Direct and Indirect Detection of Dark Matter in Heterotic Orbifold Models

P. Binétruy $^{1,3}$, Y. Mambrini $^2$, E. Nezri $^{1,3}$

$^1$ Laboratoire de Physique Théorique des Hautes Énergies
Université Paris-Sud, F-91405 Orsay.

$^2$ Departamento de Física Teórica C-XI and Instituto de Física Teórica C-XVI,
Universidad Autónoma de Madrid, Cantoblanco, 28049 Madrid, Spain.

$^3$ APC, Université Paris 7,
Collège de France, F-75231 Paris Cedex 05.

Abstract

We study the neutralino dark matter phenomenology in the context of effective field theories derived from the weakly–coupled heterotic string. We consider in particular direct detection and indirect detection with neutrino telescopes rates. The two cases of moduli dominated and dilaton dominated SUSY breaking lead to completely different phenomenologies. Even if in both cases relic density constraints can be fulfilled, moduli domination generically leads to detection rates which are much below the present and future experimental sensitivities, whereas dilaton domination gives high detection rates accessible to the next generation of experiments. This could make dark matter searches an alternative way to constrain high energy fundamental parameters.
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1 Introduction

There exists a large collection of measurements providing convincing evidence in favor of the existence of cold dark matter in the universe \cite{1}. But the exact nature of this dark matter is still an open question. One of the most promising and best motivated candidates is a Weakly Interacting Massive Particle (WIMP). Direct detection via its rare scattering with a nucleus in a detector, or indirect detection via its annihilation after gravitational storage in a massive body provide two possible experimental strategies.

In the framework of supersymmetry (SUSY), most extensions of the Standard Model (SM) predict a massive neutral weakly interacting particle in the form of a neutralino ($\chi^0$). Moreover, in the simplest versions of SUSY models such as the minimal supergravity model (mSUGRA), this particle is predicted to be stable. Recent works have constrained a large part of the parameter space available in mSUGRA \cite{2} in constrained or unified versions of the model. For small values of $M_0$, we can even find in the literature some strong consequences on the limit of the neutralino mass: $< 500 \text{ GeV}$ \cite{3} if we took into account all the recent accelerator analysis, but this can be strongly evaded when one allows high values of $M_0$ and $m_{1/2}$ \cite{4}. Up to now, several works have generalized these simple models to a non-universal framework in the higgs sector \cite{5,6,7,8}, the gaugino sector \cite{8,9,10} or the sfermion sector \cite{11}. Such analyses show that neutralino dark matter searches are sensitive to the spectrum of supersymmetric particles, but the direct connection with the supersymmetry-breaking sector is not made because it appears only through the parameters of the effective theory, the so-called soft supersymmetry breaking parameters.

If one wants to make explicit the supersymmetry breaking mechanism, one must identify the type of breaking as well as the nature of the mediation between the supersymmetry breaking sector and the observable sector. A standard example is gaugino condensation in a hidden sector which interacts only gravitationally with the sector of quarks and leptons. Since this involves gravitational interactions, it is natural to consider such a model in a string context and thus to include all fields associated with the gravitational sector, in particular the dilaton (whose vacuum expectation value determines the magnitude of the string coupling) and moduli fields (whose expectation values determine the size of the compact manifold). Both types of fields play an important role in supersymmetry breaking. The most elaborated class of models and probably the most realistic one from a phenomenological point of view is associated with the weakly coupled heterotic string compactified on a Calabi-Yau manifold or an orbifold.

Recently, the full one loop soft supersymmetry breaking terms in a large class of superstring effective theories have been calculated \cite{12} based on orbifold compactifications of the weakly-coupled heterotic string (including the so-called anomaly me-
diated contributions). Such models yield specific non-universalities which make their phenomenology significantly different from the minimal supergravity model. The parameter space in this class of models has already been severely constrained by taking into account accelerator and relic density constraints \[13\] \[14\] or benchmark models at the Tevatron \[15\] \[16\].

In what follows, we take this specific class of models to discuss how direct and indirect neutralino detection depend on the properties of supersymmetry breaking. Indeed, supersymmetry breaking induced by gaugino condensation may be transmitted through the auxiliary field vacuum expectation values of the compactification moduli or of the dilaton. We will see that the two corresponding cases have a very different behavior regarding dark matter detection.

2 Theoretical framework

2.1 Structure of heterotic orbifolds models at one loop

The task of string phenomenology is to make contact between the high energy string theory, and the low energy world. For this purpose, we need to build a superstring theory in four dimensions, able to give us the Standard Model gauge group, three generations of squarks, and a coherent mechanism of SUSY breaking. We will set here our analysis in the framework of orbifold compactifications of the heterotic string within the context of the supergravity effective theory. We concentrate on those models where the action is dominated by one loop order contributions to soft breaking terms. The key property of such models is the non–universality of soft terms, consequence of the beta–function appearing in the superconformal anomalies. This non–universality gives a particular phenomenology in the gaugino and the scalar sector, modifying considerably the predictions coming from mSUGRA. In fact, these string–motivated models show new behavior that interpolates between the phenomenology of unified supergravity models (mSUGRA) and models dominated by the superconformal anomalies (AMSB). The constraints arising from accelerator searches and relic density have been already studied in \[13\]. It is thus interesting, to see to which extent direct and indirect detection of dark matter will be able to bring us extra information on the models.

We provide a phenomenological study within the context of orbifold compactifications of the weakly–coupled heterotic string, where we distinguish two regimes. In the first one, the SUSY breaking is transmitted by the compactification moduli \( T^\alpha \), whose vacuum expectation values determine the size of the compact manifold. Generic (0,2) orbifold models contain three \( T^\alpha \) moduli fields. We considered a situation in which only an ”overall modulus \( T \)” field contributes to SUSY–breaking. The use of an overall modulus \( T \) is equivalent to the assumption that the three \( T^\alpha \) fields of generic orbifold models have similar contributions to SUSY–breaking. This is expected in the
absence of some dynamical effect that would strongly discriminate between the three moduli. In our second example, it is the dilaton field $S$ present in any four–dimensional string (whose vacuum expectation value determine the magnitude of the unified coupling constant $g_{\text{STR}}$ at the string scale), that transmits, via its auxiliary component, the SUSY breaking. We work in the context of models in which string nonperturbative corrections to the Kähler potential act to stabilize the dilaton in the presence of gaugino condensation [17, 18]. The origin of the soft breaking terms are completely different in the two scenarios. Some are coming from the superconformal anomalies and are non–universal (proportional to the beta–function of the Standard Model gauge groups), others are generated in the hidden sector (from Green–Schwarz mechanism or gaugino condensation) and are thus universal. This mixture between universality and non–universality gives the richness of the phenomenology in this type of effective string models and confirms the interest of non–universal studies in the prospect of supersymmetric dark matter detection, the non–universality being in this case connected with the basic properties of the model.

2.2 The moduli dominated scenario

In the moduli dominated scenario, the one loop order supersymmetric SUSY breaking terms at GUT scale can be written [12, 19, 20]:

$$M_a = \frac{g_a^2(\mu)}{2} \left\{ 2 \left[ \frac{\delta_{G S}}{16\pi^2} + b_a \right] G_2(T, T) F^T + \frac{2}{3} b_a \overline{M} \right\}, \quad (2.1)$$

$$A_{ijk} = -\frac{1}{3} \gamma_i \overline{M} - p \gamma_i G_2(T, T) F^T + \text{cyclic}(ijk), \quad (2.2)$$

$$M_i^2 = (1 - p) \gamma_i |M|^2 / 9. \quad (2.3)$$

where $M_a$ and $M_i$ are the soft masses for the gauginos and scalars and $A_i$, the trilinear coupling. $b_a$ is the beta–function coefficient for the gauge group $G_a$:

$$b_a = \frac{1}{16\pi^2} \left( 3 C_a - \sum_i C_a^i \right). \quad (2.4)$$

where $C_a$, $C_a^i$ are the quadratic Casimir operators for the group $G_a$ in the adjoint representation and in the representation of the field $i$ respectively. $F^S$ and $F^T$ are the auxiliary fields for the dilaton and the Kähler modulus, respectively, $\overline{M}$ is the supergravity auxiliary fields whose vacuum expectation value ($\vev$) determines the gravitino mass $m_{3/2} = -\frac{1}{3} \overline{M}$, and $\delta_{G S}$ is the Green–Schwarz coefficient which is a (negative) integer between 0 and $-90$. The function $G_2(T, T)$ is proportional to the
Eisenstein function and vanishes when $T$ is stabilized at one of its two self–dual points. From Eq. (2.1), it follows that when the moduli are stabilized at a self dual point, only the second term contributes to gaugino masses. This is precisely the "anomaly mediated" contribution. The loop contributions have been computed using the Pauli–Villars (PV) regularization procedure. The PV regular fields mimic the heavy string modes that regulate the full string amplitude. The phenomenological parameter $p$ which represents the effective modular weight of the PV fields is constrained to be no larger than 1, though it can be negative in value. Thus the scalar squared mass for all matter fields is in general non–zero and positive at one loop (only the Higgs can have a negative running squared mass). The limiting case of $p = 1$, where the scalar masses are zero at one loop level and for which we recover a sequestered sector limit, occurs when the regulating PV fields and the mass–generating PV fields have the same dependence on the Kähler moduli. Another reasonable possibility is that the PV masses are independent of the moduli, in which case we would have $p = 0$. and $\gamma_i$ is related to the anomalous dimension through $\gamma_j^i = \gamma_i \delta_j^i$. (see [12, 13, 15] for notations and conventions)

We clearly see in these formulae the competition between universal terms and non–universal ones. The scalar mass terms are all non–universal and proportional to their anomalous dimension $\gamma_i$ and thus loop suppressed. The Green-Schwarz mechanism generates universal breaking terms for the gauginos (proportional to $\delta_{GS}$) whereas superconformal anomalies introduce non–universal contributions (proportional to $b_a$). The nature of the neutralino thus depends mainly on the value of the Green–Schwarz counterterm $\delta_{GS}$, whereas the mass scale is the gravitino mass $m_{3/2}$.

### 2.3 The dilaton dominated scenario

We turn now to a scenario where the dilaton is the primary source of supersymmetry breaking in the observable sector. It is well known that if we use the standard Kähler potential derived from the tree level string theory, it is very difficult to stabilize the dilaton at acceptable weak–coupling values. We postulate in our study nonperturbative correction of stringy origin to the dilaton Kähler potential. In that case, one condensate can stabilize the dilaton at weak coupling while simultaneously ensuring vanishing expectation values at the minimum of the potential. The key feature of such models is the deviation of the dilaton Kähler metric from its tree level value. If we imagine the superpotential for the dilaton having the form $W(S) \propto e^{-3S/b_+}$, with $b_+$ being the largest beta–function coefficient among the condensing gauge groups of the hidden sector, then we are led to consider the phenomenology of models given by the following pattern of soft supersymmetry breaking terms [12, 19, 20]:
\[ M_a = \frac{g^2}{2} \left\{ \frac{2}{3} b_a M + [1 - 2b'_a K_s] F^S \right\} \tag{2.5} \]

\[ A_{ijk} = -\frac{K_s}{3} F^S - \frac{1}{3} \gamma_i M + \tilde{\gamma}_i F^S \left\{ \ln(\mu^2_{\text{UV}}/\mu^2_R) - p \ln [(t + \bar{t})|\eta(t)|^4] \right\} + (ijk) \tag{2.6} \]

\[ M^2_i = \frac{|M|^2}{9} \left[ 1 + \gamma_i - \left( \sum_a \gamma^a_i - 2 \sum_{jk} \gamma^{jk}_i \right) \left( \ln(\mu^2_{\text{UV}}/\mu^2_R) - p \ln [(t + \bar{t})|\eta(t)|^4] \right) \right] \]

\[ + \left\{ \tilde{\gamma}_i - \frac{MF^S}{6} + \text{h.c.} \right\}, \tag{2.7} \]

where \( \mu_{\text{UV}} \) is an ultraviolet regularization scale (of the order of the string scale \( M_{\text{STR}} \)) and \( \mu_R \) the renormalization scale (taken at the boundary value of \( M_{\text{GUT}} \)). Moreover

\[ F^S = \sqrt{3} m_3(K_{s\bar{s}})^{-1/2}, \quad K_{s\bar{s}} = \partial_s \partial_{\bar{s}} K. \tag{2.8} \]

and

\[ (K^{s\bar{s}})^{-1/2} = \sqrt{3} \frac{2}{3} b_+ - \frac{2}{3} b_+ K_s, \quad K_s = -g^2_{\text{STR}}/2. \tag{2.9} \]

with \( g_{\text{STR}} \) the unified constant at \( M_{\text{STR}} \) (see \cite{12, 13, 15} for the notations and conventions) to ensure a vanishing vacuum energy in the dilaton–dominated limit.

The phenomenology of a dilaton–dominated scenario is completely different from a moduli–dominated one. If we look at formulae (2.5) and (2.7), it is clear that we are in a domain of heavy squarks and sleptons (of the order of magnitude of the gravitino mass) and relatively light gauginos. Indeed, the factor \( b_+ \), as it contains a loop factor, can suppress the magnitude of the auxiliary field \( F^S \) relative to that of the supergravity auxiliary field \( M \) through the relation (2.8). The resulting gaugino soft breaking terms are less universal for low values of \( b_+ \).

\footnote{For simplicity, we assume here that \( M_{\text{STR}} \sim \mu_{\text{UV}} \sim \mu_R \sim M_{\text{GUT}} \) (the corresponding error is logarithmic and appears in a loop factor).}
3 Supersymmetric dark matter phenomenology

3.1 Relic density

We recall that, in a general supersymmetric model, the neutralino mass matrix reads in the bino, wino, higgsino basis \( (\tilde{B}, \tilde{W}^3, \tilde{H}_1^0, \tilde{H}_2^0) \):

\[
M_N = \begin{pmatrix}
M_1 & 0 & -m_Z \cos \beta \sin \theta_W & m_Z \sin \beta \sin \theta_W \\
0 & M_2 & m_Z \cos \beta \cos \theta_W & -m_Z \sin \beta \cos \theta_W \\
-m_Z \cos \beta \sin \theta_W & m_Z \cos \beta \cos \theta_W & 0 & -\mu \\
m_Z \sin \beta \sin \theta_W & -m_Z \sin \beta \cos \theta_W & -\mu & 0
\end{pmatrix}. \tag{3.10}
\]

Thus the lightest neutralino \( \chi \) is generically a superposition of these states:

\[
\chi = z_{1,1} \tilde{B} + z_{2,2} \tilde{W} + z_{3,3} \tilde{H}_1 + z_{4,4} \tilde{H}_2 \tag{3.11}
\]

In a large parameter space of the mSUGRA model, the \( \chi \) is mainly bino–like because of the renormalization group evolution of \( m_{H_u}^2 \) down to the electroweak scale. Indeed, let us recall the radiative electroweak symmetry breaking relation

\[
\mu^2 = \frac{\left( m_{H_u}^2 + \delta m_{H_u}^2 \right) - \left( m_{H_d}^2 + \delta m_{H_d}^2 \right) \tan^2 \beta - \frac{1}{2} \tan^2 \beta}{\tan^2 \beta - 1}, \tag{3.12}
\]

where \( \delta m_{H_u}^2 \) and \( \delta m_{H_d}^2 \) represent the one loop tadpole corrections to the running Higgs masses \( m_{H_u}^2 \) and \( m_{H_d}^2 \) \[21\,22\,23\]. In (3.12) all the parameters are running and set at the minimization scale. We see that, as \( m_{H_u}^2 \) is becoming negative and relatively large in absolute value at low energy, \( |\mu| \) which sets the higgsino mass scale, becomes large, in fact larger than the gaugino mass terms \( M_1 \) and \( M_2 \). Since in mSUGRA models, \( M_{1,2} \) are linked at the weak scale by the well–known relation \( M_1 = \frac{5}{3} \tan^2 \theta_W M_2 \sim 0.5 M_2 \) (with some corrections at 2–loop order running), the \( \chi \) aligns along its bino component. When \( m_{H_u}^2 \) (negative) decreases in absolute value, \( \mu \) is then smaller, increasing the higgsino content of the neutralino. In the \( m_{H_u}^2 \) decreasing direction, this leads successively to a mixed bino-higgsino neutralino and then an higgsino LSP and finally the no EWSB boundary when \( \mu = 0 \). This happens typically in the hyperbolic branch (focus point) mSugra region thanks to the heavy scalar scale.

Different processes lead to a cosmologically favoured neutralino relic density. For a bino neutralino one needs light sfermions, \( \tilde{t}(\tilde{\tau}) \) coannihilation or annihilation into pseudo-scalar \( A \). If the lightest neutralino has a dominant wino component, the relic density drops down because of efficient annihilations into gauge bosons as well as strong \( \chi \chi_1^+ \) coannihilations. For a non negligible higgsino component neutralino annihilates into gauge bosons or \( t\bar{t} \) and relic density is also decreased by \( \chi \chi_1^+ \) and \( \chi \chi_2^0 \) coannihilations.
3.2 Direct detection

Direct detection consists of measuring the energy deposited in a low background detector by the recoil of a nucleus from its elastic scattering with a Weakly Massive Interacting Particle (WIMP) [24]. In our case, the best WIMP candidate is the lightest neutralino $\chi$. The rates follow the neutralino-proton spin–independent elastic cross sections ($\sigma^{\text{scal}}_{\chi-p}$) or spin–dependent one ($\sigma^{\text{spin}}_{\chi-p}$), function on the target nucleus spin [25, 26].

$\sigma^{\text{scal}}_{\chi-p}$ is essentially driven by first generation scalar quark ($\bar{u}_i$, $\bar{d}_i$) exchanges or neutral Higgs ($h, H$) ($\chi q \xrightarrow{H, \bar{d}_i} \chi q$) and the spin–dependent one $\sigma^{\text{spin}}_{\chi-p}$ by first generation squark and $Z$ exchanges ($\chi q \xrightarrow{Z, \bar{d}_i} \chi q$). The processes involving $Z$-boson exchange being completely dependent on the neutralino higgsino fraction.

The main impact of the null searches, particularly at LEP, is in the increase in the lower limit to the LSP mass, as well as the rest of the sparticle spectrum.

In the mSUGRA case, the bino-like nature of the lightest neutralino $\chi$ implies a highly suppressed scalar cross section via heavy neutral higgs exchange because of the low couplings $\chi H \chi$ (proportional to the product of their higgsino and gaugino component) at moderate value of tan $\beta$. If we look at the spin–dependent cross section, dominated by $Z$ exchange, it becomes much larger when $\chi$ is mostly higgsino because of the enhancement due to the $\chi Z \chi$ coupling (proportional to the square of the higgsino components).

The real possibilities of an enhancement of $\sigma^{\text{scal}}_{\chi-p}$ are thus in the high tan $\beta$ regime or in models that predict low first generation squarks masses for moderate gaugino masses, like some of the string inspired models that we study in this paper. Another possibility is to find regions of parameter space that increase the higgsino component of the lightest neutralino, enhancing in the same way its coupling. This happens for large $M_0$ in mSUGRA along the “no EWSB” boundary. We will see that for a large class of heterotic orbifold models, all these constraints can be achieved. At the opposite, some of them will exhibit a complete depletion of the scalar cross–section, far below the sensitivity of the next generation of detectors. In this sense, direct detection of dark matter can become an important tool in the effort of discriminating string inspired models and constraining their parameter space.

The current experimental status may be briefly summarized as follows. Although one of the current experiments, the DAMA collaboration [27] claimed an evidence and gave a determination of the allowed maximum–likelihood region in the WIMP–mass and WIMP–nucleon cross section of $10^{-6} - 10^{-5}$ pb for a WIMP’s mass between 30 and 270 GeV, other experiments exclude almost (CDMS [28], EDELWEISS [29]) or completely (ZEPLIN I [31]) the DAMA region. But many new or upgraded versions of direct detection experiments will soon reach a significantly improved sensitivity for WIMP detection (EDELWEISS II [30], ZEPLIN(s) [31]). Our study will be placed
in the light of this next generation of detectors to see whether a SUSY dark matter candidate could be directly detected in the next years, and how this would constrain some of the fundamental SUSY breaking terms, and, as a consequence, the parameters of the more fundamental string theory.

3.3 Neutrino indirect detection

Neutralinos can also be gravitationally captured in massive astrophysical bodies like the Sun or the Earth by successive diffusions on their nuclei leading to a trapped neutralino population at their center. Then neutralinos can annihilate and the annihilation products, essentially gauge bosons and heavy quarks, decay emitting neutrinos. After conversion into muons through the Earth, these neutrinos can be observed by neutrino telescopes collecting Čerenkov light of induced muons traveling in water or in ice. The annihilation rate depends both on the capture rate \( \sigma_{\chi-p}^{\text{scal/spin}} \) (depending on the target nucleus spin) and on the neutralino annihilation cross section \( \sigma_{\chi-\chi}^{\text{ann}} \). The capture rate in the Earth depends on \( \sigma_{\chi-p}^{\text{scal}} \) because of the zero spin of the iron nucleus. This leads to muon fluxes far beyond reach of detection \([32]\). In the case of the Sun, capture rates are enhanced by \( Z \) exchange in \( \sigma_{\chi-p}^{\text{spin}} \) thanks to the non zero spin of the hydrogen. Fluxes are then maximized for a substantial higgsino fraction. Furthermore, mixed higgsino neutralino states annihilate into \( W^+W^- \), \( ZZ \) or \( t\bar{t} \) leading to more energetic neutrinos/muons than other annihilation channels. This happens along the “noEWSB” boundary, where (thanks to \( m_{H_u}^2 \) running) the neutralino gets a dominant higgsino fraction. This is the case in mSUGRA for high values of \( M_0 \) \([33, 32]\) and can be strongly favored with non universal gaugino masses by decreasing \( M_3|_{\text{GUT}} \) \([8]\).

Some experiments have constrained these fluxes (Macro \([34]\), Baksan \([35]\), Super-Kamiokande \([36]\)), but future neutrinos telescopes like Antares \([37]\) and Icecube \([38]\) will be much more efficient. They will improve current sensitivities of order \( 5 \times 10^3 \mu \text{km}^{-2} \text{yr}^{-1} \) to \( 10^3 \) – \( 10^2 \mu \text{km}^{-2} \text{yr}^{-1} \) on muon fluxes coming from neutralino annihilations in the centre of the Sun. We will compare our predictions for fluxes coming from the Sun with both current and future sensitivities.

As we will see below, in the models that we are considering in this work, the muon fluxes coming from the Sun are very dependent on the nature of SUSY breaking, i.e. whether it is dominated by moduli or dilaton \( F \) terms. In this way, a neutralino dark matter signal in a neutrino telescope would provide key information on the nature of SUSY breaking and on the fundamental underlying theory.

Works on prospects for direct and/or neutrino indirect detection of neutralino dark matter in mSUGRA/CMSSM or non universal frameworks can be found in \([4, 5, 39, 3, 4, 9, 10, 41, 42, 43, 33, 32, 11, 10]\).
3.4 Tools and experimental constraints

In this section we describe the tools that we have used and the various constraints that we have imposed to obtain a correct phenomenology at the electroweak scale for our analysis.

We have used the Fortran code SuSpect2 [44] to solve the renormalization group equations (RGEs) for the soft supersymmetry breaking parameters between the high energy boundary scale \( \mu_{\text{UV}} \) and the scale given by the Z-boson mass (electroweak scale). While the initial scale \( \mu_{\text{UV}} \) should itself be treated as a model-dependent parameter, for our purposes we have chosen for \( \mu_{\text{UV}} \) the scale of grand unification \( M_{\text{GUT}} \). We use \( \tan \beta \) and the sign of the supersymmetric \( \mu \) parameter in the superpotential as free parameters, defined at the low-energy (electroweak) scale.

The magnitude of the \( \mu \) parameter is determined by imposing electroweak symmetry breaking (EWSB) at the usual scale \( (m_{\tilde{t}_1} m_{\tilde{t}_2})^{1/2} \) [45, 46]. The one-loop corrected \( \mu \) is obtained from the condition (3.12).

The soft supersymmetry breaking parameters at the weak scale are then passed on to the C code micrOMEGAs [47] to perform the calculation of physical masses for the superpartners and various indirect constraints, to be described below.

We note that the value of the \( \mu \) parameter is fundamental in the analysis of astroparticle processes, because it determines the nature, mass and couplings of the lightest neutralino \( \chi \) which we require to be the LSP. We then estimated the detection rates using the DarkSusy package [48].

The remaining parameter space is further reduced by limits on superpartner and Higgs masses from various collider experiments. We take the most recent bounds given by the different experiments of the LEP Working Group [49, 50]. Concerning the light CP-even neutral Higgs mass \( m_h \), we assume that a 95\% confidence level (CL) lower limit on \( m_h \) is set at 111.5 GeV. The search for an invisibly decaying Higgs boson in \( hZ \) production has allowed a 95\% CL lower limit on \( m_h \) to be set at 114.5 GeV, assuming a production cross section equal to that in the Standard Model and a 100\% branching fraction to invisible decays [51]. We believe the value of 113.5 GeV will serve as a good mean. Concerning the chargino limit, we take 103.5 GeV, bearing in mind that in some degenerate cases and for light sleptons the limit can go down to 88 GeV [52]. For the squark sector the limit of 97 GeV [53] is used. For all mass bounds we should keep in mind that experimental limits are always given in the context of a particular SUSY model framework which is not generally a string motivated one. The bounds we use could possibly be weakened in some cases.

Various non-collider observations can be used to further reduce the allowed parameter space of the loop-dominated orbifold models that we consider. We will focus our attention on the three sets of data that are the most constraining for these models: the density of relic neutralino LSPs, the branching ratio for decays involving the process
$b \to s\gamma$ and the measurement of the anomalous magnetic moment of the muon. We apply:

**Relic density:**

Recent evidence suggests [54] that $\Omega_\chi \sim 0.3$ with $h^2 \sim 0.5$. We will take as a conservative favored region

$$0.03 < \Omega_\chi h^2 < 0.3.$$  \hspace{1cm} (3.13)

In addition to our conservative dark matter limit, we also take into account the recent results of WMAP [55] that give a $2\sigma$ range for the density of cold dark matter, $\Omega_{\text{CDM}} h^2 = 0.1126^{+0.0161}_{-0.0181}$. Let us stress that the requirement of (3.13) should not be treated as an exact constraint, but rather as an indication of the region preferred by cosmological considerations because of the uncertainty in mass spectrum calculation.

**$b \to s\gamma$ Constraint:**

Another observable where the SUSY particle contributions can be important and measurable is the flavor changing decay $b \to s\gamma$ [56]. In the Standard Model, this process is mediated by virtual isospin $+1/2$ quarks and $W$-bosons. In supersymmetric theories, the spectrum allows new contributions involving loops of charginos and squarks or top quarks and charged Higgs bosons. For our analysis, we use the results given by the CLEO and BELLE collaborations [57]. We adopt the procedure taken in the recent benchmark study of Battaglia et al. [58] and choose to impose the constraint

$$2.33 \times 10^{-4} < \text{BR}(b \to s\gamma) < 4.15 \times 10^{-4}.$$ \hspace{1cm} (3.14)

**The Muon Anomalous Magnetic Moment:**

Recently, the Brookhaven collaboration has given a new measurement of the anomalous magnetic moment of the muon [62]

$$\frac{(g_\mu^{\exp} - 2)}{2} = a_\mu^{\exp} = 11 659 202 (14) (6) \times 10^{-10},$$ \hspace{1cm} (3.15)

Following [63] we introduce the parameter $\delta_\mu$ to quantify the difference between theoretical and experimental determinations of $a_\mu$:

$$\delta_\mu \equiv (a_\mu - 11 659 000 \times 10^{-10}) \times 10^{10}.$$ \hspace{1cm} (3.16)

From this the current experimental determination of the parameter $\delta_\mu$ is $\delta_\mu^{\exp} = 203 \pm 8$. In our discussion, we are less conservative than the authors of [63] and consider a $2\sigma$ standard deviation region about the anomalous moment of the muon based on the $\tau$ decay analysis [59]:

$$-11.6 < \delta_\mu^{\text{new physics}} = \delta_\mu^{\exp} - \delta_\mu^{\text{SM}} < 30.4 \quad [2 \sigma].$$ \hspace{1cm} (3.17)
4 The Models

4.1 The general moduli-dominated case

We present in Figs 1–4 the parameter space allowed by experimental constraints and the corresponding values of two important observables in the context of dark matter: the spin-independent scalar cross-section $\sigma^{\text{scal}}_{\chi p}$ and the muon flux from the Sun.

Let us start with a general comment concerning the relic density. As we can see looking back at (2.1) and (2.3), the main feature of the moduli-dominated regime is to have gaugino and scalar masses mediated by one loop corrections and threshold effects. This implies a phenomenology with light neutralinos and possibly light squarks. In any case, it has been shown in [13] that the lightest neutralino $\chi$ keeps a wino nature in a broad region of parameter space, being degenerated with the lightest chargino $\chi^+_1$. The first effect of this degeneracy is a complete depletion of the relic density which barely reaches $10^{-2}$ due to the strong coannihilations channel ($\chi\chi^+_1$). The only way of splitting these two masses is by the influence of the universal negative Green–Schwarz counterterm $\delta_{\text{GS}}$ in (2.1). Increasing $|\delta_{\text{GS}}|$ decreases the ratio $M_1/M_2$ leading first to the critical value $M_1|_{\text{l.e.}}/M_2|_{\text{l.e.}} = 1$ around which the bino and wino contents give an interesting relic density. By increasing further $|\delta_{\text{GS}}|$, we can reach a bino-like region, where the relic density is enhanced up to the point where the neutralino mass reaches the lightest stau mass giving a density compatible with the last WMAP results ($\tilde{\tau}$ coannihilation corridor). This is very well illustrated on the regions denoted $0.1 < \Omega < 0.3$ and WMAP in Figs. 3 and 4.

Figures 1 and 2 (left) present the space constraints in a $(t, M_3/2)$ plane for $\delta_{\text{GS}} = 0$, $p = 0$ and $\tan \beta = 35$. With this choice of parameters, the lightest neutralino is always wino-like, independently of the value of $t \equiv (\tilde{T})$. Indeed, the gaugino masses $M_a$ in (2.1) are proportional to their respective beta–function coefficients $b_a$, which yields the standard relation, at the GUT scale:

$$\frac{M_1}{M_2} \sim \frac{b_1}{b_2} = \frac{33}{5} = 6.6$$

implying $M_2 < M_1$ at the electroweak scale. This leaves the lightest neutralino $\chi$ and chargino $\chi^+_1$ in a wino state, which depletes the relic density of the neutralino ($10^{-2}$ at its maximum level value).

If we decrease the value of $t$ from its self–dual point $t = 1$, the function $G_2(t, \tilde{t})$ will reach a point where $M_{i=1,2} = 0$, corresponding to

$$(t + \bar{t})G_2(t, \bar{t}) = 1 \rightarrow \text{Re } t = 0.523$$

giving a null value for gaugino mass terms, and, as a consequence, for the scalar neutralino-proton cross section (proportional to $m_\chi$), as shown in Fig. 4. After this
Figure 1: The spin–independent scalar cross section in the moduli parameter space for \( \delta_{GS} = p = 0 \), in the \((t, M_{3/2})\) plane, for \( \tan \beta = 35 \) (left) and \((p, \tan \beta)\) plane for \( M_{3/2} = 40 \text{ TeV} \) (right). Accelerators and cosmological constraints are given for \( \mu > 0 \). The labels on the black lines correspond to the \( \text{Log}_{10} \) value of \( \sigma_{\chi p}^{\text{vis}} \) (pb). For a description of the experimental constraints applied, see Section 3.4.

Figure 2: Muon fluxes from the sun in the moduli parameter space for \( \delta_{GS} = p = 0 \), in the \((p, M_{3/2})\) plane, for \( \tan \beta = 35 \) (left) and \((t, \tan \beta)\) plane for \( M_{3/2} = 40 \text{ TeV} \) (right). Accelerators and cosmological constraints are given for \( \mu > 0 \). The labels on the black lines correspond to the \( \text{Log}_{10} \) value of the flux \( \mu \text{ km}^{-2} \text{ yr}^{-1} \). For a description of the experimental constraints applied, see Section 3.4.
Figure 3: The spin–independent scalar cross section in the moduli parameter space in the $(\delta_{GS}, M_{3/2})$ plane, for $t = 0.25$, $p = 0$, $\mu > 0$, for $\tan \beta = 5$ (left) and 35 (right). The labels on the black lines correspond to the $\log_{10}$ value of $\sigma_{\chi p}^{\text{scal}}$ (pb). For a description of the experimental constraints applied, see Section 3.3.

Figure 4: Muon fluxes from the sun in the moduli parameter space in the $(\delta_{GS}, M_{3/2})$ plane, for $t = 0.25$, $p = 0$, $\mu > 0$, for $\tan \beta = 5$ (left) and 35 (right). The labels on the black lines correspond to the $\log_{10}$ value of the flux ($\mu$ km$^{-2}$ yr$^{-1}$). For a description of the experimental constraints applied, see Section 3.4.
pole, $M_{i=1,2}$ becomes negative, and its absolute value increases with $t$: this restores a non-vanishing cross-section. This cross-section falls with $t$ becoming small because all the sparticles (even the scalars driven by $M_3$ through the renormalization group equations) become more massive ($G_2 \to \infty$ as $t \to 0$), which increases the virtuality of the exchanged squarks or neutral Higgs in the elastic diffusion process. Generically, for a wino-like neutralino, neutrino flux coming from the Sun are small although annihilations into gauge bosons (which give more energetic neutrinos than other annihilation channels) are favoured. Indeed the capture is small due to suppression of the $Z$ exchange, proportional to the higgsino content of the neutralino. This explains the small values of the fluxes in the $(t, M_{3/2})$ plane. This is illustrated on the left panel of Fig. 2 where muon fluxes (which we will denote by $\text{flux}_\mu$) coming from the Sun are smaller than $10^{-2}$ yr$^{-1}$, well below possible detection.

We show on the right panel of Figs. 1 and 2 the effects of the parameter $p$ and $\tan \beta$ for a fixed value of the gravitino mass ($M_{3/2} = 40$ TeV) keeping $\delta_{GS} = 0$. From (2.3), we see that increasing $p$ up to one decreases scalar soft masses, increasing the scalar cross section $\sigma_{\chi-p}$: for $p=0.95$, the scalars are sufficiently light (2.3) to allow a cross section of the order of $10^{-8}$ pb in some region of the parameter space. For high values of $\tan \beta$, $m_H$ decreases and one has lower values of $\mu$ so higher higgsino content of the neutralino before the “noEWSB” boundary which also enhances the Higgs coupling in $\sigma_{\chi-p}$ (proportional to $z_{\chi,i=1,2}$). Along this boundary, although it is excluded by limits on $b \to s\gamma$ and $m_h$, fluxes coming from the Sun for neutrino indirect detection can be high. For small $\tan \beta$, the $\sigma_{\chi-p}$ enhancement comes from the light Higgs contribution $\chi q \to \chi q$.

Whereas direct detection, being driven by Higgs ($h, H$) exchange has a strong $\tan \beta$ dependance as can be seen on Fig. 1, the indirect detection ($Z$ exchange in the capture) is clearly independent of $\tan \beta$ as shown in Fig. 2.

Figures 3 and 4 show, respectively for direct detection and indirect detection, the $(\delta_{GS}, M_{3/2})$ plane for $t = 0.25$, $p = 0$ and $\tan \beta = 5$ or 35. Because of the positive sign of $b_1$, the effect of $\delta_{GS}$ (which is negative) allows the neutralino to be bino-like and splits $m_\chi$ and $m_\chi^+$: indeed, looking back at (2.1), we can see that, at grand unification,

$$M_1 \bigg|_{\text{GU}} = \frac{a b_1 - \delta_{GS}}{a b_2 - \delta_{GS}} M_2 |_{\text{GU}},$$

with $\alpha$ independent of the gauge group considered. Decreasing from zero $\delta_{GS}$ decreases the wino (versus the bino) content of the neutralino and thus increases the relic density. For $\delta_{GS} \sim -23/5$, we have the relation $M_1 |_{\text{e.e.}} / M_2 |_{\text{e.e.}} \sim 1$ at low energy giving $\Omega h^2 |_{\chi} \sim \Omega h^2 |_{\text{CDM}}^{\text{WMAP}}$. Taking larger $|\delta_{GS}|$ values means driving to the universal case, with $M_1 |_{\text{GUT}} = M_2 |_{\text{GUT}}$ at the high energy scale, and so $M_1 |_{\text{e.e.}} \sim 0.6 M_2 |_{\text{e.e.}}$ at the electroweak scale, giving back a bino-like lightest neutralino. Cosmologically favoured relic densities then come from $\chi \tilde{\tau}$ coannihilation along the $\tilde{\tau}$ LSP boundary line. However in this region, fluxes are very small as capture is disfavoured by squark (though quite light) exchange and very small higgsino content of neutralino in Higgs
Figure 5: The spin–independent scalar cross section as a function of the Green-Schwarz counterterm $\delta_{GS}$, for $t = 0.25$, $p = 0$ and different values of $M_{3/2}$, with $\tan \beta = 5$ (left) and 35 (right).

Figure 6: Muon fluxes from the sun as a function of the Green-Schwarz counterterm $\delta_{GS}$, for $t = 0.25$, $p = 0$ and different values of $M_{3/2}$, with $\tan \beta = 5$ (left) and 35 (right).
Figure 7: The spin–independent scalar cross section as a function of the neutralino mass $M_\chi$ for $t = 0.25$, $p = 0$ and $\tan\beta = 5, 20, 35$. We have scanned the moduli parameter space on $M_{3/2}$ and $\delta_{GS}$ before (left) and after (right) having applied the accelerator and cosmological constraints of Section 3.4.

Decreasing $\delta_{GS}$ first decreases $|M_3|$ down to gluino LSP points and then increases $|M_3|$ (which is now negative), scalar masses (through $M_3$ RGE effect) and $\mu$ (decreasing the neutralino higgsino fraction). In particular the lightest Higgs $h$ (through its stop radiative correction) and the heavy neutral Higgs $H$ masses increase and the coupling $z_{\chi,1(2)}z_{\chi,3(4)}$ decreases so that $\sigma_{\chi^p}^{scal}$ decreases following the higgs mass experimental limit contour (see Fig. 3). For high values of $M_{3/2}$, $\sigma_{\chi^p}^{scal}$ is actually beyond reach of experiment sensitivities ($\lesssim 10^{-11}$ pb). Moreover, for small values of $\delta_{GS}$, $\sigma_{\chi^p}^{scal}$ follows the contour of iso-$\mu$ values (the higgsino content of neutralino) along the $m_{\chi^+_1}$ limit. Higher values of $\tan\beta$ are also more favourable to direct detection because the Higgs $H$ is lighter. For low values of $M_{3/2}$, $\mu$ values are smaller and the higgsino fraction of neutralino drives the phenomenology. Relic density is closed to the WMAP range and both direct ($\sigma_{\chi^p}^{scal} \sim 10^{-8}$ pb in Fig. 3) and neutrino indirect detection ($\text{flux}_\mu^\odot \sim 10^{3-2}$ km$^{-2}$yr$^{-1}$ in Fig. 4) can be interesting but this region is excluded by limits on $m_h$ and $b \to s\gamma$. Muon fluxes coming from the Sun follow the iso-$\mu$ shape given by the limit on $m_{\chi^+_1}$. Small $|\delta_{GS}|$ can also lead to gluino LSP.

We have illustrated explicitly the $\delta_{GS}$ dependance of $\sigma_{\chi^p}^{scal}$ and $\text{flux}_\mu^\odot$ on Figs. 5.
Figure 8: Muon fluxes from the sun as a function of the neutralino mass $M_\chi$ for $t = 0.25$, $p = 0$ and $\tan \beta = 5, 20, 35$. We have scanned the moduli parameter space on $M_{3/2}$ and $\delta_{GS}$ before (left) and after (right) having applied the accelerator and cosmological constraints of Section 3.4.

As previously said, decreasing $\delta_{GS}$ increases scalar masses (through $M_3$ running effect) so $m_H$ which dominantly drive $\sigma_{\chi^0-\mu}$. It also exchanges wino and bino content leading to smaller coupling ($\tan \theta_W$ suppressed) and lower values of $\sigma_{\chi^0-\mu}$. We also see the well known strong influence of $\tan \beta$ on $m_H(m_A)$ leading to higher values for $\tan \beta = 35$.

Concerning indirect detection, bino content decreases the fluxes but small values of $|\delta_{GS}|$ can lead to gluino LSP and high fluxes nearby. This corresponds to the effect of small values of $M_{3[GUT]}$ in soft parameter running, leading to small values of $\mu$ and favoring both direct and indirect detection [6]. But in the models we consider in this section, this happens in region excluded by accelerator constraints because of light chargino contribution in $b \rightarrow s\gamma$ and too light Higgs $h$.

To conclude the discussion of the moduli-dominated SUSY breaking case, we show a wide range of models in the $(m_\chi, \sigma_{\chi^0-\mu})$ plane for direct detection (Fig. 7) and in the $(m_\chi, \text{flux}_\mu)$ plane for neutrino indirect detection (Fig. 8). We then clearly see the complementarity between dark matter and accelerator searches by comparing left panel (no accelerator constraints) and right panel (accelerator constraints applied) of these two figures.
In regions satisfying accelerator constraints, moduli dominated model give rise to either wino neutralino with strongly suppressed relic density and very low detection rates, or bino neutralino (thanks to $\delta_{GS}$) and interesting relic density via $\chi^\tau$ coannihilation or adjusted wino vs bino content but again the direct and indirect detection rates are beyond reach of detection (except sometimes for a larger than one ton size direct detection experiment such as Zeplin Max). This situation is summarized on Figs. 7 and 8.

One interesting feature of this scenario is that $M_{i=1,2}$ is becoming negative for large values of $|\delta_{GS}|$. This negative sign can allow cancellation between the up–quark and down–quark contribution to the scalar cross section $\sigma_{\chi-p}$ together as for negative values of $\mu$ pointed in [60]. This phenomenon is responsible of the "seagull" shape of the plots 5 and 7 depleting in the same way the scalar cross section down to $10^{-13}$ (for $\tan\beta=5$) or $10^{-11}$ (for $\tan\beta=35$).

Moduli models satisfying accelerator constraints and a WMAP favoured relic density are generically not detectable by dark matter searches.

4.2 Dilaton dominated case

The phenomenology of the dilaton dominated scenario is completely different from the moduli domination just discussed. If we look at Eqs (2.7) and (2.5), it is clear that we are in a domain of heavy squarks and sleptons (of the order of the gravitino scale) and light gaugino masses (determined by the dilaton auxiliary component vev).

Indeed, the beta–functions $b_a$ being of the order of $10^{-2}$, the corresponding terms are not competitive by comparison to the $F$ term of the dilaton in (2.5). In fact, looking more closely at (2.8) and (2.9), for not too large values of the universal beta–function coefficient of the first condensing group ($b_+$), we may consider that $F_S$ is a linear function of $b_+$. Increasing $b_+$ means approaching the universal case for the gaugino sector (and the scalar one, driven by $M_3/2$).

Figures 9 and 10 present the $(b_+, M_3/2)$ plane for $\tan\beta = 5$ and $35$ with experimental exclusions, neutralino relic density and respectively iso-$\sigma_{\chi-p}$ and iso-flux$^\odot$ curves. We first discuss the experimental exclusion plots.

For low values of $M_{3/2}$, scalars and gauginos are light so that accelerator constraints are strong. At fixed values of $M_{3/2}$ and decreasing values of $b_+$, we see from (2.7) that gaugino masses decrease. $M_3|_{GUT}$ being smaller, the $m^2_{H_u}$ running slope is softer and yields positive $m^2_{H_u}$ at low energy leading to lower values of $\mu$. This explains the region with too light a chargino (i.e. Higgsino) mass followed by the “no EWSB” region as one goes along a decreasing $b_+$ direction.

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2See Fig. 11 For very low values of $b_+$ around 0.1 (independent on $M_{3/2}$ value), we have a very localised range with gluino LSP due to cancellation in $M_3$ coming from the cancellation between $b_3 (< 0)$ and $b_+$.20
As one increases $M_{3/2}$ at fixed $b_+$, one decreases $M_3$ and the same focussing effect as discussed in the previous paragraph leads to a lower value of the SUSY parameter $\mu$. It is thus not surprising that the “no EWSB” extends further to the right as $M_{3/2}$ increases.

We now turn to a discussion of the direct and indirect (neutrino) rates. For low values of the gravitino mass, these rates are favoured because of the light squark contributions in $\sigma^{\text{scal}}$ and in $\sigma^{\text{spin}}$ (enhancing capture and flux$_{\mu}^\odot$).  

As one increases $M_{3/2}$, one finds a region where the lightest neutralino is a mixed higgsino–gaugino state. It can satisfy WMAP requirements on relic density through $\chi_1 \rightarrow W^+W^-, ZZ$ or $t\bar{t}$ annihilation and $\chi_1^+,\chi_2^0$ coannihilation processes. This region is safe from limits on $m_h$ and $b \rightarrow s\gamma$ thanks to the high $M_{3/2}$ values. Direct detection is favoured because the higgsino component enhances coupling in $\chi q \rightarrow H \chi q$ leading to $\sigma^{\text{scal}} \sim 10^{-7-8}$ pb. In this mixed region, indirect detection is also powerful (Fig. 10). Indeed, the higgsino component allows efficient capture via $\chi q \rightarrow Z \chi q$ and neutralino annihilations into gauge bosons or $t\bar{t}$ lead to energetic neutrinos/muons. Muon fluxes coming from the Sun are enhanced and can reach $10^1$ to $10^3$ km$^{-2}$ yr$^{-1}$.

In both cases (direct or indirect detection), the detection rates increase with tan$\beta$ because the higgsino fraction is higher, direct detection being also enhanced by a lighter $H$ Higgs.

The interesting effect of $b_+$ on both direct ($\sigma^{\text{scal}}$) and indirect (flux$_{\mu}^\odot$) detection is illustrated on Fig. 12. We see that decreasing $b_+$ can give a gain of one order of magnitude for direct detection and up to 2 orders of magnitude for indirect detection which is favoured twice by the higgsino fraction (capture plus annihilation in gauge bosons giving more energetic neutrinos/muons). $\sigma^{\text{scal}}$ being proportional to $z_{\chi_1} z_{\chi_3(4)}$, finally decreases when $b_+(\mu)$ is further decreased (too much higgsino like). The bumps on Fig. 12 (left panel) occur around the $\mu$-$M_1$ crossing (Fig. 11). This requires higher values of $b_+$ for higher values of $M_{3/2}$. The last re-increasing of $\sigma^{\text{scal}}$ (for $M_{3/2} = 1$ TeV) comes from the $\mu$ evolution (Fig. 11) and $z_{\chi_1} z_{\chi_3(4)}$ dependence.

Fluxes$_{\mu}^\odot$ have the same behaviour with $M_{3/2}$ and $b_+$, but being proportional to $z_{\chi_3(4)}^2$, when $b_+$ decreases they still increase or actually follow the (inverse) evolution of the parameter $\mu$ as can be clearly seen by comparing Figs. 11 and 12 (right panel).

Starting from high values of $b_+$, we can also notice the first bump corresponding to $m_\chi = m_t$ because neutrino coming from $\chi\chi \rightarrow t\bar{t}$ annihilation are less energetic than $\chi\chi \rightarrow W^+W^-, ZZ$. This very favourable effect on neutralino dark matter relic density and detection rates consisting in decreasing $\mu$ and scalar masses through $M_{3/2}$ running effect has been pointed out in [8] but is directly related here to the $b_+$ parameter.

To conclude this dilaton case, if we compare model predictions with experimental

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$^3b_+$ acts in the same way since we have seen that it is increasing with $\mu$. But these low $M_{3/2}$ regions are ruled out by accelerator constraints.
Figure 9: **The spin–independent scalar cross section in the dilaton parameter space**, in the \((b_+, M_{3/2})\) plane, for \(\tan \beta = 5\) (left) and \(\tan \beta = 35\) (right). Accelerators and cosmological constraints are given for \(\mu > 0\). The labels in the black lines correspond to the \(\log \) value of \(\sigma_{\chi p}^{\text{scal}}\) (pb). For a description of the experimental constraints applied, see Section 3.4.

Figure 10: **Muon fluxes from the sun in the dilaton parameter space in the \((b_+, M_{3/2})\) plane** for \(\tan \beta = 5\) (left) and 35 (right). Accelerators and cosmological constraints are given for \(\mu > 0\). The labels in the black lines correspond to the \(\log_{10}\) value of the flux \((\mu \text{ km}^{-2} \text{ yr}^{-1})\). For a description of the experimental constraints applied, see Section 3.4.
Figure 11: Low energy parameter $M_1$, $M_2$, $M_3$, $\mu$ evolution with $b_+$, for $\tan \beta = 35$ and $M_{3/2}=1000$, 3000 GeV.

sensitivities for both direct (Fig. left panel) and neutrino indirect detection (Fig. right panel), we see that dilaton models satisfying both accelerator constraints and approximate WMAP relic density give mixed higgsino-gaugino neutralino with generically high detection rates: $\Omega_\chi h^2 \sim (\Omega h^2)_{\text{CDM}}^{\text{WMAP}}$, $\sigma_{\chi-\mu}^{\text{scal}} \sim 10^{-7}$ to $9$ pb and flux$_{\odot}^{\chi} \sim 10^0$ to $3$ km$^{-2}$ yr$^{-1}$.

One of the most interesting features of the dilaton type models is its closed parameter space. Indeed, the beta function of the first condensing gauge group cannot exceed the largest one of the models ($b_{E8} \sim 0.57$). On the other hand, the No EWSB condition forbids high values of $M_{3/2}$, the upper limit depending on $\tan \beta$. Imposing WMAP constraints on the relic density, in this this closed ($M_{3/2}, b_+$) plane, neutralino mass is limited to 500 GeV (resp. 1500 GeV) for $\tan \beta = 35$ (resp. 5). Those limits can be compared to the mSUGRA case where $m_\chi < 500$ GeV for $\tan \beta < 45$ (small values of $M_0$ and $M_{1/2}$) but where this bound can be higher ($\simeq$ TeV) if one considers the hyperbolic branch/focus point region.

Dilaton models could be detected by future dark matter searches especially a ton-size direct detection experiment like Zeplin, or, for points with $m_\chi \leq 500$ GeV a km$^3$ neutrino telescope like Icecube.
5 Conclusion

In the specific context of the class of string models that we have considered, we have seen that the predictions regarding dark matter are strikingly different according to the type of supersymmetry breaking considered. In the case of moduli domination, one does not expect any signal in the forthcoming direct or indirect (neutrino) detection experiments. On the other hand, these experiments should not miss the neutralino signal in the case of dilaton domination. Thus the detection of dark matter or the absence of detection may give key information on the nature of supersymmetry breaking, at least in the context of this given class of models.

Obviously there are connections between these results and detection of the LSP at colliders. Small direct detection cross sections or small indirect detection fluxes are obviously correlated with small production cross sections at colliders. In any case, it is interesting for collider searches to note the characteristics of the regions that satisfy the criterion of satisfactory relic density. For moduli domination, we have identified two regions of interest: one where $m_\chi \sim m_{\tilde{\chi}_1^+} \sim m_{\tilde{\chi}_2^0}$ through the bino and wino content of the LSP (for sufficiently large values of $\tan \beta$), and the other one close to stau LSP region where $m_{\tilde{\tau}_1} \sim m_\chi$. In the case of dilaton domination, the cosmologically interesting region corresponds to a LSP with a proper higgsino content.
Figure 13: The spin–independent scalar cross section (left) and the muon flux from the sun (right) in the dilaton parameter space, as a function of the neutralino mass $M_\chi$ after a scan on $b_+$ and $M_{3/2}$ for $\tan \beta = 5, 20$ and $35$. Accelerators and cosmological constraints (taken for $\mu > 0$) have been included. For a description of the experimental constraints applied, see Section 3.4.

$(m_\chi \sim m_{\chi_1} \sim m_{\chi_2})$. Furthermore the parameter space being closed, this case gives an upper bound on neutralino mass: $m_\chi < 1500$ GeV.

We think that the type of conclusions that we have reached is relevant to a larger class of models. It remains however to make similar analyses on other classes of models to see if dark matter searches, whether they are positive or negative, give interesting indications on the way supersymmetry is broken, and on the way it is transmitted to the observable sector.

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