

## Particle dark matter – A theorist’s perspective

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**Abstract.** Dark matter (DM) is presumably made of some new, exotic particles that appear in extensions of the standard model. After giving a brief overview of some popular candidates, I discuss in more detail the most appealing case of the supersymmetric neutralino.

**Keywords.** Dark matter; supersymmetry.

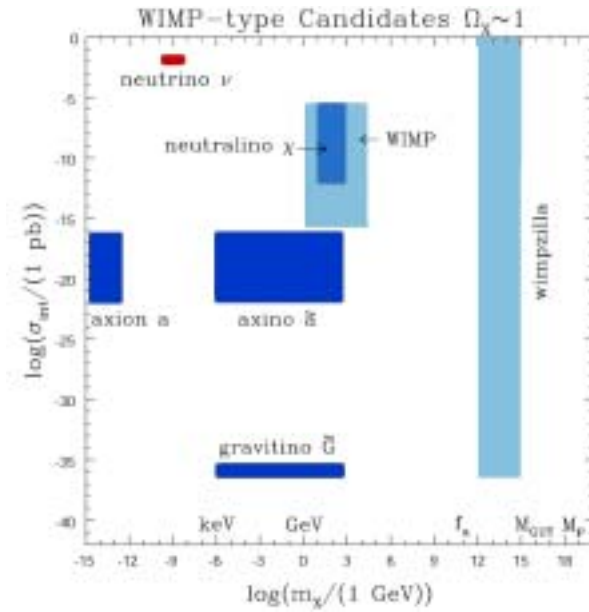
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### 1. Introduction – WIMP-type candidates for DM

While one cannot complain about a shortage of candidates for explaining the nature of the dark matter (DM) in the Universe, from the point of view of particle physics, the WIMP (weakly interacting massive particle) looks particularly attractive. In many ‘scenarios’ as well as more complete theories beyond the SM there often appear several new WIMPs and it is typically not too difficult to ensure that the lightest of them is stable by means of some discrete symmetry or topological invariant. (For example, in supersymmetry, one invokes  $R$ -parity.) In order to meet stringent astrophysical constraints on exotic relics (like anomalous nuclei), they must be electrically and (preferably) color neutral. They can however interact weakly.

Contrary to common wisdom, WIMP-type candidates are neither bound to interact with roughly weak interaction strength (in the sense of electroweak) nor does their mass have to be in the GeV to TeV regime. From a particle theorist’s point of view, it is most sensible to concentrate on the cases where the WIMP appears as a ‘by-product’ in some reasonable frameworks beyond the SM which have been invented to address some other major puzzle in particle physics. In other words, let us talk about WIMP candidates that have not been invented for the sole purpose of solving the DM problem.

One way to present this is to consider a big ‘drawing board’ as in figure 1. A plane spanned by the mass of the WIMP on the one side and by a typical strength  $\sigma_{\text{int}}$  of its interaction with ordinary matter (i.e., detectors) on the other. To a first approximation the mass range can in principle extend up to the Planck mass scale,



**Figure 1.** A schematic representation of some well-motivated WIMP-type particles for which *a priori* one can have  $\omega \sim 1$ .  $\sigma_{\text{int}}$  Represents a typical order of magnitude of interaction strength with ordinary matter. The neutrino provides hot DM which is disfavored. The box marked ‘WIMP’ stands for several possible candidates, e.g., from Kaluza–Klein scenarios.

but not above, if we are talking about elementary particles. The interaction cross-section could reasonably be expected to be of the electroweak strength ( $\sigma_{\text{EW}} \sim 10^{-38} \text{ cm}^2 = 10^{-2} \text{ pb}$ ) but could also be as tiny as that purely due to gravity:  $\sim (m_W/M_P)^2 \sigma_{\text{EW}} \sim 10^{-32} \sigma_{\text{EW}} \sim 10^{-34} \text{ pb}$ .

What can we put into this vast plane shown in figure 1? One obvious candidate is the neutrino, since we know that it exists. Neutrino oscillation experiments have basically convinced us that its mass is at least  $\sim 0.1 \text{ eV}$ . On the upper side, if it were heavier than a few  $\text{eV}$ , it would overclose the Universe. The problem of course is that such a WIMP would constitute *hot* DM which hardly anybody favors these days. While some like it hot, or warm, most like it cold.

Cold, or non-relativistic at the epoch of matter dominance (although not necessarily at freeze-out!) and later, DM particles are strongly favored by a few independent arguments. One is numerical simulations of large structures. Also, increasingly accurate studies of CMB anisotropies, most notably recent results from WMAP [1], imply a large cold DM (CDM) component and strongly suggest that most ( $\sim 90\%$ ) of it is non-baryonic.

In the SUSY world, of course, we could add a sneutrino  $\tilde{\nu}$ , which, like neutrinos, interacts weakly. From LEP its mass is  $\gtrsim 70 \text{ GeV}$  (definitely a cold DM candidate), but then  $\Omega_{\tilde{\nu}} \ll 1$ . Uninteresting and  $\tilde{\nu}$  does not appear in figure 1.

The main suspect for today is of course the neutralino  $\chi$ . Unfortunately, we still know little about its properties. LEP bounds on its mass are actually not

too strong, nor are they robust: they depend on a number of assumptions. In minimal SUSY (the so-called MSSM) ‘in most cases’  $m_\chi \gtrsim 70 \text{ GeV}$ , but the bound can also be much lower. Theoretically, because of the fine tuning argument, one expects its mass to lie in the range of several tens or hundreds of GeV. More generally,  $m_\chi \gtrsim \text{few GeV}$  from  $\Omega_\chi h^2 \lesssim 1$  (the so-called Lee–Weinberg bound [2]) and  $m_\chi \lesssim 300 \text{ TeV}$  from unitarity [3]. Neutralino interaction rates are generally suppressed relative to  $\sigma_{\text{EW}}$  by various mixing angles in the neutralino couplings. In the MSSM they are typically between  $\sim 10^{-3} \sigma_{\text{EW}}$  and  $\sim 10^{-10} \sigma_{\text{EW}}$ , although they could be even lower in more complicated models where the LSP would be dominated for example by a singlino (fermionic partner of an additional Higgs singlet under the SM gauge group). This uncertainty of the precise nature of the neutralino is reflected in figure 1 by showing both the smaller (dark blue) region of minimal SUSY and an extended one (light blue) with potentially suppressed interaction strengths in non-minimal SUSY models. Another example of a WIMP that would belong to the light blue box is the the lightest Kaluza–Klein state which is massive, fairly weakly interacting and stable in some currently popular Kaluza–Klein frameworks [4].

One can see that, while the (s)neutrinos interact weakly *sensu stricto*, this is not quite the case with the neutralino and the other WIMP candidates above. In fact, a typical strength of their interactions can be several orders of magnitude less, while still giving  $\Omega h^2 \sim 1$ .

There are some other cosmologically relevant relics out there whose interactions would be much weaker than electroweak. One well-known example is the axion – a light neutral pseudoscalar particle which is a by-product of the Peccei–Quinn solution to the strong CP problem. Its interaction with ordinary matter is suppressed by the PQ scale  $\sim (m_W/f_a)^2 \sigma_{\text{EW}} \sim 10^{-18} \sigma_{\text{EW}} \sim 10^{-20} \text{ pb}$  ( $f_a \sim 10^{11} \text{ GeV}$ ), hence extremely tiny, while its mass  $m_a \sim \Lambda_{\text{QCD}}^2/f_a \sim (10^{-6}–10^{-4}) \text{ eV}$  if  $\Omega_a \sim 1$ . The axion, despite being so light, is of CDM-type because it is produced by the non-thermal process of misalignment in the early Universe.

In the supersymmetric world, the axion has its fermionic superpartner, called axino. Its mass is strongly model-dependent but, in contrast to the neutralino, often not directly determined by the SUSY breaking scale  $\sim 1 \text{ TeV}$ . Hence the axino could be light and could naturally be the LSP, thus stable. An earlier study concluded that axinos could be *warm* DM with mass less than  $2 \text{ keV}$  [5]. More recently it has been pointed out that more massive axinos quite naturally can also be *cold* DM as well, as marked in figure 1. Relic axinos can be produced either through thermal scatterings and decays involving gluinos and/or squarks in the plasma, or in out-of-equilibrium decays of the next-to-LSP, e.g. the neutralino [6]. The first mechanism is more efficient at larger reheat temperatures  $T_R \gtrsim 10^4 \text{ GeV}$ , the second at lower ones. Axino cosmology is very interesting but I have no time to discuss it here [6].

Lastly, there is the gravitino – the fermionic superpartner of the graviton – which arises by coupling SUSY to gravity. The gravitino relic abundance can be of order one [7] but one has to also worry about the so-called gravitino problem: heavier particles, like the NLSP, will decay to gravitinos very late, around  $10^7 \text{ s}$  after the Big Bang, and the associated energetic photons may cause havoc to BBN products. The problem is not unsurmountable but more conditions/assumptions need to be

satisfied. One way is to assume that the NLSP is mostly a higgsino but this does not normally happen in the framework of grand unification, at least not in minimal models. In figure 1 the gravitino is marked in the mass range of keV to GeV and gravitational interactions only, although keV-gravitinos have actually strongly enhanced couplings via their goldstino component.

While figure 1 is really about WIMPs which arise in attractive extensions of the SM, it is worth mentioning another class of relics, popularized under the name of WIMPzillas, for which there exist robust production mechanisms (curvature perturbations) in the early Universe [8]. As the name suggests, they are thought to be very massive,  $\sim 10^{13}$  GeV or so. There are no restrictions on WIMPzilla interactions with ordinary matter, other than that they must interact at least gravitationally, as schematically depicted in figure 1.

In summary, the number of well-motivated WIMP and WIMP-type candidates for CDM is in the end not so large. On the other hand, one should remember that in the box marked with the generic name ‘WIMP’ one can accommodate not just the neutralino but also some other stable states appearing in various extensions of the SM, e.g., Kaluza–Klein type theories. One can add to this picture other candidates, like cryptons and other particles arising in the context of superstrings, or mirror DM [9].

SUSY, which to many is the most promising extension of the SM, provides three robust WIMP candidates for the CDM: the neutralino, the axino and the gravitino. Each has its virtues and weak points but I will have no time to discuss this here. I think that it is fair to say that the WIMP for today, and this decade, is the neutralino. It is therefore of primary interest here because it is present in any sensible SUSY spectrum and is testable in today’s experimental programmes.

In this talk, I will explore cosmological properties of the neutralino as the LSP in the general MSSM and in two unification-based models: the constrained MSSM (CMSSM) and in the GUT model based on the  $SO(10)$  gauge group. (See the talk by Raby [10].) I will next present ensuing predictions for the cross-section for DM neutralino elastic scattering with detector material in both models and contrast them with the case of the general MSSM. See also the complementary presentation of Nojiri [11] for a discussion of indirect detection and collider search for SUSY aspects of the neutralino.

## 2. The neutralino

### 2.1 The frameworks

There are two basic schemes in which one usually considers cosmological properties of the neutralino. One is a rather general framework of the MSSM where superpartner masses originate primarily from soft SUSY-breaking terms. Additionally, one assumes a common mass parameter  $m_{1/2}$  for all the gauginos in the spectrum which leads to the well-known relations  $M_1 = \frac{5}{3} \tan^2_{\theta_W} M_2 \simeq 0.5 M_2$  and  $M_2 = \frac{\alpha_2}{\alpha_s} m_{\tilde{g}} \simeq 0.3 m_{\tilde{g}}$ , with  $M_1/M_2/m_{\tilde{g}}$  being the soft bino/wino/gluino mass. One further imposes the  $R$ -parity to make the LSP stable.

Alternatively, it is more popular today to consider specific boundary conditions at some high scale, like the grand unification scale  $M_{\text{GUT}}$  or the string scale. In unified models one writes down the Lagrangian at the unification scale and next employs the renormalization group equations (RGEs) to compute the couplings and masses in an effective theory valid at the electroweak scale. One such popular model is the CMSSM, also known as mSUGRA, where one assumes a common soft mass scale  $m_0$  for all the scalars (sfermions and Higgs) and a common trilinear soft SUSY-breaking parameter  $A_0$ . These parameters are run using their respective RGEs down from  $M_{\text{GUT}}$  to some appropriately chosen low-energy scale  $Q_0$  where the Higgs potential (including full one-loop corrections) is minimized while keeping the usual ratio  $\tan\beta = v_t/v_b$  of the Higgs VEVs fixed. The Higgs/higgsino mass parameter  $\mu$  and the bilinear soft mass term  $B\mu$  are then computed from the conditions of radiative electroweak symmetry breaking (EWSB), and so are the Higgs and superpartner masses. The CMSSM thus *a priori* has only the usual  $\tan\beta, m_{1/2}, m_0, A_0, \text{sgn}(\mu)$  as input parameters. However, in the case of large  $m_{1/2}, m_0 \gtrsim 1$  TeV and/or large  $\tan\beta \sim \mathcal{O}(m_t/m_b)$  some resulting masses will in general be highly sensitive to the assumed physical masses of the top and the bottom (as well as the tau).

Another interesting scenario is a fully realistic effective SUSY model which derives from a minimal GUT with the  $SO(10)$  gauge group (MSO<sub>10</sub>SM [10,12]). It involves a different set of boundary conditions at  $M_{\text{GUT}}$  which leads to a distinctively different set of phenomenological and cosmological predictions [10,13]. One starts with a well-defined model at the GUT scale: the sfermions of all the three families have a unified (soft) mass  $m_{16}$ , while for the Higgs (which come in  $\mathbf{10}_H$  and its conjugate) the analogous unified quantity is  $m_{10}$ . After running the parameters down to  $m_Z$ , one finds that experimental constraints require  $A_0 \sim -2m_{16}$ ,  $m_{10} \sim \sqrt{2}m_{16}$ ,  $m_{16} \geq 2$  TeV  $\gg \mu, m_{1/2}$ , and  $\Delta m_H^2 \sim 10\%$ , where  $m_{(H_t, H_b)}^2 = m_{10}^2(1 \mp \Delta m_H^2)$  [12].  $\tan\beta$  is necessarily large ( $\sim 50$ ) because of  $t$ - $b$ - $\tau$  Yukawa unification.

The case of large  $\tan\beta$  requires one to treat the top, bottom and tau masses with special care. These (especially  $m_b$ ) receive large radiative corrections from SUSY and one has to be careful in extracting from them the corresponding Yukawa couplings which in turn have an important effect on the running of the RGEs at large  $\tan\beta$ . The masses of the top and tau are treated with a similar accuracy although corrections to their masses are typically smaller.

Despite a small number of independent parameters in unified models, their analysis is often technically rather involved, especially at large  $\tan\beta$ , as mentioned above. Furthermore, in order to reduce the scale dependence of the Higgs sector and related conditions for the EWSB it is important to include full one-loop corrections to the Higgs potential *and* also minimize the Higgs potential at the scale  $Q_t \sim \sqrt{m_{t_1} m_{t_2}}$  with  $m_{t_1} (m_{t_2})$  denoting the physical masses of the stops. This is because, at this scale, the role of the otherwise large log-terms  $\sim \log(m_t^2/Q^2)$  from the dominant stop-loops will be reduced. At this scale one evaluates the one-loop conditions for the EWSB which determine  $\mu^2$  as well as the bilinear soft mass parameter  $B\mu$ . Alternatively, one can run down all the mass parameters down to  $m_Z$  (or their physical values, if they are above  $m_Z$ ) and include all the threshold corrections to the  $\beta$ -functions due to the decoupling of states.

The mass of the pseudoscalar  $m_A$  plays an important role in evaluating  $\Omega_\chi h^2$ , especially at very large  $\tan\beta \sim 50$ . This is so for three reasons: (i)  $m_A$  decreases

with increasing  $\tan\beta$  [14] due to the increased role of the bottom Yukawa coupling which at large enough  $\tan\beta$  opens up a wide resonance  $\chi\chi \rightarrow f\bar{f}$ ; (ii)  $A$ -resonance is dominant due to the coupling  $Af\bar{f} \sim \tan\beta$  for down-type fermions; and (iii) in contrast to the heavy scalar  $H$ , this channel is not  $p$ -wave suppressed [15]. One needs to stress however that there still remains considerable uncertainty (of the order of  $\sim 5$ – $10\%$ ) in the procedure for computing Higgs masses and conditions for EWSB.

## 2.2 The bino is the winner

The lightest neutralino is the lowest-lying mass eigenstate of the two gauginos ( $\tilde{B}$  and  $\tilde{W}_3$ ) and the two higgsinos ( $\tilde{H}_b^0, \tilde{H}_t^0$ ) which are the fermionic superpartners of the neutral gauge and Higgs boson states, respectively. Remarkably, despite *a priori* many free parameters present in SUSY theories, the state that appears to most naturally give  $\Omega_\chi h^2 \sim 1$  is the nearly pure bino [16]. This is true in both the MSSM and in basically all unified models. Indeed, for the higgsino one typically finds  $\Omega_\chi h^2 \ll 1$  due to efficient annihilation to  $WW, ZZ, t\bar{t}$  and coannihilation. In unified models, the LSP neutralino is often a nearly pure bino [17–19] because of the requirement of radiative EWSB which typically gives  $|\mu| \gg M_1$ . This often allows one to impose strong constraints from  $\Omega_\chi h^2 < \mathcal{O}(1)$  on *most* of the region of  $m_{1/2}$  and  $m_0$  in the ballpark of 1 TeV, as originally shown in [18,19] and later confirmed by many studies. There are a few exceptions to this rough upper bound which play an important role in unified SUSY models. The bino as the cosmologically favored LSP has now become part of the standard lore.

## 2.3 Experimental constraints

First I will summarize the relevant experimental constraints.

1.  $m_{\chi_1^\pm} > 104 \text{ GeV}$  from LEP. This constraint is fairly robust, except for some fairly degenerate cases.
2.  $m_h > 111 \text{ GeV}$  in the case of light SM-like Higgs which holds in unified models. The actual limit from LEP is  $m_h > 114.1 \text{ GeV}$  but theoretical uncertainty is still about 2–3 GeV. In the MSSM, there remains a narrow corridor of  $m_h \simeq m_A$  down to some 90 GeV.
3.  $\text{BR}(B \rightarrow X_s \gamma) = (3.34 \pm 0.68) \times 10^{-4}$  [20,21]. This constraint is very important in the light of a rather good agreement between experiment and the SM prediction ( $\text{BR}(B \rightarrow X_s \gamma)_{\text{SM}} = (3.70 \pm 0.30) \times 10^{-4}$  [22]). However, it is also very sensitive to underlying theoretical assumptions and has to be applied with much care. If one makes the usual simplifying assumption that the mass mixing in the down-type squark sector is the same as in the corresponding quark sector (the so-called minimal flavor violation, or MFV, scenario), then the constraint provides a very strong lower bound on  $m_{1/2}$  and  $m_0$ , especially at large  $\tan\beta$  where SUSY contributions are strongly enhanced. However, by allowing for even a mild relaxation of the MFV assumption, one finds that

the constraint from  $b \rightarrow s\gamma$  may become very much weaker, or even disappear altogether [21]. We will show its effect in the (C)MSSM but not in the MSO<sub>10</sub>SM where the MFV scenario does not necessarily hold. The constraint favors  $\mu > 0$  although there are ways to overcome this [21].

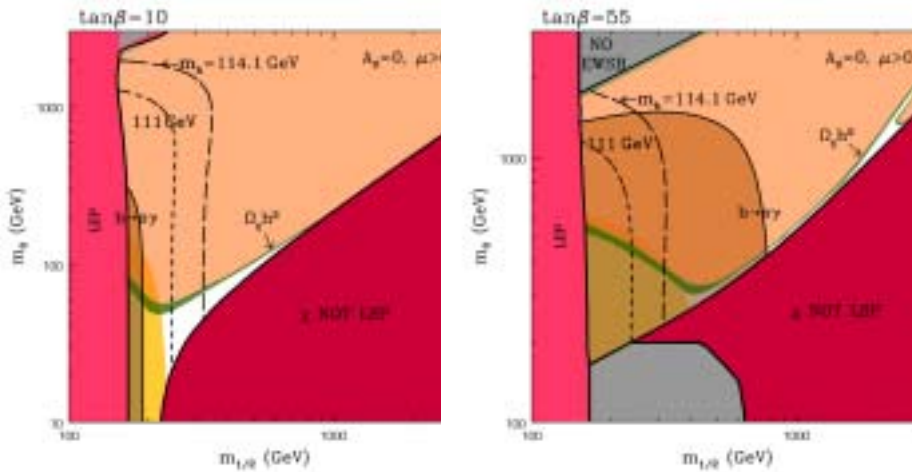
4.  $a_\mu^{\text{expt}} - a_\mu^{\text{SM}} = (22.1 \pm 11.3) \times 10^{-10}$  ( $e^+e^-$  data) and  $a_\mu^{\text{expt}} - a_\mu^{\text{SM}} = (7.4 \pm 10.5) \times 10^{-10}$  ( $\tau$  data), where  $a_\mu = (g_\mu - 2)/2$  is the anomalous magnetic moment of the muon. These numbers are quoted from the latest analysis [23]. This reduces previous discrepancy with the SM value:  $1.9\sigma$  and  $0.7\sigma$ , respectively, and also between the analyses using the  $e^+e^-$  and  $\tau$ -data. It also favors  $\mu > 0$ .
5.  $0.095 < \Omega_\chi h^2 < 0.13$  ( $2\sigma$ ), a stringent range determined recently by WMAP (+CBI and ACBAR) [1]. Larger values are excluded. Smaller ones are in principle allowed but would imply that the neutralino would only be a subdominant component of the CDM. This constraint will have an extremely strong effect on the parameter space of SUSY and allowed ranges of  $m_\chi$ , but not on the detection cross-sections, as we will see later.

#### 2.4 Allowed parameter space

I will now present the implications of experimental and cosmological constraints, as described above, on the parameter space of the CMSSM and the MSO<sub>10</sub>SM. Next we will discuss the ensuing ranges for direct detection elastic scattering cross-sections.

*The CMSSM.* In the CMSSM there are two distinct cases: low to moderate  $\tan\beta$  regime and the case  $\tan\beta \sim 50$  [20]. We will illustrate this in figure 2 in the plane  $(m_{1/2}, m_0)$  for a representative choice of  $\tan\beta$  for each class. In the left and right panel of figure 2 the experimental and cosmological bounds are presented for  $\tan\beta = 10$  and  $50$ , respectively. We also fix  $A_0 = 0$ ,  $\mu > 0$ ,  $m_t \equiv m_t^{\text{pole}} = 175$  GeV and  $m_b \equiv m_b(m_b)_{\overline{\text{MS}}} = 4.25$  GeV. We can see many familiar features. At  $m_{1/2} \gg m_0$  there is a dark red wedge where the  $\tilde{\tau}_1$  is the LSP. On the other side, at  $m_0 \gg m_{1/2}$  we find large gray regions where the EWSB is not achieved. Just below the region of no-EWSB the parameter  $\mu^2$  is small but positive which allows one to exclude a further (light red) band by imposing the LEP chargino mass bound. As one moves away from the wedge of no-EWSB,  $\mu^2$  increases rapidly. That implies that, just below the boundary of the no-EWSB region, the LSP neutralino very quickly becomes the usual nearly pure bino. This causes the relic abundance  $\Omega_\chi h^2$  to accordingly increase rapidly from very small values typical for light higgsinos, through the narrow strip ( $\Delta m_{1/2} \sim 5$  GeV) of the cosmologically expected (green) range ( $0.095 < \Omega_\chi h^2 < 0.13$ ) up to larger values which are excluded (light orange). In particular, in the whole region allowed by the chargino mass bound the LSP is mostly bino-like.

A general pattern of the cosmologically favored regions does not change much until  $\tan\beta \sim 45$ – $50$  (depending somewhat on  $m_b$ ,  $m_t$ , etc.). Generally one finds a robust (green) region of expected  $\Omega_\chi h^2$  at  $m_{1/2} \sim m_0$  in the range of a few hundred GeV [19]. In addition, at  $m_{1/2} \gg m_0$ , just above the wedge where the LSP is the  $\tilde{\tau}_1$ , the coannihilation of the neutralino LSP with  $\tilde{\tau}_1$  opens up a very



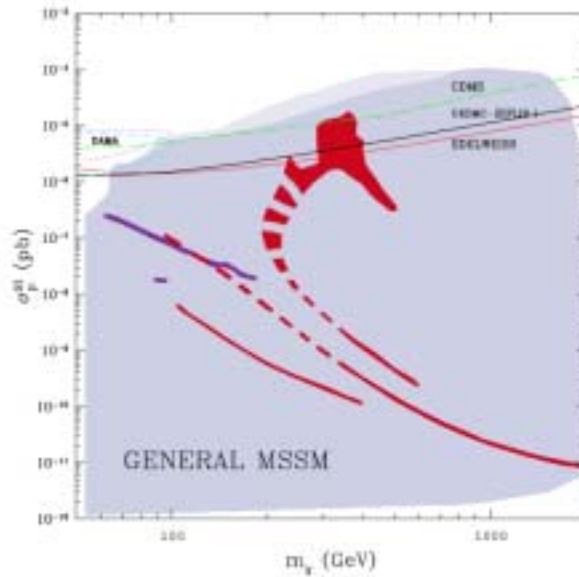
**Figure 2.** The plane  $(m_{1/2}, m_0)$  for  $A_0 = 0$ ,  $\mu > 0$  and for  $\tan\beta = 10$  (left) and  $\tan\beta = 55$  (right). We fix  $m_t \equiv m_t^{\text{pole}} = 175$  GeV and  $m_b \equiv m_b(m_b)_{\text{SM}}^{\text{MS}} = 4.25$  GeV. The light red bands on the left are excluded by chargino searches at LEP. In the gray regions the electroweak symmetry breaking conditions are not satisfied or  $m_A^2 < 0$  while in the dark red region denoted ‘ $\chi$  NOT LSP’ the LSP is the lighter stau. The large light orange regions of  $\Omega_\chi h^2 > 0.13$  are excluded by cosmology while the narrow green bands correspond to the expected range  $0.095 < \Omega_\chi h^2 < 0.13$  ( $2\sigma$ ). The region to the left of the lightest Higgs scalar mass  $m_h = 111$  GeV and 114.1 GeV are excluded. The light brown region is excluded by  $b \rightarrow \gamma$  (assuming MFV). Also shown is the semi-oval (yellow) region which is excluded at  $2\sigma$  by the anomalous magnetic moment of the muon measurement.

long and narrow corridor of  $m_{1/2}$  and  $m_0$  favored by  $0.095 < \Omega_\chi h^2 < 0.13$ . At  $m_0 \gg m_{1/2}$ , very close to the region of no-EWSB, again one finds a very narrow range of  $\Omega_\chi h^2$  consistent with observations.

The new feature that appears at large  $\tan\beta$  is the effect of a very distinct, wide pseudoscalar Higgs resonance in the annihilation process  $\chi\chi \rightarrow A \rightarrow f\bar{f}$ . Since, as mentioned above,  $m_A$  decreases with increasing  $\tan\beta$ , at some point, this opens up a corridor in the plane of  $(m_{1/2}, m_0)$  along  $m_A = 2m_\chi$ . Clearly, at large  $\tan\beta$  cosmological constraints on  $\Omega_\chi h^2$  permit much larger superpartner masses, not only in the very narrow strips close to the regions of no-EWSB and/or  $\tilde{\tau}_1$ -LSP, but especially because of the resonance.

The existence of the resonance and of the region of no-EWSB are quite generic but their exact positions at large  $\tan\beta$  are rather sensitive to the relative values of the top and bottom masses. Generally, at fixed  $\tan\beta$ , increasing(decreasing) the top mass relative to the bottom mass causes the region of no-EWSB to move up(down) considerably because of the diminishing(growing) effect of the bottom Yukawa coupling on the loop correction to the conditions of EWSB. At fixed top and bottom masses, as  $\tan\beta$  decreases, the region of no-EWSB moves toward somewhat larger  $m_0$  and smaller  $m_{1/2}$  but the overall effect is not very significant.





**Figure 3.** Scalar interaction cross-section  $\sigma_p^{SI}$  vs.  $m_\chi$ . The lowest (intermediate) red and blue bands correspond to the cosmologically favored regions of figure 2 with  $\tan\beta = 10$  ( $\tan\beta = 55$ ), which are allowed by experimental constraints, except  $(g-2)_\mu$ , and by the favored range  $0.095 < \Omega_\chi h^2 < 0.13$  ( $2\sigma$ ). The top band is for the non-unified Higgs model case with  $\tan\beta = 60$ ,  $A_0 = 0$  and  $\delta_u = 1$ . In the dashed red bands  $b \rightarrow s\gamma$  is not satisfied (within the MFV scenario). The blue (light blue) region is predicted by the general MSSM by varying  $10 < \tan\beta < 65$  and assuming cosmological and collider constraints including (except for)  $b \rightarrow s\gamma$ . See text for more details.

Regarding other constraints, the lightest Higgs mass excludes a sizeable region of smaller  $m_{1/2}$ . In the figures we plot the contours of  $m_h = 114.1$  GeV and  $m_h = 111$  GeV to show the earlier mentioned effect of the uncertainties in computing  $m_h$ . Larger values of  $m_h$  are given by contours which are shifted along the  $m_{1/2}$  axis with roughly equal spacings but diverge somewhat at larger  $m_0$ . The (light brown) region excluded by  $\text{BR}(B \rightarrow X_s \gamma)$  grows significantly because the dominant chargino-squark contribution to the branching ratio grows linearly with  $\tan\beta$ . On the other hand, one has to remember that the constraint has been derived in the MFV scenario and can easily be relaxed (or strengthened) by even a small departure from scheme. Finally,  $(g-2)_\mu$  robustly excludes an oval-shape (yellow) region of small  $m_{1/2}$  and  $m_0$ . An upper bound still exists at  $1\sigma$  but is now much weaker than before, and it disappears completely at  $2\sigma$ .

One of the most promising strategies to detect WIMPs in the galactic halo is to look for the effect of their elastic scattering from a target material in an underground detector. How do the ‘theory plots’ of figure 2 translate into the more familiar (to the general community) language of WIMP mass vs. cross-section ones? And how do they compare with predictions following from less constrained SUSY models, like the general MSSM?

A relevant quantity is a scalar, or spin-independent (SI), interaction cross-section on a free proton at zero momentum transfer  $\sigma_p^{\text{SI}}$ . First, in figure 3 we show the case of the CMSSM for the same choices of parameters as in figures 2. We plot allowed ranges of  $\sigma_p^{\text{SI}}$  vs. the neutralino mass. For comparison, the wide (blue) region is predicted by the general MSSM assuming rather generous ranges of SUSY parameters [24]. In particular, we take  $5 < \tan \beta < 65$  and  $A_0 = 0, \pm 1$  TeV.

The two lower narrow (red) strips correspond to the allowed configurations of the CMSSM of figure 2, with the lower (middle) one corresponding to  $\tan \beta = 10$  ( $\tan \beta = 55$ ). In deriving them, the collider bounds mentioned above and  $0.095 \lesssim \Omega_\chi h^2 \lesssim 0.13$  were applied, as well as the  $1\sigma$  constraint from  $b \rightarrow s\gamma$ . Since the last one is highly sensitive to theoretical assumptions at the unification scale and excludes large regions of smaller  $m_{1/2}$  and  $m_0$ , the effect of removing it is shown by a dashed line. One can see that, at large  $\tan \beta$ , ranges of much larger  $\sigma_p^{\text{SI}}$  (at smaller  $m_\chi$ ) are allowed. On the other hand, we do not apply the constraint from  $a_\mu$ . Its effect in the right window would be to exclude (at  $2\sigma$ ) the part with  $m_\chi \lesssim 150$  GeV in the main (red) band, but *not* in the (violet) focus point band. For  $\tan \beta = 10$  instead the Higgs mass bound remains stronger.

For each  $\tan \beta$  a general pattern is that of two distinct ‘branches’. The lower (red) one corresponds for the most part to the  $m_A$ -resonance and/or  $\tilde{\tau}$ -coannihilation region at  $m_{1/2} \gg m_0$  in figures 2. On the other hand, the upper (violet) branch comes from the focus point region of large  $m_0 \gg m_{1/2}$ . In this region, even though the masses of squarks and sleptons are large, in the TeV range, the mass of the lightest neutralino remains moderate, since roughly  $m_\chi \sim 0.4m_{1/2}$ . In both branches, the dominant contribution to  $\sigma_p^{\text{SI}}$  comes from the exchange of the heavy scalar Higgs.

One can clearly see that predictions of the CMSSM are rather definite (thus fully justifying the name the model bears). In particular, note a remarkably narrow band of  $\sigma_p^{\text{SI}}$  as a function of  $m_\chi$ . Generally larger values of  $\sigma_p^{\text{SI}}$  and smaller  $m_\chi$  are predicted by the extremely narrow focus point region of  $m_0 \gg m_{1/2}$ , while the long (red) tail comes from the region of the neutralino-slepton coannihilation and (at large  $\tan \beta \gtrsim 50$ ) of the wide resonance. Generally, as  $\tan \beta$  increases, so does the cross-section at fixed  $m_\chi$ .

This implies that, for both small and large  $\tan \beta$  there is also an upper bound on  $m_\chi$  because the cosmologically allowed regions in the plane  $(m_{1/2}, m_0)$  eventually end at large  $m_{1/2} \gg m_0$ . As  $m_{1/2}$  grows, the neutralino-slepton coannihilation eventually ceases to remain effective and  $\Omega_\chi h^2$  grows beyond 0.13. At very large  $\tan \beta \gtrsim 50$  the  $A$ -resonance extends in the plane  $(m_{1/2}, m_0)$  up to very large values of  $m_{1/2}$  but eventually ends, too. Thus follows a lower bound on  $\sigma_p^{\text{SI}}$  of about  $\sim 8 \times 10^{-11}$  pb and an upper limit  $m_\chi \lesssim 1950$  GeV, as can be (barely) seen in the figure.

The top red region (with an amusing shape) has been added to show a rather exceptional case of boundary conditions for which the predicted  $\sigma_p^{\text{SI}}$  is unusually high, in fact already partly excluded by experiment. It corresponds to the scenario where the soft mass terms of the Higgs doublet at the GUT scale is not unified with the other scalars. The case presented in the figure is for  $\tan \beta = 60$ ,  $A_0 = 0$  and  $\delta_u = 1$  where  $m_{H_2}^2 = m_0^2 (1 + \delta_u)$ . This specific example is somewhat unusual because for other choices of  $\tan \beta$  and/or  $\delta_u < 1$  the predicted ranges of  $\sigma_p^{\text{SI}}$  are

typically  $\lesssim 10^{-7}$  pb, like in the CMSSM, while on the other hand usually giving patterns distinctively different from those predicted in the CMSSM. The hashed (full) red region corresponds to relaxing (imposing) the constraint from  $b \rightarrow s\gamma$ .

In contrast to the specific patterns resulting from different unification assumptions, in the general MSSM one is struck by the enormity of the *a priori* allowed ranges, extending down to  $\sim 10^{-12}$  pb, or, at lower  $m_\chi$ , even below [24]. The lowest values of  $\sigma_p^{\text{SI}}$  generally correspond to one or more SUSY mass parameters being very large, above 1 TeV, and thus can be considered more fine-tuned, and therefore (hopefully!) less likely. Requiring SUSY mass parameters to be less than a few TeV allows one to put a lower bound on  $\sigma_p^{\text{SI}}$ .

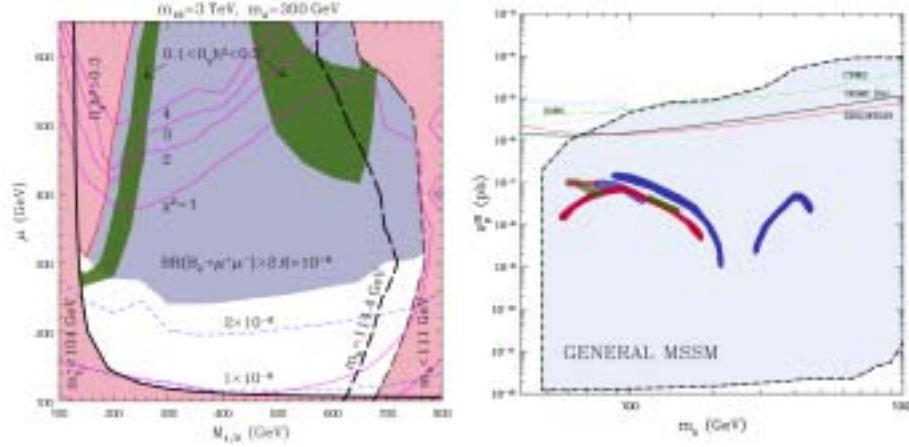
At lower  $m_\chi$  the lower bound on  $\sigma_p^{\text{SI}}$  can actually extend to much lower values. This is because  $\Omega_\chi h^2$  is determined there primarily by neutralino coannihilation with sleptons. By selecting a slepton mass not too much above  $m_\chi$ , one can make coannihilation reduce  $\Omega_\chi h^2$  to the favored range and have at the same time very low  $\sigma_p^{\text{SI}}$ . This requires some fine-tuning but three or four orders of magnitude for  $\sigma_p^{\text{SI}}$  (or even smaller) are allowed by allowing for SUSY mass parameters in the multi-TeV range. The coannihilation with sleptons is very effective at  $m_\chi$  in the range of several hundred GeV (with the upper limit growing with  $\tan\beta$ ) but at some point ceases to be effective enough and  $\Omega_\chi h^2$  grows to too large values. One can still reduce it to acceptable values by either fine-tuning  $m_A$  to twice  $m_\chi$  or by taking  $|\mu|$  very (!) large.

That is for the lower ranges, which hopefully will never have to be probed! On the upper side, values as large as  $\sim 10^{-5}$  pb are allowed. At lower  $m_\chi$  they are limited from above by a lower bound on light Higgs mass and by the lower limit on  $\Omega_\chi h^2$ . At larger  $m_\chi$  the limit from  $b \rightarrow s\gamma$  is the main constraint. The effect of lifting it is marked by a lighter shade of blue. The constraint from  $a_\mu$  puts an upper bound (at  $1\sigma$ ) of  $m_\chi \gtrsim 450$  GeV but not on  $\sigma_p^{\text{SI}}$  [24], and not at  $2\sigma$ .

The current experimental sensitivity is also shown for comparison. Denoted are the regions claimed by DAMA to be consistent with the annual modulation signal allegedly present in their data when one ignores (magenta dots) or includes (blue dash) the Collaboration's previous limits. Some other experiments (CDMS, UKDMC, Edelweiss) have by now excluded much of the DAMA region.

*The  $\text{MSO}_{10}\text{SM}$ .* A typical example of the parameter space is presented in the left panel of figure 4. (For other cases see [13] and the talk by Raby [10].) The mass of the pseudoscalar is fairly low which allows for efficient depletion of the number of WIMP neutralinos in the early Universe through the  $A$ -resonance, as in the case of the CMSSM. The vertical position of the pole is clearly visible in the panel at  $m_{1/2} \simeq 350$  GeV. Away from it,  $\Omega_\chi h^2$  is too large, close to it, it is too small. In between one finds sizeable cosmologically favored (green) regions. Impact of other relevant constraints is also shown. We do not apply the constraint from  $b \rightarrow s\gamma$  because in this model one is not bound by the MFV framework. On the other hand, all the sfermions are heavy and the model predicts basically the SM value for  $(g_\mu - 2)/2$ .

Another aspect of the  $\text{MSO}_{10}\text{SM}$  which is of much interest to phenomenology is the process  $B_s \rightarrow \mu^+ \mu^-$ . Because the process involves a flavor-changing exchange of the pseudoscalar, which is fairly light, the resulting  $\text{BR}(B_s \rightarrow \mu^+ \mu^-)$  can be large, often well within the reach of Run II at the Tevatron. At lower  $m_A$ , it can even



**Figure 4.** Left panel: The cosmologically favored (green) and allowed (white) regions in  $MSO_{10}SM$  for  $m_{16} = 3\text{ TeV}$  and  $m_A = 300\text{ GeV}$ . Red and blue regions are excluded, as marked in the panel. Right panel: Allowed regions of  $\sigma_p^{SI}$  in  $MSO_{10}SM$  for different choices of parameters. See text for more details.

exceed the current bound  $2.6 \times 10^{-6}$  from CDF which is expected to be improved by a some two orders of magnitude. In short, the process offers very good prospects for the Tevatron.

In the right panel of figure 4 we present  $\sigma_p^{SI}$  vs.  $m_\chi$  for several typical choices of  $m_A = 300\text{ GeV}$  (light shade) and  $500\text{ GeV}$  (dark shade), and for  $m_{16} = 2.5\text{ TeV}$  (green),  $3\text{ TeV}$  (red) and  $5\text{ TeV}$  (blue) [13]. The patterns are somewhat different from those of the CMSSM (thus possibly allowing for discriminating between the two models should a WIMP signal be detected and  $\sigma_p^{SI}$  and  $m_\chi$  measured with some accuracy). One also finds  $\sigma_p^{SI} \lesssim 10^{-7}\text{ pb}$  although also  $\sigma_p^{SI} \gtrsim 10^{-9}\text{ pb}$  which is encouraging, compared to the CMSSM.

In summary, it is clear that the current experimental sensitivity is sufficient to already probe a part of the parameter space of the general MSSM. On the other hand, it is still typically at least one order of magnitude above the preferred ranges of cross-sections that are predicted by unified models as the examples of the CMSSM and the  $MSO_{10}SM$  demonstrate. For these ranges to be explored a new generation of detectors will be required and is actually already being constructed. On the other hand, theoretically lower ranges of WIMP mass and therefore larger  $\sigma_p^{SI}$  are more natural (fine-tuning and  $a_\mu$ ). It is therefore quite possible that a nice surprise may come before long. Finally, since the specific ranges of  $\sigma_p^{SI}$  and  $m_\chi$  are typically very narrow and model-dependent, once a WIMP signal is detected, one may hope to be able to discriminate among different unification scenarios.

## References

- [1] D N Spergel *et al*, *Astrophys. J. Suppl.* **148**, 175 (2003)
- [2] B W Lee and S Weinberg, *Phys. Rev. Lett.* **39**, 165 (1977)

- For more references, see E W Kolb and M S Turner, *The early universe* (Addison Wesley, Redwood City, 1990)
- [3] K Griest and M Kamionkowski, *Phys. Rev. Lett.* **64**, 615 (1990)
  - [4] G Servant and T Tait, *Nucl. Phys.* **B650**, 391 (2003); *New J. Phys.* **4**, 99 (2002)
  - [5] K Rajagopal, M S Turner and F Wilczek, *Nucl. Phys.* **B358**, 447 (1991)
  - [6] L Covi, J E Kim and L Roszkowski, *Phys. Rev. Lett.* **82**, 4180 (1999)  
L Covi, H B Kim, J E Kim and L Roszkowski, *J. High Energy Phys.* **0105**, 033 (2001)  
L Covi, L Roszkowski and M Small, *J. High Energy Phys.* **0207**, 023 (2002)
  - [7] H Pagels and J R Primack, *Phys. Rev. Lett.* **48**, 223 (1982)  
S Weinberg, *Phys. Rev. Lett.* **48**, 1303 (1982)  
J Ellis, A D Linde and D V Nanopoulos, *Phys. Lett.* **B118**, 59 (1982); *Phys. Lett.* **B443**, 209 (1998); and several more recent papers
  - [8] D J H Chung, E W Kolb and A Riotto, *Phys. Rev. Lett.* **81**, 4048 (1998)  
V Kuzmin and I Tkachev, hep-ph/9809547
  - [9] R N Mohapatra and V L Teplitz, *Phys. Rev.* **D62**, 063506 (2000)
  - [10] S Raby, *Pramana – J. Phys.* **62**, 523 (2004)
  - [11] M Nojiri, *Pramana – J. Phys.* **62**, 335 (2004)
  - [12] T Blažek, R. Dermíšek and S Raby, *Phys. Rev. Lett.* **88**, 111804 (2002); *Phys. Rev.* **D65**, 115004 (2002)
  - [13] R Dermíšek, S Raby, L Roszkowski and R Ruiz de Austri, *J. High Energy Phys.* **0304**, 037 (2003)
  - [14] M Drees and M Nojiri, *Phys. Rev.* **D47**, 376 (1993)
  - [15] See, e.g., G Jungman, M Kamionkowski and K Griest, *Phys. Rep.* **267**, 195 (1996)
  - [16] L Roszkowski, *Phys. Lett.* **B262**, 59 (1991)  
See also J Ellis, D V Nanopoulos, L Roszkowski and D N Schramm, *Phys. Lett.* **B245**, 251 (1990)
  - [17] P Nath and R Arnowitt, *Phys. Lett.* **B289**, 368 (1992)
  - [18] R G Roberts and L Roszkowski, *Phys. Lett.* **B309**, 329 (1993)
  - [19] G L Kane, C Kolda, L Roszkowski and J Wells, *Phys. Rev.* **D49**, 6173 (1994)
  - [20] L Roszkowski, R Ruiz de Austri and T Nihei, *J. High Energy Phys.* **0108**, 024 (2001)
  - [21] K Okumura and L Roszkowski, *J. High Energy Phys.* **0310**, 024 (2003); hep-ph/0208101
  - [22] A J Buras, A Czarnecki, M Misiak and J Urban, *Nucl. Phys.* **B631**, 219 (2002)
  - [23] M Davier, S Eidelman, A Hocker and Z Zhang, hep-ph/0308214
  - [24] Y G Kim, T Nihei, L Roszkowski and R Ruiz de Austri, *J. High Energy Phys.* **0212**, 034 (2002)