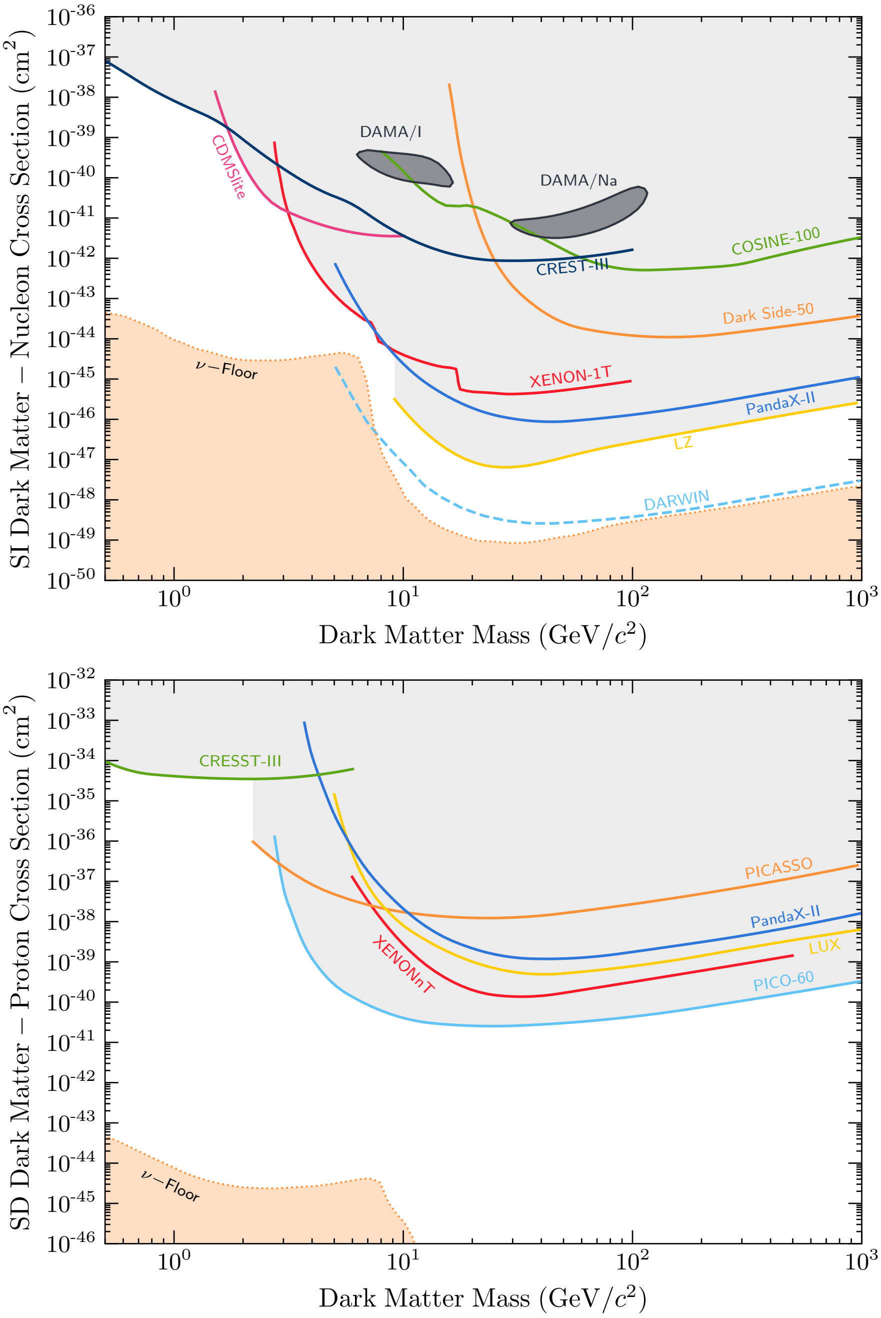

Figure 1.5: Current status of direct detection searches for dark matter. **Top:** Spin-independent dark matter-nucleon scattering. **Bottom:** Spin-dependent dark matter-proton scattering.

The current leading constraints on the dark matter-nucleon scattering cross-section are shown in Fig. 1.5, with SI in the top panel and SD in the bottom. The SI limits are set by liquid noble gas experiments (LZ [50], XENON-1T [51], PandaX-II [52], and DarkSide-50 [53]), solid-state cryogenic detectors (CRESST-III [54], CDMSlite [55], with projected DARWIN sensitivities [56]), and room temperature crystals (DAMA/LIBRA [57], and COSINE-100 [58]).

The SD experiments require their targets to carry non-zero spin for the dark matter to couple to. ^{19}F is the favourable choice proton scattering, as it has an unpaired proton giving it its overall spin. The leading constraints come from superheated liquid experiments such as the PICO-60 [59] as well as PICASSO [60] In terms of the SD proton scattering shown in Fig 1.5, These interactions are also searched for by many of the same experiments in the SI case, with the inclusion of LZ's predecessor LUX [61].

The orange dashed line represents the neutrino floor[10], a theoretical lower limit on the discoverability of WIMP-like dark matter. In this region of parameter space, doctors will become sensitive to the irreducible background from neutrino scattering, which will produce signals almost indistinguishable from a true dark matter interaction. A significant amount of effort is being put toward overcoming this hindrance, with the main strategy being to take advantage of the directionality of dark matter flux [63].

Many experiments begin to lose sensitivity to low-mass dark matter ($m_{\rm DM} \lesssim 10\,{\rm GeV}$) as the targets recoil with energies below the detector threshold. Current energy thresholds can reach as low as $\sim \mathcal{O}(100\,{\rm eV})$, which is on the same order of magnitude as the recoil energy due to a 1 GeV dark matter collision. The sensitivity also falls off at a slower rate at larger masses, though this is due to the number of dark matter particles that pass through the detector given $N_{\rm DM} = \rho_{\rm DM}/m_{\rm DM}$, and the dark matter density is known to be $0.4\,{\rm GeV\,cm^{-3}}$.

Direct detection limits also assume that the scattering cross-section is independent of the dark matter velocity and momentum transfer in the interaction. Given that the local dark matter dispersion velocity is predicted to be $v_d = 270\,{\rm km\,s^{-1}} \approx 10^{-3}c$, a back-of-the-envelope estimation for the momentum transfer gives $q_{\rm tr} \lesssim 100\,{\rm MeV}$. Therefore, cross-sections proportional to $v_{\rm DM}$ or $q_{\rm tr}$ will result in significantly lower event rates and hence much weaker limits than the unsuppressed interactions.

This leads us to indirect detection methods, which can provide complementary probes to direct detection while also exploring interactions that are difficult, if not impossible, for terrestrial-based detectors to observe.

[10]Calling this the "neutrino fog" rather than floor has been gaining traction in recent years [62]

1.3.3 Indirect Detection

Indirect detection experiments aim to infer the presence of dark matter through its annihilation or decay into Standard Model states. These searches look for anomalies in astrophysical data, though dark matter accumulating within the Earth's core can also produce a detectable signal [64, 65]. The signals searched for include:

- Gamma-rays at terrestrial-based telescopes such as HESS [66–68], VERITAS [69–71], MAGIC [72, 73] and HAWC [74–77] as well as the Fermi-LAT [78–82] satellite;
- Neutrino signals at IceCube [83, 84], ANTARES [65, 85, 86], Super-K [87–89], and will be searched for at the upcoming Hyper-K [90–92], JUNO [93] experiments.
- Cosmic-Rays by the AMS-02 experiment [94, 95]

Signals from dark matter annihilation are best searched for by looking at regions where the dark matter density is expected to be high, boosting the annihilation rate. Natural places to look include the Galactic Centre [96, 97], dwarf-spheroidal galaxies [98], and celestial bodies where dark matter can accumulate over time.

1.4 Dark Matter Signals from the Sun

Stars have a rich history of being used as astrophysical laboratories to help in the search for dark matter. Depending on the type of dark matter being searched for, there are various signals one can look for. Light bosonic dark matter, such as ALPs and dark photons, can be produced within the plasma of stars, altering the energy transport properties within them. This can ultimately lead to deviations in the evolution of the star, which can be used to place some of the strongest constraints on these models [99–101]. WIMP-like dark matter in the halo that couples to visible matter can scatter within the stars as they pass through. If the dark matter loses enough energy in these interactions, it can become gravitationally bound to the object and a population of dark matter will be accumulated within the star over time [23, 102–105].

This idea of WIMPs accumulating within the cores of stars has been applied extensively to the star closest to us, the Sun. The formalism for stellar capture of dark matter was set up by Gould [103, 104, 106] in the late 80's, and has remained quite successful to this day, with many authors continuing to build upon these foundations over time [105, 107, 108].

Once the dark matter is captured, it will continue to scatter with the stellar constituents until it thermalises within the core of the Sun, collecting with an isothermal sphere. The radius of this sphere can be found by applying the virial theorem, with the gravitational potential given by

$$\Phi(r) = -\int_r^\infty \frac{GM_\odot(r')}{r'^2} dt', \tag{1.25}$$

$$\approx \frac{2}{3}\pi G\rho_{\odot,c} r^2, \tag{1.26}$$

assuming the density of the Sun within this region is constant, $\rho_{\odot,c}$. The resulting radius is

$$r_{\rm iso}^2 = \frac{3T_\odot}{2\pi G m_{\rm DM} \rho_\odot}. \tag{1.27}$$

The dark matter number density will follow a Gaussian profile,

$$n_{\rm iso}(r) = n_0 \exp\left(-\frac{m_{\rm DM}\Phi(r)}{T_\odot}\right), \tag{1.28}$$

$$= n_0 \exp(-r^2/r_{\rm iso}^2), \tag{1.29}$$

where n_0 is a normalistion constant fixed by requiring that the total number of dark matter particles is

$$N_{\rm DM} = \int d^3r\, n_{\rm iso}(r). \tag{1.30}$$

In addition, dark matter velocity will follow a Maxwell-Boltzmann distribution,

$$f_{\rm MB}(v) = 4\pi \left(\frac{m_{\rm DM}}{4\pi T_\odot}\right)^{3/2} v^2 \exp\left[-\frac{m_{\rm DM} v^2}{4T_\odot}\right]. \tag{1.31}$$

The time evolution of the total number of dark matter particles within the Sun is governed by three processes. Capture acts to increase the number over time, while annihilation and evaporation will reduce this number over time. This is described by the differential equation

$$\frac{dN_{\rm DM}}{dt} = C - EN_{\rm DM} - AN_{\rm DM}^2, \tag{1.32}$$

where C and E are the capture and evaporation rates respectively, with A being related to the annihilation rate, $\Gamma_{\rm ann}$ through

$$\Gamma_{\rm ann} = \frac{1}{2}\int dr^3 n_{\rm iso}^2(r)\langle\sigma_{\rm ann} v\rangle \tag{1.33}$$

$$\equiv \frac{1}{2} AN_{\rm DM}^2, \tag{1.34}$$

where the factor of 1/2 accounts for each annihilation removing two dark matter particles from the Sun.

In this context, evaporation refers to the process in which dark matter can be ejected back out of the Sun by up-scattering[11] off an energetic constituent. This becomes increasingly important for lighter dark matter masses, as less energy will be required to boost the dark matter above the local escape velocity. Below a certain mass, the capture and evaporation come into equilibrium, and a net-zero amount of dark matter is contained within the Sun. This critical mass places a lower bound on the dark matter mass that can be probed through stellar capture and is named the evaporation mass, $m_{\rm evap}$.

There are three regimes we are interested in solving this equation for. First, if both evaporation and annihilation are negligible, or at least have characteristic timescales longer than the age of the Sun, the solution is simply,

$$N_{\rm DM}(t) = Ct. \tag{1.35}$$

This simply tells us that the dark matter will continue to accumulate over time.

Next, assume that annihilations are negligible ($A = 0$) and that the capture and evaporation rates are constant in time. The result is

$$N_{\rm DM}(t) = Ct\left(\frac{1-e^{-Et}}{Et}\right), \tag{1.36}$$

where the first factor is the number of captured dark matter if evaporation is negligible. From this, we can estimate the evaporation mass by asking when the evaporation rate is large enough to cause a significant reduction in the number of captured dark matter particles relative to the $E = 0$ case. This can be expressed formally as [107]

$$\frac{1}{N_{\rm DM}(m_{\rm evap})}\left|N_{\rm DM}(m_{\rm evap}) - \frac{C(m_{\rm evap})}{E(m_{\rm evap})}\right| = \alpha, \tag{1.37}$$

where α is the fraction of evaporated dark matter.

The second regime is when evaporation can be neglected, i.e. $m_{\rm DM} \gtrsim m_{\rm evap}$. The solution in this case is

$$N_{\rm DM}(t) = \sqrt{\frac{C}{A}}\tanh(\sqrt{CA}t) \xrightarrow{t\to\infty} \sqrt{\frac{C}{A}} \tag{1.38}$$

The signals searched for depend on whether the dark matter can annihilate or not. If the dark matter is asymmetric, it cannot annihilate, and we can set

[11] Up-scattering refers to the interactions in which the dark matter gains rather than loses energy.

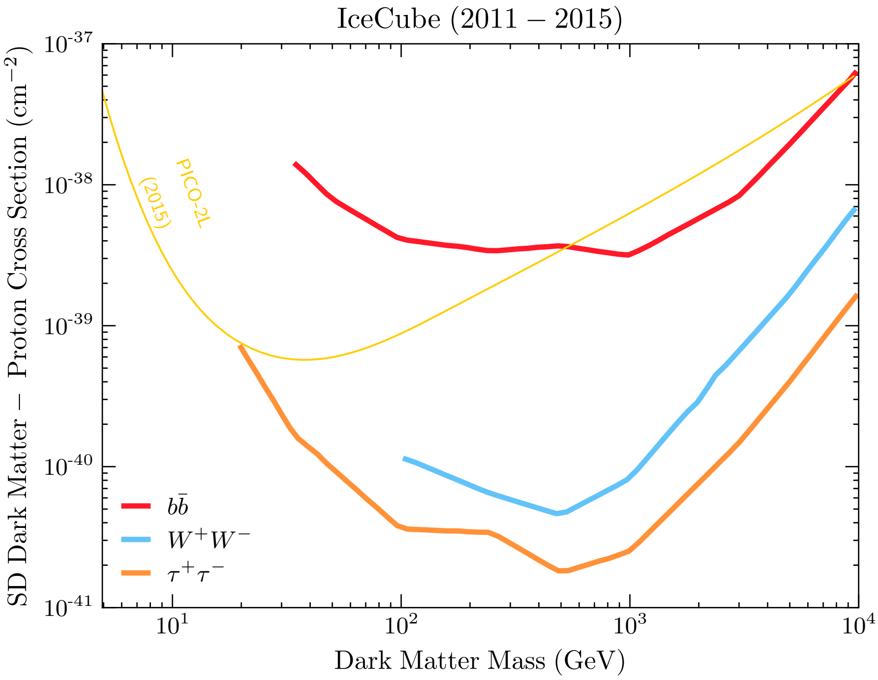

Figure 1.6: Limits on the SD dark matter-proton cross-section from the IceCube collaboration assuming 100% branching fraction to $b\bar{b}$ (red), W^+W^- (blue) or $\tau^+\tau^-$ (orange) final states. Also shown is the result from the PICO-2L DD experiment. This plot was recreated with data taken from Ref. [83].

$A = 0$ in Eq. 1.32. This leads to the population continuing to grow over time, with $N_{\mathrm{DM}}(t) = Ct$ if evporation is negligible. A large enough population of captured dark matter can alter the energy transport within the Sun, leading to the modifications of the solar neutrino flux, or even the solar structure itself [93, 109–111].

Instead, if the dark matter can annihilate an equilibrium will eventually be reached between the capture and annihilation rates, and the total number of dark matter particles will be constant. If the annihilation products can escape the Sun, they can be searched for by various experiments depending on the nature of the final states. These could be neutrinos produced from the decays of other charged annihilation products [85–87, 112, 113], or to some other long-lived state that can escape the Sun and decay into visible states [114–118].

In comparison to DD searches, interpretation of indirect detection data will require additional model-dependent assumptions, namely the relevant annihilation channels of the dark matter. The most general limits can be placed by assuming that the dark matter only has a single annihilation channel, i.e. annihilation to a $\tau^+\tau^-$ final state 100% of the time. Under these assumptions, limits on the SD dark matter-proton cross-section have been placed that exceed current DD constraints, due to the rather large abundance of Hydrogen within the Sun. Constraints from the IceCube collaboration are shown in Fig. 1.6

Change

Overcoming the first issue requires either a colder star or one that is much heavier. The second requires dark matter to scatter with the constituent material at relativistic energies to overcome the suppression in the cross-sections. Fortunately, there exists objects that meet all these criteria, allowing for a wider variety of dark matter models to be explored then direct detection or traditional indirect detection experiments: compact objects.

1.5 Compact Objects as Dark Matter Probes

The main goal behind this work is to explore how compact objects can be used to probe a wide variety of dark matter interactions that terrestrial direct detection experiments are insensitive to. By compact objects, we are referring to Neutron Stars (NSs) and White Dwarfs (WDs), and not Black Holes that also fall into this category.

Compact objects offer a unique laboratory for studying dark matter and its interactions with the Standard Model in environments unachievable anywhere else in the Universe. They generate strong gravitational fields and are composed of incredibly dense matter, with NSs reaching super-nuclear densities in their central cores. The capture rate within these objects is therefore enhanced due to these properties, with benefits over solar capture including:

- **Gravitational focusing of the DM flux:** The strong gravitational field will increase the impact parameter of the infalling dark matter. This increases the effective size of the capturing body, increasing the flux of dark matter passing through it.

- **Realtivistic Interaction Energies:** In general, the infalling dark matter will be accelerated to (semi-)relativistic velocities ($\sim 0.2 - 0.7c$). Moreover, the stellar constituents will also have relativistic energies. As such, interactions that are momentum/velocity dependent will suffer far less suppression than in DD experiments.

- **Large Number of Targets:** The extremely high densities of these objects correspond to a considerable number of targets for scattering to occur. This allows these objects to probe very small scattering cross-sections, with NSs in particular expected to reach as low as $\sim 10^{-45}\,\mathrm{cm}^2$.

- **Forgot the last point....**

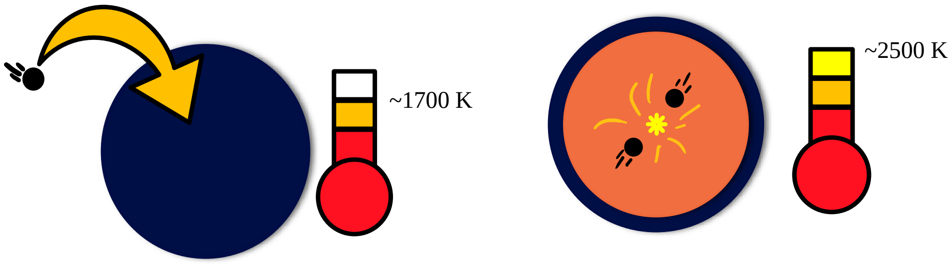

Figure 1.7: Illustration of DM-induced heating of compact objects. **Left:** kinetic heating due to DM scattering, raising the temperature to $\sim 1700\,\mathrm{K}$. **Right:** Annihilation heating contributes an additional $\sim 800\,\mathrm{K}$. This image is inspired by Ref. [136].

In the past, capture in NSs has been applied primarily in the context of sending gravitational collapse into black holes [119–125], and the modifications of NS merger rates as well as the gravitational wave signatures of these mergers [126–129]. Capture in WDs has also been considered, with a variety of different applications of the capture process [130–135].

In recent years, dark matter induced heating of NSs has reemerged as a potential detection frontier [136–143]. It was shown that dark matter could reheat old, isolated NSs in our local neighbourhood back up to temperatures that would cause them to radiate as blackbody peaked in the near-infrared. The aim is to locate the NSs with radio telescopes such as the Square-Kilometer-Array (SKA), and determine their age through their spindown rate. Once located, the star's temperature can be determined through observations from infrared telescopes such as the James Webb Space Telescope (JWST).

This heating occurs in two stages. The dark matter will first deposit its kinetic energy into the star through the scatterings required for capture and its subsequent thermalisation within the NS core, with this process called *kinetic heating*. If the dark matter can annihilate, it will deposit its mass energy, assuming the products are trapped within the star, termed *annihilation heating*. These processes are illustrated in Fig. 1.7. Assuimng a NS in our local neighbourhood, i.e., within

In order to accurately determine the limits on dark matter interactions that such an observation could place, one first requires an accurate calculation of the capture rate. However, all previous calculations relied on the formalism set up by Gould for capture in the Sun, with only minor modifications made to accommodate the extreme nature of the compact objects.

1.6 Thesis Outline

Chapters **??** and **??** of this thesis are devoted to reformulating Gould's capture formalism to account for the physics specific to compact objects in a self-consistent manner. These include a relativistic treatment of the kinematics, using General Relativity to calculate the correct dark matter flux passing through the star, and accounting for Pauli blocking of the final state target using Fermi-Dirac statistics for the stellar constituents. In addition, we incorporate the internal structure of these objects by calculating the radial profiles for the relevant microscopic quantities (e.g., chemical potentials and number densities) via the adoption of a realistic equation of state.

Further considerations are required when considering dark matter interactions with the baryonic matter inside NSs. Due to the high density of the NS interior, the baryonic matter undergoes strong interactions amongst themselves and should not be treated as a free Fermi gas. Instead, adopting an equation of state that accounts for these interactions is required. These interactions modify the mass of the baryons, leading them to obtain an effective mass smaller than their vacuum mass. Furthermore, as we will see, the dark matter may interact with the baryons with momentum transfers on the order of $10\,\mathrm{GeV}$. This is high enough that the dark matter will begin to resolve the internal structure of the baryon. To account for this, the momentum dependence of the baryon form factors that are typically neglected in direct detection and solar capture must be reintroduced.

This formalism is made in preparation for a thorough analysis of the timescales involved in the dark matter heating of compact objects. The enrgy deposited in both the kinetic and annihilation heating stages does not occur instantaneously, and the timescales involved in them need to be compared to the age of the star in question. We will define kinetic heating timescale as the time required for dark matter to deposit 99% of its initial kinetic energy into the star. For annihilation heating to occur, the dark matter must reach a state of capture-annihilation equilibrium within the stellar core. In standard calculations of this timescale, the dark matter must first become thermalised with the star. Only then can annihilations occur efficiently enough to heat the star.

We will work with the EFT operators in Table 1.1 that describe Dirac fermion dark matter interacting with Standard Model leptons. Each operator will be studied in isolation, i.e., by considering a Lagrangian that contains only one of the operators rather than a linear superposition of multiple. This way, we can analyse specific types of interactions independently, allowing us to take as model-independent an approach to the phenomenology as possible.

Definition of Symbols and Abbreviations

ALP Axion Like Particle

BBN Big Bang Nucleosynthesis

$\boldsymbol{C_{\mathrm{geo}}}$ Geometric Capture Rate
CMB Cosmic Microwave Background

DD Direct Detection
DM Dark Matter
$\boldsymbol{K_\chi}$ Dark Matter Kinetic Energy
$\boldsymbol{\rho_\chi}$ DM halo density
$\boldsymbol{m_\chi}$ Dark Matter Mass

EFT Effective Field Theory
EoS Equation of State

$\boldsymbol{f_{\mathrm{FD}}}$ Fermi-Dirac Distribution
$\boldsymbol{\varepsilon_{F,i}}$ Fermi kinetic energy of target species

JWST James Webb Space Telescope

$\boldsymbol{|\overline{\mathcal{M}}|^2}$ Spin-averaged squared matrix element
$\boldsymbol{\mu}$ DM-Target mass ratio, m_χ/m_i

nEMD Neutron ELectric Dipole Moment
NS Neutron Star

PB Pauli Blocking
PQ Peccei-Quinn

QMC Quark-Meson-Coupling EoS

$\boldsymbol{\sigma_{\mathrm{th}}}$ Threshold Cross-Section
SM Standard Model

$\boldsymbol{T_{\mathrm{eq}}}$ Equilibrium Temperature
$\boldsymbol{t_{\mathrm{eq}}}$ Capture-Annihilation equilibrium time
$\boldsymbol{T_\star}$ Temperature of the star
$\boldsymbol{t_{\mathrm{therm}}}$ Thermalisation time

UV Ultraviolet

$\boldsymbol{v_d}$ DM halo dispersion velocity
$\boldsymbol{v_\star}$ Star velocity

WD White Dwarf
WIMP Weakly Interacting Massive Particle

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