DESY Theoretical
Physics Department

DESY Summer School Project

Dark Forces and the Hidden Sector

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Contents

1 Introduction

The Standard Model (SM) is one of the most highly tested theories in science, providing some of the most accurate predictions ever made by a physical theory [1]. However, over the past few years it has become increasingly apparent that the theory is incomplete (besides the exclusion of gravity). Experimental evidence of phenomena such as 'neutrino oscillations' [2] and 'CP violation' [3] as well as the indirect evidence of 'dark matter' [4], [5] all indicate that the SM needs to be extended. The most popular theory for this extension is 'Supersymmetry' (SUSY). The theory of SUSY was developed in order to cure the divergences in the Higgs boson mass m_H which arise due to the quantum corrections of every particle that couples to the Higgs field [6]. SUSY predicts that for every boson and fermion in the SM there exists a respective fermion and boson partner that differs in spin by $\frac{1}{2}$. These SUSY particles are predicted to be naturally heavier than their SM counterparts and hence provide a solution to the so-called 'hierarchy problem'. This problem concerns the fact that the Higgs boson energy scale is considerably smaller than that of the Planck energy scale M_P, the scale at which all of the fundamental forces are expected to be unified in strength. Not only do the SUSY particles help to bridge the void between these energy scales but they also provide a natural dark matter candidate - the lightest SUSY particle (LSP). However, although SUSY provides solutions to many of the problems in the SM, there are still several important mechanisms that are not well understood, the most of important of which is how the particles acquire their mass. In the SM the fermions and vector bosons acquire mass via the 'Higgs mechanism', which 'breaks' the gauge symmetry. However, in order for SUSY particles to become massive, one must break Supersymmetry. There are several proposed theories that perform this breaking, most of which agree that there must exist a separate *hidden* world of particles, detached from the visible world. This is often referred to as the 'hidden sector'. The existence of a hidden sector has not only been proposed in the context SUSY, but also in string theory [7], [8] and theories devised to explain the origin of dark matter in the universe [9]- [17]. Hidden sector theories help to explain many of the unsolved problems in physics and are an interesting and rich area of research. This review aims to provide both a brief overview of the most popular hidden sector theories, as well as a more in depth discussion about the phenomenology of one such model [9] and its relevance to dark matter searches.

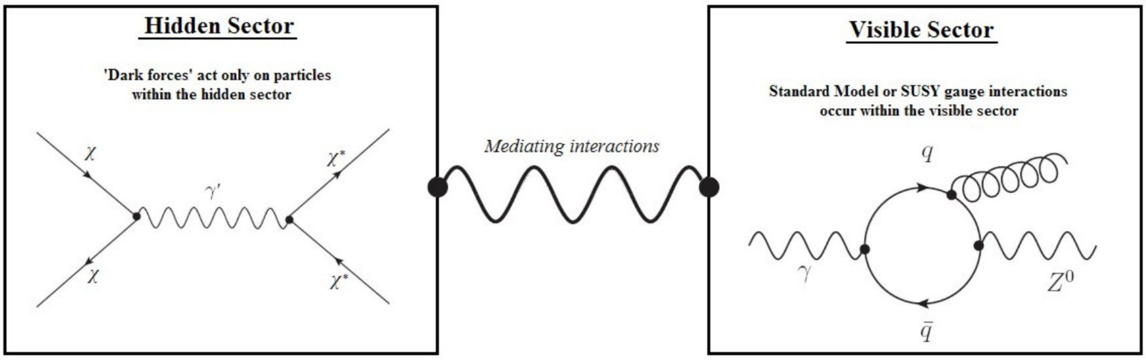

Figure 1: The schematic structure of a hidden sector.

2 The Hidden Sector

Every hidden sector theory is defined by a corresponding gauge group. This group determines the interactions that can occur within the hidden sector as well as with the visible Standard Model sector. The forces that are mediated by the hidden sector gauge bosons are often referred to as 'dark forces'. There are various ideas as to how the interactions between the hidden and visible sectors are mediated, most of which will be discussed in the proceeding section. However, the general schematic structure is shown in figure 1. Although hidden sector theories are discussed in many contexts in physics, they are mostly proposed in order to explain either dark matter or Supersymmetry breaking. The following sections will provide an overview of how hidden sectors have been applied to these areas and in particular will discuss the interactions and predictions made by such theories.

3 Supersymmetry breaking via the Hidden Sector

As discussed in the introduction, supersymmetric theories were devised in order to cure the quadratic divergences in the Higgs boson's mass. In order to do this, it was suggested that every spin-1 boson and spin-$\frac{1}{2}$ fermion in the Standard Model must have a corresponding spin-$\frac{1}{2}$ fermion and spin-0 (scalar) boson partner respectively. Conventionally, the SUSY boson partners to the fermions were called 'sfermions', and the SUSY fermion partners to the gauge bosons were called 'gauginos'. This phenomenology was achieved by introducing a new symmetry to the system called 'Supersymmetry'. Transformations that respected this symmetry allowed bosonic states to be turned into fermionic states and vice versa, creating a duality been the SM and SUSY particles. Each of these SM-SUSY paired states were then grouped together into a single object called a 'supermultiplet'. This was defined to be an 'irreducible representation' of the SUSY symmetry group in much the same way that the 'weak isospin doublet' was for $SU(2) \times U(1)$. Fermion-sfermion states were called 'chiral supermultiplets' and gauge boson-gaugino states called 'gauge supermultiplets'. As well as all of the known fermions and gauge bosons being part of supermultiplets it was shown that in the simplest possible SUSY model there must also exist two (chiral) Higgs supermultiplets (and hence 5 Higgs bosons!). This model is called the 'Minimum Supersymmetric Model' (MSSM)[1].

All theories of SUSY assume that Supersymmetry is an *exact* symmetry. This has the nice feature of ensuring that the quadratic divergences to the Higgs mass will be cancelled to *any order* in *any* MSSM process. A direct consequence of this is that every SM-SUSY pair of particles must be mass degenerate. However, this clearly cannot be the full story because otherwise SUSY particles would have been easily observed at high-energy colliders. This problem alone has led physicists to widely accept that any realistic SUSY model *must* contain a mechanism that causes Supersymmetry to become broken, thus allowing the SUSY particles to acquire different masses to their SM counterparts. Although there are many different ideas as to the origin of SUSY breaking, most theories agree that the symmetry must be 'spontaneously' broken [6]. What this means is that in the vacuum state of the system the Lagrangian for the theory is no longer invariant under SUSY transformations. This 'breaking' of Supersymmetry allows the SM and SUSY particles to gain different masses in much the same way as the $W^{\pm}$ and Z^0 bosons do after electroweak symmetry breaking. More explicitly, what symmetry breaking does is it introduces new terms into the Lagrangian which destroy the Lagrangian's invariance under the symmetry transformations. In the context of SUSY, these terms are called the 'soft Supersymmetry-breaking terms'. Even without knowledge of the specific breaking mechanism, there are only a few types of allowed 'soft' terms that will both break SUSY and ensure the theory remains renormalisable. Just like the scalar Higgs potential in the SM acquires a non-zero 'vacuum expectation value' (VEV) and breaks electroweak symmetry, it is the *scalar* potential in SUSY which is the source of symmetry breaking. It can be shown [6] that the scalar potential for *any* general SUSY theory with scalar fields $\{\phi_i\}$ can be written:

$$V(\{\phi_i\},\{\phi_i^*\}) = F^{*i}F_i + \frac{1}{2}\sum_a D^a D^a \tag{1}$$

$$= \frac{\delta W}{\delta \phi_i}\frac{\delta W^*}{\delta \phi^{*i}} + \frac{1}{2}\sum_a g_a^2(\phi^{*i}T^a\phi_i)^2 \tag{2}$$

where $F^{*i}F_i$ and D^aD^a are called the 'F-terms' and 'D-terms' in the potential[2]. The index[3] a labels the generators T^a and coupling constant g_a^2 of the scalar field's gauge group. W is a specific function of the scalar fields called the 'superpotential'. It can be seen directly from equation 2 that if V is to have a non-zero VEV, it requires that either the D-term or F-term have non-zero VEVs. The first case is called 'Fayet-Iliopoulos (D-term) Supersymmetry breaking' and the second case is called 'O'Raifeartaigh' (F-term) Supersymmetry breaking'.

[1]See pages 8-9 of [6] for a full list of all the MSSM supermultiplets and an in depth discussion of how the general characteristics of SUSY were developed.

[2]These terms are called the 'auxiliary fields' and are introduced in order to ensure that the SUSY Lagrangian remains supersymmetrically invariant 'off shell' (i.e. when the fields do not obey the classical equations of motion) for each gauge interaction.

[3]Technically a runs over the adjoint representation of the corresponding gauge group. a= 1,...,8 for $SU(3)_C$ colour; a = 1,2,3 for $SU(2)_L$ weak isospin; a=1 for $U(1)_Y$ weak hypercharge.

D-term SUSY breaking occurs when the scalar fields of the theory are all charged under a $U(1)$ gauge group. When this is the case, it implies that a linear term $-\kappa D$ can be introduced into the scalar potential without destroying the supersymmetric invariance of the Lagrangian. This is called the 'Fayet-Iliopoulos' or FI term. The D-term part of the potential V_D is then:

$$V_D = \kappa D - \frac{1}{2}D^2 - gD\sum_i q_i|\phi_i|^2 \tag{3}$$

Here q_i are the charges of the scalar fields ϕ_i. At the minimum of the potential D now acquires a non-zero VEV $\langle D\rangle = \kappa - g\sum_i q_i|\phi_i|^2$, which means that in the vacuum the potential has the form [6]:

$$V = \sum_i |m_i|^2|\phi_i|^2 + \frac{1}{2}(\kappa - g\sum_i q_i|\phi_i|^2)^2 \tag{4}$$

Clearly this implies that even if $\langle\phi_i\rangle = 0$, then the VEV of the potential is non-zero ($\frac{1}{2}\kappa^2$). Now, since massive scalar particles ϕ with masses m always have mass terms $m^2\phi^2$ in the potential, one can see directly from V that the scalar particles have acquired squared-masses $|m_i|^2 - gq_i\kappa$. The fermions however, are unaffected by the D-terms and will remain with squared-masses $|m_i|^2$. This inequality in fermion and scalar (sfermion) particle masses explicitly shows that Supersymmetry has been broken.

F-term breaking occurs when the supermultiplets and superpotential, W are such that the equations $F_i = -\frac{\delta W^*}{\delta\phi^{*i}} = 0$ have *no* simultaneous solutions. This implies that the F-term part of the potential ($\sum_i |F_i|^2$) *must* be positive, and hence V will have a non-zero VEV. In a similar manner to D-term breaking this leads the scalar particles to acquire an additional contribution to their masses, resulting in the breaking of SUSY.

However, it has been established that models that incorporate either F or D-term breaking within a MSSM framework soon run into difficulties. In D-term models, allowing D to obtain a VEV for $U(1)_Y$ leads to an unacceptable spectrum and in F-term models it has been found to be difficult to incorporate a supermultiplet whose F-term could develop a VEV. From these findings it is now largely accepted that *the MSSM is incomplete and must be extended by an additional 'hidden sector' in order to accommodate for SUSY breaking.* As outlined in figure 1, it is thought that SUSY breaking occurs indirectly in the MSSM as a consequence of the interactions between the hidden and visible sectors. Although there are many theories regarding how this SUSY breaking is communicated between the sectors, the most popular theories are; 'Planck-scale mediated SUSY breaking' (PMSB), 'gauge mediated SUSY breaking' (GMSB), 'anomaly-mediated SUSY breaking' (AMSB) and 'extra-dimensional mediated SUSY breaking'.

In PSMB the hidden sector is connected to the visible MSSM sector primarily through gravitational interactions. These interactions are added into the MSSM by demanding that SUSY is a *local* symmetry and incorporating the spin-2 gravitational gauge boson (the 'graviton') and it's spin-$\frac{3}{2}$ superpartner the 'gravitino'. This locally supersymmetric theory is called 'supergravity' [6]. The massive gravitino (with mass $m_{3/2}$) plays a special

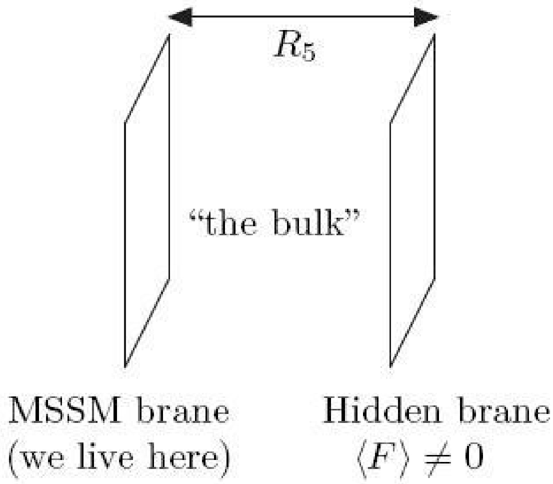

Figure 2: The schematic structure of 'gaugino mediation'. The figure is from page 62 of [6].

role in this theory because it's inclusion automatically implies the existence of F 'soft mass' terms: $m_{3/2} \sim \langle F \rangle / M_P$ in the MSSM Lagrangian, where M_P is the Planck mass and $\langle F \rangle$ is an F-term for some chiral supermultiplet in the hidden sector. Clearly the size of $\langle F \rangle$ will determine whether the effects of PSMB are dominant compared to other breaking effects. If one assumes that $\sqrt{\langle F \rangle} \sim 10^{10} - 10^{11}$GeV then soft terms of a few hundred GeV will be introduced into the Lagrangian, implying the existence of SUSY particles with a mass range of this order.

GMSB is another popular theory that connects the visible and hidden sectors. This theory works by introducing a new set of chiral supermultiplets called 'messenger fields' which couple both to the hidden sector and indirectly to specific MSSM particles via the ordinary $SU(3) \times SU(2)_L \times U(1)_Y$ interactions. The idea is that some unknown process in the hidden sector causes one (or more) of the messenger fields to acquire a non-zero F-term. As a result this splits the masses of the particles in the messenger sector. Different messenger particles then interact with different MSSM particles, via 'radiative corrections', and this causes the MSSM particles to acquire different masses from one another. In other words, the mass splitting in the messenger sector is transferred to the visible sector resulting in the breaking of SUSY.

It has been suggested[4] that extra dimensions may be the source of SUSY breaking. If the MSSM and hidden sectors resided on different dimensional 'branes' separated by a distance R_5 (as shown in figure 2), and the gauge multiplets could propagate in the bulk this would introduce soft gaugino mass terms $M \sim \frac{\langle F \rangle}{R_5 M_5^2}$ into the Lagrangian that would break SUSY (in a similar manner to the PSMB scenario). This mechanism is often called 'gaugino mediation'. AMSB works in a similar way but instead the gauge multiplets are also confined to the MSSM brane and supergravity effects dominate the SUSY breaking.

[4]See [6] for more details.

4 Hidden Sector Dark Matter

4.1 An Overview of Hidden Sector Dark Matter Theories

There exist two mains types of hidden sector dark matter theories; abelian $U(1)$, and non-abelian. Within these theories there are also many variations such as the inclusion of supersymmetric interactions, 'kinetic mixing', and different mediation mechanisms between the sectors. The main aim of these theories have been to provide a consistent dark matter theory that agrees with both cosmological constraints as well as the results from direct and indirect dark matter searches. Table 4.1 summarises the details of some of these experiments and anomalies that they observed.

Experiment Name	**Purpose of Experiment**	**Anomaly**
PAMELA (The Payload for Antimatter Matter Exploration and Light-nuclei Astrophysics)	To detect cosmic rays from the upper atmosphere.	A sharp upturn in the positron fraction from 10-100 GeV was detected which did not agree with what was expected from cosmic rays interacting with the interstellar medium.
ATIC (Advanced Thin Ionisation Calorimeter)	To detect cosmic rays up to an energy of 1 TeV.	Large excesses of electrons and positrons were detected with energies of $\sim$ 300-800 GeV.
WMAP (Wilkinson Microwave Anisotropy Probe)	To detect the cosmic background radiation.	High-energy microwave radiation was detected from the galactic centre which was not correlated with any known galaxy.
EGRET (Energetic Gamma-Ray Experiment Telescope)	To detect high-energy gamma rays in the range 30 MeV-30 GeV.	Gamma-ray measurements from the galactic centre showed excesses at 10-50 GeV.
INTEGRAL (International Gamma-Ray Astrophysics Laboratory)	To detect low-energy gamma rays in the range 15 keV-10 MeV.	A higher than expected 511 keV positron signal was detected.
DAMA/LIBRA (Dark Matter observation experiments)	To directly detect photons which have been emitted from DM-nucleon scattering.	An annual modulation of photons from recoiling nuclei was detected.

Abelian $U(1)$ hidden sector theories provide the simplest extension of the SM. By definition these types of hidden sector must include the minimum of a photon-like spin-1 gauge boson, however, in all DM theories additional fields are also included. In models which consider non-SUSY sectors [13], [18] a defining feature is whether the hidden and visible sectors are 'kinetically mixed', an idea first suggested by Holdem et al [19]. The simplest case of kinetic mixing is when there exists a mixed hidden-$U(1)$-visible-$U(1)_Y$ gauge field term in the Lagrangian of the form:

$$\mathcal{L}_{mixed} = -\frac{\epsilon}{2} F'_{\mu\nu} F^{\mu\nu} \tag{5}$$

where $F'_{\mu\nu}$ and $F^{\mu\nu}$ denote the field strength tensors of the hidden and visible sectors respectively. A main element of kinetic mixing is that it allows simple indirect interactions between the visible and hidden sectors, an attractive feature for any dark matter candidate theory aiming to explain the results in table 4.1. SUSY $U(1)$ hidden sector models work in much the same way as non-SUSY models except that there exist many more possible interactions. A downside to these types of theories is that they inevitably introduce additional free parameters. However, if these parameters can be fixed (by kinetic mixing effects for example), SUSY models offer more potential dark matter candidate states. As with SM $U(1)$ hidden sector theories, most studies tend to consider the simplest possible models either with $W \supset \mu H H^c$ or $W \supset \lambda S H H^c$ [14], [9], [10], [12]. These are often referred to as MSSM and NMSSM abelian hidden sectors. The latter is often considered to be the most realistic SUSY hidden sector theory because the inclusion of a singlet supermultiplet S (which cannot directly couple to either the visible or hidden sectors), allows the Higgs mass scale to be generated naturally without the need for 'fine tuning' [6].

Non-abelian hidden sector theories have also been developed to describe dark matter interactions. However, these theories have been investigated far less extensively. Some of the models that have been suggested include $SU(2) \times U(1)$ [16] and $SU(3) \times SU(2) \times U(1)$ [17], [11]. General non-abelian hidden sector gauge interactions have also been discussed [15]. As in the abelian case, there are models which both include and exclude SUSY and kinetic mixing. Although non-abelian hidden sectors are far more complicated in their structure, they have the distinct advantage in that they can easily accommodate for excited dark matter states [16], which in turn helps to explain the results from both INTEGRAL and DAMA.

4.2 $U(1)_x$ Hidden Sector NMSSM with Kinetic Mixing

As has just been discussed, there are many different hidden sector models that aim to provide a theoretically consistent theory of dark matter that matches with the indirect detection signals observed at the experiments listed in table 4.1. With each of these models there are both advantages and drawbacks, however it seems that in particular, models that involve 'kinetic mixing' have had the strongest predictive success. This section aims to focus on one such model [9] deriving; the particles that are present in the theory, the allowed interactions and also the cross section for a dark matter candidate annihilation process.

Consider a $U(1)_x$ hidden sector with gauge coupling g_x, hypercharge[5] x_H and superpotential W where:

$$W \supset \lambda s h h^c \tag{6}$$

[5]This is analogous to hypercharge Y in QED.

and λ is a positive real Yukawa coupling constant[6]. The hidden sector contains two chiral supermultiplets H and H^c with scalar components h, h^c and fermionic components $\tilde{h}$, $\tilde{h^c}$ respectively. There is also a singlet supermultiplet with components s, $\tilde{s}$ and an intrinsic gauge supermultiplet G with gauge boson A and gaugino $\tilde{A}$. The most important aspect of this theory is that there exists a small kinetic mixing ϵ between the gauge fields of the hidden $U(1)_x$ and visible $U(1)_Y$ sectors. This indirect connection between the sectors implies [20] that there must exist a term of the form $-\epsilon D_Y D_x$ in the scalar potential, where D_Y and D_x are the D-terms introduced for $U(1)_Y$ and $U(1)_x$ interactions. As a result, this means that when the Higgs fields acquire a VEV $\langle\phi_{Yi}\rangle$ in the MSSM, this automatically generates an FI term[7]:

$$V_{FI} \supset \epsilon\langle D_Y\rangle D_x \sim \epsilon\langle\phi_{Yi}\rangle D_x \tag{7}$$

in the hidden sector potential, and hence SUSY breaking (see section A.2). In other words, electroweak symmetry breaking in the visible sector implies that D-term SUSY breaking will occur in the hidden sector. Using the form of the superpotential W (equation 6) and equation 2, the full tree-level scalar potential for the hidden sector is:

$$V = |\lambda|^2|h|^2(|h^c|^2 + |s|^2) + |\lambda|^2|h^c|^2|s|^2 + \frac{g_x^2}{2}(x_H|h|^2 - x_H|h^c|^2 - \xi)^2 \tag{8}$$

where ξ here is the D-term VEV induced by kinetic mixing (derived in section A.2 of the appendix). If one assumes that $\xi > 0$, $x_H > 0$ and $\lambda < \sqrt{2}x_H g_x$, then at the supersymmetric global minimum of this potential h acquires a non-zero VEV but s and h^c do not ($\langle h\rangle \equiv \eta = \sqrt{\xi/x_H}$, $\langle s\rangle = \langle h^c\rangle = 0$). In order to determine the mass spectrum of this theory one must consider the form of the Lagrangian (see section B) after the supermultiplets have acquired their VEVs. This is done by looking at an object called the 'fermion mass matrix' M_f which contains all of the fermion mass terms in the Lagrangian. By defining ψ_f as a vector in the complete fermion basis ($\tilde{A}$, $\tilde{h}$, $\tilde{h^c}$, s), one can write the (tree-level) fermionic part of the Lagrangian as:

$$\mathcal{L}_{fermion} = -\frac{1}{2}(\psi_f)^T M_f \psi_f = -\frac{1}{2}\begin{pmatrix} \tilde{A}^\dagger & \tilde{h}^\dagger & \tilde{h^c}^\dagger & \tilde{s}^\dagger \end{pmatrix} \begin{pmatrix} 0 & \sqrt{2}x_H g_x\eta & 0 & 0 \\ \sqrt{2}x_H g_x\eta & 0 & 0 & 0 \\ 0 & 0 & 0 & \lambda\eta \\ 0 & 0 & \lambda\eta & 0 \end{pmatrix} \begin{pmatrix} \tilde{A} \\ \tilde{h} \\ \tilde{h^c} \\ \tilde{s} \end{pmatrix}$$

By determining the eigenvalues of M_f one finds that there are 2 mass eigenstates ($\frac{1}{\sqrt{2}}(\tilde{A} \pm \tilde{h})$) with a mass of $\sqrt{2}g_x x_H\eta$ and 2 mass eigenstates ($\frac{1}{\sqrt{2}}(\tilde{h^c} \pm \tilde{s})$) with a mass of $\lambda\eta$. Physically this corresponds to 2 Dirac fermions with masses $\sqrt{2}g_x x_H\eta$ and $\lambda\eta$. In a similar way one finds that the squared-mass matrix of the scalar fields M_s^2 (in the basis (A, h, h^c, s)) is given by:

$$M_s^2 = \begin{pmatrix} 2x_H^2 g_x^2\eta^2 & 0 & 0 & 0 \\ 0 & 2x_H^2 g_x^2\eta^2 & 0 & 0 \\ 0 & 0 & \lambda^2\eta^2 & 0 \\ 0 & 0 & 0 & \lambda^2\eta^2 \end{pmatrix}$$

[6]This form for W can only be enforced if one imposes an additional discrete or continuous global symmetry such as Z_3. Such a symmetry prevents the singlet S from coupling to the visible and SUSY breaking sector.

[7]In this case $\kappa = -\epsilon\langle D_Y\rangle$ when written in the general form of an FI term $-\kappa D$.

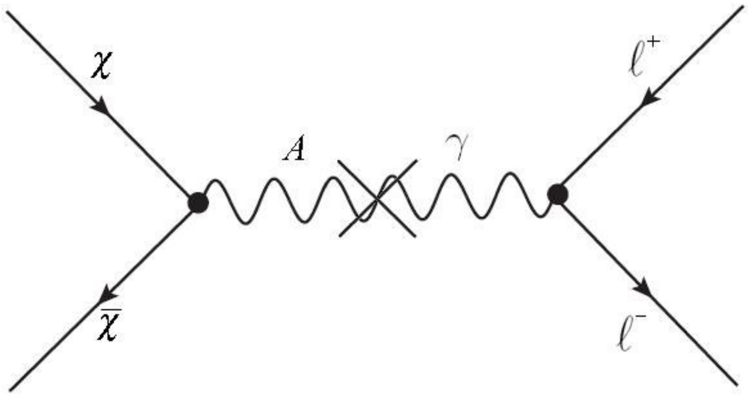

Figure 3: The annihilation mode of the LHP into visible sector leptons.

Unlike the fermion states, the scalar states remain unmixed at tree-level. One can see directly that the $U(1)_x$ gauge boson A and the scalar boson h will acquire masses $\sqrt{2}g_x x_H \eta$, and both the h^c and s bosons will acquire masses $\lambda\eta$.

This hidden sector has a fairly simple structure until Supersymmetry is broken as a result of soft terms in the Lagrangian. It has been found [9] that the phenomenology of this broken hidden sector depends solely on whether gravity mediated breaking or gauge mediated breaking dominates. These regimes effectively correspond to $m_{3/2} \gtrsim m_{hid}$ and $m_{3/2} \ll m_{hid}$ respectively. It turns out that when $m_{3/2} \ll m_{hid}$, the model can not produce a viable dark matter candidate, whereas when $m_{3/2} \gtrsim m_{hid}$ such a candidate does exist. When Supersymmetry is broken in the hidden sector the scalar potential will have soft mass terms: $\{m_h^2|h|^2,\ m_{h^c}^2|h^c|^2,\ m_s^2|s|^2\}$ and the fermion sector will have a gaugino soft mass term $-\frac{1}{2}m_{\tilde{A}}\tilde{A}^2$. This shifts the VEV of h to: $\sqrt{\xi/x_H - m_H^2/(x_H g^x)^2} \equiv \eta'$ but $\langle h^c\rangle = \langle s\rangle = 0$ as before. Although the scalar particle spectrum is altered[8], it is the fermion spectrum which has the most relevance in providing a dark matter candidate. The inclusion of the soft gaugino mass alters the form of the fermion mass matrix:

$$M_f = \begin{pmatrix} m_{\tilde{A}} & \sqrt{2}x_H g_x \eta' & 0 & 0 \\ \sqrt{2}x_H g_x \eta & 0 & 0 & 0 \\ 0 & 0 & 0 & \lambda\eta' \\ 0 & 0 & \lambda\eta' & 0 \end{pmatrix}$$

The fermion spectrum now consists of a Dirac fermion with mass $\lambda\eta'$, and a pair of gaugino-Higgsino[9] Majorana states with masses $\frac{1}{2}(m_{\tilde{A}}+\sqrt{m_{\tilde{A}}^2 + 8x_H^2 g_x^2 \eta'^2})$ and $\frac{1}{2}(\sqrt{m_{\tilde{A}}^2 + 8x_H^2 g_x^2 \eta'^2} - m_{\tilde{A}})$. Depending on the size of $m_{\tilde{A}}$ the lightest of these Majorana states χ could potentially be the 'lightest hidden sector particle' (LHP) and hence a dark matter candidate. This LHP would not have any hidden sector annihilation modes and would only be able to annihilate through s-channel $U(1)_x$ gauge boson exchange into the visible sector as shown in figure 3. The annihilation cross-section for the process in figure 3 is given by

[8]Most notably the scalar particles acquire different masses to their fermion superpartners.
[9]These are asymmetrically mixed $\tilde{A}/\tilde{h}$ states

(derived in section C of the appendix):

$$\langle \sigma v \rangle \simeq \frac{g_x^2 x_H^2 e^2 c_W^2}{12\pi} \epsilon^2 |U_{fh}|^4 \frac{m_\chi^2}{(4m_\chi^2 - m_A^2)^2 + \Gamma_A^2 m_A^2} v_\chi^2 \tag{9}$$

where c_W is the vector coupling constant of the MSSM leptons, m_χ is the mass of the LHP, v_χ is the initial velocity of the LHPs and $|U_{fh}|$ is the Higgsino $\tilde{h}$ fraction of the LHP state. It has been shown that for a small gauge boson mass m_A or large gauge coupling g_x this cross section can explain the electron and positron excesses at PAMELA and ATIC as well as the 511 keV gamma-ray line at INTEGRAL [4].

5 Experimental Signatures of Hidden Sectors

Experimental signatures of hidden sectors have already been discussed in the context of dark matter in section 4.1 (table 4.1), however, this section aims to briefly touch on effects that could be observed at high and low energy colliders, fixed target experiments, 'Light shining through a wall' (LSW) experiments, and in low energy spectroscopy.

Depending on the structure of the hidden sector, many studies [21], [18], [22] agree that there is potential to observe the direct effects of these sectors at low and high energy colliders such as the LHC, BABAR and the TEVATRON. For abelian hidden sectors in particular, events such as $\pi^0 \to \gamma A$ or $K^+ \to \pi_+ A$ could well be observed with a high enough luminosity. In fact the anomalous hyperon decays observed by the HyperCP collaboration could already be a direct signature of a hidden sector photon [18]. Fixed target and beam dump experiments also offer a window of opportunity for both constraining hidden sector parameters and directly observing hidden sector processes [23].

LSW experiments, like the name suggests, work on the principle of shining light at a wall and detecting if any light re-emerges at the other end. LSW experiments such as ALPS [24] and OSQAR [25] have all produced null results so far in observing hidden sector processes such as $\gamma \to A \to \gamma$. However, like the fixed target and beam dump experiments, these observations have allowed a significant amount of the parameter space (mostly for simple hidden $U(1)$ models) to be ruled out. Low energy spectroscopy has also been a useful probe of hidden sector theories because of precision to which certain atomic characteristics can be measured. Measuring the 'Lamb shift'[10] in particular, has put significant bounds on the hidden sector parameter space [26]. Some of the bounds imposed by experiment on the hidden sector photon mass and it's kinetic mixing can be seen in figure 4.

6 Conclusion

Hidden sector theories are a fascinating and extremely relevent field of research in physics. In particular, what characterises these theories is their tremendous scope of application.

[10] This is the difference in the $2S_{\frac{1}{2}}$ and $1P_{\frac{1}{2}}$ energy levels of Hydrogen or Muonium caused by fluctuations in the vacuum.

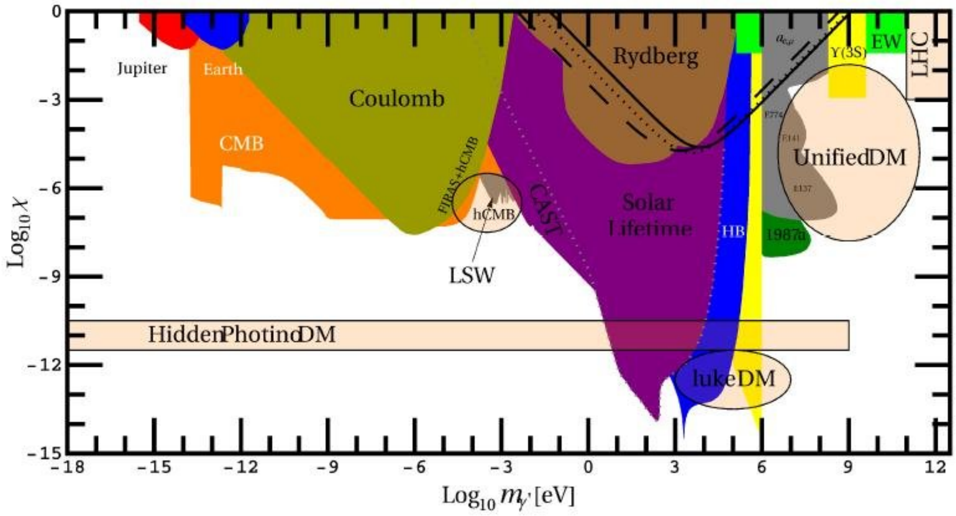

Figure 4: Some of the experimental restrictions imposed on the mass $m_{\gamma'}$ and kinetic mixing χ of the hidden photon. This figure has been taken from [26].

The first main area of interest with hidden sector models is their application to extensions of the Standard Model such as Supersymmetry and String theory. In Supersymmetry it is now pretty well established [6] that SUSY breaking is not possible to achieve within the MSSM alone, and requires an external 'breaking sector' that can communicate this breaking to the MSSM. Hidden sectors nicely fit the criteria for these sectors because they are neutral (or very weakly coupled) with respect to SM interactions. This means that the effects of SUSY breaking are able to occur while at the same time providing a justification as to why direct hidden sector effects have not yet been observed at experiments. In String theory it has also been shown [7], [8] that hidden sectors are actually a fairly natural occurrence. The other main interest with hidden sectors is their application in helping to answer one of the biggest questions in modern-day physics: *What is Dark Matter?* Hidden sector theories have been used to not only describe the interactions that dark matter could have, but also describe how these dark matter interactions could be mediated to the visible Standard Model sector and the signatures that one could detect in experiments. For these reasons alone one can be sure that hidden sectors will play a key role in theoretical physics developments for years to come.

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A Kinetic Mixing

Kinetic mixing provides a natural way of generating the mass-scale of the hidden sector (i.e. the soft mass terms) and breaking Supersymmetry via D-term breaking. The following derivations show how both of these features arise.

A.1 Mass-scale Generation from Kinetic Mixing

When a hidden $U(1)_x$ gauge field is kinetically mixed by a factor ϵ with the SM $U(1)_Y$ gauge field this contributes a 'gauge mixing term' $-\frac{\epsilon}{2}F^{'}_{\mu\nu}F^{\mu\nu}$ to the Lagrangian of the theory. The full kinetic part of the Lagrangian can then be written:

$$\mathcal{L}_{kinetic} = -\tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} - \tfrac{\epsilon}{2}F^{'}_{\mu\nu}F^{\mu\nu} - \tfrac{1}{4}F^{'}_{\mu\nu}F^{'\mu\nu} \tag{10}$$

$F^{'}_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$;$F_{\mu\nu} = \partial_\mu A_{Y\nu} - \partial_\nu A_{Y\mu}$ denote the gauge field strength tensors for $U(1)_x/U(1)_Y$, and $A_\mu/A_{Y\mu}$ are the corresponding gauge boson fields. By making the gauge field transformations (suppressing the 4-vector indices):

$$A \mapsto cA = \frac{1}{\sqrt{1-\epsilon^2}}A \tag{11}$$

$$A_Y \mapsto A_Y - bA = A_Y - \frac{\epsilon}{\sqrt{1-\epsilon^2}}A \tag{12}$$

this eliminates the gauge mixing term from $\mathcal{L}_{kinetic}$:

$$\mathcal{L}_{kinetic} \mapsto -\tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} - \tfrac{1}{4}(b^2+c^2-2\epsilon bc)F^{'}_{\mu\nu}F^{'\mu\nu} + \tfrac{1}{2}(b-\epsilon c)F^{'}_{\mu\nu}F^{\mu\nu} \tag{13}$$

$$= -\tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} - \tfrac{1}{4}(b^2+c^2-2\epsilon bc)F^{'}_{\mu\nu}F^{'\mu\nu} + 0 \tag{14}$$

Because the D-terms are proportional to the equation of motion of the boson fields (by construction), it means that the D-term part of the potential V_D will have a similar structure to $\mathcal{L}_{kinetic}$. In fact V_D is given by:

$$V_D = \frac{1}{2}D_Y^2 + \frac{1}{2}D_x^2 + \epsilon D_Y D_x + g_Y D_Y \sum_i Y_i|\phi_{Yi}|^2 + g_x D_x \sum_i x_i|\phi_{xi}|^2 \tag{15}$$

where ϕ_{Yi} and ϕ_{xi} denote the visible and hidden sector scalar (Higgs) fields respectively, and Y_i and x_i denote their visible and hidden sector hypercharges. It also follows that the transformations in equation 12 will automatically transform the D-terms in an analogous manner: $D_x \mapsto cD_x$;$D_Y \mapsto D_Y - bD_x$. Just as in equation 14 the gauge mixing term will cancel leaving a potential of the form:

$$V_D = \frac{1}{2}D_Y^2 + \frac{1}{2}(b^2+c^2-2bc\epsilon)D_x^2 + g_Y D_Y \sum_i Y_i|\phi_{Yi}|^2 \tag{16}$$

$$+ g_x D_x[c\sum_i x_i|\phi_{xi}|^2 + \sum_i(-\tfrac{bg_Y Y_i}{g_x})|\phi_{Yi}|^2] \tag{17}$$

The last term in this expression shows that the visible sector scalar fields have acquired an effective $U(1)_x$ charge: $x_i^{eff} = -\frac{bg_Y Y_i}{g_x}$. This means that these fields will act as 'gauge messengers' between the visible and hidden sectors, generating soft terms in the hidden sector of the form $m_{soft}^{hid} \sim \epsilon m_{soft}^{vis}$ [9]. These terms, depending on the size of ϵ and the visible sector soft terms m_{soft}^{vis}, will then naturally set the mass scale of the hidden sector particles. This mechanism is often called 'little gauge mediation'.

A.2 D-term SUSY Breaking from Kinetic Mixing

As discussed in the previous section, the D-term part of the scalar potential will have the form of equation 15. Because this potential contains an extra linear term in D_Y ($\epsilon D_Y D_x$), when electroweak symmetry is broken and the visible sector Higgs obtains a VEV, D_Y itself will subsequently obtain a non-zero VEV (via D-term breaking):

$$\langle D_Y \rangle = \epsilon D_x - g_Y \langle \phi_{Yi} \rangle^2 \tag{18}$$

Therefore in the visible sector Higgs vacuum the scalar potential becomes (ignoring constant and $\mathcal{O}(\epsilon^2)$ terms):

$$\begin{aligned} V_D &= \frac{1}{2}\langle D_Y^2 \rangle + \frac{1}{2} D_x^2 + \epsilon \langle D_Y \rangle D_x + g_Y \langle D_Y \rangle \langle \phi_{Yi} \rangle^2 + g_x D_x \sum_i x_i |\phi_{xi}|^2 && (19) \\ &= \frac{1}{2} D_x^2 - \epsilon g_Y \langle \phi_{Yi} \rangle^2 D_x + g_x D_x \sum_i x_i |\phi_{xi}|^2 && (20) \end{aligned}$$

Now again because there is an FI term $-\epsilon g_Y \langle \phi_{Yi} \rangle^2 D_x \equiv \xi D_x$ in the potential this will now cause D_x to obtain a non-zero VEV (by D-term breaking):

$$\langle D_x \rangle = \xi - g_x \sum_i x_i |\phi_{xi}|^2 \tag{21}$$

and hence break Supersymmetry in the hidden sector. The overall D-term potential in this sector is therefore (from equation 2):

$$V_D = \frac{1}{2}\langle D_x \rangle^2 = \frac{1}{2}(g_x \sum_i x_i |\phi_{xi}|^2 - \xi)^2 \tag{22}$$

B The Full Lagrangian of the $W \supset \lambda s h h^c$ Theory

A $U(1)_x$ supersymmetric hidden sector with superpotential $W \supset \lambda s h h^c$ will contain a single gauge boson A_μ and have a Lagrangian:

$$\begin{aligned} \mathcal{L} = &- \tfrac{1}{4} F'_{\mu\nu} F'^{\mu\nu} - \tfrac{\epsilon}{2} F'_{\mu\nu} F^{\mu\nu} + i[\tilde{A}^\dagger \bar{\sigma}^\mu D_\mu \tilde{A} + \tilde{s}^\dagger \bar{\sigma}^\mu D_\mu \tilde{s} + \tilde{h}^\dagger \bar{\sigma}^\mu D_\mu \tilde{h} + \tilde{h}^{c\dagger} \bar{\sigma}^\mu D_\mu \tilde{h}^c] \\ &- [D^\mu s^* D_\mu s + D^\mu h^* D_\mu h + D^\mu h^{c*} D_\mu h^c] - \tfrac{1}{2}\lambda [h \tilde{h}^c \tilde{s} + s \tilde{h} \tilde{h}^c + H^c \tilde{S} \tilde{H}] \\ &- \tfrac{1}{2}\lambda^* [h^* \tilde{h}^{c\dagger} \tilde{s}^\dagger + s^* \tilde{h}^\dagger \tilde{h}^{c\dagger} + h^{c*} \tilde{h}^\dagger \tilde{s}^\dagger] - \tfrac{\sqrt{2}}{2} g_x x_H [s^* \tilde{s} \tilde{A} + \tilde{A}^\dagger \tilde{s}^\dagger s + h^* \tilde{h} \tilde{A} \\ &+ \tilde{A}^\dagger \tilde{h}^\dagger h + h^{c*} \tilde{h}^c \tilde{A} + \tilde{A}^\dagger \tilde{h}^{c\dagger} h^c] - V(h, h^*, h^c, h^{c*}, s, s^*) \end{aligned}$$

where $D_\mu = \partial_\mu - i g_x x_H A_\mu$ is the covariant derivative for the $U(1)_x$ group.

C Dark Matter Annihilation Cross section

The Feynman diagram for the s-channel annihilation of the LHP into leptons is shown in figure 3. Let the incoming LHPs χ have momenta p_1, p_2 and the outgoing lepton anti-lepton pair have momenta p_3, p_4. In the limit where one considers m_ℓ to be negligibly small, the most significant contribution to the cross section comes from S-wave initial states in which the LHP has an *axial* coupling to A and the leptons have a *vector* coupling (c_W) to A. From figure 3 it is clear that there will be two separate propagators involved: the propagator for the hidden $U(1)_x$ boson $\frac{-ig_{\mu\nu}}{q^2-m_A}$, and the propagator for the MSSM photon $\frac{-ig_{\mu\nu}}{q^2}$. However, one can re-write these combined propagators as a single 'Breit-Wigner' propagator [27] $\frac{-ig_{\mu\nu}}{q^2-m_A+im_A\Gamma_A}$ which represents the virtual propagation of a single unstable boson A with decay width Γ_A and couplings $g_x x_H|U_{fh}|$, $ec_W|U_{fh}|\epsilon$ to the hidden and visible sectors respectively. The quantity $\langle\sigma v\rangle$, where v is the relative velocity of the interacting particles, is often of more interest in dark matter physics than simply the cross section. This is because it can be shown [16] that the so-called 'relic abundance' of dark matter in the universe today Ω is given by:

$$\Omega \simeq 0.1 \times \frac{1}{H_0^2} \times \left(\frac{3\times 10^{-26} cm^3 s^{-1}}{\langle\sigma v\rangle}\right) \tag{23}$$

where H_0 is Hubble's constant. However, since $v \simeq 2$ [28], $\langle\sigma v\rangle$ can simply be derived by computing the cross-section of the process. Now, the complete amplitude $\mathcal{M}$ for the process in figure 3 is:

$$\mathcal{M} = \bar{\chi}^{s'}(p_2)(-ig_x x_H|U_{fh}|\gamma^\mu\gamma^5)\chi^s(p_1)\left(\frac{-ig_{\mu\nu}}{q^2-m_A+im_A\Gamma_A}\right)\bar{\ell}^r(p_3)(-iec_W|U_{fh}|\epsilon\gamma^\nu\gamma^5)\bar{\ell}^{r'}(p_4) \tag{24}$$

where $\{s, s', r, r'\}$ denote the spins of the incoming and outgoing particles, and $\{\chi, \bar{\chi}, \ell, \bar{\ell}\}$ are Dirac spinors. After summing over initial and final state spins and setting $q = 2E = 2m_\chi$ the squared-amplitude $|\mathcal{M}|^2$ becomes:

$$\begin{aligned}
|\mathcal{M}|^2 &= \frac{g_x^2 x_H^2 e^2 c_W^2 \epsilon^2 |U_{fh}|^4}{4[(4m_\chi^2 - m_A^2)^2 + \Gamma_A^2 m_A^2]} Tr[(\not{p}_1 + m_\chi)\gamma_\mu\gamma_5(-\not{p}_2 + m_\chi)\gamma_\nu\gamma_5] Tr[(\not{p}_3 + m_\ell)\gamma^\mu(-\not{p}_4 + m_\ell)\gamma^\nu] \\
&= \frac{8g_x^2 x_H^2 e^2 c_W^2 \epsilon^2 |U_{fh}|^4}{(4m_\chi^2 - m_A^2)^2 + \Gamma_A^2 m_A^2}[(p_1\cdot p_3)(p_2\cdot p_4) + (p_1\cdot p_4)(p_2\cdot p_3) - (p_3\cdot p_4)m_\chi^2 \\
&\quad + (p_1\cdot p_2)m_\ell^2 - 2m_\chi^2 m_\ell^2] \\
&= \frac{16g_x^2 x_H^2 e^2 c_W^2 \epsilon^2 |U_{fh}|^4}{(4m_{LHP}^2 - m_A^2)^2 + \Gamma_A^2 m_A^2}[p_\chi^2 p_\ell^2(1 + cos^2\theta) + 2p_\chi^2 m_\ell^2]
\end{aligned}$$

After averaging over angles, setting $m_\ell \ll 1$ and multiplying by the phase space integration factor $\frac{1}{256\pi}\frac{1}{4m_\chi^2}$ one obtains the result:

$$\langle\sigma v\rangle = \frac{g_x^2 x_H^2 e^2 c_W^2}{12\pi}\epsilon^2|U_{fh}|^4 \frac{m_\chi^2}{(4m_\chi^2 - m_A^2)^2 + \Gamma_A^2 m_A^2} v_\chi^2 \tag{25}$$

which is valid for S-wave tree-level processes.

D References

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