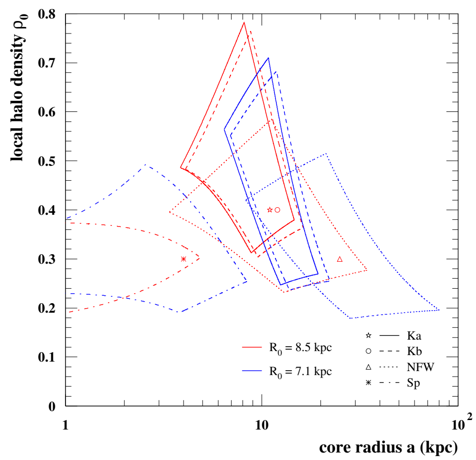

Figure 9: The range of local dark matter densities acceptable with observations of rotation curves for a variety of halo profiles and galactocentric distances. Densities in the range of $0.2 - 0.8\,\mathrm{GeV/cm^3}$ are shown to be acceptable. From Ref. [75].

3.1 The Non-Baryonic Candidate Zoo

There is no shortage of candidates for non-baryonic dark matter. In this section we briefly describe some of these candidates.

- *Standard Model neutrinos*

  Neutrinos have been considered, until recently, excellent dark matter candidates for their "undisputed virtue of being known to exist" [81]. However, a simple calculation shows that, if we call m_i the mass of the i-th neutrino, their total relic density is predicted to be

$$\Omega_\nu h^2 = \sum_{i=1}^{3} \frac{m_i}{93\ \mathrm{eV}} \ . \tag{56}$$

  The best laboratory constraint on neutrino masses comes from tritium β-decay experiments at Troitsk and Mainz [468], pointing to the following upper limit on the neutrino mass

$$m_\nu < 2.05\ \mathrm{eV} \quad (95\%\ \mathrm{C.L.})\ , \tag{57}$$

  while next-generation experiments are expected to reach a sensitivity of approximately 0.2 eV (see Ref. [468] and references therein). The above upper limit applies to all three mass eigenvalues [69], since the mass differences among them must be very small to explain solar ($\Delta m^2 \approx$

$7 \cdot 10^{-5}$ eV^2) and atmospheric ($\Delta m^2 \approx 3 \cdot 10^{-3}$ eV^2) neutrino anomalies (see *e.g.* Ref. [253]). This implies an upper bound on the total neutrino relic density of

$$\Omega_\nu h^2 \lesssim 0.07 \ , \tag{58}$$

which means that neutrinos are simply not abundant enough to be the dominant component of dark matter. A more stringent consraint on the neutrino relic density comes from the analysis of CMB anisotropies, combined with large-scale structure data, suggesting $\Omega_\nu h^2 < 0.0067$ (95% confidence limit). For three degenerate neutrino species this implies $m_\nu < 0.23$ eV.

Furthermore, being relativistic collisionless particles, neutrino free streaming erases (moving from high to low density regions) fluctuations below a scale of ~ 40 Mpc ($m_\nu/30$ eV), called the *free-streaming* length [117]. This would imply a *top-down* formation history of structure in the Universe, where big structures form first. The fact that our galaxy appears to be older than the Local Group [392], and the discrepancy between the predicted late formation of galaxies, at redshift $z \lesssim 1$, against observations of galaxies around $z > 4$ [118], is a further argument against neutrinos as a viable dark matter candidate.

- *Sterile neutrinos*

  These hypothetical particles are similar to Standard Model neutrinos, but without Standard Model weak interactions, apart from mixing. They were proposed as dark matter candidates in 1993 by Dodelson and Widrow [183]. Stringent cosmological and astrophysical constraints on sterile neutrinos come from the analysis of their cosmological abundance and the study of their decay products (see Ref. [1] and references therein).

  Light neutrinos, with masses below a few keV, would be ruled out as dark matter candidates. In fact, if the WMAP result for the reionization optical depth is correct, then dark matter structures were in place to form massive stars prior to redshift $z > 20$, which is is simply not possible if the dark matter particle mass is smaller than ~ 10 keV [476]. An alternative explanation for the WMAP optical depth is reionization by decaying particles, such as sterile neutrinos (see Ref. [274] and references therein). Sterile neutrinos could also be cold dark matter, if there is a very small lepton asymmetry, in which case they are produced resonantly with a non-thermal spectrum [423].

- *Axions*

  Introduced in an attempt to solve the problem of CP violation in particle physics, axions have also often been discussed as a dark matter candidate.

  Laboratory searches, stellar cooling and the dynamics of supernova 1987A constrain axions to be very light ($\lesssim 0.01$ eV). Furthermore, they are expected to be extremely weakly interacting with ordinary particles, which implies that they were not in thermal equilibrium in the early universe.

The calculation of the axion relic density is uncertain, and depends on the assumptions made regarding the production mechanism. Nevertheless, it is possible to find an acceptable range where axions satisfy all present-day constraints and represent a possible dark matter candidate (see *e.g.* Ref. [428]).

- *Supersymmetric candidates*
  - *Neutralinos*

    Neutralinos in models of R-parity conserving supersymmetry are by far the most widely studied dark matter candidates. We devote Sec. 3.2 to their presentation.
  - *Sneutrinos*

    The superpartners of the Standard Model neutrinos in supersymmetric models have long been considered as dark matter candidates. It has been shown that sneutrinos will have a cosmologically interesting relic density if their mass is in the range of 550 to 2300 GeV. However, the scattering cross section of a sneutrino with nucleons is easily calculated and is much larger than the limits found by direct dark matter detection experiments [216].
  - *Gravitinos*

    Gravitinos are the superpartners of the graviton in supersymmetric models. In some supersymmetric scenarios, gauge mediated supersymmetry for example, gravitinos can be the lightest supersymmetric particle and be stable. Gravitinos are thus very strongly theoretically motivated. With only gravitational interactions, however, gravitinos are very difficult to observe [224].

    It has been known for some time that long lived gravitinos can pose problems for cosmology [205, 207, 373, 389, 466, 477]. In particular, their presence can destroy the abundances of primordial light elements in some scenarios [368, 162, 209, 223]. Gravitinos may also be overproduced in the early universe if the temperature of the reheating epoch is not sufficiently low [368]. In some scenarios, however, these problems can be circumvented [126, 145, 237].
  - *Axinos*

    Axinos, the superpartner of the axion, were believed until recently to only be capable of acting as a warm, or hot, dark matter candidate [119, 256]. It has been shown, however, that for quite low reheating temperatures, cold axino dark matter may be possible [157, 158, 159]. In many ways, axinos and gravitinos share similar phenomenological properties.
- *Light scalar dark matter*

  Considering fermionic dark matter candidates with standard Fermi interactions, Lee and Weinberg concluded that relic density arguments preclude such a WIMP with a mass less than a few GeV [340]. If the dark matter is made up other types of particles, however, this limit could be evaded. For example, a 1-100 MeV scalar candidate has been proposed [112, 113].

Such a candidate, although somewhat ad hoc from a particle physics perspective, has recently become experimentally motivated. In Ref. [115], it has been suggested that the 511 keV gamma-ray line emission observed by the INTEGRAL satellite from the galactic bulge could be the product of light dark matter particles annihilating into positrons which then annihilate producing the observed gamma-ray line. To confirm this hypothesis, more tests are needed. In particular, a similar signature could be expected from dwarf spheroidal galaxies [291].

Very recently, light decaying dark matter particles such as axinos with R-parity violation [298] or sterile neutrinos [394] have been suggested as the source of the observed 511 keV emission.

- *Dark matter from Little Higgs models*

As an alternative mechanism (to supersymmetry) to stabilize the weak scale, the so-called "little Higgs" models have been proposed and developed [36, 37, 38, 39]. In these models, the Standard Model Higgs is a pseudo-Goldstone boson with its mass protected by approximate non-linear global symmetries. The divergences to the Higgs mass which remain are present only at the two-loop level and, therefore, the weak scale can be stabilized in an effective field theory which is valid up to ~ 10 TeV. Recall that in supersymmetry, the divergences to the Higgs mass are exactly cancelled at all orders.

At least two varieties of little Higgs models have been shown to contain possible dark matter candidates. One of these classes of models, called "theory space" little Higgs models, provide a possibly stable, scalar particle which can provide the measured density of dark matter [105]. In Ref. [100], the detection prospects for such a candidate were found to be not dissimilar to WIMPs predicted in models of supersymmetry or universal extra dimensions.

Cheng and Low [140] have developed another variety of little Higgs model, motivated by the problem of the hierarchy between the electroweak scale and the masses of new particles constrained by electroweak precision measurements. They solve this problem by introducing a new symmetry at the TeV scale which results in the existence of a stable WIMP candidate with a $\sim$TeV mass.

For a potential dark matter candidate from a little Higgs model to be stable, we must assume that the discrete symmetry which protects it from decay is fundamental and is not broken by the operators in the UV completion.

- *Kaluza-Klein states*

Kaluza-Klein excitations of Standard Model fields which appear in models of universal extra dimensions have also been discussed a great deal recently as a candidate for dark matter. They are discussed in Sec. 3.3. Additionally, a dark matter candidate has been proposed in the framework of "warped" universal extra-dimensions: an exotic particle with gauge quantum numbers of a right-handed neutrino, but carrying fractional baryon-number [14].

- *Superheavy dark matter*

  Superheavy dark matter particles, sometimes called *Wimpzillas*, have interesting phenomenological consequences, including a possible solution to the problem of cosmic rays observed above the GZK cutoff. These are discussed in Sec. 3.4.

- *Q-balls [336, 335], mirror particles [285, 227, 301, 366], CHArged Massive Particles (CHAMPs) [174], self interacting dark matter [433, 165], D-matter [424], cryptons [206, 203], superweakly interacting dark matter [224], brane world dark matter [135], heavy fourth generation neutrinos [406, 307], etc.* Although some of these candidates or classifications present some intriguing features we will not discuss them here. We refer the interested reader to the wide literature on the subject, *e.g.* the reviews of non-baryonic candidates by Ellis [199] and Bergstrom [81].

We stress that it is by no means assured that the dark matter is made of a single particle species. On the contrary, we already know that Standard Model neutrinos contribute to dark matter, but cannot account for all of it. Even in supersymmetry models for dark matter, $N = 2$ supersymmetry allows the possibility of two stable dark matter relics (see, for example, Ref. [114]).

In what follows, we will assume that the abundance of our candidates satisfy the limits provided by the analysis of the CMB discussed in Sec.2.3, but we stress that, although the upper bound is a strict limit, the lower bound can be relaxed, assuming that our candidate is a sub-dominant component of dark matter. The interested reader will find in Ref. [192] a detailed discussion on the detection prospects of a subdominant density component of dark matter.

3.2 Supersymmetry

It would be impossible to review in only a few pages the history and theory of Supersymmetry (SUSY). Instead, we prefer here to review the motivations that led to its introduction and to briefly present the concepts and the notations that we will use in the following chapters. Furthermore, we present a few of the supersymmetric models discussed in the literature (we reserve the word "scenario" for a specific choice of parameters in the framework of a given model) and discuss the consequences of various assumptions, involved in the process of model-building, on SUSY phenomenology. For further discussions of supersymmetry, we refer the interested reader to Refs. [96, 150, 469, 182, 305, 355, 383, 146].

3.2.1 Basics of supersymmetry

As we saw in Sec.1.3, in the Standard Model of particle physics there is a fundamental distinction between bosons and fermions: while bosons are the mediators of interactions, fermions are the constituents of matter. It is therefore natural to ask whether a symmetry exists which relates them, thus providing a sort of "unified" picture of matter and interactions.

Another way to state the problem is to ask whether a Lie group exists mixing internal (Isospin, *etc.*) and space-time (Lorentz) symmetries [269]. Although apparently uncorrelated to the differing behaviour of bosons and fermions, this problem led to the study of the same algebraic structures. Early attempts to find a broad Lie group including the Poincaré and internal symmetry groups had to face the limitations imposed by the so–called *no-go* theorem of Coleman and Mandula. Such limitations were finally circumvented with the introduction of *graded* Lie algebras, *i.e.* algebras involving fermionic generators satisfying anticommutation relations (see below).

For those who are not convinced by these symmetry arguments, there are other major reasons for interest in supersymmetry. One reason is its role in understanding the *hierarchy problem*. The hierarchy problem is linked to the enormous difference between the electroweak and Planck energy scales. This problem arises in the radiative corrections to the mass of the Higgs boson.

All particles get radiative corrections to their mass, but while fermion masses increase only logarithmically, scalar masses increase quadratically with energy, giving corrections at 1-loop of

$$\delta m_s^2 \sim \left(\frac{\alpha}{2\pi}\right)\Lambda^2, \tag{59}$$

where Λ is a high-energy cut-off where new physics is expected to play an important role. The radiative corrections to the Higgs mass (which is expected to be of the order of the electroweak scale $M_W \sim 100$ GeV) will destroy the stability of the electroweak scale if Λ is higher than $\sim$ TeV, *e.g.* if Λ is near the Planck mass.

An appealing, though not the only, solution to this problem is to postulate the existence of new particles with similar masses but with spin different by one half. Then, since the contribution of fermion loops to δm_s^2 have opposite sign to the corresponding bosonic loops, at the 1-loop level, Eq. 59 becomes

$$\delta m_s^2 \sim \left(\frac{\alpha}{2\pi}\right)\left(\Lambda^2+m_B^2\right)-\left(\frac{\alpha}{2\pi}\right)\left(\Lambda^2+m_F^2\right)=\left(\frac{\alpha}{2\pi}\right)\left(m_B^2-m_F^2\right). \tag{60}$$

Furthermore, the supersymmetric algebra insures that (provided $|m_B^2-m_F^2|\lesssim 1$ TeV) the quadratic divergence to the Higgs mass is cancelled at all orders of perturbation theory. The algebra of supersymmetry naturally guarantees the existence of new particles, with the required properties, associating to all of the particles of the Standard Model superpartners with the same mass and opposite spin-type (boson or fermion).

Another reason for interest in supersymmetric theories comes from the unification of gauge couplings at a scale $M_U \sim 2\times 10^{16}$ GeV (see Fig. 10). Although

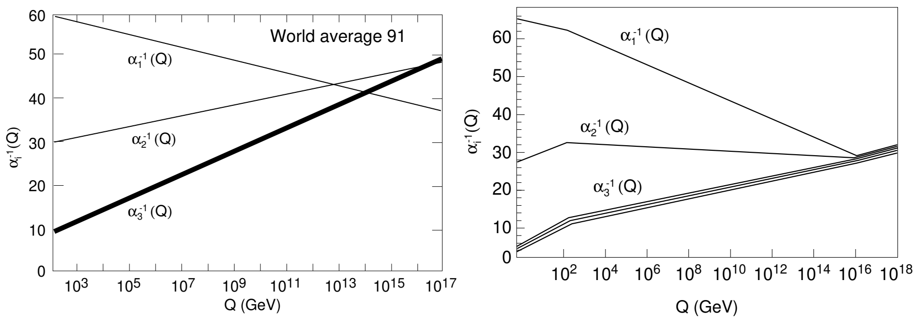

Figure 10: The measurements of the gauge coupling strengths at LEP do not (left) evolve to a unified value if there is no supersymmetry but do (right) if supersymmetry is included [210].

extrapolation of the coupling constants using only Standard Model particles fails to unify them to a common value (left frame of Fig. 10), by introducing supersymmetry at the TeV scale, it was shown [27] that these forces naturally unify at a scale $M_U \sim 2 \times 10^{16}$ GeV (right frame of Fig. 10). This has been taken as a strong hint in favor of a Grand Unified Theory (GUT) which predicts gauge coupling unification below the Planck scale.

The new generators introduced with supersymmetry change fermions into bosons and vise versa, *i.e.*

$$Q|\text{fermion}\rangle = |\text{boson}\rangle; \quad Q|\text{boson}\rangle = |\text{fermion}\rangle. \tag{61}$$

Because of their fermionic nature, the operators Q must carry spin 1/2, which implies that supersymmetry must be a spacetime symmetry. The question then arises of how to extend the Poincaré group of spatial translations and Lorentz transformations to include this new boson/fermion symmetry. The structure of such a group is highly restricted by the Haag-Lopuszanski-Sohnius extension of the Coleman and Mandula theorem cited above. For realistic theories, the operators, Q, which we choose by convention to be Majorana spinors, must satisfy

$$\{Q_a, \overline{Q}_b\} = 2\gamma^\mu_{ab} P_\mu \tag{62}$$
$$\{Q_a, P_\mu\} = 0 \tag{63}$$
$$[Q_a, M^{\mu\nu}] = \sigma^{\mu\nu}_{ab} Q^b \tag{64}$$

where

$$\overline{Q}_a \equiv \left(Q^\dagger \gamma_0\right)_a \tag{65}$$

and

$$\sigma^{\mu\nu} = \frac{i}{4}[\gamma^\mu, \gamma^\nu] \tag{66}$$

are the structure constants of the theory.

Just as Lorentz invariance is manifest in Minkowski space-time, supersymmetry is manifest in the so-called *superspace* formalism, where a superspace is defined as a set of coordinates $\{x, \theta, \overline{\theta}\}$, where $x = x^\mu$ are the usual coordinate of Minkowski spacetime, and $\theta, \overline{\theta}$ are anti-commuting Weyl spinors.

A *superfield* is then a function, $\Phi(x, \theta, \overline{\theta})$, defined on a superspace; it is common to introduce *chiral fields* representing matter and *vector fields* representing gauge fields.

3.2.2 Minimal supersymmetric Standard Model

To continue our brief introduction to SUSY, we consider the *minimal* supersymmetric extension of the Standard Model (MSSM, for Minimal Supersymmetric Standard Model). The MSSM is minimal in the sense that it contains the smallest possible field content necessary to give rise to all the fields of the Standard Model. This can be done as follows:

- We associate fermionic superpartners to all gauge fields. Gluons, $W^\pm$ and B bosons then get fermionic partners called *gluinos* ($\tilde{g}$), *winos* ($\tilde{W}^i$) and *binos* ($\tilde{B}$), respectively. The common name for all partners of gauge fields is the *gaugino*.
- We associate scalar partners to the fermions, *i.e.* quarks and leptons get scalar partners called *squarks* and *sleptons*.
- We introduce one additional Higgs field (for a total of two Higgs doublets, corresponding to five physical Higgs states) and associate one spin 1/2 *Higgsino* to each Higgs boson. This is done to give masses to both up and down-type quarks upon electroweak symmetry breaking and also preserve supersymmetry (therefore, we cannot use the conjugate of the Higgs as is done in Standard Model). Introducing another Higgs doublet also makes the theory anomaly free.

The resulting particle content of the theory is shown in tables 3 and 4.

The MSSM is then specified through the *superpotential*, defined as

$$W = \epsilon_{ij}\left[y_e H_1^j L^i E^c + y_d H_1^j Q^i D^c + y_u H_2^i Q^j U^c\right] + \epsilon_{ij}\mu H_1^i H_2^j \tag{67}$$

where i and j are SU(2) indices, and y are Yukawa couplings. Color and generation indices have been suppressed in the above expression. The superpotential represents a supersymmetrization of the Standard Yukawa couplings plus a bilinear Higgs term. The superpotential enters the Lagrangian of the theory through the terms

$$\mathcal{L}_{\text{SUSY}} = -\frac{1}{2}(W^{ij}\psi_i\psi_j + W^*_{ij}\psi^{i\dagger}\psi^{j\dagger}) - W^i W^*_i \tag{68}$$

Superfield	SM particles	Spin	Superpartners	Spin
Q	$\begin{pmatrix} u_L \\ d_L \end{pmatrix}$	1/2	$\begin{pmatrix} \tilde{u_L} \\ \tilde{d_L} \end{pmatrix}$	0
U^c	$\bar{u}_R$	1/2	$\tilde{u}_R^*$	0
D^c	$\bar{d}_R$	1/2	$\tilde{d}_R^*$	0
L	$\begin{pmatrix} \nu_L \\ e_L \end{pmatrix}$	1/2	$\begin{pmatrix} \tilde{\nu_L} \\ \tilde{e_L} \end{pmatrix}$	0
E^c	$\bar{e}_R$	1/2	$\tilde{e}_R^*$	0
H_1	H_1	0	$\tilde{H}_1$	1/2
H_2	H_2	0	$\tilde{H}_2$	1/2
G^a	g	1	$\tilde{g}$	1/2
W_i	W_i	1	$\tilde{W}_i$	1/2
B	B	1	$\tilde{B}$	1/2

Table 3: Field content of the MSSM.

Standard Model particles and fields		Supersymmetric partners			
		Interaction eigenstates		Mass eigenstates	
Symbol	Name	Symbol	Name	Symbol	Name
$q = d, c, b, u, s, t$	quark	$\tilde{q}_L$, $\tilde{q}_R$	squark	$\tilde{q}_1$, $\tilde{q}_2$	squark
$l = e, \mu, \tau$	lepton	$\tilde{l}_L$, $\tilde{l}_R$	slepton	$\tilde{l}_1$, $\tilde{l}_2$	slepton
$\nu = \nu_e, \nu_\mu, \nu_\tau$	neutrino	$\tilde{\nu}$	sneutrino	$\tilde{\nu}$	sneutrino
g	gluon	$\tilde{g}$	gluino	$\tilde{g}$	gluino
$W^\pm$	W-boson	$\tilde{W}^\pm$	wino		
H^-	Higgs boson	$\tilde{H}_1^-$	higgsino	$\tilde{\chi}_{1,2}^\pm$	chargino
H^+	Higgs boson	$\tilde{H}_2^+$	higgsino		
B	B-field	$\tilde{B}$	bino		
W^3	W^3-field	$\tilde{W}^3$	wino		
H_1^0	Higgs boson			$\tilde{\chi}_{1,2,3,4}^0$	neutralino
H_2^0	Higgs boson	$\tilde{H}_1^0$	higgsino		
H_3^0	Higgs boson	$\tilde{H}_2^0$	higgsino		

Table 4: Standard Model particles and their superpartners in the MSSM (adapted from Ref. [194]).

where we have used $W^i \equiv \partial W/\partial\phi_i$, $W_i^* \equiv \partial W/\partial\phi^{i*}$, and $W^{ij} \equiv \partial^2 W/\partial\phi_i\partial\phi_j$. ϕ_i and ψ_i are scalar and fermion fields, respectively.

One additional ingredient of the MSSM is the conservation of R-parity. R-parity is a multiplicative quantum number defined as

$$R \equiv (-1)^{3B+L+2s} \; . \tag{69}$$

All of the Standard Model particles have R-parity $R = 1$ and all sparticles (*i.e.* superpartners) have $R = -1$. Thus, as a consequence of R-parity conservation, sparticles can only decay into an odd number of sparticles (plus Standard Model particles). The lightest sparticle (dubbed the LSP, for Lightest Supersymmetric Particle) is, therefore, stable and can only be destroyed via pair annihilation, making it an excellent dark matter candidate [244, 204]. Note that this not the original motivation for R-parity. In fact, R-parity was first introduced to suppress the rate of proton decay [218, 175, 467, 411, 177].

The nature of the LSP in the MSSM is constrained by many observations. It cannot have a non-zero electric charge or color, or it would have condensed with baryonic matter to produce heavy isotopes, in conflict with observations. Among the neutral candidates, a possibile LSP could be the sneutrino. Sneutrino LSPs have, however, been excluded by direct dark matter detection experiments (see sections 4.1 and 5). Although axinos and gravitinos cannot be *a prori* excluded, they arise only in a subset of supersymmetric scenarios and have some unattractive properties (see section 3.1). In particular, gravitinos and axinos have very weak interactions and would be practically impossible to detect, making them less interesting from a phenomenological perspective. The lightest *neutralino* remains an excellent dark matter candidate, and is further discussed in the next section.

To determine the identity of the LSP (or other characteristics) in a given supersymmetric scenario, we have to specify how supersymmetry is *broken*. If supersymmetry were not broken, then each superpartner would have a mass identical to its Standard Model counterpart, which is clearly not the case. Thus, new terms which break supersymmetry must be added to the Lagrangian. These terms, however, should be added carefully, in order not to destroy the hierarchy between Planck and electroweak scales. The possible forms for such terms are

$$\begin{aligned}\mathcal{L}_{soft} &= -\frac{1}{2}M_\lambda^a \lambda^a \lambda^a - \frac{1}{2}(m^2)^i_j \phi_i \phi^{j*} \\ &\quad -\frac{1}{2}(BM)^{ij}\phi_i\phi_j - \frac{1}{6}(Ay)^{ijk}\phi_i\phi_j\phi_k + h.c.,\end{aligned} \tag{70}$$

where the M_λ^a are gaugino masses, m^2 are soft scalar masses, B is a bilinear mass term, and A is a trilinear mass term. We will discuss some specific supersymmetry breaking scenarios later in this section.

3.2.3 The lightest neutralino

In the MSSM, the superpartners of the B, W_3 gauge bosons (or the photon and Z, equivalently) and the neutral Higgs bosons, H_1^0 and H_2^0, are called binos

($\tilde{B}$), winos ($\tilde{W}_3$), and higgsinos ($\tilde{H}_1^0$ and $\tilde{H}_2^0$), respectively. These states mix into four Majorana fermionic mass eigenstates, called neutralinos. The four neutralino mass eigenstates are typically labelled $\tilde{\chi}_1^0$, $\tilde{\chi}_2^0$, $\tilde{\chi}_3^0$ and $\tilde{\chi}_4^0$, ordered with increasing mass. In the following we will refer to $\tilde{\chi}_1^0$, *i.e.* the lightest of the four neutralinos, as *the* neutralino, and denote it simply as, $\chi \equiv \tilde{\chi}_1^0$.

In the basis $(\tilde{B}, \tilde{W}_3, \tilde{H}_1^0, \tilde{H}_2^0)$, the neutralino mass matrix can be expressed as

$$\mathcal{M}_N = \begin{pmatrix} M_1 & 0 & -M_Z \cos\beta \sin\theta_W & M_Z \sin\beta \sin\theta_W \\ 0 & M_2 & M_Z \cos\beta \cos\theta_W & -M_Z \sin\beta \cos\theta_W \\ -M_Z \cos\beta \sin\theta_W & M_Z \cos\beta \cos\theta_W & 0 & -\mu \\ M_Z \sin\beta \sin\theta_W & -M_Z \sin\beta \cos\theta_W & -\mu & 0 \end{pmatrix}, \tag{71}$$

where M_1 and M_2 are the bino and wino mass parameters, respectively, θ_W is the Weinberg angle and $\tan\beta$ is the ratio of the vacuum expectation values of the Higgs bosons. μ is the higgsino mass parameter. As we have seen, the (lightest) neutralino is a linear combination of $\tilde{B}$, $\tilde{W}_3$, $\tilde{H}_1^0$ and $\tilde{H}_2^0$,

$$\chi = N_{11}\tilde{B} + N_{12}\tilde{W}_3 + N_{13}\tilde{H}_1^0 + N_{14}\tilde{H}_2^0. \tag{72}$$

We then define the *gaugino fraction*, f_G, and the *higgsino fraction*, f_H, as

$$f_G = N_{11}^2 + N_{12}^2 \tag{73}$$

and

$$f_H = N_{13}^2 + N_{14}^2 \,. \tag{74}$$

For the analytic expressions used to diagonalize the neutralino mass matrix, see Appendix A.

The neutralino interactions most relevant for the purposes of dark matter are self annihilation and elastic scattering off of nucleons. Neutralinos are expected to be extremely non-relativistic in the present epoch, allowing us to safely keep only the a-term in the usual expansion of the annihilation cross section,

$$\sigma v = a + bv^2 + \mathcal{O}\left(v^4\right). \tag{75}$$

The b-term must be included in performing calculations of the neutralino relic density, however.

At low velocities, the leading channels for neutralino annihilation are annihilations to fermion-antifermion pairs (primarily heavy fermions, such as top, bottom and charm quarks and tau leptons), gauge bosons pairs (W^+W^- and Z^0Z^0) and final states containing Higgs bosons. In appendix B, we give the most important neutralino annihilation diagrams, amplitudes and cross sections (in the low velocity limit). For a complete list of all tree level processes, diagrams, amplitudes and cross sections, see the excellent review of Jungman, Kamionkowski and Griest [305].

3.2.4 Supersymmetric models

Although relatively simple in many respects, the MSSM has a huge number of free parameters. Most of these parameters represent masses and mixing angles, much as in the case of the Standard Model. To allow for the practical phenomenological study of the MSSM, the number of parameters which are considered must be reduced. This can be done by making (theoretically well motivated) assumptions which reduce the free parameters from more than 100 to a more tractable quantity. Depending on the assumptions used, one obtains different supersymmetric models. In the following section, we will describe a few of the most widely considered supersymmetric scenarios, including mSUGRA (often called the constrained MSSM) and a phenomenologically simplified MSSM (called the phenomenological, or, pMSSM). We also discuss the phenomenological features of the MSSM in anomaly, gauge and gaugino mediated scenarios.

mSUGRA The mSUGRA, or constrained MSSM, scenario is a simple phenomenological model based on a series of theoretical assumptions (see *e.g.* Kane *et al.* [313]). The number of free parameters is reduced in this scenario by assuming that the MSSM parameters obey a set of boundary conditions at the Grand Unification scale:

- Gauge coupling unification

$$\alpha_1(M_U) = \alpha_2(M_U) = \alpha_3(M_U) \equiv \alpha_U \tag{76}$$

  with $\alpha_i = g_i^2/4\pi$

- Unification of the gaugino masses

$$M_1(M_U) = M_2(M_U) = M_3(M_U) \equiv m_{1/2} \tag{77}$$

- Universal scalar [sfermion and Higgs boson] masses

$$\begin{aligned} M_{\tilde{Q}}(M_U) = M_{\tilde{u}_R}(M_U) = M_{\tilde{d}_R}(M_U) = M_{\tilde{L}}(M_U) = M_{\tilde{l}_R}(M_U) \\ = M_{H_u}(M_U) = M_{H_d}(M_U) \equiv m_0 \end{aligned} \tag{78}$$

- Universal trilinear couplings:

$$A_u(M_U) = A_d(M_U) = A_l(M_U) \equiv A_0 \tag{79}$$

By requiring the minimization of the Higgs potential (in order to recover electroweak symmetry breaking), we are left with five (four continuous and one discrete) free parameters:

$$\tan\beta \ , \ m_{1/2} \ , \ m_0 \ , \ A_0 \ , \ \ \mathrm{sign}(\mu), \tag{80}$$

where $\tan\beta$ is the ratio of the vacuum expectation values of the two Higgs fields and μ is the higgsino mass parameter.

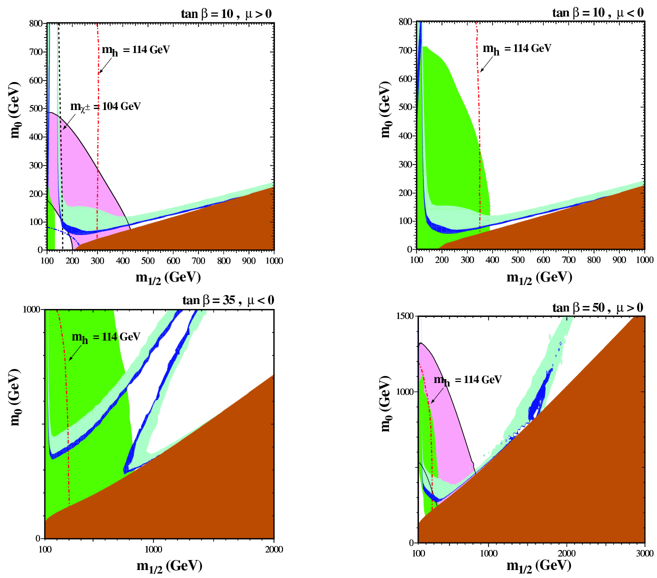

Figure 11: The $(m_{1/2}, m_0)$ planes for mSUGRA with (a) $\tan\beta = 10, \mu > 0$, (b) $\tan\beta = 10, \mu < 0$, (c) $\tan\beta = 35, \mu < 0$, and (d) $\tan\beta = 50, \mu > 0$. In each panel, the region allowed by the older cosmological constraint $0.1 \leq \Omega_\chi h^2 \leq 0.3$ has cyan shading, and the region allowed by the newer cosmological constraint $0.094 \leq \Omega_\chi h^2 \leq 0.129$ has dark blue shading. For more details, see Ref. [208].

A recent study of mSUGRA parameter space in light of the WMAP measurement of the dark matter relic density can be found in Ref. [208]. We show in Fig. 11 and Fig. 12 the regions of the $(m_{1/2}, m_0)$ plane consistent with CMB and accelerator data.

In addition to constraints on models in mSUGRA which come from the WMAP measurements, strong constraints can also be placed by collider data. In particular, constraints arise from the absence of new particles at LEP below ≈ 100 GeV and the agreement of $b \to s\gamma$ decays with predictions of the Standard Model. Measurements of the anomalous magnetic momentum of the muon, $g_\mu - 2$, also provide a possible constraint. The interested reader will find a discussion of mSUGRA parameters and the definition of SUSY *benchmarks* points in Ref. [61]. For more on collider constraints, see section 3.5.

The phenomenological MSSM The scenario we present in this section is not necessarily motivated by any theoretical arguments, but rather is justified by focusing on the aspects of supersymmetric phenomenology which are the most interesting for neutralino dark matter. The phenomenological MSSM, or pMSSM, is an adaptable framework which can be described by as many as tens of parameters, or as few as five or seven. It is NOT a model, but rather a convenient description of the phenomenology most relevant to supersymmetric

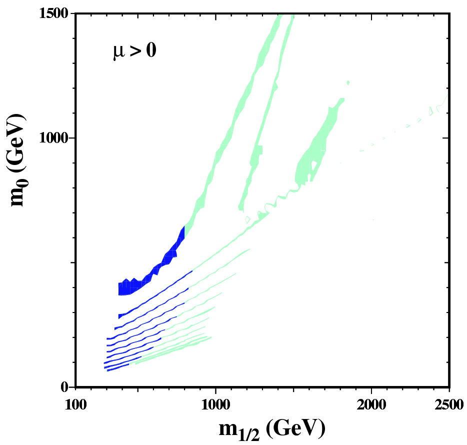

Figure 12: Regions of the $(m_{1/2}, m_0)$ plane in mSUGRA that are compatible with $0.094 < \Omega_\chi h^2 < 0.129$ and laboratory constraints for $\mu > 0$ and $\tan\beta = 5, 10, 15, 20, 25, 30, 35, 40, 45, 50, 55$. The parts of the strips compatible with $g_\mu - 2$ at the 2-σ level have darker shading. From Ref. [208].

dark matter. Common choices in defining a phenomenological MSSM include a) no new sources of CP violation (all the phases in the SUSY breaking terms are set to zero), b) no flavor-changing neutral currents and c) first and second generation universality.

One example of a phenomenological MSSM is used in the DarkSusy program package [252]. In this scheme, in addition to the common features described above, gaugino unification is assumed (similar to Eq. 77). The remaining inputs are defined by seven free parameters:

$$\mu, M_2, \tan\beta, \ M_A, \quad m_0, \ A_b \text{ and } A_t \tag{81}$$

where M_A is the mass of the pseudo-scalar Higgs boson, m_0 is the common scalar mass, and $A_{b,t}$ are trilinear couplings appearing in SUSY breaking terms. Unlike in the case of the mSUGRA scenario, the input parameters are chosen at the electroweak scale without making use of renormalization group equations. The inputs used in DarkSusy can be expanded beyond these seven to include other parameters, thus representing a more general MSSM.

We note that these scenarios are less theoretically motivated in comparison to mSUGRA. Various combinations of theoretically and phenomenologically-based descriptions for supersymmetry are often considered in the literature, often maintaining some of the theoretically motivated constraints of mSUGRA while relaxing other requirement (for example, see Refs. [85, 104]).

The focus point region of mSUGRA In most of the parameter space of mSUGRA or other similar scenarios, the lightest neutralino is a gaugino-like neutralino with a mass of a couple hundred GeV or less. In the so-called "focus point" region of mSUGRA, however, the lightest neutralino can have a considerable higgsino content, and be significantly more heavy [186, 190, 219, 220, 221].

In the focus point region, very large scalar masses are possible without violating naturalness constraints. This occurs because the soft masses squared of the Higgs boson, $m_{H_u}^2$, have pseudo fixed-point behavior, and can start with a wide range of input values and run to a similar negative value at the low scale. This is interesting because it indicates that, in the focus point region, electroweak symmetry breaking does not require fine-tuning in the high energy input values.

A typical feature of the focus point region are large scalar masses (usually $\sim$ TeV). The main reason for a larger higgsino content in the LSP is the larger input value of the soft scalar mass. The tree level electroweak symmetry breaking condition gives

$$\frac{1}{2}m_Z^2 \sim -m_{H_u}^2 - \mu^2. \tag{82}$$

In the typical mSUGRA scenarios, $m_{H_u}^2$ is driven to some large negative value due to the running of the renormalization group equations. This requires a large value of μ to give the correct Z mass. In the focus point region, however, it is possible that the large input value of the scalar soft mass makes $m_{H_u}^2$ less negative. Hence, a smaller value of μ is possible, which leads to a larger higgsino content in the LSP.

Anomaly mediated SUSY breaking Anomaly Mediated Supersymmetry Breaking (AMSB) is an attractive alternative to general gravity mediated scenarios as it provides an elegant solution to the so-called flavor problem through an elegant decoupling mechanism. The resulting soft parameters are "UV insensitive". In this scenario, the SUSY parameters can all be expressed in terms of low energy parameters such as the Yukawa and gauge couplings. Although the details of AMSB are quite technical, and are beyond the scope of this work, we will here describe some of the phenomenological features of this scenario which are most relevant for dark matter.

In AMSB, the gaugino spectrum is given by

$$M_{\lambda^a} \quad = \quad \frac{\beta_{g_a}}{g_a} m_{3/2}, \tag{83}$$

where β are beta-functions computed in the supersymmetric limit [401, 243, 238]. The proportionality of the β-function to the low energy masses leads to a specific relationship between the gaugino masses: $M_1 : M_2 : M_3 = 2.8 : 1 : 7.1$. This is very different than would be predicted by the GUT relations used in mSUGRA, for example, with the wino and bino mass hierarchy reversed. When the neutralino mass matrix is diagonalized in AMSB, these ratios result in an LSP which is almost purely (neutral) wino [350, 452]. Additionally, a charged wino, with a mass only a few hundred MeV heavier than the LSP, is predicted.

This leads to a long lived chargino with distinctive collider signatures. Also, in AMSB, with such large values of M_3, the gluino and squarks are predicted to be considerably heavier than in mSUGRA or other similar scenarios.

The heterotic orbifold model The weakly coupled heterotic string with orbifold compactification is among the earliest and best understood string models that can accomodate in four dimensions the Standard Model gauge group, three generations of squarks and a coherent mechanism of supersymmetry breaking. These models show a behavior that interpolates between the phenomenology of mSUGRA and models dominated by superconformal anomalies (AMSB) [95].

Recently, the full one loop soft supersymmetry breaking terms in a large class of superstring effective theories have been calculated [97] based on orbifold compactifications of the weakly coupled heterotic string (including the so-called anomaly mediated contributions). The parameter space in this class of models has already been severely constrained by taking into account accelerator and relic density constraints, as well as direct and indirect searches (see Refs. [90, 95, 94, 312, 103].

Gauge mediated SUSY breaking Another alternative SUSY breaking mechanism is mediated by gauge interactions [178, 374, 26, 179, 180, 181]. In Gauge Mediated Supersymmetry Breaking (GMSB), we have the following approximate relationship between the low energy SUSY masses and the gravitino mass

$$\frac{m_{3/2}}{m_{\rm SUSY}} \sim \frac{1}{\alpha_a}\frac{M_S}{M_{Pl}} \ll 1, \tag{84}$$

where M_S is some typical supersymmetry breaking scale. Therefore, generically, we will have a very light gravitino as the LSP [176]. Such a scenario provides a dark matter candidate which is very difficult to observe. We will not discuss gravitino dark matter or GMSB further for this reason.

Gaugino mediated SUSY breaking Gaugino mediated supersymmetry breaking [316, 139] represents another class of SUSY breaking mediation motivated by the brane-world scenario. In this scenario, unwanted supersymmetry breaking effects, such as the flavor violating couplings, are suppressed by the separation of the observable and hidden sectors via the separation of their respective branes. Gauginos are allowed to propagate off of the branes (in the bulk) in this scenario, communicating SUSY breaking from the hidden sector.

The most important phenomenological feature of this mechanism is that the sfermion masses are suppressed relative to the gaugino masses. This is because sfermion masses can only be generated from the 1-loop diagrams in which a gaugino is emitted, travels through the bulk to the supersymmetry breaking brane, gets the information of SUSY breaking and then returns to join the sfermion propagator again. Generically, these masses are suppressed relative to the gaugino mass by a loop factor, $m_{\tilde{f}}^2 \sim M_\lambda^2/(16\pi^2)$.

3.3 Extra Dimensions

Although our world appears to consist of 3+1 (three space and one time) dimensions, it is possible that other dimensions exist and appear at higher energy scales.

From the physics point-of-view, the concept of extra dimensions received great attention after the idea of Kaluza, in 1921, to unify electromagnetism with gravity by identifying the extra components of the metric tensor with the usual gauge fields. More recently, it has been realized that the hierarchy problem (see section 3.2.1) could be addressed, and possibly solved, by exploiting the geometry of spacetime.

In many extra-dimensional models, the 3+1 dimensional spacetime we experience is a structure called a *brane*, which is embedded in a $(3+\delta+1)$ spacetime called the *bulk*. The hierarchy problem can then addressed by postulating that all of the extra dimensions are compactified on circles (or other topology) of some size, R, as has been done in the Arkani-Hamed, Dimopoulos and Dvali (ADD) scenario [40], thus lowering the fundamental Planck scale to an energy near the electroweak scale. Alternatively, this could be accomplished by introducing extra dimensions with large curvature (warped extra dimensions) as has been suggested by Randall and Sundrum [400]. The extra dimensional scenario which we will focus on throughout the remainder of this review (universal extra dimensions) does not share the features of the ADD or RS scenarios. Rather, it introduces flat extra dimensions which are much smaller than those in the ADD framework.

In addition to the hierarchy problem, motivation for the study of theories with extra dimensions comes from *string theory* and *M-theory*, which today appear to be the best candidates for a consistent theory of quantum gravity and a unified description of all interactions. It appears that such theories may require the presence of six or seven extra-dimensions.

A general feature of extra-dimensional theories is that upon compactification of the extra dimensions, all of the fields propagating in the bulk have their momentum quantized in units of $p^2 \sim 1/R^2$. The result is that for each bulk field, a set of Fourier expanded modes, called Kaluza–Klein (KK) states, appears. From our point of view in the four dimensional world, these KK states appear as a series (called a tower) of states with masses $m_n = n/R$, where n labels the mode number. Each of these new states contains the same quantum numbers, such as charge, color, etc.

In many scenarios, the Standard Model fields are assumed to be confined on the brane, with only gravity allowed to propagate in the bulk. Nevertheless, if the extra-dimensions are small, it would be possible for all fields to freely propagate in the extra dimensions. Such is the case in models with universal extra dimensions, which we discuss in the next section.

3.3.1 Universal extra dimensions

Scenarios in which all fields are allowed to propagate in the bulk are called Universal Extra Dimensions (UED) [33]. Following Ref. [421], we note that there is significant phenomenological motivation to have all Standard Model fields propagate in the bulk, including:

- Motivation for three families from anomaly cancellation
- Attractive dynamical electroweak symmetry breaking
- Prevention of rapid proton decay
- *Provides a viable dark matter candidate.*

In the case of one extra dimension, the constraint on the compactification scale in UED models from precision electroweak measurements is as low as $R^{-1} \gtrsim 300$ GeV [33]. Recently, it was shown that this bound can be weakened to $R^{-1} \gtrsim 280$ GeV if one allows a Higgs mass as heavy as $m_H \gtrsim 800$ GeV [34]. This is to be contrasted with another class of models where Standard Model bosons propagate in extra dimensions while fermions are localized in 4 dimensions. In such cases, the constraint on the compactification scale is much stronger, requiring $R^{-1} \gtrsim$ several TeV [144].

The prospect of UED models providing a viable dark matter candidate is indeed what motivates us in our discussion here. The existence of a viable dark matter candidate can be seen as a consequence of the conservation of momentum in higher dimensional space. Momentum conservation in the compactified dimensions leads to the conservation of KK number. This does not stabilise the lightest KK state, however. To generate chiral fermions at the zero mode, the extra dimensions must be moded out by an orbifold, such as S/Z_2 for one extra dimension or T^2/Z_2 for two. This orbifolding results in the violating of KK number, but can leave a remnant of this symmetry called KK-parity (assuming that the boundary terms match). All odd-level KK particles are charged under this symmetry, thus ensuring that the lightest (first level) KK state is stable. In this way, the Lightest Kaluza-Klein Particle (LKP) is stabalized in a way quite analogous to the LSP in R-parity conserving supersymmetry.

In the next section, we will discuss some of the characteristics of the LKP in models of UED.

3.3.2 The lightest Kaluza–Klein particle

The study of the Lightest Kaluza-Klein Particle (LKP) as a dark matter candidate dates back to the work of Kolb and Slansky in 1984 [324], where the KK excitations were referred to as *pyrgons*, from the Greek $\pi\upsilon\rho\gamma o\varsigma$ for "scale" or "ladder". The LKP has since been reconsidered in the framework of universal extra dimensions, in which it is likely to be associated with the first KK excitation of the photon, or more precisely the first KK excitation of the hypercharge gauge boson [142]. We will refer to this state as $B^{(1)}$.

Channel	Branching ratio
quark pairs	35%
charged lepton pairs	59%
neutrino pairs	4%
Higgs bosons	2%

Table 5: Branching ratios for the annihilation of the $B^{(1)}$ particle. Note that small variations from these results can occur with variation in the KK spectrum.

A calculation of the $B^{(1)}$ relic density was performed by Servant and Tait [421], who found that if the LKP is to account for the observed quantity of dark matter, its mass (which is inversely proportional to the compactification radius R) should lie in the range of 400 to 1200 GeV, well above any current experimental constraint.

We show in Fig. 13 the relic density of the $B^{(1)}$ particle versus its mass, including coannihilations (see section 1.5) with the next-to-lightest KK particle, which in the case shown is $e_R^{(1)}$, the first KK excitation of the right-handed electron. This figure is a new version of Fig. 3 in Ref. [421], updated to include the new WMAP constraints on the cold dark matter relic density [4]. We also show in Fig. 14 the effect of increasing the number of extra dimensions to 2.

Note that the results of the LKP relic density calculation can vary depending on the spectrum of other first level KK states. Unlike in the case of supersymmetry, the density of KK dark matter is *increased* through coannihilations with other KK particles. The spectrum of first level KK states has been calculated to one loop by Cheng, Matchev and Schmaltz [142], although higher dimensional operators localized at the boundary may change the details of the spectrum (without affecting KK parity). Variations in this spectrum can result in variations for the predicted LKP relic abundance.

The $B^{(1)}$ annihilation cross section has been studied in Ref. [421], and is given by

$$\sigma v = \frac{95 g_1^4}{324 \pi m_{B^{(1)}}^2} \simeq \frac{0.6\,\mathrm{pb}}{m_{B^{(1)}}^2[\mathrm{TeV}]} \,. \tag{85}$$

The branching ratios for $B^{(1)}$ annihilation (see table 5) are almost independent of the particle mass. Unlike in the case of supersymmetry, the bosonic nature of the LKP means that there will be no chirality suppression in its annihilations, and thus can annihilate efficiently to fermion-fermion pairs. In particular, since the annihilation cross section is proportional to hypercharge^4 of the final state, a large fraction of LKP annihilations produce charged lepton pairs.

Direct detection of the LKP via its elastic scattering with nuclei was investigated in Refs. [143, 422]. It was emphasized in Ref. [422] that a one-ton detector is needed to probe the expected heavy masses as indicated by the relic density calculation [421] of the LKP. One must, therefore, wait for the next generation

[4]see section 2.3 for a discussion of the CMB and, in particular, the recent WMAP data

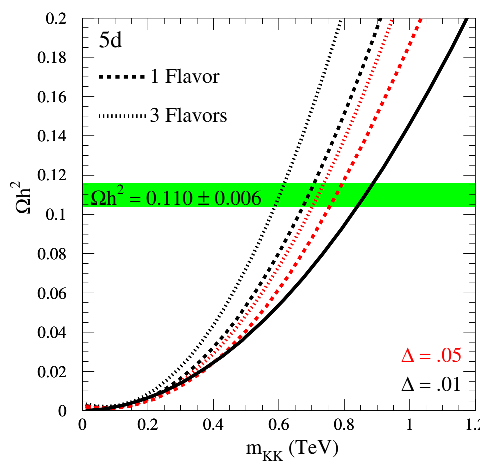

Figure 13: Relic density versus mass of the $B^{(1)}$. The solid line is the case for $B^{(1)}$ alone, dashed and dotted lines are for one (three) flavors of nearly degenerate $e_R^{(1)}$. For each case, black curves (upper of each pair) are for $\Delta = 0.01$ and red curves (lower of each pair) for $\Delta = 0.05$. Figure kindly provided by G. Servant.

of direct detection experiments such as GENIUS [319] or XENON [35] (see section 4.1). Simultaneously, the LHC should probe most of the relevant KK mass parameter space (up to $R^{-1} \sim 1.5$ TeV [141]) and confirm or rule out UED at the TeV scale.

3.4 Superheavy Candidates

Dark matter particles are usually assumed to be relatively "light", meaning lighter than a few hundred TeV. This "limit" is a consequence of the existence of a maximum annihilation cross section, σv, for a particle of a given mass, m_{DM}, set by the so-called *unitarity bound* (see *e.g.* Ref. [465]). Griest and Kamionkowski [264] applied this bound and the constraint on the relic density to infer an upper limit on the dark matter particle mass:

$$m_{DM} \lesssim 340 \text{ TeV} . \tag{86}$$

We note that nowadays, using the WMAP constraint on $\Omega_{DM} h^2$, such a constraint can be made ten times stronger,

$$m_{DM} \lesssim 34 \text{ TeV} . \tag{87}$$

The assumption behind this argument is that the dark matter particle is a thermal relic of the early Universe, otherwise we could not have applied the

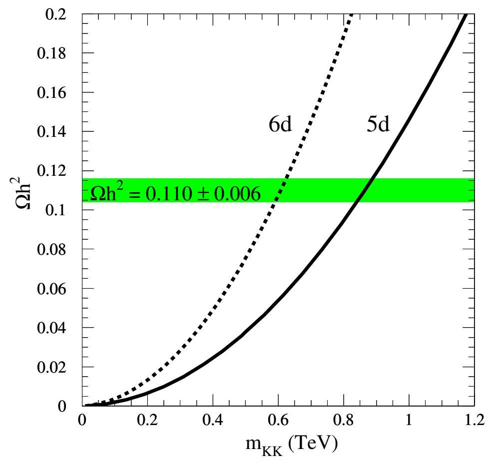

Figure 14: Relic density versus mass of the $B^{(1)}$. Shown is the effect of increasing by 1 the number of extra dimensions. Figure kindly provided by G. Servant.

relation between $\Omega_{DM}h^2$ and σv. In this section, we consider superheavy dark matter candidates, defined as candidates with mass $m_{DM} > 10^{10}$ GeV, that we call generically *wimpzillas* [325]. Thus the first condition for this scenario is that *wimpzillas must not have been in thermal equilibrium during freeze-out.* Since they are not in thermal equlibrium during freeze-out, their relic abundance does not depend on their annihilation cross section, but rather is a function of the wimpzilla's production cross section. Furthermore, we want them to be sufficiently stable against decay and annihilation to significantly contribute to the present day matter density.

There are many ways to produce wimpzillas in the early Universe. Among the most studied is *gravitational production* at the end of inflation, resulting from the expansion of the background spacetime (for details on this and other scenarios see *e.g.* Ref. [325] and references therein).

Natural mass scales for wimpzillas include the inflaton or grand unified masses, which are usually assumed to be roughly 10^{11} GeV or 10^{16} GeV, respectively. Alternatively, D-matter provides a good candidate for wimpzillas with a somewhat larger mass [424]. The interaction cross sections with ordinary matter for such particles can vary from very weak to strong (in the latter case supermassive particles are sometimes called *simpzillas*).

A common motivation for superheavy dark matter comes from the observation of cosmic rays at ultra-high energies[15], above the so-called GZK (Greisen-Zatsepin-Kuzmin) cutoff [260, 478]. Above this cutoff, which occurs at $\sim 5 \times 10^{19}$ eV, protons interact at resonance with CMB photons with a center-of-mass energy nearly equal to the mass of the Δ-hadron (1.232 GeV). The cross section for this interaction is quite large, thus making the Universe opaque to ultra-high energy protons over cosmological distances ($\gtrsim$ 50 Mega-parsecs). Since no astrophysical sources of ultra-high energy protons are known within this range, more exotic scenarios have been developed to account for these observed events. Such scenarios include ultra-high energy cosmic-ray production via the decay or annihilation of superheavy dark matter particles, called top-down cosmic-ray models (see, for example, Refs. [108, 73, 419, 417, 106, 480, 154, 426]).

3.5 Collider Constraints

The constraints which can be placed on a dark matter candidate from collider experiments are highly model dependent in nature. It is, unfortunately, impossible to completely or simply describe the reach of colliders in their search for dark matter in any kind of general way. We will here, rather, review several of the most important collider searches which have been carried out for dark matter particles and for particles associated with a dark matter candidate.

3.5.1 Current collider constraints

- *Invisible Z width*

  If a dark matter candidate is sufficiently light, Z bosons may decay invisibly to such particles with a non-zero branching fraction. Of course, there is a substantial background to such events, namely $Z \to \nu\bar{\nu}$ decays. Presently, to contribute less than one standard deviation to the measured neutrino contribution, the analysis of LEP2 finds that a decay width of $\Gamma_{Z\to XX} < 4.2$ MeV is required (X denotes a dark matter particle).

  Similarly, single photon events can be an interesting search channel for light dark matter particles. At LEP2, the Standard Model background process for this signature is $e^+e^- \to Z \to \nu\bar{\nu}$ with an additional photon radiated off of the initial state. The total cross section for γZ production at LEP2 is less than 31 pb with a minimum 1 GeV transverse momentum cut for the photon. The contribution to this final state (single photon) from particle dark matter, in addition to $Z \to XX$ with an additional photon radiated off of the initial state, are t-channel dark matter producing processes in which a photon is radiated off of a charged propagator (a selectron in supersymmetry, for example).

- *Searches for new charged particles*

  LEP2 has placed very stringent bounds on charged particles lighter than about 100 GeV. In e^+e^- colliders, cross sections for the direct pair production of charged particles are quite large, allowing for limits to be placed at or slightly below half of the center-of-mass energy of the collision. For LEP2, which reached a center-of-mass energy of 209 GeV, limits of 87-103 GeV have been placed for such particles; in particular for charginos

($m_{\chi^\pm} > 103$ GeV) and charged sleptons ($m_{\tilde{e}} > 99$ GeV, $m_{\tilde{\mu}} > 96$ GeV, $m_{\tilde{\tau}} > 87$ GeV) in supersymmetric models. If the LSP is only slightly less massive than the charged particle, however, this limit may be substantially lower [56, 3, 11, 280].

Limits on charged particles can only indirectly constrain dark matter, however. In supersymmetry, chargino masses and neutralino masses are, in some models, related by the unification of gaugino masses. Although such a relationship is often assumed, it is quite possible that the pattern of gaugino masses is not so simple. If gaugino mass unification is assumed chargino limits can translate to neutralino mass limits of about half of the chargino mass limit ($m_{\chi^0} \gtrsim 50$ GeV). Without such a relationship, the LSP could be much lighter [295, 121].

- *Sneutrino limits*

  Limits for charged sleptons can be used to indirectly limit the possible masses for sneutrinos beyond the invisible Z width constraints. Such a bound is the result of a basic SU(2) symmetry between the supersymmetry breaking masses of the left handed slepton and the sneutrino of a given lepton flavor. Limits somewhat lower than for charged particles ($m_{\tilde{\nu}} \gtrsim 85$ GeV) can be placed on sneutrinos if such theoretical assumptions are made [3, 11, 280].

- *Searches for colored particles*

  Hadron colliders, such as the Tevatron, can place the strongest limits on colored particles (squarks and gluinos in supersymmetry or KK excitations of quarks and gluons in models with universal extra dimensions, for example). Such particles would most likely undergo a series of cascades upon decaying, possibly producing dark matter candidates among other particles. Combinations of squarks and gluinos are searched for using jets and missing energy signatures. This leads to exclusion contours in the squark and gluino mass plane.

  In supersymmetry, the spectrum of neutralinos, charginos and sleptons lighter than the decaying squark/gluino is very important in placing limits on squark and gluino masses. Similar ambiguities are present in other models as well, such as universal extra dimensions, etc. Typically, limits of ~200 GeV are obtained for new colored particles, unless there exists an invisible final state particle with a mass close to the new colored particle's mass [4, 5, 6, 9, 397].

- *New gauge bosons*

  Heavy gauge bosons appear in many models of particle physics beyond the Standard Model. Heavy charged gauge bosons (called W''s) and heavy neutral gauge bosons (called Z''s) have been excluded below about 600-800 GeV, depending on the details of the analysis [13, 8]. These limits assume that these particles have couplings equal to their Standard Model counterparts. If their couplings were smaller, the resulting limits could be considerably weaker. Electroweak precision measurements can also constrain heavy gauge bosons considerably (see below).

- *Higgs searches*

  In supersymmetric models, the Higgs mass is increased from its tree level mass (below m_Z) by loop processes involving superparticles, most importantly top squarks. Current bounds on the (lightest) Higgs mass, therefore, constrain the masses of top squarks and other superparticles. Furthermore, if supersymmetry is manifest below 1 TeV, as is normally expected, the Higgs mass must be less than about 130 GeV, not very far above current limits from LEP2 ($m_h < 114.1$ GeV) [341, 342, 281, 282, 2, 10]. Note that this bound is somewhat lower for cases with very large values of $\tan\beta$.

  Searches for charged Higgs bosons can also provide constraints on models of physics beyond the Standard Model. More sensitive, however, may be the impact of charged Higgs bosons in the branching fraction of $b \to s\gamma$.

- *Flavor changing neutral currents*

  Many models of physics beyond the Standard Model introduce flavor changing neutral currents, often at tree level. To avoid the corresponding flavor constraints, either the masses of new particles involved must be quite large, or symmetries must be imposed to solve the "flavor problem". For example, the squark and slepton mass matrices are flavor diagonal in the constrained MSSM (mSUGRA) scenario, thus suppressing such processes. Flavor changing neutral currents in models with universal extra dimensions have also been explored [128].

- $b \to s\gamma$

  The branching fraction for $b \to s\gamma$ [127, 20, 7, 344, 43], measured at CLEO and BELLE, is of particular interest for supersymmetry and other beyond the Standard Model phenomenology. In many scenarios, the contributions to this process from new physics can add substantially to the Standard Model prediction. In particular, light charged Higgs bosons and/or charginos can be quite important for this decay. In supersymmetry, the constraint is considerably stronger if $\mu < 0$, but also relevant for $\mu > 0$, especially for large values of $\tan\beta$ [132]. $b \to s\gamma$ is also an important constraint in models of universal extra dimensions [128].

- $B_s \to \mu^+\mu^-$

  The branching fraction for $B_s \to \mu^+\mu^-$ is quite small in the Standard Model ($\simeq 3.5 \times 10^{-9}$) [32]. The contribution from supersymmetry scales as $\tan^6\beta$, and thus becomes quite large for models with large values of $\tan\beta$. In run I of the Tevatron, a value consistent with the Standard Model was found. The sensitivity of run II of the Tevatron to this quantity will be considerably greater.

- *The anomalous magnetic moment of the muon, $g_\mu - 2$*

  In 2001, the E821 experiment at the Brookhaven National Laboratory reported a measurement of the muon's magnetic moment which was 2.6 standard deviations from the Standard Model prediction [124]. Since then, however, an important error in the theoretical calculation was discovered which reduced the significance of this anomaly to about 1.6 standard deviations [322, 323, 277, 110, 111]. With the reduction of statistical error which has been achieved more recently, the deviation from the Standard Model prediction of this measurement is again about 3σ using e^+e^- data (although the significance is somewhat less using τ decay data) [66, 166, 270, 302]. These measurements, although somewhat difficult to interpret, can be used to constrain TeV-scale particle physics beyond the Standard Model.

- *Electroweak precision measurements*

  In addition to the useful direct particle searches at LEP2, the Tevatron and other experiments, impressively accurate electroweak measurements have been made. Various limits on the scale of new physics and associated particle masses have been inferred from these measurements. Given these constraints, models with universal extra dimensions are limited to the scale $\sim$ 300 GeV or higher [33]. These measurements also yield particularly important bounds for models without a custodial SU(2) symmetry, such as many little Higgs models [161].

Together, these constraints can be very powerful, often providing very tight bounds for specific models. For example, in Fig. 15, we show the impact of collider and cosmological constraints on the constrained Minimal Supersymmetry Standard Model (or mSUGRA). We find that over the parameter space shown, constraints from LEP2 searches (Higgs, charginos and selectrons), along with $g_\mu - 2$ and relic density constraints, leave only a small region near $m_{1/2} \cong 300 - 400$ GeV and $m_0 \cong 80 - 150$ GeV. Although the power of these and other constraints is quite model dependent, they are often very useful in supersymmetry and other classes of models.

3.5.2 The reach of future collider experiments

- *Future reach of the Tevatron*

  The reach of the Tevatron extends to higher energies than any other accelerator until the time at which the Large Hadron Collider (LHC) becomes operational. The range of masses which can be searched for colored particles (squarks, gluinos and KK quarks, for example), heavy gauge bosons and other new physics will be increased significantly at the Tevatron IIb [311, 9].

- *The Large Hadron Collider*

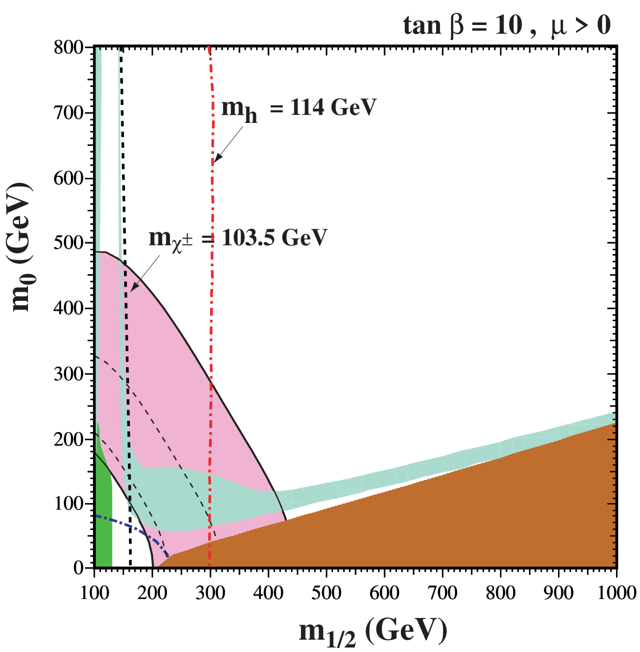

Figure 15: An example of the impact of collider (and cosmological) constraints on a model of particle dark matter [383]. The model shown is the constrained Minimal Supersymmetric Standard Model (cMSSM or mSUGRA) with $\tan\beta = 10, \mu > 0$ and $A_0 = 0$. The almost vertical lines represent the limits on chargino (left) and Higgs (right) masses from LEP2. The blue dot-dash curve in the bottom left corner follows the 99 GeV selectron mass contour, excluded by LEP2. In the dark red region in the lower right, the LSP is a stau and is not, therefore, a viable dark matter candidate. The green region in the lower left corner is excluded by the $b \to s\gamma$ constraint. The long and often narrow turquoise region provides a relic density of $0.1 \leq \Omega h^2 \leq 0.3$, near the observed quantity. The pink region extending over much of the lower left is the region within the 2σ range for $g_\mu - 2$. The two dashed curves within this region are the 1σ bounds.

The Large Hadron Collider (LHC) is expected to begin operation around 2007 with proton-proton collisions at 14 TeV center-of-mass energy. A luminosity of 300 inverse femtobarns is expected to be achieved, making the prospects for discovering new physics at the LHC excellent. Numerous classes of models which provide interesting dark matter candidates will be tested at this very important experiment, searching at scales of up to several TeV. In addition to the Higgs boson(s), the LHC will be sensitive to most supersymmetry scenarios, models with TeV-scale universal extra dimensional, little Higgs models, etc.

For a few examples of studies which discuss the sensitivity of the LHC to new physics, see Refs.[42, 148, 168, 23, 24, 46, 284, 398, 102]. In Fig. 16, an example of such a study is shown [46]. It is interesting to note that in the region of the MSSM which is the most difficult to probe at the LHC, direct dark matter detection rates are very high [45].

- *Beyond the LHC*

  After the LHC, other collider experiments are likely to follow. Although no specific post-LHC program is certain at this time, a 500-1000 GeV linear collider is a possibility, perhaps followed by a Very Large Hadron Collider (VLHC). These or other post-LHC colliders will, of course, have great value to particle dark matter studies [464, 63].

4 Experiments

4.1 Direct Detection Experiments

Direct detection experiments appear today as one of the most promising techniques to detect particle dark matter. The idea is very simple: if the galaxy is filled with WIMPs, then many of them should pass through the Earth, making it possible to look for the interaction of such particles with matter, e.g. by recording the recoil energy of nuclei, as WIMPs scatter off them [191, 254].

The key ingredients for the calculation of the signal in direct detection experiments are the density and the velocity distribution of WIMPs in the solar neighborhood and the WIMP-nucleon scattering cross section. With this information, it is then possible to evaluate the rate of events expected in an experiment (*i.e.* WIMP–nucleon scattering events) per unit time, per unit detector material mass.

The rate is approximately given by

$$R \approx \sum_i N_i n_\chi < \sigma_{i\chi} >, \tag{88}$$

where the index, i, runs over nuclei species present in the detector

$$N_i = \frac{\text{Detector mass}}{\text{Atomic mass of species } i}$$