Supersymmetry with light higgsinos

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Contents

Preface

This Report has been written in partial fulfilment of the requirements for obtaining a "habilitation à diriger des recherches" at Montpellier University. As such, it is not intended to be an unbiased review (if there ever was such a thing) of the field, but its main purpose is to point out my own contributions to it, while embedding them in their proper context. Much of its material has been drawn from Refs. [1–10] where further technical background, additional details and many more literature references can be found; it moreover touches upon the subjects of [11, 12] without reviewing them in detail. Amongst the few previously unpublished results contained in this document, the one which is the most likely to be of interest to some readers is probably the compilation of exclusion bounds on higgsino-like charginos and neutralinos shown in Fig. 3.9.

Over the recent years I have had the privilege to work on these and closely related subjects together with several knowledgeable colleagues, from whom I have learned many profound and fascinating things and whom I would like to thank warmly: Emanuele Bagnaschi, Mikael Berggren, Aoife Bharucha, Sergei Bobrovskyi, Wilfried Buchmüller, Sylvain Fichet, Jan Hajer, Arthur Hebecker, Masahiro Ibe, Rolf Kappl, Sabine Kraml, Suchita Kulkarni, Jenny List, Gudrid Moortgat-Pick, Michael Ratz, Tanja Robens, Krzysztof Rolbiecki, Ronan Ruffault, Kai Schmidt-Hoberg, Hale Sert, Ritesh Singh, Alexander Voigt, Georg Weiglein and Tsutomu Yanagida.

I have also benefited greatly from discussions on the topic of light higgsinos and the μ problem with Howard Baer, Riccardo Barbieri, Nishita Desai, Matthew Dolan, Jack Gunion, Jan Kalinowski, Jörn Kersten, Valery Khoze, Andrea Romanino, Kazuki Sakurai, Veronica Sanz, Pedro Schwaller, Jamie Tattersall, Susanne Westhoff, Robert Ziegler and José Zurita, among many others.

Chapter 1

Introduction

Since the discovery of supersymmetry as a possible enlargement of Poincaré space-time symmetry in the 1970s, and since the potential implications of supersymmetry for elementary particle physics were realized in the early 1980s, supersymmetric extensions of the Standard Model (SM) have been extensively studied. The appeal of supersymmetry for theoretical high-energy physics has many reasons: First, it is the unique nontrivial extension of Poincaré symmetry in four dimensions which allows for particle interactions and masses. This statement does require adopting a generalized notion of symmetry, since the generators of other symmetries governing particle physics form Lie algebras defined by certain commutation relations, while the supersymmetry generators form a Lie superalgebra defined by anticommutation relations. However, in view of the fact that all symmetry generators transform in some representation of the Lorentz algebra, and that fermionic representations are linked to anticommutators by the spin-statistics theorem, this generalization of what defines a symmetry very naturally imposes itself when considering fermionic conserved charges, i.e. supersymmetry charges. History has shown as that, when exploring shorter and shorter distance scales, or higher and higher energies, physics appears to be governed by more and more symmetries. It seems not too bold a proposition that we should encounter supersymmetry at some sufficiently large energy scale.

A second motivation comes from superstring theory. String theory includes a prescription for calculating gravitational scattering amplitudes perturbatively at arbitrary energies in a first-quantized setting. Despite all the difficulties in finding a non-perturbative definition (on which much progress has been made during the last decades), this alone is still a tremendous achievement, and makes string theory an excellent candidate for a perturbative limit of the elusive quantum theory of gravity. To obtain a well-defined ground state, string theory has to be supplemented with supersymmetry; consequently, a low-energy effective field theory deriving from it will also be supersymmetric. Supersymmetry may be partly or entirely broken by the gravitational background used to compactify the resulting ten-dimensional supergravity theories to four dimensions, but unless one demands that (at least) one copy of the four-dimensional supersymmetry algebra is unbroken by compactification, superstring models are prone to instabilities and loss of perturbative control. Hence most recent superstring model building has focussed on compactifications whose effective field theories are $\mathcal{N} = 1$ supersymmetric in four dimensions, with the remaining supersymmetry then to be broken dynamically at a lower scale.

A third argument to consider supersymmetry one of the most promising, if not the single most promising possible extension of the SM is its ability to solve the electroweak hierarchy problem. The hierarchy problem can be phrased as follows: The SM contains a single fundamental mass scale, which is given by the mass-squared parameter m^2 of the Higgs field, or equivalently by the scale of electroweak symmetry breaking. In the presence of any further states with masses $\Lambda \gg m$ and with non-negligible couplings to the SM Higgs boson, direct or loop-induced, the Higgs mass-squared parameter receives quantum corrections of the order Λ^2 which need to precisely cancel amongst each other in order to yield a parametrically smaller electroweak scale. There are good reasons to believe that such states exist at high scales, either as part of field-theoretical ultraviolet embeddings of the Standard Model or as part of an ultimate ultraviolet-complete theory of quantum gravity such as superstring theory. The relative smallness of the SM Higgs mass parameter would then be the result of a delicate cancellation between the a priori uncorrelated fundamental parameters of the underlying UV theory. While finely tuned fundamental parameters would be perfectly consistent from the purely mathematical point of view, they do conflict with common physical experience from other systems, and therefore the need for fine-tuning is usually taken as a strong indication that we do not yet understand the physics of the electroweak scale.

If the world were supersymmetric, with supersymmetry softly broken not too far from the electroweak scale, this conundrum would be solved. Quantum corrections in a supersymmetric theory can arise at most from wave-function renormalization, and are therefore much milder and much less sensitive to ultraviolet dynamics than those in a generic quantum field theory. In particular, any states entering the theory at high energy scales Λ would appear in the form of complete supersymmetric multiplets, with the Λ^2 corrections to the electroweak scale cancelling between their fermionic and bosonic components, not as a result of fine-tuning but as a consequence of supersymmetry enforcing certain relations between their masses and couplings. Indeed, at the heart of this cancellation is that exact supersymmetry requires all components of a single supermultiplet to be degenerate in mass, and that the masses of the fermionic components are protected from additive renormalization by chiral symmetry, thus the masses of the bosonic components such as the Higgs mass are protected as well.

A fourth motivation for supersymmetry is the observation that a significant part of the universe's energy budget is constituted by dark matter. Supersymmetric models of TeV-scale particle physics are typically supplemented by additional discrete symmetries to ensure the absence of baryon- and lepton-number violating interactions, which are unobserved in Nature. The most common example is R-parity or equivalently matter parity. While these symmetries are introduced somewhat ad hoc and with a purely phenomenological motivation, they do lead to consistent quantum field theories. Moreover, they force the lightest supersymmetric particle (LSP — or more precisely the lightest particle transforming under these symmetries) to be stable. This has a significant impact on the characteristic collider signatures. Moreover, if the LSP is electrically and colour-neutral, it could account for the observed dark matter abundance in the universe. Several candidates have been studied, the most prominent being the lightest neutralino (a superposition of the fermionic superpartners of the electroweak gauge bosons and the Higgs fields). A neutralino LSP is, in fact, the prime example of a Weakly Interacting Massive Particle (WIMP) dark matter candidate, which is thermally produced in the early universe, ceases annihilating once the cross-section drops below the Hubble rate, and leaves a thermal relic density which is calculable as a function of its interaction strength and mass. It turns out that the order of magnitude of the observed thermal relic abundance is roughly reproduced for electroweak interactions and electroweak-scale masses. Other possible dark matter candidates in supersymmetric models include the superpartners of the right-handed neutrino, of the graviton (with a different production mechanism where the relic abundance depends on the reheating temperature), or of other possible particles added to the Standard Model such as a QCD axion.

A fifth motivation is specific to a particular set of supersymmetric models which includes the Minimal Supersymmetric Standard Model (MSSM). When embedding the SM particles into supersymmetry representations, and adding a second Higgs supermultiplet as required by anomaly cancellation, the beta functions of the gauge couplings change in a way such that the three gauge couplings unify at a scale $M_{\rm GUT} \approx 10^{16}$ GeV (provided that the additional states have masses not too far above the electroweak scale). This can be seen as supporting the hypothesis of a grand-unified theory (GUT), i.e. a common origin of the three Standard Model gauge group factors from a single simple gauge group spontaneously broken at $M_{\rm GUT}$. The gauge couplings will continue to unify if the MSSM is extended by further complete representations of the grand-unified group, but will fail to do so when adding light supermultiplets in incomplete GUT representations.

The last three of these arguments favour a mass scale M_S of the supersymmetric partners of the SM particles which is close to the scale of electroweak symmetry breaking. This should, in particular, be true if supersymmetry is to solve the electroweak hierarchy problem, since the SM Higgs potential is sensitive to threshold corrections from supersymmetry breaking terms. At the very least, those states with sizeable couplings to the Higgs sector (direct or through loops) should not be much heavier than a TeV at most. Taking the MSSM as an example, the Higgs sector consists of two complex scalar doublets h_u and h_d as well as their supersymmetric partners, the fermionic higgsino fields $\tilde{h}_u$ and $\tilde{h}_d$. Given that the electroweak scale is of the order of 100 GeV, this should also be the natural mass scale for the states constituting the MSSM Higgs sector. For other states coupling strongly to the Higgs and higgsino fields, the supersymmetry-breaking masses should not be larger than the electroweak scale by about a loop factor. In the MSSM this concerns mostly the scalar superpartners of the third-generation quarks, since their Yukawa couplings are largest, but also the fermionic superpartner of the SU(2) gauge bosons. Finally, the supersymmetric partners of the SM gluons affect the electroweak scale through two-loop corrections which can be sizeable due to the large α_s and large group-theoretic factors, and may further be enhanced by large logarithms. The MSSM with these supersymmetric partners within their naturalness limits, and all other states potentially much heavier, is sometimes called "natural supersymmetry". It is already severely constrained by the LHC, specifically by the null results in stop, sbottom and gluino searches.

If one gives up on the MSSM as a solution to the hierarchy problem, but insists on preserving unification and WIMP dark matter, one may also consider a different limit in which the mass scales for scalar and fermionic supersymmetric particles are hierarchically different. In this "split supersymmetry" scenario the higgsinos are at the electroweak scale, together with the "gaugino" fermionic superpartners of the gauge bosons, while the scalar superpartners of quarks and leptons as well as the additional Higgs bosons are parametrically heavier. This does not spoil unification because the heavy states form complete GUT representations, apart from one heavy Higgs doublet whose impact is small.

The topic of the present report is the physics of higgsinos with electroweak-scale masses. These figure prominently in both natural and split supersymmetry, which are two of the reasons for studying them. They can however be motivated independently by several other arguments:

- In the limit where the other electroweak gauginos as well as the third-generation squarks and sleptons are decoupled, the higgsino sector becomes an extremely simple and predictive extension of the SM, with characteristic collider signatures that are typical for exotic long-lived particles. It therefore constitutes a model system in which these signatures can be analysed.
- Likewise, an almost pure neutral higgsino furnishes a simple and predictive example of a WIMP dark matter candidate, if its mass is about 1.1 TeV. While the LHC is not sensitive to such massive higgsinos, almost-pure higgsino dark matter can be probed by future direct detection experiments.
- A higgsino-like neutralino with a sizeable bino component is a good dark matter candidate at a mass which is closer to the electroweak scale, of the order of a few 100 GeV (although direct detection limits are already constraining parts of the parameter space).
- From a purely theoretical point of view, pure higgsinos are motivated by the observation that the higgsino mass parameter μ is the only dimensionful MSSM parameter which does not break supersymmetry. It is constrained from below by direct searches to be larger than about 100 GeV, and from above by naturalness if this is a concern, but in the absence of a mechanism connecting its origins to those of the supersymmetry-breaking parameters, it has no a priori reason to be of the same size. In particular, when ignoring naturalness issues, it might well be parametrically smaller than the typical supersymmetry-breaking masses.
- Related to the latter point, there exist UV-scale models (some of which we will review) predicting that the typical higgsino mass should be about an order of magnitude below the typical supersymmetry-breaking masses. The study of their phenomenology therefore has to include the physics of higgsino-like states, since they might be the only ones kinematically accessible to experiments.

The main part of this Report is divided into two chapters, with Chapter 2 treating topics of more theoretical interest, while Chapter 3 is focussed on phenomenology. In Section 2.1 we will begin with a brief review of the particle content and the parameters characterizing the MSSM, since this is the low-scale model we will be mostly concerned with. We will subsequently review the MSSM Higgs sector and the chargino-neutralino sector, to the extent that is needed for the subsequent discussion, exhibiting in particular some characteristics of higgsino-like neutralinos and charginos. This is followed by a discussion of the μ problem in Section 2.2, i.e. the problem of generating a supersymmetric higgsino mass of the order of (or below) the scale M_S which characterizes supersymmetry breaking within the MSSM supermultiplets. We will review several aspects of this problem in global and local supersymmetry, and also briefly discuss it in the context of gauge-mediated models and of the MSSM. Section 2.3 follows Refs. [1,3] in discussing a solution to the μ problem motivated by higher-dimensional grand unified theories, where the structure of the MSSM Higgs sector at the grand-unified scale is dictated by a shift symmetry. In Section 2.4 we will discuss the generation of a μ parameter in gravity-mediated models with approximate R-symmetries, following Ref. [2]. Section 2.5 deals with the possibility to generate a μ parameter independently of supersymmetry breaking as originally proposed in Ref. [9], while the subject of Section 2.6 is the generation of the μ parameter in the hybrid gauge-gravity mediated models of Refs. [4, 6, 7]. Finally in Section 2.7 we discuss light higgsinos in the context of an extremely split scenario where almost all superpartners have masses close to the Planck scale [10].

We begin chapter 3 by reviewing the production (in Section 3.1) and decays (in Section 3.2) of higgsino-like neutralinos and charginos at colliders. Present-day collider constraints are collected in Section 3.3. We subsequently discuss the discovery potential for light higgsinos in Section 3.4, both at the LHC (drawing upon Ref. [5] among others) and at a future linear collider, recapitulating the ILC analysis of Ref. [8]. Section 3.5 contains a brief discussion of higgsino dark matter.

Finally, conclusions and some possible future directions are given in Chapter 4.

Chapter 2

The μ problem and possible origins of the higgsino mass

2.1 The Minimal Supersymmetric Standard Model

2.1.1 Symmetries, particle content and parameters

In this Section we will review the essential features of the MSSM in so far as they are needed for the subsequent discussion; see e.g. [13–15] for more complete and pedagogical introductory texts. The particle content of the SM is summarized in Table 2.1. The

field	spin-	SU(3)	SU(2)	U(1)
q_I	$\frac{1}{2}$	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$
u^c_I	$\frac{1}{2}$	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$
d^c_I	$\frac{1}{2}$	$\bar{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$
ℓ_I	$\frac{1}{2}$	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$
e^c_I	$\frac{1}{2}$	$\mathbf{1}$	$\mathbf{1}$	1
B_μ	1	$\mathbf{1}$	$\mathbf{1}$	0
W_μ	1	$\mathbf{1}$	$\mathbf{3}$	0
G_μ	1	$\mathbf{8}$	$\mathbf{1}$	0
h	0	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$

Table 2.1: Standard Model field content, indicating spin and gauge quantum numbers. All fermions are written in terms of left-handed Weyl spinors. The index $I = 1, 2, 3$ is a generation index. Right-handed neutrinos are omitted.

MSSM is obtained by embedding each Standard Model state in a $\mathcal{N} = 1$ supermultiplet: Weyl fermions in a chiral supermultiplet and gauge bosons in a vector supermultiplet. The corresponding extra states are called squarks and sleptons (for the scalar superpartners of quarks and leptons) and gauginos (for the Majorana fermion superpartners of the gauge fields). The SM Higgs boson would in principle have the correct gauge quantum numbers to furnish the scalar superpartner of one of the lepton doublets, but this turns out to be difficult to implement while respecting all phenomenological constraints. The SM Higgs field is therefore embedded in a chiral supermultiplet on its own by adding to the SM a Weyl fermion superpartner, the higgsino. A second chiral supermultiplet with the conjugate quantum numbers is needed for anomaly cancellation and to allow for holomorphic Yukawa couplings to down-type quarks and leptons in the superpotential. The resulting particle content is given in Table 2.2. The most general renormalizable

superfield	type	spin-1	spin-$\frac{1}{2}$	spin-0	SU(3)	SU(2)	U(1)
Q_I	chiral	—	q_I	$\tilde{q}_I$	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$
U_I	chiral	—	u^c_I	$\tilde{u}_I$	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$
D_I	chiral	—	d^c_I	$\tilde{d}_I$	$\bar{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$
L_I	chiral	—	ℓ_I	$\tilde{\ell}_I$	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$
E_I	chiral	—	e^c_I	$\tilde{e}_I$	$\mathbf{1}$	$\mathbf{1}$	1
B	vector	B_μ	$\widetilde{B} \equiv \lambda_1$	—	$\mathbf{1}$	$\mathbf{1}$	0
W	vector	W_μ	$\widetilde{W} \equiv \lambda_2$	—	$\mathbf{1}$	$\mathbf{3}$	0
G	vector	G_μ	$\widetilde{G} \equiv \lambda_3$	—	$\mathbf{8}$	$\mathbf{1}$	0
H_u	chiral	—	$\tilde{h}_u$	h_u	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$
H_d	chiral	—	$\tilde{h}_d$	h_d	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$

Table 2.2: Superfield content of the Minimal Supersymmetric Standard Model, listing supermultiplet type, propagating component fields by spin, and gauge quantum numbers.

superpotential allowed by gauge invariance is

$$\begin{aligned} W_{\text{MSSM}} = {} & y^{(u)}_{IJ}\, H_u Q_I U_J + y^{(d)}_{IJ}\, H_d Q_I D_J + y^{(e)}_{IJ}\, H_d L_I E_J + \mu\, H_u H_d \\ & + \lambda_{IJK}\, L_I L_J E_K + \lambda'_{IJK}\, L_I Q_J D_K + \lambda''_{IJK}\, U_I D_J D_K + \mu'_I\, H_u L_I\,. \end{aligned} \tag{2.1}$$

where I, J, K are flavour indices. The first line gives rise to the usual SM Yukawa couplings, additional Yukawa couplings involving the higgsinos, and the corresponding quartic interactions between Higgs, squark and slepton fields. The last term on the first line gives rise to a Dirac mass μ for the higgsino fields and to a corresponding contribution to the Higgs masses. The second line represents baryon-number and lepton-number violating interactions. It is usually discarded by imposing a discrete symmetry such as R-parity (a $\mathbb{Z}_2$ under which all superpartners are odd), although small $\mathbb{Z}_2$-violating effects can be of phenomenological interest. In the following we will adopt R-parity as part of the definition of the MSSM and thus keep only the first line of Eq. (2.1).

Supersymmetry breaking is parameterized by the following soft breaking terms in the Lagrangian:

$$\begin{aligned} \mathcal{L}_{\text{soft}} = & -\frac{1}{2}\sum_{a=1}^{3} M_a\ \text{tr}\,\lambda_a\lambda_a + \text{ h.c.} \\ & - m^2_{QIJ}\,\tilde{q}^\dagger_I\tilde{q}_J - m^2_{UIJ}\,\tilde{u}^\dagger_I\tilde{u}_J - m^2_{DIJ}\,\tilde{d}^\dagger_I\tilde{d}_J - m^2_{LIJ}\,\tilde{\ell}^\dagger_I\tilde{\ell}_J - m^2_{EIJ}\,\tilde{e}^\dagger_I\tilde{e}_J \\ & - a_{UIJ}\,h_u\tilde{q}_I\tilde{u}_J - a_{DIJ}\,h_d\tilde{q}_I\tilde{d}_J - a_{EIJ}\,h_d\tilde{\ell}_I\tilde{e}_J + \text{ h.c.} \\ & - m^2_{H_u}\,|h_u|^2 - m^2_{H_d}\,|h_d|^2 - (B\mu\,h_u h_d + \text{ h.c.})\ . \end{aligned} \tag{2.2}$$

The terms in Eq. (2.2) are the most general set of soft supersymmetry-breaking terms which are allowed by R-parity and which can be generated, at the leading order, by spontaneous supersymmetry breaking in a hidden sector. Their origins become apparent when using a spurionic hidden-sector chiral superfield

$$X = F\theta^2 \tag{2.3}$$

for an effective description of spontaneously broken supersymmetry. If $M > \sqrt{|F|}$ is the mediation scale, i.e. the scale of interactions between the MSSM and the hidden sector, then the leading-order interaction terms are schematically

$$
\begin{aligned}
\mathcal{L}_{\rm spurion} = &\int \mathrm{d}^2\theta\, \frac{X}{M} W_a^\alpha W_{a\alpha} + \text{h.c.} \\
&+ \int \mathrm{d}^2\theta\, \mathrm{d}^2\bar\theta \left(\frac{X^\dagger + X}{M} + \frac{X^\dagger X}{M^2}\right)\left(Q^\dagger Q + U^\dagger U + D^\dagger D + L^\dagger L + E^\dagger E\right) \\
&+ \int \mathrm{d}^2\theta\, \frac{X}{M}\left(H_u Q U + H_d Q D + H_d L E\right) + \text{h.c.} \\
&+ \int \mathrm{d}^2\theta\, \mathrm{d}^2\bar\theta \left(\frac{X^\dagger + X}{M} + \frac{X^\dagger X}{M^2}\right)\left(H_u^\dagger H_u + H_d^\dagger H_d + (H_u H_d + \text{h.c.})\right) .
\end{aligned} \tag{2.4}
$$

Here we have suppressed generation indices and dimensionless coefficients. The component expansion of Eq. (2.4) yields that the first line leads to gaugino masses. The second line gives squark and slepton soft masses, as well as scalar trilinear terms (a-terms) which are proportional to the respective MSSM Yukawa couplings. The third line gives a-terms which are in general independent of the MSSM Yukawa matrices. The fourth line gives rise to Higgs soft masses $m_{H_u^2}$ and $m_{H_d}^2$, as well as to soft Higgs mass mixing $B\mu$. Moreover, because of the $(X^\dagger H_u H_d + \text{h.c.})$ coupling it contributes to the effective higgsino mass parameter μ. All masses induced by these operators are of the order $\frac{F}{M}$, up to dimensionless coupling constants.

To supplement Eq. (2.2), occasionally also the terms

$$
\mathcal{L}'_{\rm soft} = -c_{UIJ}\, h_d^\dagger \tilde q_I \tilde u_J - c_{DIJ}\, h_u^\dagger \tilde q_I \tilde d_J - c_{EIJ}\, h_u^\dagger \tilde\ell_I \tilde e_J - \tilde\mu \tilde h_u \tilde h_d + \text{h.c.} \tag{2.5}
$$

are included in the list of soft terms. They are, however, not generated at the leading order in $\frac{F}{M}$ in the above parametrization of supersymmetry breaking by a spurion field, and thus generically subdominant. For example, a non-supersymmetric contribution $\tilde\mu$ to the higgsino mass can be generated by the operator [16]

$$
\mathcal{L}_{\tilde\mu} = \int \mathrm{d}^2\theta\, \mathrm{d}^2\bar\theta \frac{X^\dagger X}{M^3} D^\alpha (H_u e^{-V}) D_\alpha (e^V H_d) + \text{h.c.} \tag{2.6}
$$

where D_α is the supercovariant derivative and $V = g_2\, W + \frac{g_1}{2} B$ is the combination of $\mathrm{SU}(2) \times \mathrm{U}(1)$ gauge superfields minimally coupled to the Higgs. However, the resulting contribution is of the order $\frac{F^2}{M^3}$ and therefore generically subdominant.

2.1.2 The Higgs potential in the MSSM

The scalar components of the superfields H_u and H_d furnish two scalar Higgs doublets,

$$
h_u = \begin{pmatrix} h_u^+ \\ h_u^0 \end{pmatrix}, \qquad h_d = \begin{pmatrix} h_d^0 \\ h_d^- \end{pmatrix}. \tag{2.7}
$$

By a choice of gauge, h_u^+ can be set to zero; one finds that then also h_d^- vanishes in the vacuum, i.e. electric charge is unbroken.

The μ term in the superpotential Eq. (2.1) gives rise to supersymmetric masses for h_u^0 and h_d^0. In addition, there are soft supersymmetry-breaking mass parameters $m_{H_u}^2$ and

$m_{H_d}^2$ in Eq. (2.2), as well as a soft mass mixing parameter $B\mu$. Finally, there is a quartic Higgs self-interaction from the SU(2) × U(1) D-term potential. Altogether the scalar potential for h_u^0 and h_d^0 is, when setting the squarks and sleptons to zero,

$$\begin{aligned}\mathcal{V} = & \left(|\mu|^2 + m_{H_u}^2\right) |h_u^0|^2 + \left(|\mu|^2 + m_{H_d}^2\right) |h_d^0|^2 + \left(B\mu\, h_u^0 h_d^0 + \text{h.c.}\right) \\ & + \frac{1}{8}\left(g^2 + g'^2\right)\left(|h_u^0|^2 - |h_d^0|^2\right)^2 .\end{aligned} \tag{2.8}$$

Since $B\mu$ is the only parameter which depends on the phases of h_u^0 and h_d^0, one may absorb any phase by a field redefinition into the Higgs fields and choose $B\mu > 0$ without loss of generality. A stable non-trivial vacuum exists provided that (in the tree-level approximation)

$$\left(m_{H_u}^2 + |\mu|^2\right)\left(m_{H_d}^2 + |\mu|^2\right) - (B\mu)^2 < 0\,, \qquad m_{H_u}^2 + m_{H_d}^2 + 2\,|\mu|^2 - 2\,B\mu > 0\,. \tag{2.9}$$

The first of these conditions is that the Higgs mass-squared matrix has a negative eigenvalue, hence the point $h_u^0 = h_d^0 = 0$ is unstable and electroweak symmetry is broken. The second condition serves to obtain a potential which is bounded from below along the direction $h_u^0 = h_d^0$, along which the quartic part in Eq. (2.8) vanishes. If these conditions are satisfied, the Higgs fields acquire vacuum expectation values $\langle h_u^0 \rangle = v_u$ and $\langle h_d^0 \rangle = v_d$. The known electroweak gauge boson masses imply that this breaking of the electroweak symmetry takes place at a scale

$$v^2 = v_u^2 + v_d^2 = (174\,\text{GeV})^2\,. \tag{2.10}$$

As in the Standard Model, the Higgs vacuum expectation value gives masses to three of the electroweak gauge bosons, while the photon remains massless. Defining $\tan\beta = \frac{v_u}{v_d}$, one has the following relations between $\tan\beta$, the observed electroweak scale, and the fundamental parameters in the Higgs potential:

$$\frac{\sin 2\beta}{2} = \frac{1}{\tan\beta + \cot\beta} = \frac{B\mu}{m_{H_u}^2 + m_{H_d}^2 + 2\,|\mu|^2} \tag{2.11}$$

and

$$m_Z^2 = \frac{\left|m_{H_d}^2 - m_{H_u}^2\right|}{|\cos 2\beta|} - m_{H_u}^2 - m_{H_d}^2 - 2\,|\mu|^2\,. \tag{2.12}$$

There are five physical Higgs bosons: Two CP-even neutral scalars h^0 and H^0, one CP-odd pseudoscalar A^0, and a complex charged scalar $H^\pm$. Their mass eigenvalues are

$$\begin{aligned} m_{A^0}^2 &= \frac{2\,B\mu}{\sin 2\beta} = 2|\mu|^2 + m_{H_u}^2 + m_{H_d}^2\,, \\ m_{h^0,H^0}^2 &= \frac{1}{2}\left(m_{A^0}^2 + m_Z^2 \mp \sqrt{\left(m_{A^0}^2 - m_Z^2\right)^2 + 4\,m_Z^2 m_{A^0}^2 \sin^2 2\beta}\right)\,, \\ m_{H^\pm}^2 &= m_{A^0}^2 + m_W^2\,. \end{aligned} \tag{2.13}$$

The implied inequality

$$m_{h^0} < m_Z \tag{2.14}$$

is a result of the tree-level approximation. Quantum corrections can in fact lift the lightest Higgs boson mass from its tree-level bound m_Z to 125 GeV, although they need to be quite sizeable, of the same order as the tree-level value itself. At the one-loop level, the dominant corrections in the decoupling limit $m_{h^0} \ll m_{A^0}$ are given by

$$\delta m_{h^0}^2 = \frac{3}{4\pi^2} \frac{m_t^4}{v^2} \left(\log \frac{M_S^2}{m_t^2} + \frac{X_t^2}{M_S^2} \left(1 - \frac{X_t^2}{12\, M_S^2} \right) \right) . \tag{2.15}$$

Here m_t is the running top mass at the scale m_t, $M_S^2 = m_{\tilde{t}_1} m_{\tilde{t}_2}$ with $m_{\tilde{t}_{1,2}}$ the stop masses, and X_t is the stop mixing parameter, defined at the scale M_S as

$$X_t = a_t / y_t - \mu \cot \beta \, . \tag{2.16}$$

Higher-order contributions to $\delta m_{h^0}^2$ in the general case are known up to the dominant three-loop corrections [17]. Evidently, obtaining large contributions to the Higgs mass requires either heavy stop squarks or large trilinear a-terms in the stop sector (the implications of the latter having been studied in detail in Ref. [11]).

2.1.3 The chargino and neutralino sector

The fermionic superpartners of the electroweak gauge and Higgs bosons can mix after electroweak symmetry breaking. For the electrically neutral states $\widetilde{B}$, $\widetilde{W}^0$, $\tilde{h}_d^0$, $\tilde{h}_u^0$ the mass matrix is

$$M_{\chi^0} = \begin{pmatrix} M_1 & 0 & -c_\beta s_w m_Z & s_\beta s_w m_Z \\ 0 & M_2 & c_\beta c_w m_Z & -s_\beta c_w m_Z \\ -c_\beta s_w m_Z & c_\beta c_w m_Z & 0 & -\mu \\ s_\beta s_w m_Z & -s_\beta c_w m_Z & -\mu & 0 \end{pmatrix} \tag{2.17}$$

where M_1 and M_2 are the gaugino masses, $c_\beta \equiv \cos\beta$ and $s_\beta \equiv \sin\beta$, and c_w and s_w are the cosine and sine of the weak mixing angle respectively. Without loss of generality we can take M_2 to be real and positive, since only relative phases between μ, M_2 and M_1 are physical. We assume that no new sources of CP violation are present in the visible sector, so M_1 and μ are real (but can be of either sign).

The neutralino mass matrix is diagonalized by a complex symmetric matrix $\mathcal{N}$ such that

$$\mathcal{N}^T M_{\chi^0} \mathcal{N} = \mathrm{diag}\,(m_{\chi_1^0}, m_{\chi_2^0}, m_{\chi_3^0}, m_{\chi_4^0}) \tag{2.18}$$

with the positive neutralino masses $m_{\chi_i^0}$ ordered according to their size. Of particular interest for the present Report is the higgsino limit, i.e. the case $|\mu| \ll |M_{1,2}|$, in which the lightest two neutral mass eigenstates χ_1^0 and χ_2^0 are predominantly higgsino-like (see e.g. [18] for an early study). Their masses are

$$m_{\chi_{1,2}^0} = |\mu| \mp \frac{m_Z^2}{2} \left(1 \pm s_{2\beta}\, \mathrm{sign}(\mu)\right) \left(\frac{s_w^2}{M_1} + \frac{c_w^2}{M_2} \right) , \tag{2.19}$$

up to terms suppressed by higher powers of M_1 or M_2. In terms of gauge eigenstates, the neutralinos are given by

$$\begin{aligned} \chi_1^0 &= \frac{1}{\sqrt{2}} \left(\tilde{h}_d^0 - \tilde{h}_u^0 \right) + \frac{s_\beta + c_\beta}{\sqrt{2}} \frac{m_Z}{M_1} s_w \widetilde{B} - \frac{s_\beta + c_\beta}{\sqrt{2}} \frac{m_Z}{M_2} c_w \widetilde{W}^0 , \\ \chi_2^0 &= \frac{1}{\sqrt{2}} \left(\tilde{h}_d^0 + \tilde{h}_u^0 \right) - \frac{|s_\beta - c_\beta|}{\sqrt{2}} \frac{m_Z}{M_1} s_w \widetilde{B} - \frac{|s_\beta - c_\beta|}{\sqrt{2}} \frac{m_Z}{M_2} c_w \widetilde{W}^0 , \end{aligned} \tag{2.20}$$

for $\mu > 0$; similar expressions can be derived for negative μ.

The chargino mass matrix is given by

$$M_{\chi^+} = \begin{pmatrix} M_2 & \sqrt{2}\, m_Z c_w s_\beta \\ \sqrt{2} m_Z c_w c_\beta & \mu \end{pmatrix} . \tag{2.21}$$

It is diagonalized by unitary matrices U and V such that

$$U^* M_{\chi^+} V^\dagger = \text{diag}\,(m_{\chi_1^+}, m_{\chi_2^+}) \tag{2.22}$$

with the chargino masses $m_{\chi_{1,2}^+}$ real and positive. In the higgsino limit, the lighter chargino is predominantly higgsino-like and its mass is given by

$$m_{\chi_1^+} = |\mu| - s_{2\beta}\, \text{sign}(\mu) c_w^2 \frac{m_Z^2}{M_2} . \tag{2.23}$$

For $\mu > 0$ the higgsino-like charginos correspond to the following combination of gauge eigenstates:

$$\begin{aligned} \chi_1^+ &= \tilde{h}_u^+ - \sqrt{2}\, s_\beta \frac{m_W}{M_2} \widetilde{W}^+ , \\ \chi_1^- &= \tilde{h}_d^- - \sqrt{2}\, c_\beta \frac{m_W}{M_2} \widetilde{W}^- . \end{aligned} \tag{2.24}$$

In the deep higgsino limit where the electroweak gauginos decouple completely, the spectrum exhibits the following features:

- The neutralinos χ_1^0 and χ_2^0 are exactly degenerate in mass. This is a consequence of an accidental "higgsino number" U(1) symmetry in this limit which forbids Majorana masses for χ_1^0 and χ_2^0.
- The chargino is degenerate with the neutralinos at the tree level, but electroweak corrections lift this degeneracy, see Fig. 2.1. While the loop corrections to both

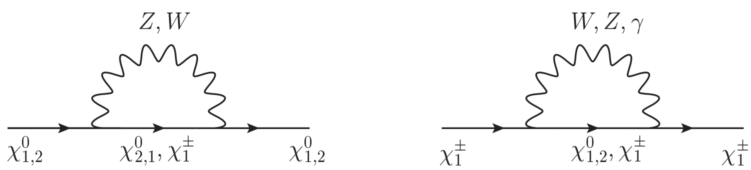

Figure 2.1: One-loop diagrams inducing a mass splitting between neutral and charged higgsinos. (All Feynman diagrams in this Report were generated using the `JaxoDraw` package [19].)

  neutralino and chargino masses are individually divergent, the mass difference is finite. At one loop,

$$\Delta m^{(\text{1-loop})}_{\chi_1^+ - \chi_1^0} = \frac{g^2}{16\pi^2} m_{\chi_1^+} s_w^2\, f\left(\frac{m_Z}{m_{\chi_1^+}}\right) , \tag{2.25}$$

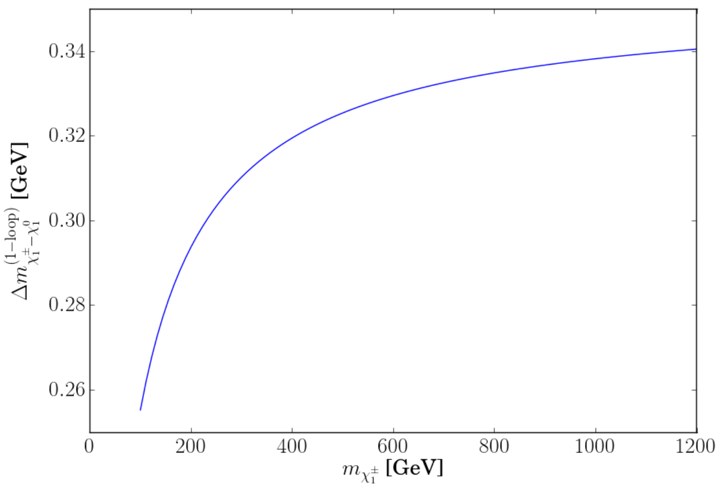

Figure 2.2: One-loop induced mass splitting between χ_1^0 and $\chi_1^\pm$ in the limit of decoupled wino- and bino-like charginos and neutralinos according to Eq. (2.25).

where

$$f(x) = \frac{x}{2}\left(2\,x^3\log x - 2\,x + \sqrt{x^2-4}(x^2+2)\log\frac{x^2 - x\sqrt{x^2-4}-2}{2}\right) . \qquad (2.26)$$

This mass difference is shown in Fig. 2.2 as a function of the chargino mass.

- We finally note that the case of light higgsinos with completely decoupled gauginos and heavy Higgs bosons is fine-tuned, since these states contribute to μ with one-loop threshold corrections. This can be explained by noticing that, once supersymmetry is broken with large gaugino and heavy Higgs boson masses, neither the R-symmetry nor the Peccei-Quinn symmetry which formerly protected μ (see Section 2.2) remain intact [20].

2.2 The μ problem

2.2.1 The μ problem in global supersymmetry

All particle masses in the MSSM have their origins either in electroweak symmetry breaking or in supersymmetry breaking, i.e. they arise either through Yukawa couplings or scalar quartic couplings to the Higgs fields, or from couplings to the hidden sector. The only exception to this rule is the supersymmetric higgsino mass parameter μ (which also contributes to the scalar Higgs masses).

As far as squarks, sleptons and gauginos are concerned, it is quite fitting that unbroken electroweak symmetry permits mass terms for these extra states, since this can explain why all of their masses are above the electroweak scale (as apparent from their non-discovery so far). Note however that the electroweak scale itself is given by the typical mass scale in the Higgs potential, whose supersymmetry-breaking contributions receive quantum corrections from the mass parameters of the other MSSM states. It is therefore rather unnatural to have a significant mass hierarchy between the Higgs sector on the one hand and the squarks, sleptons and gauginos on the other.

By contrast, when considering only the higgsinos, there is no a priori reason why the supersymmetric parameter μ should be of the order of the supersymmetry-breaking soft masses. Instead, by common effective field theory reasoning one should expect it to be either zero (in case it were forbidden by some symmetry) or of the order of the UV cutoff Λ, to be identified with the mediation scale M or eventually with the Planck mass M_P. Neither of these two options is acceptable: A vanishing higgsino masses would imply a chargino mass below the electroweak scale, in contradiction with direct search bounds from the LEP experiment among others. On the other hand, a very large μ parameter would require large cancellations between the parameters entering the Higgs potential, as is evident from Eq. (2.12). The first guise of the so-called "μ problem" [21] is therefore the need for an explanation for the approximate coincidence between μ and the electroweak scale. The usual approach is to forbid a bare μ parameter in the UV embedding of the MSSM by symmetry, and to have an effective μ parameter generated by breaking that symmetry with the same dynamics which induces supersymmetry-breaking masses.

For instance, one may postulate a Peccei-Quinn (PQ) U(1) symmetry under which the Higgs superfield bilinear $H_u H_d$ carries a charge $q \neq 0$. This evidently forbids a bare μ parameter in the superpotential. In the spurion superfield Lagrangian of Eq. (2.4), imposing PQ symmetry and assigning a PQ charge q to the SUSY breaking spurion X will allow for a Kähler term [22]

$$\mathcal{L}_{\rm GM} = \int \mathrm{d}^2\theta\, \mathrm{d}^2\bar{\theta} \left(\frac{X^\dagger}{M} H_u H_d + \text{h.c.} \right) \tag{2.27}$$

which becomes an effective μ term once the F-component of X takes its vacuum expectation value. The resulting higgsino mass is generically of the same order as the soft supersymmetry breaking masses, since PQ symmetry is spontaneously broken at the same scale as supersymmetry. This is known as the Giudice-Masiero mechanism. Note that this symmetry forbids also the gaugino mass term $\int \mathrm{d}^2\theta\, X W^\alpha W_\alpha + \text{h.c.}$, and therefore some model building effort is necessary to generate gaugino masses at the same order.

Alternatively, a U(1) R-symmetry may be used to forbid a bare μ parameter. Assigning opposite R-charges to the Higgs superfields[1] and $R[X] = 2$, the μ term is forbidden while the Giudice-Masiero term Eq. (2.27) is allowed, and an effective μ term of the order of the soft supersymmetry breaking masses F/M is again generated by simultaneous supersymmetry and $\mathrm{U}(1)_R$ breaking.

2.2.2 The μ problem in supergravity

Supergravity is a more appropriate framework to discuss concrete implementations of the above mechanism for generating a phenomenologically acceptable μ parameter in high-scale UV completions of the MSSM. Supergravity is needed, in particular, in models where the mediation of supersymmetry breaking cannot be described in terms of a

[1] We follow the usual conventions of identifying the R-charge of a supermultiplet with that of its lowest component, and of assigning the superpotential an R-charge of $R[W] = 2$.

renormalizable low-energy effective field theory, but involves Planck-scale suppressed interactions between the hidden and visible sectors in an essential way. Since in some of the following Sections will concern just such scenarios, we will now briefly review the origin of the μ term in supergravity following [2] (see also [23] for an earlier review).

We consider a four-dimensional $\mathcal{N} = 1$ supergravity theory containing some chiral superfields Φ^i, comprising hidden-sector chiral superfields X^i as well as the MSSM Higgs superfields H_u and H_d (plus, eventually, MSSM matter as well as hidden- and visible-sector gauge fields, which will however play no role here). The X^i could, for instance, represent the moduli of some superstring compactification, which remain massless at the perturbative level before supersymmetry breaking; all Planck-scale massive states are assumed to be integrated out for the purposes of an effective field theory description. We recall that hidden-sector fields, by definition, do not share their quantum numbers with any of the visible-sector fields.

Barring higher-derivative terms, and ignoring gauge interactions for the moment, the theory is characterized by a real Kähler potential K and a holomorphic superpotential W. The scalar potential is[2]

$$V = e^K \left(D_i W \overline{D}_{\bar{\jmath}} \overline{W} K^{i\bar{\jmath}} - 3|W|^2 \right) \tag{2.28}$$

where we have set $M_P = 1$ momentarily, $D_i = \partial_i + (\partial_i K)$ with $\partial_i = \frac{\partial}{\partial \Phi^i}$, and $K^{i\bar{\jmath}} = (\partial_i \partial_{\bar{\jmath}} K)^{-1}$. We recall that in a supersymmetry breaking vacuum, some of the X^i obtain F-term vacuum expectation values, $F_i \neq 0$, where $F_i = D_i W$. The positive definite part of the scalar potential $F_i \overline{F}_{\bar{\jmath}} K^{i\bar{\jmath}}$ must be cancelled by a nonzero vacuum expectation value of the superpotential $W_0 \neq 0$ in order to yield a Minkowski vacuum. The gravitino mass in this vacuum is given by

$$m_{3/2}^2 = e^K |W_0|^2 \,, \tag{2.29}$$

and therefore serves as an order parameter for supersymmetry breaking. To some extent it can also serve as an order parameter for R-symmetry breaking in R-symmetric models, since the superpotential carries $R[W] = 2$; however, while $W_0 \neq 0$ implies that R-symmetry is spontaneously broken, the converse is not necessarily true.

Expanding K and W up to quadratic order in H_u and H_d, one finds

$$\begin{aligned} K &= K_0 + Y_u \,|H_u|^2 + Y_d \,|H_d|^2 + \;(Z \, H_u H_d + \text{ h.c.}) + \dots \,, \\ W &= W_0 + \hat{\mu} H_u H_d + \dots \end{aligned} \tag{2.30}$$

where K_0, Y_u, Y_d, Z, W_0 and $\hat{\mu}$ are functions of the X_i (K_0 and $Y_{u,d}$ being real, and W_0 and $\hat{\mu}$ being holomorphic).

We define the μ parameter as the supersymmetric mass parameter of Higgs and higgsino fields after canonical field normalization. It receives three separate and a priori independent contributions:

$$\mu = \frac{1}{(Y_u Y_d)^{1/2}} \left(e^{\frac{K}{2}} W_0 Z - \overline{F}^{\bar{\imath}} \frac{\partial Z}{\partial \overline{X}^{\bar{\imath}}} + e^{\frac{K}{2}} \hat{\mu} \right) . \tag{2.31}$$

The first term in Eq. (2.31) is the gravitino mass up to a factor $Z/\sqrt{Y_u Y_d}$ which is generically $\mathcal{O}(1)$. The gravitino mass, when supersymmetry is mediated by Planck-suppressed interactions as we are assuming here, is necessarily of the order as the soft

[2] By a common abuse of notation we use the same symbols for chiral superfields and their lowest components, and regard K and W as functions of either depending on the context.

supersymmetry-breaking masses, and therefore this first contribution to the μ parameter does not entail a μ problem.

The second term results from the supergravity generalization of the Giudice-Masiero term Eq. (2.27) (in the simple spurion Lagrangian Eq. (2.4) there was only one hidden sector field $X = F\theta^2$, and $Z = \overline{X}/M$). It is likewise of the order of the soft supersymmetry-breaking masses.

The last term in Eq. (2.31), finally, is a bare μ parameter from the superpotential, whose presence constitutes the actual μ problem. As in global supersymmetry, there is no reason why $\hat{\mu}$ should be correlated with supersymmetry breaking, so to avoid an unnatural coincidence of scales, one should find a mechanism to forbid or to suppress $\hat{\mu}$.

We finally remark that, in supergravity, the superpotential and Kähler potential are defined only up to Kähler-Weyl transformations $K \to K + f + \bar{f}$, $W \to We^{-f}$ with f holomorphic; physical quantities can depend only on the Kähler-Weyl invariant G-function

$$G = K + \log |W|^2 \,. \tag{2.32}$$

In particular, it is always possible to absorb $\hat{\mu}$ in the Kähler potential by choosing $f = \hat{\mu} H_u H_d / W_0$. However, this of course does not solve the μ problem but rather induces a correspondingly large contribution $\hat{\mu}/W_0$ to the Z function : the resulting μ parameter will still be of the order of $\hat{\mu}$ rather than of the order of the soft supersymmetry breaking masses or, equivalently, the gravitino mass.

2.2.3 The $\mu/B\mu$ problem in gauge-mediated supersymmetry breaking

The μ problem is particularly severe in models of gauge-mediated supersymmetry breaking (see e.g. [24] for a review), or more generally models where supersymmetry breaking is mediated at a lower scale than M_P. In gauge-mediated supersymmetry breaking, the visible sector (the MSSM or some extension thereof, containing all states which are relevant to TeV-scale physics) and hidden the sector (presumably, some supersymmetric field theory with spontaneous and eventually dynamical supersymmetry breaking) are coupled, directly or via intermediate "messenger" states, through the Standard Model gauge couplings [25]. In models of direct gauge mediation, some subgroup of the hidden-sector global symmetry group is identified with the SM gauge group. Messenger models, by contrast, contain a separate messenger sector with states of mass M which are vector-like and charged under the SM gauge interactions, and coupled to the hidden sector by superpotential couplings.

Defining gauge mediation in this way, a purely gauge-mediated model can never generate a nonzero μ term, since the SM gauge interactions do not break the PQ symmetry protecting μ. The hidden sector, therefore, needs to be directly coupled to the visible-sector Higgs fields. In principle it is straightforward to do so, and to generate a μ term which is of the correct order. For example, consider a messenger model with messenger superfields P, $\widetilde{P}$ transforming in the $\mathbf{10} \oplus \overline{\mathbf{10}}$ of SU(5) $\supset$ SU(3) $\times$ SU(2) $\times$ U(1), and a supersymmetry-breaking spurion X with $\langle X \rangle = F\theta^2$. The gauge and superpotential interactions of the messenger fields with the hidden and visible sector,

$$W = W_{\text{hidden}} + W_{\text{MSSM}} + (X + M) P\widetilde{P} + \kappa_u H_u PP + \kappa_d H_d \widetilde{P}\widetilde{P} \,, \tag{2.33}$$

(with projections on $H_u \subset \overline{\mathbf{5}}$ and $H_d \subset \mathbf{5}$ understood) will induce soft terms of the order

$$M_S \sim \frac{g^2}{16\pi^2}\frac{F}{M} \tag{2.34}$$

through messenger and messenger-gaugino loops, but also a μ term of the order

$$\mu = \frac{\kappa_u \kappa_d}{16\pi^2}\frac{F}{M} \tag{2.35}$$

through the superfield graph depicted on the left of Fig. 2.3. Unfortunately, the presence of these superpotential couplings will also lead to a $B\mu$ term through the superfield graph on the right of Fig. 2.3 which is

$$B\mu = \frac{\kappa_u \kappa_d}{16\pi^2}\frac{|F|^2}{M^2} \,. \tag{2.36}$$

Despite being a mass dimension-2 parameter, $B\mu$ is thus generated at the one-loop level

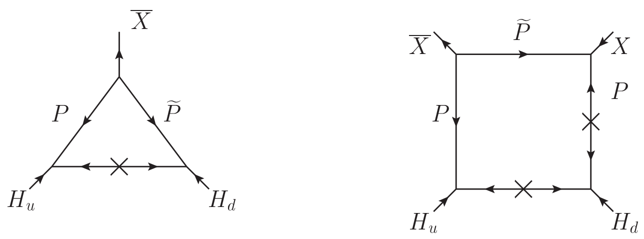

Figure 2.3: Left: Superfield graph inducing an effective μ term at one loop. Right: Effective $B\mu$ term which is also generated at one loop.

and therefore too large to lead to realistic electroweak symmetry breaking. This is a generic problem in gauge-mediated models: If one introduces superpotential couplings between the Higgs and messenger fields, and these lead to a Giudice-Masiero-like term

$$K_{\text{eff}} \supset \frac{1}{16\pi^2}\frac{X^\dagger}{M} H_u H_d + \text{h.c.} \tag{2.37}$$

in the one-loop effective Kähler potential after integrating out the messenger fields and thus to a realistic μ term, then they will typically also give rise to a $B\mu$ parameter which is too large by a loop factor [26].

Several possible solutions to the $\mu/B\mu$ problem in gauge mediation have been discussed in the literature. In Section 2.6 we will discuss one possible approach, viable for models where the mediation scale is lower than but close to the Planck scale, where the interactions inducing μ and $B\mu$ are gravitationally suppressed but appear at the tree level. Thus, μ and $B\mu$ are effectively induced by gravity mediation, whereas the other soft supersymmetry-breaking terms are predominantly gauge mediated.

While well motivated in that particular context by the existence of suitable messenger states in string-theoretic UV completions, a general model of gauge mediation has no a priori relation between the Planck scale and the messenger scale, and the latter could in principle be as low as 100 TeV (in which case the gravity-induced μ term, while still of the order of the gravitino mass, would be tiny since the gravitino mass is much lower than the typical soft mass scale in low-scale mediation models. To cure the μ problem, other approaches are therefore needed. For example, enlarging the messenger sector by at least two singlets, one can generate μ through a higher supercovariant derivative operator rather than through the effective Kähler potential [26] or by an appropriately engineered superpotential [27]. One may also enlarge the visible sector by adding a singlet to the MSSM field content, see Section 2.2.4. Finally, it has been suggested that the $\mu/B\mu$ problem might be resolved by hidden sector dynamics [28]: Even if a too large $B\mu$ is generated at the mediation scale, a strongly coupled and near-conformal hidden sector with the right properties could in principle suppress the corresponding effective operator through renormalization at a lower scale, without affecting the effective μ term, although a calculable example would be difficult to construct.

2.2.4 The μ problem beyond the MSSM

In extensions of the MSSM the μ problem can be addressed without the need for extensive hidden-sector model building. Notably, the next-to-minimal supersymmetric Standard Model (NMSSM, see e.g. [29] for a review) contains a singlet superfield S with scalar and fermionic components s and $\tilde{s}$. Imposing a Z_3 symmetry under which all visible sector particles carry charge 2, a μ term is forbidden but the following superpotential terms are allowed:

$$W_{\rm NMSSM} = W_{\rm Yukawa(MSSM)} + \lambda\, S H_u H_d + \frac{\kappa}{3} S^3 \,. \tag{2.38}$$

The soft supersymmetry-breaking terms permitted by $\mathbb{Z}_3$ and involving the singlet are

$$\mathcal{L}_{\rm soft} = -m_S^2 |s|^2 - a_\lambda\, s h_u h_d + \text{h.c.} - \frac{1}{3} a_\kappa\, s^3 + \text{h.c.} \tag{2.39}$$

Note that $\mathbb{Z}_3$ enforces $B\mu = 0$, while all other MSSM soft terms are allowed. For suitable parameter choices the singlet field s will take a vacuum expectation value, thus spontaneously breaking $\mathbb{Z}_3$ and providing a mass to the higgsinos by the λ term in the superpotential, as well as a supersymmetry-breaking mass mixing term for the scalar Higgs doublets. Since the only mass scale in this model is that of the supersymmetry-breaking soft terms, this is also the scale of the effective higgsino mass. With either gauge- or gravity-mediated supersymmetry breaking, the μ problem can thus be solved in this way. As the singlet carries no Standard Model gauge interactions, using this model for gauge mediation requires additional superpotential couplings between the visible sector and the messengers (see e.g. [24] for details).

The neutralino sector of this model is more complicated than that of the MSSM since there is now also a singlino gauge eigenstate $\tilde{s}$, giving rise to a fifth neutralino mass eigenstate which can mix with the higgsino (as well as the gaugino-like neutralinos). Reviewing the details of the resulting NMSSM phenomenology is beyond the scope of the present report; see [29] and references therein, or [30] for a recent study.

2.3 A μ parameter from a shift symmetry

A particularly interesting class of gravity-mediated models where μ is generated by the Giudice-Masiero mechanism (the $\overline{F}^{\bar{\imath}} \partial_{\bar{\imath}} Z$ term in Eq. (2.31)) are models where, at some high scale M, the Higgs sector is subject to an approximate shift symmetry

$$H_u \to H_u + i\alpha\,, \qquad H_d \to H_d + i\alpha\,, \qquad \alpha \in \mathbb{R}^2\,. \tag{2.40}$$

In four-dimensional field theory, this shift symmetry can be the consequence of an approximate global symmetry G which is spontaneously broken to a subgroup $H \supset \mathrm{SU}(2)_L$ at a scale $f > M$, with the combination of Higgs superfields $\overline{H}_u - H_d$ representing the Goldstone direction. The shift symmetry then corresponds to a non-linear realization of the symmetry G. In five-dimensional field theory, compactified on an interval to four dimensions, the origin of the shift symmetry can be a gauge symmetry G of the five-dimensional bulk, which is broken explicitly to H on the boundaries; the gauge symmetry of the effective four-dimensional theory below the compactification scale M will be then be H. The Goldstone superfield can emerge from the 5d gauge multiplet in the bulk (this is known as supersymmetric gauge-Higgs unification [31]) or from a boundary field which breaks $G \to H' \supset H$ on the corresponding boundary (as in the "holographic GUT" models of [32]). The four-dimensional and five-dimensional mechanisms can be related explicitly, if the bulk geometry is a slice of AdS_5, by the AdS/CFT dictionary [33], since bulk gauge symmetries correspond to global CFT symmetries in the holographic picture.

To further illustrate the possible 5D origins of the model, we depict these two 5D realizations of the Higgs fields and the shift symmetry in Fig. 2.4. In both cases there is a

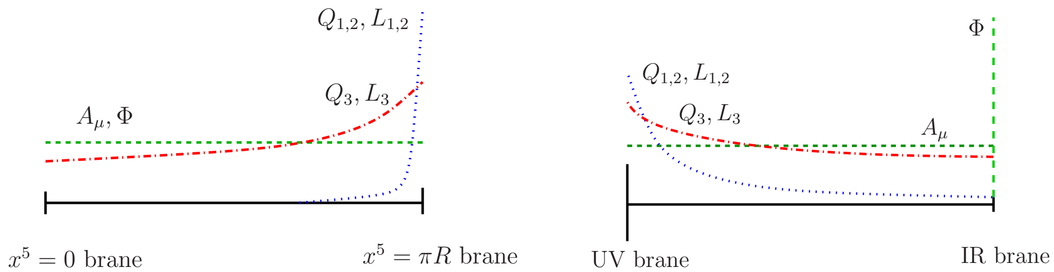

Figure 2.4: Sketch of two 5D models giving rise to a shift-symmetric Higgs sector at the compactification scale, taken from [2]. Left: a flat bulk metric with bulk gauge symmetry G and a bulk gauge field containing the chiral adjoint Φ. The 5D gauge symmetry is broken to the SM by boundary conditions. Right: an AdS bulk metric with bulk gauge symmetry G which is broken by boundary conditions on the UV brane and by the VEV of the brane field Φ on the IR brane.

chiral superfield Φ transforming in the adjoint of the bulk gauge symmetry G, which is assumed to split as

$$\Phi \to H_u \oplus H_d \oplus \ldots \tag{2.41}$$

under the breaking $G \to H \supset \mathrm{SU}(2)_L$. The simplest realizations of this mechanism have $G = \mathrm{SU}(6) \supset \mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$. In the gauge-Higgs unified model depicted on the left of Fig. 2.4 the bulk gauge symmetry is broken to $\mathrm{SU}(5) \times \mathrm{U}(1)$ and to $\mathrm{SU}(4) \times \mathrm{SU}(2)$ on the two branes respectively; their intersection gives the SM gauge group (up to an extra $\mathrm{U}(1)$ factor which is subsequently broken at a high scale in the 4D effective field theory). Third-generation matter fields are only weakly localized to account for their large Yukawa couplings, while first- and second-generation matter fields are effectively localized on one of the boundaries. In the holographic model on the right, the bulk gauge symmetry is likewise SU(6), broken explicitly by boundary conditions to SU(5) $\times$ U(1) on the UV brane, but spontaneously by the VEV of a chiral adjoint superfield Φ on the IR brane. In the first model, the field Φ containing the MSSM Higgs fields (and with $\Phi^\dagger - \Phi$ representing the Goldstone directions) forms part of the 5D gauge supermultiplet, which as a $5D$ $\mathcal{N} = 1$ vector multiplet contains a vector and a chiral adjoint Φ in $4D$ $\mathcal{N} = 1$ language. In the second model Φ is put in by hand as a brane superfield living on the 4D IR boundary.

Let us now turn to the implications of having a shift symmetry Eq. (2.40) at the compactification scale M, which we can identify roughly with the four-dimensional unification scale $M_{\rm GUT}$. At the scale M, the Kähler potential can only depend on the combination $\overline{H}_u + H_d$, i.e.

$$K = Z(X, \overline{X}) \left(H_u + \overline{H}_d\right) \left(\overline{H}_u + H_d\right) + \dots \tag{2.42}$$

where X denotes some hidden-sector fields or compactification moduli. Likewise, the superpotential cannot depend on the Higgs bilinear $H_u H_d$, so there is no μ term from the superpotential. When some of the hidden sector fields develop F-term vacuum expectation values, the Higgs mass matrix resulting from Eq. (2.42)

$$M_H^2 = \begin{pmatrix} m^2 & m^2 \\ m^2 & m^2 \end{pmatrix} \tag{2.43}$$

is therefore degenerate with a flat direction: one has

$$|\mu|^2 + m_{H_u}^2 = |\mu|^2 + m_{H_d}^2 = B\mu\,. \tag{2.44}$$

In other words, the stability conditions Eq. (2.9) are only marginally satisfied. On the other hand, any μ parameter generated in this class of models is automatically of the correct order of magnitude, namely of the order of the soft masses.

The shift symmetry is explicitly broken by the Yukawa couplings in the superpotential, to the extent that they derive from bulk-boundary couplings. The third-generation Yukawa couplings in gauge-Higgs unification derive mostly from 5D gauge couplings, which of course respect the 5D gauge symmetry. However, below the compactification scale the states in the effective theory no longer furnish complete representations of the bulk gauge group, since those parts of the bulk hypermultiplets that do not correspond to MSSM fields have been projected out. Therefore, in the effective theory below the scale M the shift symmetry is no longer respected, and the Kähler potential in Eq. (2.42) is subject to renormalization: the flat direction will eventually be lifted by radiative corrections.

The conditions on the soft terms for obtaining realistic electroweak-scale spectra in such models were studied numerically in [1, 3, 34]. To this end one needs to solve the renormalization group equations between the grand-unified scale, taking into account the boundary conditions Eq. (2.44), and the electroweak scale, matching to the Standard Model observables. A subtlety lies in the sign of the $B\mu$ parameter, which is conventionally defined to be positive at the electroweak scale but may change sign during its renormalization group evolution. Hence the sign in front of the off-diagonal terms in Eq. (2.43), or in front of $B\mu$ in Eq. (2.44), may need to be flipped by a superfield redefinition $H_u \to -H_u$, $U_i \to -U_i$, in order to ensure $B\mu|_{M_S} > 0$.

A particularly simple and predictive soft term pattern results from radion-mediated supersymmetry breaking in five dimensions, where the compactification radius R of the fifth dimension is embedded in a chiral superfield T. The F-term expectation value F^T of this radion superfield can then be used together with the F-term expectation value F^φ of the chiral compensator (the non-dynamical scalar in the 4D gravitational multiplet, whose VEV parametrizes SUSY breaking in the gravitational background) to calculate the soft masses. In the gauge-Higgs unified model, one obtains the Higgs kinetic function

$$Z(T,\overline{T}) = \frac{\pi R}{g_5^2}\left(1 + c'\frac{2R}{T+\overline{T}}\right)\frac{2R}{T+\overline{T}}, \tag{2.45}$$

where g_5 is the 5D gauge coupling and c' is the coefficient of the Chern-Simons term for the 5D gauge field [35]. This leads to the Higgs masses [1, 34, 35]

$$\pm\,\mu = \overline{F}^{\bar\varphi} - \frac{\overline{F}^T}{2R}\frac{1+2c'}{1+c'}\,, \tag{2.46}$$

$$m_{H_u}^2 + |\mu|^2 = m_{H_d}^2 + |\mu|^2 = \pm B\mu = |F^\varphi|^2 - \frac{(F^\varphi\overline{F}^T + \text{h.c.})}{2R}\frac{1+2c'}{1+c'} + \frac{|F^T|^2}{(2R)^2}\frac{2c'^2}{(1+c')^2}\,. \tag{2.47}$$

The 4D gauge-kinetic term is

$$S \supset \frac{\pi R}{g_5^2}\int d^4x \int d^2\theta\left(\frac{T}{R} + c'\right)\text{tr}\, W^\alpha W_\alpha \,+\, \text{h.c.}\,, \tag{2.48}$$

giving for the gaugino masses

$$M_{1/2} = \frac{\overline{F}^T}{2R}\frac{1}{1+c'} \tag{2.49}$$

and the 4d gauge coupling

$$\frac{1}{g_4^2} = \frac{2\pi R}{g_5^2}(1+c')\,. \tag{2.50}$$

The soft masses and trilinear terms for the matter multiplets depend on the corresponding kinetic functions, defined in analogy to those of the Higgs fields Eq. (2.30):

$$K = Y_U(T,\overline{T})\,|U|^2 + Y_Q(T,\overline{T})|Q|^2 + Y_D(T,\overline{T})|D|^2 + Y_E(T,\overline{T})\,|E|^2 + Y_L(T,\overline{T})|L|^2 + \ldots \tag{2.51}$$

These give rise to the soft masses and trilinear couplings according to according to

$$\begin{aligned} m_X^2 &= \,-\,|F^T|^2\frac{\partial^2}{\partial T\partial\overline{T}}\log Y_X(T,\overline{T})\,, \\ A_{U,D} &= F^T\frac{\partial}{\partial T}\log\left(Y_H Y_Q Y_{U,D}\right)\,, \qquad A_E = F^T\frac{\partial}{\partial T}\log\left(Y_H Y_L Y_E\right)\,. \end{aligned} \tag{2.52}$$

Assuming that the first two generations are effectively brane-localized, their soft terms at the compactification scale will vanish. For the third generation, this is no longer a good approximation, so the kinetic functions Y_X should be chosen such as to reproduce the corresponding Yukawa couplings. The result will be model dependent. For the model of Ref. [31] they have been studied in Ref. [1] in detail, numerically solving the renormalization group equations with appropriate boundary conditions at both the high and the low scale, with the result that the higgsino in realistic benchmark points tends to be among the heavier superpartners. This scenario thus constitutes a predictive implementation of the Giudice-Masiero mechanism where the higgsinos are "light" with respect to the cutoff scale, with masses of the order of the gravitino mass, but nevertheless the heaviest of the electroweakinos.

Taking a more model-independent point of view, one may investigate whether it is possible to obtain lighter higgsinos from the boundary conditions of Eq. (2.44) without assuming the above structure for the other soft terms. In Ref. [3] the parameter space of more general models was sampled using a Markov Chain Monte Carlo method, the free parameters being μ, $\tan\beta$, a universal GUT-scale gaugino mass $M_{1/2}$, and the GUT-scale squark and slepton soft masses and scalar trilinear soft terms. For the latter, the study was carried out assuming either (i) universal values m_0 and A_0 for the soft masses and trilinears at the GUT scale, or (ii) that the first two generations do not couple to supersymmetry breaking at leading order, setting their GUT-scale soft terms to zero, while leaving those of the third generation as free parameters.

For the latter class of models it is indeed possible to obtain higgsinos as the lightest supersymmetric particles. The lightest values of μ one can reach are around 500 GeV, and thus still somewhat above the weak scale. Fig. 2.5 shows the Bayesian posterior probability contours for a set of parameter points computed by Markov Chain parameter sampling. The prior used for generating these plots favours points with low Barbieri-

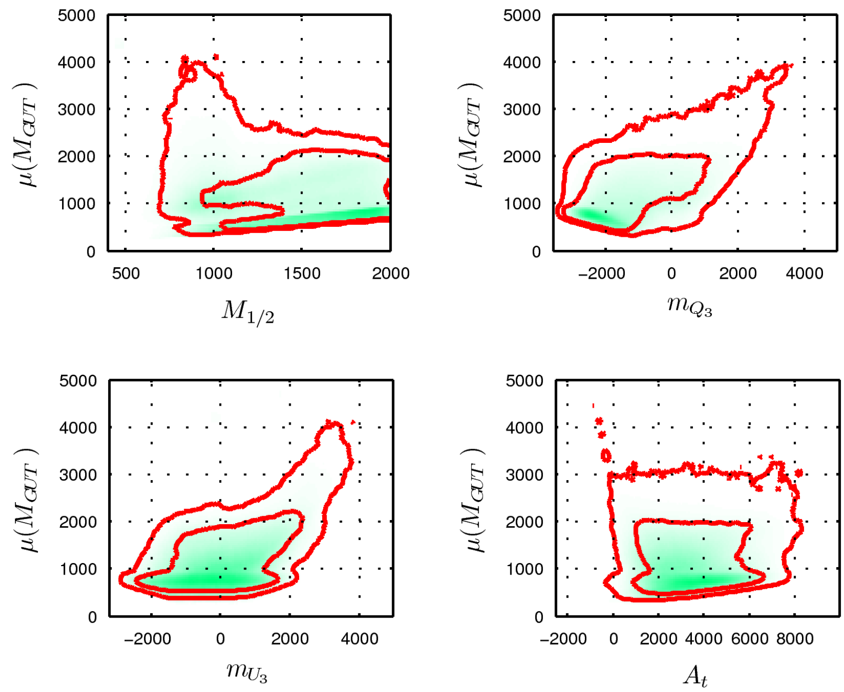

Figure 2.5: Contours of 95% and 68% Bayesian posterior probability, showing correlations between the GUT-scale μ parameter and the most important other supersymmetry-breaking parameters, from [1]. All masses are in units of GeV. The green shading corresponds to the normalized Bayesian likelihood. Note that some parts of this parameter space are by now excluded by LHC limits and by the Higgs mass measurement.

Giudice fine-tuning measure [36], but the results were checked to be reasonably prior-independent. Care should be taken in interpreting these pre-LHC plots, because the observables used obviously do not include any LHC results, neither the lightest Higgs boson mass nor any direct exclusion limits beyond those from LEP. Using up-to-date mass limits, the allowed parameter space would shrink significantly, although some of the points with large gluino and squark masses would certainly survive. We note moreover that this scan required the parameter points to reproduce the observed dark matter relic density, which implies that any higgsino-like lightest neutralino states below a TeV would need to have some sizeable bino component, see Section 3.5. The indicated GUT-scale higgsino mass parameter $\mu = \mu|_{M_{\rm GUT}}$ does not change much in running to the electroweak scale, and therefore provides a good approximation for the physical higgsino mass.

To conclude, models with an approximate shift symmetry in the Higgs sector at the grand unification scale cannot only solve the μ problem, via a variant of the Giudice-Masiero mechanism, but even give rise to MSSM spectra where the lightest supersymmetric states are higgsino-like.

2.4 A μ parameter from an approximate R-symmetry

2.4.1 Suppressing $m_{3/2}$ and $\hat{\mu}$ with $\mathrm{U}(1)_R$

A different approach to solving the μ problem can be taken in models which possess an approximate R-symmetry. An R-symmetry in $\mathcal{N} = 1$ supersymmetry is a U(1) symmetry which does not commute with supersymmetry, and under which, consequently, the different components of a superfield carry different charges. By convention one identifies the R-charge of a supermultiplet with that of its lowest component. For the Lagrangian to be invariant, the superpotential W then needs to carry nonzero R-charge, which as usual we normalize to be $R[W] = 2$.

Continuous global symmetries are believed to be broken in realistic theories of quantum gravity. However, a theory may still be exactly invariant under some discrete symmetry which becomes an approximate continuous symmetry at the level of renormalizable operators, or even at the level of higher-dimensional operators $\sim \Phi^N/M_P^{N-3}$ for some given N. In Ref. [37] an approximate R-symmetry was argued to potentially explain the hierarchy between the Planck scale and the scale of supersymmetry breaking. Notably, if R-symmetry is explicitly broken in W by terms of order at least Φ^N, and if the typical field expectation values $\langle\Phi\rangle$ are mildly suppressed with respect to M_P (as is necessary for an effective field theory description), then $\langle W\rangle \sim \langle\Phi\rangle^N$ can easily be very small in Planck units (which we adopt from now on for the remainder of this Section). As explained in Section 2.2.2, in a Minkowski vacuum the scale of supersymmetry breaking, which is the gravitino mass in the Planck-scale mediated models we are considering here, is given by $\langle W\rangle$. Therefore a small ratio between the gravitino mass and the Planck scale could originate from an approximate R-symmetry (with the understanding that the actual supersymmetry-breaking F- and D-terms will still need to be tuned against $\langle W\rangle$ in order to obtain an approximately Minkowski vacuum in the first place, i.e. to solve the cosmological constant problem).

Following [2] we will now show that an approximate R-symmetry could at the same time be responsible for generating a small higgsino mass parameter μ. As before we consider a model of chiral superfields Φ_i which we separate into visible-sector fields (for concreteness, MSSM matter and Higgs fields) and hidden-sector fields X_i. We assume that none of the hidden-sector fields carries the same quantum numbers as either of the Higgs fields, and that the Higgs bilinear $H_u H_d$ is a singlet under all selection rules. Then the superpotential is

$$W = \sum_a c_a \, M_a(X_i) + H_u H_d \sum_a c'_a \, M_a(X_i) + \ldots \tag{2.53}$$

where the M_a represent some normalized monomials which are also singlets under all selection rules (except for a possible R-symmetry, under which they carry charge 2, and for possible further discrete symmetries) and c_a and c'_a are some coefficients which are generically $\mathcal{O}(1)$. We have omitted any terms depending on higher power of the Higgs bilinear or on MSSM matter fields. After setting all hidden-sector fields to their expectation values, one has, in the notation of Section 2.2.2,

$$\begin{aligned} W_0 &= \sum_a c_a \, M_a(X_i) \,, \\ \hat{\mu} &= \sum_a c'_a \, M_a(X_i) \,. \end{aligned} \tag{2.54}$$

In certain string compactifications, the numerical coefficients c_a and c'_a coincide, to leading order, up to a common $\mathcal{O}(1)$ factor λ, and hence $\hat{\mu} = \lambda W_0$ [38]. This solves the μ problem since μ is guaranteed to be of the order of the gravitino mass. However, in models in which the smallness of the superpotential $W_0 \ll 1$ is due to an approximate R-symmetry, there is no need to impose any such relation between the c_a and the c'_a. Instead, in such models W_0 is suppressed because the expectation values of all monomials M_a are independently small (rather than being suppressed as a result of an approximate cancellation between several large contributions), as we will show in the following. As a consequence, $\hat{\mu}$ is likewise suppressed, and the μ problem is solved.

It is nontrivial that all M_a are individually suppressed in the presence of an approximate R-symmetry. One may prove this statement as follows: Consider a generic superpotential. i.e. one where all terms allowed by the symmetries of the model are present in Eq. (2.53), and all coefficients c_a and c'_a are $\mathcal{O}(1)$ and uncorrelated.[3] In the exact R-symmetric limit, the expectation value of the superpotential in a supersymmetric vacuum is $W_0 = 0$ since

$$2\, W = \sum_i R[\Phi_i] \, \Phi_i \partial_i W \tag{2.55}$$

and $\partial_i W = 0$ in a globally supersymmetric vacuum. (In supergravity one has $0 = D_i W = \partial_i W + W \partial_i K$ as a condition for unbroken supersymmetry, so instead of $W = 0$ one could also have $W \neq 0$ with $\sum_i R[\Phi_i] \Phi_i \partial_i K = -2$, but we will not consider any vacua of the latter kind.) We will show that $W_0 = 0$ because all monomials M_a vanish separately, rather than cancelling against one another (see also [39] for a similar mathematical argument in a different context). To do so, we write

$$W = \sum_a c_a^0 M_a \tag{2.56}$$

[3]In the presence of additional discrete symmetries this may not be the case, but even then the argument still holds when applying it to symmetry invariants instead of individual terms.

where (c_a^0) is a generic set of coefficients. Suppose that the expectation values of the lowest components of chiral superfields are $\langle\Phi_i\rangle$ in some supersymmetric vacuum, hence W has a critical point at these field values:

$$\partial_i W\left(\langle\Phi_1\rangle, \ldots, \langle\Phi_n\rangle\right) = 0\,. \tag{2.57}$$

We consider an open neighbourhood $\mathcal{U}$ around (c_a^0) in the space of coefficients, giving rise to a family of superpotentials

$$W = \sum_a c_a M_a\,, \qquad (c_a) \in \mathcal{U}\,. \tag{2.58}$$

Since (c_a^0) is generic, we can choose $\mathcal{U}$ such that there exists a corresponding family of supersymmetric vacua with expectation values $\langle\Phi_i(c_a)\rangle$ which smoothly depend on the (c_a). Since $W = 0$ in each supersymmetric vacuum, W vanishes identically on $\mathcal{U}$ when regarded as a function of the (c_a) via

$$W(c_a) = W\left(\langle\Phi_1(c_a)\rangle, \ldots, \langle\Phi_n(c_a)\rangle\right)\,. \tag{2.59}$$

Hence

$$0 = \frac{\mathrm{d}W}{\mathrm{d}c_a} = \left(M_a + \sum_i \frac{\partial W}{\partial\Phi_i}\frac{\partial\Phi_i}{\partial c_a}\right)\Bigg|_{\langle\Phi_1(c_a)\rangle,\ldots,\langle\Phi_n(c_a)\rangle} = M_a\left(\langle\Phi_1(c_a)\rangle, \ldots, \langle\Phi_n(c_a)\rangle\right)\,. \tag{2.60}$$

Thus, to the extent that supersymmetry is unbroken and that the R-symmetry is exact rather than approximate (but may be broken or unbroken in the vacuum, see [2] for a detailed discussion), $\hat{\mu} = 0$. If the R-symmetry is merely approximate, with higher-order terms inducing a suppressed $W_0 \sim m_{3/2} \neq 0$, then also $\hat{\mu}$ will be suppressed by the same amount and take a value of order of the gravitino mass.

It is a priori not obvious that the main features of the preceding analysis will persist also in the presence of supersymmetry breaking. That is, so far our analysis applies to only supersymmetric vacua, which in the presence of small R-breaking will be anti-de Sitter with a cosmological constant $\sim -|W_0^2|$. Subsequent uplifting to a supersymmetry-breaking Minkowski (or slightly de Sitter) vacuum should not perturb the vacuum expectation value of W too much if our arguments are to remain valid.

2.4.2 Examples

We illustrate these ideas with two simple examples, taken from [2], for superpotentials which exhibit an approximate R-symmetry and, as a consequence, a suppressed gravitino mass and μ term. The superpotential for the first example will not be generic in the above sense but still serves to illustrate the point. Consider two chiral superfields X and Y with $R[X] = 2$ and $R[Y] = 0$. In the presence of an exact R-symmetry the hidden-sector superpotential is thus

$$W = X\,f(Y) \tag{2.61}$$

with f any holomorphic function. With the non-generic choice

$$f = \lambda Y^2 + \frac{1}{M}Y^3 + \text{ (higher powers)} \tag{2.62}$$

there is a supersymmetric vacuum with non-vanishing $\langle Y \rangle$, approximately located at

$$X = 0\,, \qquad Y \approx -\lambda M\,, \tag{2.63}$$

where we have assumed that λ is somewhat smaller than 1 and that any terms omitted in Eq. (2.62) are suppressed by powers of M with $\mathcal{O}(1)$ coefficients. If the R-symmetry is merely approximate and broken at higher order by a term κY^N, then the minimum will slightly shift but remain a minimum, with the vacuum expectation value W_0 given by

$$W_0 \approx \kappa(-\lambda M)^N\,. \tag{2.64}$$

In the presence of couplings between the hidden sector and the Higgs bilinear, a $\hat{\mu}$ of the order of W_0 will be induced.

A more elaborate example contains three hidden-sector chiral superfields X, Y and Z with $R[X] = 2$ and $R[Y] = -R[Z] = 3$. The most general hidden-sector superpotential is

$$W = X\, f(YZ, X^3\, Z^2)\,, \tag{2.65}$$

or, up to order 10 in the fields,

$$W = X\, P(YZ) + X^4 Z^2\, Q(YZ) + \ldots \tag{2.66}$$

where P and Q are polynomials of degree 4 and 2 respectively. There are supersymmetric vacua at the roots of P with $X = 0$. Assume that P has an isolated zero at some real value $\langle YZ \rangle = v^2 < 1$, so that it is self-consistent to treat higher-order terms as small perturbations. This corresponds to an R-symmetry breaking but supersymmetry-preserving vacuum. The Nelson-Seiberg theorem [40], stating that for a certain class of generic superpotentials spontaneously broken R-symmetry implies spontaneously broken supersymmetry, is not applicable because our superpotential does not fall into that class [2]. The rescaling $Y \to \alpha Y$ and $Z \to \frac{1}{\alpha} Z$ is a complex flat direction corresponding to the Goldstone superfield of spontaneously broken R-symmetry. Despite the fact that R-symmetry is broken, one has $W_0 = 0$ since the expectation value of X vanishes.

We now introduce small R-symmetry breaking terms. More precisely, to justify the absence of R-breaking at lower orders, we take the R-symmetry to be discrete rather than continuous, imposing invariance under $\mathbb{Z}_{16} \subset \mathrm{U}(1)_R$ for concreteness. This allows for the following R-breaking superpotential, up to order 10 in the fields:

$$\begin{aligned} W_{\not R} = {} & \lambda_1\, Y^6 + \lambda_2\, Y^7 Z + \lambda_3\, Y^8 Z^2 + \lambda_4\, Z^{10} \\ & + \kappa_1\, X^9 + \kappa_2\, X^6 Y^2 + \kappa_3\, X^6 Y^3 Z + \kappa_4\, X^3 Y^4 + \kappa_5\, X^3 Y^5 Z + \kappa_6\, X^2 Z^6 + \kappa_7\, X^2 Y Z^7 \end{aligned} \tag{2.67}$$

The terms in the first line will stabilize the flat direction, giving

$$\langle Y \rangle \approx \left(\frac{5\lambda_4}{3\lambda_1}\right)^{1/16} v^{5/4}\,, \qquad \langle Z \rangle \approx \left(\frac{3\lambda_1}{5\lambda_4}\right)^{1/16} v^{3/4}\,. \tag{2.68}$$

Using these values one finds, by solving the F-term equations for X at leading order,

$$\langle X \rangle \approx -\frac{2(3\,\lambda_1)^{5/8}(5\,\lambda_4)^{3/8}}{P'(v^2)} v^{11/2}\,. \tag{2.69}$$

This leads to a W_0 scaling as $v^{15/2}$. Put differently, if $v \approx 0.01$ is moderately small, then W_0 is suppressed with respect to the fundamental scale by 15 orders of magnitude, of the correct order to lead to the hierarchy between the Planck scale and the TeV scale.

Finally, one may also introduce couplings to the Higgs bilinear:

$$\begin{aligned} W_{\hat{\mu}} = \; & H_u H_d \Big(X\,\widehat{P}(YZ) + X^4 Z^2\,\widehat{Q}(YZ) + \hat{\lambda}_1\,Y^6 + \hat{\lambda}_2\,Y^7 Z + \hat{\lambda}_3\,Y^8 Z^2 + \hat{\lambda}_4\,Z^{10} \\ & + \hat{\kappa}_1\,X^9 + \hat{\kappa}_2\,X^6 Y^2 + \hat{\kappa}_3\,X^6 Y^3 Z + \hat{\kappa}_4\,X^3 Y^4 + \hat{\kappa}_5\,X^3 Y^5 Z + \hat{\kappa}_6\,X^2 Z^6 + \hat{\kappa}_7\,X^2 Y Z^7 \Big) \\ & + \ldots \end{aligned} \tag{2.70}$$

While these will not affect the hidden-sector vacuum expectation values at leading order, they will induce a $\hat{\mu}$ parameter which is of the same order of the gravitino mass, likewise scaling as $v^{15/2}$ asymptotically as $v \to 0$.

In summary, in this example model both the gravitino mass and the effective μ parameter are exponentially suppressed as a result of a discrete R-symmetry, which manifests itself as an approximate continuous R-symmetry at low orders. This continuous R-symmetry is also spontaneously broken, but in keeping with the preceding discussion, W_0 and $\hat{\mu}$ still vanish in the absence of explicit breaking since the vacuum is supersymmetric. Indeed, at the level of this model supersymmetry is unbroken and the vacuum is anti-de Sitter, and should be uplifted in a realistic extension by some dynamics which breaks supersymmetry without significantly perturbing W_0 and $\hat{\mu}$.

2.5 A μ parameter from a supersymmetrically broken Peccei-Quinn symmetry

In this Section we outline a third possibility for generating a μ parameter which is nonzero but parametrically below the fundamental scale. The starting point is the observation that μ is forbidden if there is an exact Peccei-Quinn symmetry $\mathrm{U}(1)_{\mathrm{PQ}}$, i.e. a non-$R$ U(1) symmetry under which the Higgs bilinear is charged. Such a symmetry would evidently forbid a bare μ term. If it is spontaneously broken, a μ term of the order of the $\mathrm{U}(1)_{\mathrm{PQ}}$ breaking scale will be generated. It is conceivable that $\mathrm{U}(1)_{\mathrm{PQ}}$ breaking is either linked to supersymmetry breaking, such that the resulting μ term is linked to the soft mass scale (see Section 2.2), or that $\mathrm{U}(1)_{\mathrm{PQ}}$ is spontaneously broken even in the supersymmetric limit, such that the μ term will be independent of the soft mass terms but its magnitude accidentally coincides with them, to within 1–2 orders of magnitude.

Let us elaborate further on the second possibility, following [9], in a globally supersymmetric model. We parametrize supersymmetry breaking as in Eq. (2.4) by a spurion superfield $X = F\theta^2$, and supersymmetric $\mathrm{U}(1)_{\mathrm{PQ}}$ breaking by a spurion Y taking a vacuum expectation value in its lowest component. The Lagrangian contains the terms

$$\mathcal{L} \supset \int \mathrm{d}^2\theta\,\frac{Y^p}{M^{p-1}}\left(1+\frac{X}{M}\right) H_u H_d + \text{h.c.} + \int \mathrm{d}^2\theta\,\mathrm{d}^2\bar{\theta}\frac{X+\overline{X}}{M}\left(H_u^\dagger H_u + H_d^\dagger H_d\right) \tag{2.71}$$

where p depends on the PQ charges of Y and of $H_u H_d$. By contrast, the terms

$$\int \mathrm{d}^2\theta\;(\hat{\mu}+X)\;H_u H_d + \text{h.c.} + \int \mathrm{d}^2\theta\,\mathrm{d}^2\bar{\theta}\,\left(\frac{X^\dagger + X}{M} + \frac{X^\dagger X}{M^2}\right)\left(H_u H_d + H_d^\dagger H_u^\dagger\right) \tag{2.72}$$

are forbidden by PQ symmetry. Thus, the effective μ parameter is given by

$$\mu \sim \frac{Y^p}{M^{p-1}}, \tag{2.73}$$

and can be independent of the typical soft mass scale (provided that the vacuum expectation values F and Y arise from independent dynamics in the underlying theory).

It is instructive to study the implications of this scenario for electroweak symmetry breaking. The $B\mu$ parameter is

$$B\mu \sim \frac{Y^p}{M^p} = \mu \frac{F}{M}. \tag{2.74}$$

In a gravity-mediated model the typical soft terms resulting from Eq. (2.4) are of the order $F/M \equiv M_S$, and thus

$$B\mu \sim \mu M_S. \tag{2.75}$$

We will now assume that $\mu \ll M_S$, which is technically natural, and that $\tan\beta$ is at least moderately large. Note that $|\mu|$ is bounded from below by negative results for chargino searches at LEP, and starting to be constrained by the LHC, which imply $|\mu| \gtrsim 100$ GeV (see Chapter 3). Requiring $\mu \ll M_S$ therefore implies that the superpartner mass scale is much larger than 100 GeV, purely due to phenomenological constraints.

Remarkably, in that case the electroweak scale is parametrically given by $|\mu|$ rather than by M_S, at least in the limiting case of an accidentally small up-type Higgs soft mass $|m^2_{H_u}| \ll M^2_S$. The latter is not a technically natural condition; indeed the very origin of the "little hierarchy problem" are the large radiative corrections which the Higgs soft masses receive from other states with large soft masses, see Section 2.1. Nevertheless, taking into account these corrections it is generically possible to fine-tune the resulting effective $m^2_{H_u}$ parameter small, of the order of $|\mu|^2$ or smaller. Then the effective Higgs mass matrix of Eq. (2.8) becomes parametrically, in terms of the respective dominant contributions,

$$M^2_H \sim \begin{pmatrix} \mu^2 & \mu M_S \\ \mu M_S & M^2_S \end{pmatrix} \tag{2.76}$$

and electroweak symmetry can be broken by the off-diagonal terms $B\mu \sim \mu M_S$ at a scale $\sim |\mu|$. That is to say, even though both diagonal entries of Eq. (2.76) are positive, the first condition of Eq. (2.9) can nevertheless be satisfied at large $B\mu$, indicating that one of the eigenvalues of M^2_H is negative and that the trivial vacuum is therefore unstable. This eigenvalue, corresponding to the lighter Higgs mass-squared parameter and setting the scale of electroweak symmetry breaking, is evidently of the order μ^2. Similar patterns of electroweak symmetry breaking with $m^2_S \gg B\mu \gg \mu^2$ have been investigated in [41, 42] in the context of gauge-mediated supersymmetry breaking.

We emphasize again that this limiting case is fine-tuned since the natural value of $|m^2_{H_u}|$ is of the order of M^2_S, while we have assumed here that it is instead subdominant. However, once a little hierarchy between the electroweak scale and the soft mass scale is forced upon us, as seems to be indicated by the absence of superpartners at the LHC so far, it is no more fine-tuned than the usual scenario where a large and negative $m^2_{H_u}$ almost cancels an equally large and positive μ^2. Moreover, since now the scale of electroweak symmetry breaking is parametrically given by $|\mu|$, the usual coincidence problem of explaining why a supersymmetrically generated μ parameter should be of the order of the electroweak scale is now absent, or more precisely merged into the little hierarchy problem.

The generic implication of this scenario for the spectrum is once more that the higgsinos will be relatively light, with masses which can be of the order of the electroweak scale, whereas the remaining superpartners tend to be much heavier. The detailed particle masses will, of course, depend on its concrete implementation in a more complete model.

2.6 A gravity-mediated μ parameter in models of high-scale gauge mediation

2.6.1 Light higgsinos in a heterotic orbifold model

Certain UV-scale models motivate a solution of the μ problem within a hybrid mediation mechanism, combining the features of gauge-mediated and gravity-mediated supersymmetry breaking. As reviewed in Section 2.2, in gravity mediation a μ term of the order of the gravitino mass is quite naturally induced by e.g. the Giudice-Masiero mechanism, while gauge mediation does not lead to a μ term. Moreover, in gauge-mediated models with messengers the gravitino mass is not of the scale of the soft masses M_S but parametrically given by

$$m_{3/2} \sim 16\pi^2 \frac{M}{M_P} M_S \tag{2.77}$$

where the scale M corresponds to the mass of some intermediate-scale messenger states which carry Standard Model gauge interactions and couple to the hidden sector via the superpotential. It is often assumed that both the scale of hidden-sector dynamical supersymmetry breaking and the messenger scale are low; messenger scales as low as 100 TeV can still give rise to TeV-scale soft masses. In that case a μ parameter of the order of the gravitino mass is clearly in conflict with phenomenology, as it would imply the existence of light charginos with masses down to and below a keV. As discussed in Section 2.2.3, models of low-scale gauge mediation therefore need some other ingredient to generate μ, and to avoid generating a too large $B\mu$ along with it.

However, there is no fundamental reason for the messenger scale to be small. For instance, it was observed in [4] that a certain class of heterotic string compactifications contains a number of vector-like exotic states with just the right properties to act as gauge mediation messenger fields. These states have masses which are naturally of the order of the scale of grand unification $M_{\rm GUT} \approx 10^{16}$ GeV, which happens to be suppressed with respect to the Planck scale by about a perturbative loop factor. Therefore, each messenger multiplet will induce a gauge-mediated contribution to the soft masses which is of the order of a generic gravity-mediated μ term $\mu \sim m_{3/2}$. Because the multiplicities of these messenger fields tend to be large, μ will in fact be somewhat suppressed with respect to the total gauge-mediated soft masses, thus naturally allowing for an electroweak-scale higgsino mass with TeV-scale superpartners.

As an example we will present a model analysed in [4], building on the spectrum of a heterotic superstring compactification on a 6D orbifold constructed in [43]. A particularly interesting limit of the underlying string compactification is that of an anisotropic orbifold grand-unified theory in six dimensions, where four of the compactification radii are Planck-sized and the other two are parametrically larger. The compactification from six to four dimensions at a radius around the GUT scale breaks the bulk gauge symmetry of the extra-dimensional model to the Standard Model. The massless spectrum contains,

superfield	SM representation	multiplicity
D'	$(\mathbf{3},\mathbf{1})_{-1/3}$	4
$\widetilde{D}'$	$(\overline{\mathbf{3}},\mathbf{1})_{1/3}$	4
L'	$(\mathbf{1},\mathbf{2})_{1/2}$	4
$\widetilde{L}'$	$(\mathbf{1},\mathbf{2})_{-1/2}$	4
M	$(\mathbf{1},\mathbf{2})_{0}$	8
S^+	$(\mathbf{1},\mathbf{1})_{1/2}$	16
S^-	$(\mathbf{1},\mathbf{1})_{-1/2}$	16

Table 2.3: The messenger content of a heterotic orbifold model [43] whose gauge-gravity mediated soft mass spectrum was studied in [4].

besides the three generations of SM quarks and leptons and a pair of Higgs doublets, several SM singlets, as well as vector-like exotic fields which become massive when some of the singlet fields acquire vacuum expectation values. The vector-like exotics are listed in Table 2.3. The minimal couplings of these messenger fields to the hidden sector fields are given by the superpotential

$$W = X_1\, D'\widetilde{D}' + X_1\, L'\widetilde{L}' + X_2\, MM + X_2\, S^+S^- \tag{2.78}$$

where X_1 and X_2 are Standard Model singlets with different quantum numbers under hidden sector selection rules. The natural expectation for the vacuum expectation values, just below the four-dimensional unification scale of 10^{16} GeV, is therefore also the mass scale of the messengers. Moreover, assuming that the goldstino is given by a superposition of X_1 and X_2, gauge mediation will induce soft masses for Standard Model gauginos, squarks, sleptons and Higgs bosons. For concreteness we set

$$\langle X_1\rangle = M + \cos\phi\, F\theta^2\,, \qquad \langle X_2\rangle = M + \sin\phi\, F\theta^2 \tag{2.79}$$

which renders the gauge-mediated contributions to the soft terms calculable in terms of F, M and the goldstino mixing angle ϕ:

$$\begin{aligned} M_1 &= \frac{g^2}{16\pi^2}\frac{F}{M}\left(4\cos\phi + \frac{24}{5}\sin\phi\right), \\ M_2 &= \frac{g^2}{16\pi^2}\frac{F}{M}\left(4\cos\phi + 4\sin\phi\right), \\ M_3 &= \frac{g^2}{16\pi^2}\frac{F}{M}\,4\cos\phi\,, \end{aligned} \tag{2.80}$$

$$\begin{aligned} m_Q^2 &= 2\left(\frac{g^2}{16\pi^2}\right)^2\left(\frac{F}{M}\right)^2\left(\frac{287}{50}+\frac{133}{50}\cos 2\phi\right), \\ m_U^2 &= 2\left(\frac{g^2}{16\pi^2}\right)^2\left(\frac{F}{M}\right)^2\left(\frac{96}{25}+\frac{64}{25}\cos 2\phi\right), \\ m_D^2 &= 2\left(\frac{g^2}{16\pi^2}\right)^2\left(\frac{F}{M}\right)^2\left(\frac{74}{25}+\frac{66}{25}\cos 2\phi\right), \\ m_L^2 = m_{H_u}^2 = m_{H_d}^2 &= 2\left(\frac{g^2}{16\pi^2}\right)^2\left(\frac{F}{M}\right)^2\left(\frac{183}{50}-\frac{3}{50}\cos 2\phi\right), \\ m_E^2 &= 2\left(\frac{g^2}{16\pi^2}\right)^2\left(\frac{F}{M}\right)^2\left(\frac{66}{25}-\frac{6}{25}\cos 2\phi\right). \end{aligned} \tag{2.81}$$

In contrast to standard gauge mediation, the gaugino masses are not universal at the messenger scale M, since the messengers do not form complete GUT multiplets. The scalar masses, likewise, do not satisfy the standard relations of minimal gauge mediation (see e.g. [24]) and are somewhat suppressed with respect to their gaugino counterparts, since they scale as $\sim \sqrt{N}$ with the messenger multiplicity while the gaugino masses scale as $\sim N$.

It should be noted that gravity-mediated contributions to the soft masses cannot be neglected if the mediation scale is as high as M_{GUT}. As a consequence, the above gauge-mediated contributions to the soft masses should be supplemented by subdominant but non-negligible gravity-mediated contributions $\sim \frac{F}{M_P}$, hence the details of the low-energy spectrum depend on many more parameters which are uncalculable in our simplified setting. Moreover, as opposed to the gauge-mediated soft masses, the gravity-mediated soft terms have no reason to be flavour universal. This model will therefore suffer from the usual supersymmetric flavour problem of gravity-mediated supersymmetry breaking, so one needs to suppose that some mechanism (such as wave-function localization in the extra dimensions or a discrete symmetry) will suppress flavour-changing neutral currents in the gravity-mediated sector.

This model does allow, on the other hand, to solve the μ problem, providing gravity-mediated μ and $B\mu$ terms

$$\mu \sim \frac{F}{M_P}, \qquad B\mu \sim \frac{F^2}{M_P^2}. \tag{2.82}$$

Choosing F such that the gravitino mass $m_{3/2} = F/\sqrt{3}M_P$ is of the order of the electroweak scale, it is straightforward to find realistic electroweak-scale spectra by appropriately choosing the mixing angle ϕ and some flavour-universal gravity-mediated contribution to the soft masses of the order of $m_{3/2}$. They will be characterized by large $\tan\beta$, since $B\mu$ is small in comparison to $m_{H_d}^2$ (see Eq. (2.11); this feature is robust under renormalization group evolution), squark masses $\gtrsim 1.5$ TeV to provide the necessary radiative corrections to uplift the lightest Higgs mass to 125 GeV, wino and gluino masses around 2 TeV with the bino somewhat lighter, and higgsinos and gravitinos as the only new particles with electroweak-scale masses. The gravitino can be the lightest supersymmetric particle and account for the dark matter relic density. Higgsino-like neutralinos and charginos will be long-lived on collider timescales, thus collider searches for almost-pure higgsinos provide a promising complementary way to probe this model experimentally (see Chapter 3), beyond standard supersymmetry searches for hard jets and missing energy. In standard cosmology, however, higgsino late-time decays will be in conflict with the successful predictions of light element abundances from primordial nucleosynthesis. This problem may be resolved by introducing e.g. small R-parity violating couplings or a mechanism for late-time entropy production.

2.6.2 Light higgsinos from the gaugino focus point

Models of hybrid gauge-gravity mediation with messengers in incomplete GUT multiplets can be motivated by concrete UV completions, as we have discussed in the preceding Section. As previously mentioned, an alternative motivation is the fine-tuning problem, or little hierarchy problem, posed by the MSSM. In a generic supersymmetric model one would expect the soft masses of superpartners to be of the order of the electroweak scale, and the lightest Higgs boson mass to be below or at most slightly above the Z boson mass. However, by now LHC data excludes coloured superparticles with sub-TeV masses under fairly mild assumptions, and we know the Higgs mass to be significantly larger than m_Z — which requires large quantum corrections, originating again from soft terms which are larger than m_Z by at least an order of magnitude. A large soft mass scale, in turn, entails a fine-tuning problem since the electroweak scale is predicted by a combination of soft masses according to Eq. (2.12), with the individual contributions needing to be tuned to approximately cancel.

More precisely, following [6] we can express the predicted Z boson mass $\widehat{m}_Z$ as a function of GUT-scale soft terms as follows: We have, at large $\tan\beta$ and at the soft mass scale $M_S = \sqrt{m_{\tilde{t}_1} m_{\tilde{t}_2}}$,

$$-\frac{\widehat{m}_Z^2}{2} = \left(|\mu|^2 + m_{H_u}^2\right)\big|_{M_S} \tag{2.83}$$

which leads to the semi-analytic relation

$$\begin{aligned}\widehat{m}_Z^2 = \Big(&2.25\,M_3^2 - 0.45\,M_2^2 - 0.01\,M_1^2 + 0.19\,M_2 M_3 + 0.03\,M_1 M_3\\ &+ 0.74\,m_U^2 + 0.65\,m_Q^2 - 0.04\,m_D^2 - 1.32\,m_{H_u}^2 - 0.09\,m_{H_d}^2\\ &+ 0.19\,A_0^2 - 0.40\,A_0 M_3 - 0.11\,A_0 M_2 - 0.02\,A_0 M_1\\ &- 1.42\,|\mu|^2\Big)\Big|_M\,.\end{aligned} \tag{2.84}$$

The coefficients in this expression have been obtained by resumming two-loop renormalization group equations between the scale $M = 10^{16}$ GeV and the scale $M_S = 3.5$ TeV at $\tan\beta = 50$. Terms with coefficients < 0.01 have been omitted. It is remarkable that, among all soft parameters, the electroweak scale is most sensitive to the gluino mass, although the gluino couples to the Higgs sector only through two-loop effects (but with a large coupling α_s and large group-theoretic factors).

Suppose now that what generates these soft terms is a model of messenger gauge mediation, with N_3 messenger pairs in the $(\mathbf{3},\mathbf{1})_0 \oplus (\bar{\mathbf{3}},\mathbf{1})_0$, N_2 messenger pairs in the $(\mathbf{1},\mathbf{2})_0 \oplus (\mathbf{1},\mathbf{2})_0$, and N_1 messenger pairs in the $(\mathbf{1},\mathbf{1})_1 \oplus (\mathbf{1},\mathbf{1})_{-1}$ representations of the SM gauge group, universally coupled in the superpotential to the singlet goldstino $X = M + F\theta^2$. This gives the leading-order soft masses at the messenger scale as

$$M_1 = \frac{6}{5} N_1\, m_{\rm GM}\,, \qquad M_2 = N_2\, m_{\rm GM}\,, \qquad M_3 = N_3\, m_{\rm GM}\,, \tag{2.85}$$

$$\begin{aligned} m_Q^2 &= \left(\frac{8}{3}N_3 + \frac{3}{2}N_2 + \frac{1}{25}N_1\right) m_{\rm GM}^2,\\ m_U^2 &= \left(\frac{8}{3}N_3 + \frac{16}{25}N_1\right) m_{\rm GM}^2\,,\\ m_D^2 &= \left(\frac{8}{3}N_3 + \frac{4}{25}N_1\right) m_{\rm GM}^2\,,\\ m_E^2 &= \left(\frac{36}{25}N_1\right) m_{\rm GM}^2\,,\\ m_{H_{u,d}}^2 = m_L^2 &= \left(\frac{3}{2}N_2 + \frac{9}{25}N_1\right) m_{\rm GM}^2\,,\end{aligned} \tag{2.86}$$

where

$$m_{\rm GM} = \frac{g^2}{16\pi^2}\frac{F}{M} \tag{2.87}$$

assuming a unified gauge coupling g at the messenger scale. Hence, using Eq. (2.84) this scenario predicts

$$\begin{aligned}\widehat{m}_Z = \big(&2.25\, N_3^2 - 0.45\, N_2^2 - 0.01\, N_1^2 + 0.19\, N_2 N_3 + 0.04\, N_1 N_3 \\ &+ 3.80\, N_3 - 1.16\, N_2 - 0.01\, N_1\,\big)\, m_{\rm GM}^2\,.\end{aligned} \tag{2.88}$$

Evidently, certain favourable ratios of messenger indices N_1, N_2 and N_3 predict a partial cancellation between the different contributions to the effective electroweak scale. For example, if $N_1 = N_2$ then for $\frac{N_3}{N_2} =\approx \frac{2}{5}$ the predicted electroweak scale is parametrically smaller than the soft mass scale by an order of magnitude. A more precise calculation yields that this cancellation is realized e.g. for $N_1 = N_2 = 23$ and $N_3 = 9$ [6] (see also [7, 44, 45] for related studies).

This observation is similar in spirit to so-called focus point supersymmetry [46]: In minimal supergravity models with a universal GUT-scale soft mass m_0 and electroweak-scale gaugino masses and μ term, it can be shown that the predicted electroweak scale is largely insensitive to m_0 because the radiative contributions which the $m_{H_u}^2$ parameter receives during its renormalization group evolution tend to cancel (as can be read of from the second line of Eq. (2.84)). In the present example, we find a similar behaviour for suitable choices of messenger indices, although here the scalar masses are not universal and the gaugino contributions are not at all negligible. [4]

Table 2.4 shows an example spectrum of electroweak-scale soft masses. Here we have set $N_1 = N_2 = 23$, $N_3 = 9$ and $m_{\rm GM} = 200$ GeV, which corresponds to $F = (2.5\times10^{10}\ {\rm GeV})^2$ and is of the same order as $m_{3/2} \approx 150$ GeV and as m_Z. To generate this spectrum we have included universal gravity-mediated contributions to the soft masses of the order of $m_{3/2}$ and, likewise, chosen suitable values for $\mu \simeq \sqrt{B_\mu} \simeq m_{3/2}$. One finds the following features:

- large $\tan\beta \approx 50$, as a consequence of the fact that $B\mu$ is generated by gravity mediation but $m_{H_{u,d}}^2$ are dominated by gauge mediation;
- a very heavy gluino, $M_3 \approx 3.8$ TeV;
- very heavy squarks, the lightest of which is the $\tilde{t}_1$ at about 2.5 TeV, while the first-generation squarks are all heavier than 3 TeV;
- the remaining Higgs bosons $H^\pm$, H^0 and A at intermediate masses, at about 1.5 TeV in our benchmark point;
- a right-handed stau as the lightest scalar superparticle;
- and, most notably, three light higgsinos whose mass scale is once more set by the gravity-mediated μ parameter.

Note that, due to the multi-TeV coloured superpartners, this model will not be probed by the LHC. The only kinematically accessible states are the light higgsinos, whose mass splittings are however in a region which is likewise extremely difficult to test at a hadron collider. A linear collider may provide the best possibility to exclude this and similar benchmark points, see Chapter 3.

[4] Focus point-like behaviour involving gauginos has previously been studied in [47].

particle	mass [GeV]
h_0	123(±3)
χ_1^0	205
$\chi_1^\pm$	207
χ_2^0	208
$\tilde{\tau}_1$	1530
H^0	1470
A	1480
$H^\pm$	1480
χ_3^0	2500
χ_4^0	3800
$\chi_2^\pm$	3800
$\tilde{g}$	3800
$\tilde{t}_1$	2500
$\tilde{u}_1$	3700
$\tilde{d}_1$	3400

Table 2.4: Some selected masses in GeV, computed with `SOFTSUSY` [48], for a hybrid gauge-gravity mediated model with messenger indices $(N_1, N_2, N_3) = (23, 23, 9)$, $m_{\rm GM} = 200$ GeV, $\mu = 240$ GeV, and $\tan\beta = 50$.

There are conflicting views about whether or not a focus point-like cancellation between radiative corrections actually constitutes a solution to the little hierarchy problem. While the prediction for the little hierarchy in this model is insensitive to variations of the continuous dimensionful parameter $m_{\rm GM}$, the cancellation does depend on the Standard Model gauge and Yukawa couplings (in particular y_t) taking their actual observed values. The coefficients in Eq. (2.88) are also logarithmically sensitive to the choice of $M_{\rm GUT}$. It seems fair to ask if this focus point model merely relegates the fine-tuning to these parameters [49].

However, in a hypothetical complete and calculable superstring model, there are no free continuous parameters; in the String Landscape, vacua are parametrized by discrete numbers, hence the fine-tuning problem as it is usually phrased will disappear. The focus point model should be regarded as an attempt to realize this, as far as possible, already in effective field theory, which however offers no guidance as to how to fix the gauge and Yukawa couplings to their measured values without introducing additional free continuous parameters. In the absence of a calculable UV completion these parameters have to be taken as they are measured, yet it is still remarkable that with these data one can construct a plausible explanation for the origin of the little hierarchy from particular combinations of messenger indices.

We finally remark that messenger indices favourable for focus point-like behaviour can also be motivated by F-theory GUT models (which generally predict vector-like exotic fields in split multiplets, see [50] for a review), or by models of product-group unification in field theory [7]. In the latter class of models the grand-unified gauge group is $\mathrm{SU}(5) \times \mathrm{U}(3)$, which allows for natural doublet-triplet splitting [51], but also requires the messenger index N_2 to be larger than N_3 when combined with high-scale gauge mediation in order to avoid a Landau pole too close to the GUT scale. Detailed example models are studied in [7]; they again exhibit the typical spectrum of hybrid GUT-scale gauge-gravity mediation at large messenger indices, notably, multi-TeV soft terms for all superpartners except the higgsinos, and higgsino masses around the electroweak scale.

2.7 Light higgsinos in high-scale supersymmetry

Given the null results in supersymmetry searches at the LHC thus far, one may speculate that the supersymmetry breaking scale in Nature might be much higher than the electroweak scale, and perhaps as large as $M_{\rm GUT}$ or M_P. This would of course imply that the electroweak hierarchy is not stabilized by supersymmetry, and require the resulting hierarchy problem to be resolved by some unknown dynamics taking place, presumably, in the UV completion.

When adopting this point of view, it becomes an interesting question if other MSSM states besides the SM Higgs boson could have electroweak-scale masses (despite not being protected by symmetry), and, depending on the light spectrum, how high the supersymmetry scale can be. For instance, split supersymmetry [20, 52] (defined by electroweak-scale gauginos and higgsinos, with all other MSSM states being parametrically heavier) has been studied extensively. It was shown in [53] that the maximal matching scale for split supersymmetry, assuming that the UV completion is the MSSM, is around 10^8 GeV, while higher UV completion scales would result in effective theory couplings at the matching scale which are in conflict with the constraints imposed by supersymmetry. In particular, the Higgs boson quartic coupling becomes negative at high scales, which is in conflict with the positive definite scalar potential required by globally supersymmetric theories. A similar if less severe problem is encountered by high-scale supersymmetry (defined by the effective theory being given by just the Standard Model) with a maximal matching scale around 10^{12} GeV [53]. For a theory with only the Standard Model and a pair of electroweak-scale higgsinos as the light degrees of freedom, the maximal matching scale will lie somewhere in between.[5]

By adding scalar degrees of freedom to the set of "light" particles with electroweak-scale masses, the maximal matching scale can be increased. For instance, when giving electroweak-scale masses to the entire Higgs sector of the MSSM, such that the effective theory above the electroweak scale is a two-Higgs doublet model, it is possible to choose the UV completion scale as high as M_P [55]. In that case, however, vacuum stability becomes an issue: Even though the scalar potential of globally supersymmetric theories is stable by construction, additional vacua besides the electroweak one may develop in an effective two-Higgs doublet model below the SUSY-breaking scale. It is therefore of interest to study in what range of parameters these models predict that the electroweak vacuum is stable, or at least cosmologically long-lived. We will summarize the results of [10] concerning vacuum stability here, as far as they pertain to models which include light higgsinos besides the two Higgs doublets.

Consider the MSSM with supersymmetry spontaneously broken at a very high scale, such that the soft masses are of the order of $M_S \sim 10^{14-17}$ GeV. Assume however that the entries of the Higgs mass matrix are parametrically smaller than M_S, and that likewise

[5] It should be mentioned that the UV scale where the quartic coupling becomes negative is subject to considerable uncertainty, resulting mostly from the uncertainty in the top Yukawa coupling. It has even been (controversially) claimed that the Standard Model can be reconciled with $\lambda = 0$ at M_P within 1.3σ [54].

$|\mu| \ll M_S$. These conditions on the spectrum are technically unnatural, not only for the Higgs bosons as previously discussed, but also for the higgsinos which receive one-loop threshold corrections from the gauginos, see Section 2.1. We will not further justify these assumptions but merely speculate that they may be find an explanation in the underlying theory of quantum gravity UV-completing the MSSM at the Planck scale, whose nature is unknown and may well defy our effective field theory intuition about fine tuning. We further assume, for simplicity, that both the Higgs and higgsino mass parameters will end up being of the order of the electroweak scale. Generalizations to models with intermediate thresholds might nevertheless be interesting to study in their own right.

At energies between the electroweak scale and M_S, the theory is therefore described by a two-Higgs doublet model (of type II, to a good approximation) with an additional Dirac fermion doublet, the higgsino. Neglecting terms which are not generated at the tree level by matching to the MSSM, and which are therefore subdominant, the Higgs potential reads

$$V_{\text{Higgs}} = V_{\text{quadratic}} + \frac{\lambda_1}{2}(h_d^\dagger h_d)^2 + \frac{\lambda_2}{2}(h_u^\dagger h_u)^2 + \lambda_3 (h_d^\dagger h_d)(h_u^\dagger h_u) + \lambda_4 |h_d^\dagger h_u|^2 \,, \tag{2.89}$$

where the tree-level matching conditions to the MSSM read, at the scale M_S,

$$\begin{aligned} \lambda_1 &= \frac{1}{4}\left(g^2 + g'^2\right) \,, \\ \lambda_2 &= \frac{1}{4}\left(g^2 + g'^2\right) \,, \\ \lambda_3 &= \frac{1}{4}\left(g^2 - g'^2\right) \,, \\ \lambda_4 &= -\frac{1}{2}g^2 \,. \end{aligned} \tag{2.90}$$

The tree-level vacuum stability conditions [56]

$$\lambda_1 > 0 \,, \qquad \lambda_2 > 0 \qquad \lambda_3 + (\lambda_1\lambda_2)^{1/2} > 0 \,, \tag{2.91}$$

$$\lambda_3 + \lambda_4 + (\lambda_1\lambda_2)^{1/2} > 0 \tag{2.92}$$

are satisfied by construction at the scale M_S, however they may be violated at intermediate scales in the renormalization-group improved potential. Numerically, one finds that Eqns. (2.91) are always satisfied at all scales, but that Eq. (2.92) may be violated. If this is the case, then the electroweak vacuum is not absolutely stable. The model may still be phenomenologically acceptable if it is metastable with a lifetime exceeding the age of the universe. The fact that at most one of the four conditions is violated allows to derive a criterion for metastability analytically (although relying on numerical solutions of the renormalization group equations): Along the direction ϕ in field space where it decreases most steeply, the quartic potential is

$$V_{\text{eff}}(\phi) = \frac{\lambda}{4}\phi^4 \,, \qquad \lambda = \frac{4\,(\lambda_1\lambda_2)^{1/2}\left(\lambda_3 + \lambda_4 + (\lambda_1\lambda_2)^{1/2}\right)}{\lambda_1 + \lambda_2 + 2\,(\lambda_1\lambda_2)^{1/2}} \,. \tag{2.93}$$

One may now apply the formalism for a single scalar field to calculate the decay probability of the electroweak vacuum along this direction (with respect to which all possible others are exponentially suppressed). According to [57] the probability for tunnelling from one vacuum into another during cosmic time τ is

$$p = \frac{\tau^4}{R^4} e^{-S_B} \tag{2.94}$$

where the "bounce action" S_B is the euclidean action of the field configuration interpolating between the vacua, and R is the length scale characterizing the radius of the nucleating bubble of true vacuum. To ensure the longevity of our vacuum we should have $p \ll 1$ for $\tau = 10^{10}$ yr. In ϕ^4 theory with a negative λ, the bounce action of the instanton for tunnelling between $\phi = 0$ and large ϕ is

$$S_B = \frac{8\pi^2}{3|\lambda|}, \tag{2.95}$$

while R is undetermined at the classical level since the action is scale invariant. In quantum theory, this scale invariance is broken by the β function of the coupling, and p can be calculated as [58]

$$p = \max_R \frac{\tau^4}{R^4} e^{-S_B(R)} \qquad S_B(R) = \frac{8\pi^2}{3\ |\lambda(\frac{1}{R})|} + \Delta S \tag{2.96}$$

where ΔS are one-loop corrections which are negligible for our purposes, and $\lambda(Q)$ denotes the running quartic coupling at the scale Q.

For the potential of the two-Higgs doublet model along the direction of Eq. (2.93), this yields the metastability condition

$$\lambda(Q) \gtrsim -\frac{2.82}{41.1 + \log_{10} \frac{Q}{\text{GeV}}} \tag{2.97}$$

which must be satisfied at all RG scales Q for the decay probability to be $\ll 1$.

Fig. 2.6 shows a section of the parameter space of this model, indicating the (meta)stability constraints. It has been derived without prescribing the lightest Higgs mass, and assuming negligible high-scale threshold corrections, which have little influence on the metastability conditions, using two-loop renormalization group equations generated by SARAH [59] and one-loop (partial two-loop) matching with FlexibleSUSY [60]. Demanding $m_h = 125 \pm 3$ GeV, with a large uncertainty to take into account the neglected thresholds, reduces the viable parameter space to rather large values of the pseudoscalar Higgs mass m_A (which sets the scale also for m_{H^0} and $m_{H^\pm}$) and to very small $\tan\beta$. A region at large $\tan\beta \approx 50$ and small $m_A \approx 150$, not shown on the plot, would be metastable rather than unstable but is ruled out by limits on $B \to s\gamma$ and $H, A \to \tau\tau$.

The vacuum stability constraints have likewise been studied in [10] for a two-Higgs doublet model with the entire gaugino and higgsino sector at the electroweak scale (or equivalently, split supersymmetry with a second light Higgs doublet). In that case, however, one cannot match this model to the Standard Model when taking the SUSY breaking scale close to M_P.

In conclusion, a two-Higgs doublet model with light higgsinos can be a viable effective field theory which allows for a UV completion by the MSSM at very high scales, provided that $\tan\beta$ is small and the extra Higgs bosons are somewhat heavy.

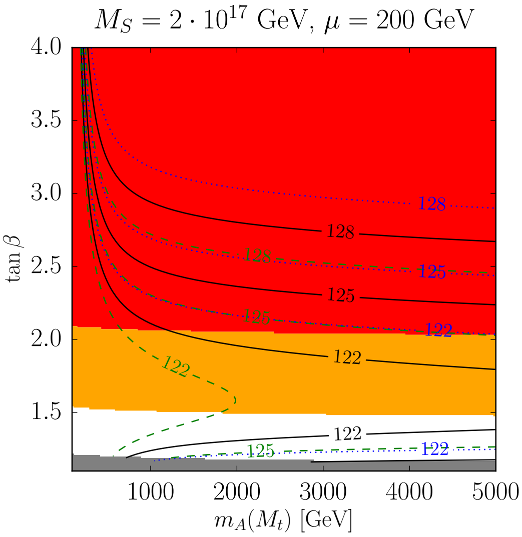

Figure 2.6: Vacuum stability and lightest Higgs mass m_h as a function of the pseudoscalar Higgs mass m_A and $\tan\beta$, using a higgsino mass of $\mu = 200$ GeV and a matching scale of $M = 2 \cdot 10^{17}$ GeV. The green dashed (blue dotted) line represents the result for a top mass at $1\sigma = 0.76$ GeV above (below) the central value of $m_t = 173.34$ GeV. Red regions are excluded by vacuum stability, in the orange regions the electroweak vacuum is metastable.

Chapter 3

Phenomenology of light higgsinos

3.1 Higgsinos production at colliders

The dominant mechanism for higgsino production at collider experiments depends to a large extent on the details of the particle spectrum. To be specific, let us assume that the particle content relevant for TeV-scale physics is that of the MSSM or some subset thereof, that R-parity is conserved, and that the lightest supersymmetric particle is a higgsino-like neutralino (apart from possible light and very weakly interacting states such as a light gravitino or axino). With these assumptions, superpartners are pair-produced at colliders, and any decay chains will end up in a higgsino and Standard Model particles.

At hadron colliders, higgsino-like neutralino and chargino pairs will thus appear in cascade decays of coloured superpartners, most notably squarks, which benefit from a strong production cross-section. This, however, obviously relies on the squarks being light enough to be kinematically accessible. Some example processes at leading order are sketched in Figs. 3.1 – 3.3. Secondly, at both hadron and electron-positron colliders, higgsinos can be pair-produced in the electroweak interactions, most importantly by the Drell-Yan process sketched in Fig. 3.4. At an e^+e^- machine, t-channel and u-channel selectron or sneutrino exchange could also play a role if these states are relatively light and if there is significant mixing between the higgsino and the electroweak gauginos (for pure higgsinos, these processes are suppressed by small Yukawa couplings).

For higgsino production from cascade decays of heavier coloured particles, which one among the possible processes will eventually dominate depends on various aspects of the spectrum:

- If first- or second-generation squarks and gluinos are within kinematic reach, they can be abundantly produced in quark fusion and quark-gluon fusion, and their cascade decays will eventually terminate in a pair of higgsino-like LSPs, producing the standard supersymmetry signature of hard jets, potentially leptons, and missing energy. Fig. 3.1 shows an example process out of many possible ones. Since there is very little to these events that is specific to the case of a higgsino LSP, we will not discuss them here (for details, see e.g. [14, 15] and references therein). We do remark, however, that the couplings between first- and second-generation squarks and higgsinos are suppressed by small couplings and mixing angles, which may cause the squarks to preferentially decay into heavier wino- or bino-like charginos and neutralinos instead of directly into higgsinos, thus softening the decay products while increasing their multiplicities.

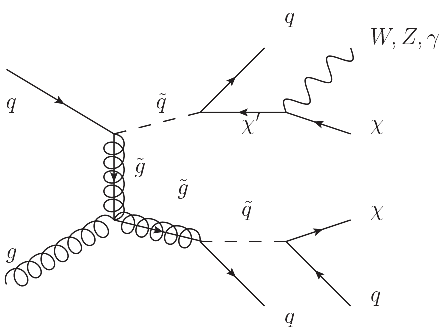

Figure 3.1: An example for strong production of a pair of supersymmetric particles (in this case, a squark and a gluino) which subsequently decay through a cascade, with a final state containing multiple quarks, possibly leptons and a higgsino pair.

- In case that first- and second-generation squarks are too heavy to be accessible, the third-generation squarks may yet be within reach. Gluino-assisted stop and sbottom pair production can give rise to events with multiple b-jets and missing energy, see Fig. 3.2. If the gluino is too heavy, third-generation squarks can still be (less efficiently) pair-produced in gluon fusion or through an s-channel gluon, as shown in Fig. 3.3. In that case the decay chains are very short, with stops and sbottoms immediately decaying into higgsino and a top or bottom quark, and could provide useful information in identifying the higgsino nature of the LSP, see Section 3.4.

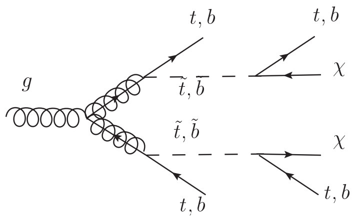

Figure 3.2: Gluino-assisted stop or sbottom pair production, with the stops and sbottoms decaying into third-generation quarks and higgsino-like neutralinos and charginos.

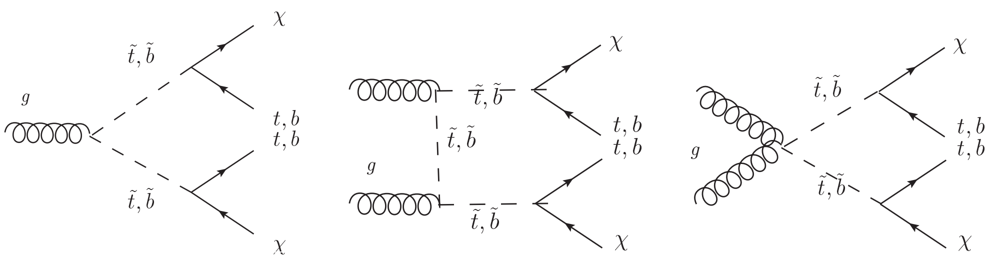

Figure 3.3: Direct stop or sbottom pair production with correspondingly short decay chains.

- In case that all squarks are too heavy to be produced, the only accessible coloured superpartner is the gluino, and it can only decay into three-body final states [61] or radiatively into gluon-neutralino [62]. Therefore, it may live long enough to hadronize. The resulting phenomenology of R-hadrons is again not specific to having a higgsino-like LSP, so we will not review it here.

Let us now turn to direct higgsino pair production through the Drell-Yan process, see Fig. 3.4. At hadron colliders the final state can be $\chi_1^0\chi_2^0$, $\chi_1^+\chi_1^-$, $\chi_1^0\chi_1^\pm$ or $\chi_2^0\chi_1^\pm$, while at e^+e^- colliders the gauge boson in the s-channel is necessarily a photon or a Z, hence the final state is either a neutralino or a chargino pair. Note that the Z coupling to $\chi_1^0\chi_1^0$ or to $\chi_2^0\chi_2^0$ is suppressed by the higgsino-wino mixing angle, and absent in the pure higgsino limit.

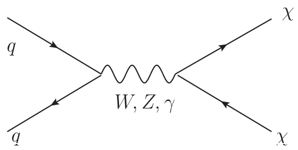

Figure 3.4: Higgsino production via the Drell-Yan process.

The cross-section for the various final states at the 14 TeV LHC are shown in Fig. 3.5.

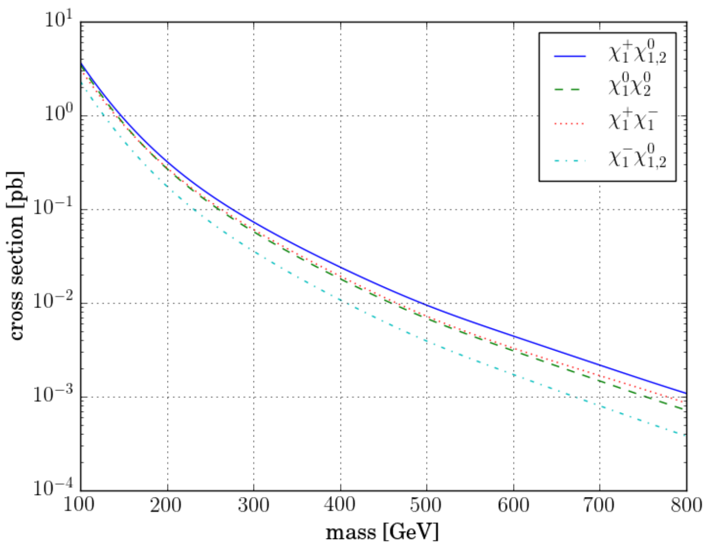

Figure 3.5: The NLO cross-section for higgsino pair production through the Drell-Yan process at the 14 TeV LHC, as a function of the higgsino mass. Generated with `PROSPINO 2.1` [63].

It is often extremely challenging to extract any useful information from direct higgsino pair production, especially in the almost-pure case of negligible mixing with other electroweakinos, since the higgsino-like charginos and neutralinos are nearly degenerate in mass. Then χ_2^0 and $\chi_1^\pm$ will typically decay into χ_1^0 and additional particles which are too soft to observe, and one is left with an effectively back-to-back χ_1^0 pair leaving the detector without any trace. In order to produce at least a missing transverse energy signature, it is convenient to study also direct higgsino pair production with an additional gluon or photon from initial- or final-state radiation. This will decrease the cross-section by a factor of approximately α_s or $\alpha_{\rm EM}$ respectively (for a jet or photon energy of the order of the higgsino mass), but provide a hard object for the neutralinos to recoil from. An initial-state jet is also frequently used in the case of somewhat milder mass degeneracy to trigger on and for better background discrimination.

3.2 Higgsino decays

In the MSSM (or the split MSSM, with a subset of the supersymmetric particles decoupled) with preserved R-parity, the decay modes of higgsino-like neutralinos χ_2^0 and charginos $\chi_1^\pm$ depend mainly on the degree of purity of the higgsino, i.e. the higgsino-wino and higgsino-bino mixing angles. This is because the mass differences $\Delta m_{\chi_2^0-\chi_1^0}$ and $\Delta m_{\chi_1^\pm-\chi_1^0}$ are induced partly by mixing according to Eqs. (2.19) and (2.23). In addition, the chargino-neutralino mass splitting receives radiative corrections after electroweak symmetry breaking which are given by Eq. (2.25) at one loop.

Neutralino decays $\chi_2^0 \to \chi_1^0 X$ for mass splittings above about 10 GeV proceed via virtual Z exchange, producing quark or lepton pairs. For smaller mass splittings $\Delta m_{\chi_2^0-\chi_1^0}$ roughly between 1 GeV and 10 GeV, the radiative decay $\chi_2^0 \to \chi_1^0\gamma$ [65] gains in relative importance as shown in Fig. 3.6. Moreover, for mass splittings smaller than about 1–2 GeV, decays into hadronic final states are no longer accurately described by a virtual Z decaying into quarks which subsequently hadronize. Hence, the region of small Δm in Fig. 3.6 is affected by a large theory uncertainty and should be taken with a grain of salt.

Similar observations can be made about the chargino decays $\chi^\pm \to \chi_1^0 X$, although there is obviously no photon channel available. While for mass splittings around 10 GeV the decays are well modelled as $\chi_1^\pm \to \chi_1^0 W^{*\pm}(\to \ell\bar{\nu}_\ell, qq')$, at mass splittings smaller than about 1–2 GeV one should instead couple to the hadronic current directly. To this end one may employ the formalism of [66], originally developed to model semileptonic tau decays. The resulting branching fractions are shown in Fig. 3.7, and the resulting decay length (which can be macroscopic for small Δm) in Fig. 3.8.

A particularly interesting limiting case is that of completely decoupled winos and binos, in which case the mass splitting Eq. (2.25) between $\chi_1^\pm$ and χ_1^0 is numerically of the order of 300 MeV, see Fig. 2.2; in this case, the chargino decays almost exclusively into a single charged pion and the neutralino (or, more precisely, into either of the neutralinos, which are themselves mass-degenerate in that limit).

Note that it is in principle possible to obtain an even smaller mass splitting, by tuning tree-level corrections to the chargino and neutralino masses due to mixing with the wino and bino, see Eqs. (2.19) and (2.23), against the one-loop correction of Eq. (2.25). One may thus obtain a $\chi_1^\pm$ which can be stable on collider scales, cf. Fig. 3.8.

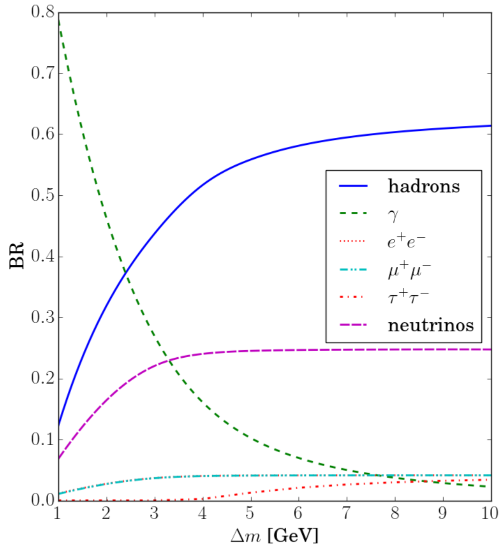

Figure 3.6: Branching fractions for the decays $\chi_2^0 \to \chi_1^0 X$ for various X and mass splittings $\Delta m_{\chi_2^0 - \chi_1^0}$ between 1 and 10 GeV, assuming $m_{\chi_1^0} = 160$ GeV. We have omitted all channels with branching ratio $< 1\%$, which includes all decays into $\chi_1^{\pm} X$. Generated with `SUSY-HIT` [64].

3.3 Collider signatures and current constraints

3.3.1 Higgsinos at the LHC

It is well known how to search for coloured supersymmetric particles at hadron colliders, and these searches have been used to put stringent bounds on standard supersymmetry scenarios such as the pMSSM (see e.g. [67, 68] for some recent results obtained with 36 fb^{-1} of 13 TeV data). Even though it is often assumed for these searches that the lightest supersymmetric particle is a bino-like neutralino, typically there is little change if it is replaced by a higgsino-like neutralino. (A possible approach to distinguish a bino LSP scenario from a higgsino LSP, in the case of a relatively simple spectrum, will be discussed in Section 3.4.) In fact, these searches rely more on the characteristics of the coloured superparticle spectrum and less on the nature of the chargino-neutralino sector, so we will not discuss them here in detail. Instead we will focus on the case where the higgsinos are effectively isolated, and where the resulting collider signatures are therefore specific to the higgsino sector.

In order to constrain directly produced higgsinos at hadron colliders, the following possibilities arise:

- In case that the mass degeneracy between the higgsino-like chargino and neutralinos are at least of the order of $5-10$ GeV, chargino decays can produce charged leptons which can be resolved and searched for. This final state becomes particularly useful in combination with an initial-state radiation jet, which provides a hard object for the higgsino to recoil from and for the experiment to trigger upon. The resulting

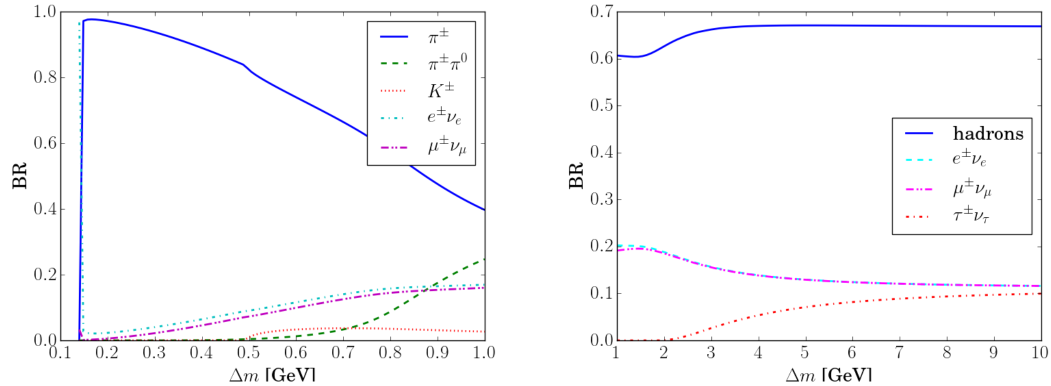

Figure 3.7: Left: Branching fractions for the decays $\chi_1^\pm \to \chi_1^0 X$ for $m_{\chi_1^0} = 160$ GeV and small mass splittings $\Delta m_{\chi_1^\pm - \chi_1^0} < 1$ GeV, where $X = \pi^\pm, \pi^\pm\pi^0, K^\pm, e^\pm\nu_e, \mu^\pm\nu_\mu$. We have omitted all channels with branching ratio $< 1\%$, and assumed that the decays into χ_2^0 are forbidden or negligible. Right: Branching fractions for larger $\Delta m_{\chi_1^\pm - \chi_1^0}$, with the hadronic BR calculated from the partial widths into quarks.

signature is a hard jet, missing transverse energy and isolated leptons.

- For the case of sufficiently small mass splittings $\Delta m_{\chi_1^\pm - \chi_1^0} \lesssim 300$ MeV (which includes the case of decoupled winos and binos), the chargino decay length becomes macroscopic and it may be searched for using a "disappearing track" signature. More precisely, in that case a chargino produced via the Drell-Yan process will decay, with a decay length which can be in the mm to cm range (see Fig. 3.8), into a very soft charged pion and a neutralino, both of which leave the detector unseen. The ionization track of the chargino effectively ends within the tracker. At the extreme end of this limiting case, for mass splittings below the charged pion mass $m_{\pi^\pm} = 140$ GeV, the chargino becomes a quasi-stable charged particle, and as such is tightly constrained by dedicated searches.

- For intermediate mass splittings, such that the heavier higgsinos decay promptly but the decay product is still too soft to be resolved, no viable approach has been found so far to probe this system at a hadron collider (but see Section 3.4.1 for some proposals). Searches for a monojet (from ISR) plus missing transverse energy may seem the most obvious option, but these searches suffer from too large irreducible backgrounds to be useful. Higgsino production in vector boson fusion or central exclusive production might significantly reduce the backgrounds, but at least at the LHC the cross-section for either process would be too small to expect enough signal events, even at high luminosity.

The strongest limits on higgsinos from missing transverse energy, an ISR jet and soft leptons are presently due to a CMS search using 36 fb^{-1} of 13 TeV data [69]. This analysis assumes $\chi_2^0\chi_1^\pm$ production, neglecting chargino pair production and neutralino-neutralino production. It requires a lepton with transverse momenta 5 GeV $< p_T <$ 30 GeV and a second isolated lepton with opposite sign, as well as missing transverse energy $E_{\mathrm{T}}^{\mathrm{miss}} > 125$ GeV and a jet with $p_T > 25$ GeV. The main backgrounds are $t\bar{t}$ production with semileptonic top decays, Drell-Yan production of leptons with an ISR jet, and

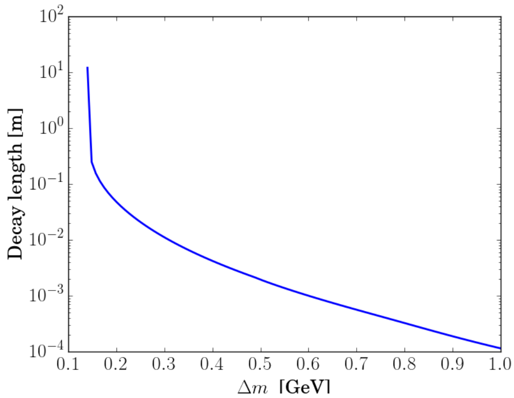

Figure 3.8: Chargino decay length for $m_{\chi_1^0} = 160$ GeV as a function of the mass splitting $\Delta m_{\chi_1^\pm - \chi_1^0}$, assuming that the only open channels are those of Fig. 3.7, i.e. $\chi_1^\pm \to \chi_1^0 X$. For mass splittings below the charged pion mass $m_{\pi^\pm} = 140$ MeV, the only remaining open channels are three-body decays into lepton-neutrino-neutralino, and the chargino becomes long-lived on collider scales.

production of an on-shell pair of electroweak gauge bosons. This analysis is sensitive to mass splittings between about 7.5 GeV $\leq \Delta m_{\chi_1^\pm - \chi_1^0} \leq 35$ GeV. It is carried out for a production cross section which corresponds to a wino-like χ_2^0 and $\chi_1^\pm$, providing limits up to 230 GeV on the wino mass, but can be recast to provide limits on higgsinos when taking into account

- the additional events obtained from $\chi_2^0 \chi_1^0$ and $\chi_1^\pm \chi_1^\mp$ production
- and the lower cross-section for higgsino-like charginos and neutralinos $\chi_1^\pm \chi_2^0$.

The resulting bound was estimated by a previous CMS analysis [70], using only 13 fb^{-1} of data and now superseded by [69], to exclude a higgsino-like chargino mass of $m_\chi^\pm = 100$ GeV for $\Delta m_{\chi_1^\pm - \chi_1^0} = 20$ GeV, hence being not competitive yet with the corresponding LEP limit yet. For the more current analysis [69] there is no interpretation in terms of higgsinos given by the experimental collaboration. We will therefore reinterpret the excluded cross-sections, given as functions of $m_{\chi_2^0}$ and of $\Delta m_{\chi_2^0 - \chi_1^0}$, in order to put bounds on higgsinos. To do so, first note that neutralino pair production and chargino pair production will provide additional events, but these will need to be weighted with respect to $\chi_2^0 \chi_1^\pm$ production, because of the possibility to obtain one of the leptons from the chargino in the latter process, and because of the different leptonic branching ratios. For simplicity we have used a flat lepton efficiency $\epsilon = 0.5$ for the reinterpretation, which is of the order of the efficiencies reported by the experiment. Moreover, we have assumed $m_{\chi_2^0} = m_{\chi_1^\pm}$ and the higgsino values for the cross-sections (implying a pure higgsino) but a non-negligible mass difference $\Delta m_{\chi_2^0 - \chi_1^0}$ (implying, contrariwise, some nonzero wino or bino admixture, for which the cross section will be correspondingly larger or smaller and the chargino will in general not be exactly degenerate with the next-to-lightest neutralino). This recast should therefore be interpreted with some caution. The result is the blue curve in Fig. 3.9. In conclusion, the CMS dilepton searches are sensitive to higgsino-like neutralinos and charginos up to masses of 155 GeV at mass splittings around 13 GeV, and can probe mass splittings down to about 8 GeV.

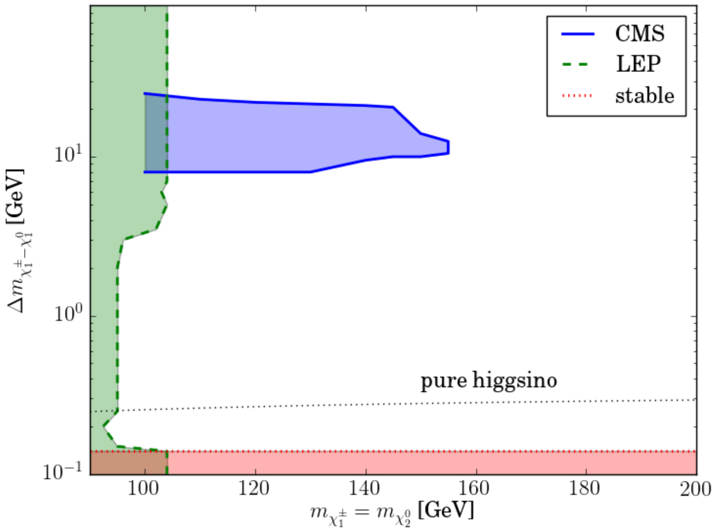

Figure 3.9: Present (2017) exclusion bounds on higgsino-like neutralinos and charginos as a function of the chargino mass and the mass splitting between the chargino and the lightest neutralino. In red, for chargino-neutralino mass splittings below the charged pion mass, the chargino is stable on collider time scales. In green, the exclusion limits given by the LEP SUSY working group [71]. In blue, the recast of the CMS soft dilepton limits [70], which assumes $m_{\chi_2^0} = m_{\chi_1^\pm}$.

For somewhat larger mass splittings of the order of 25 – 90 GeV, a CMS search for three leptons [72] can also exclude wino-like chargino masses up to 230 GeV, but the result is not directly applicable to the almost-pure higgsino case (given that such large mass splittings would imply relatively light binos or winos, and hence sizeable mixing angles). Finally, the corresponding ATLAS bounds [73] are presently not competitive.

Disappearing tracks have successfully been used by the ATLAS and CMS experiments at the LHC to constrain almost-pure wino-like charginos and neutralinos [74, 75], since the decay length of a wino-like chargino is larger than that of a higgsino-like one. More precisely, the radiative mass splitting in the pure wino system is only of the order of 160 MeV [76], leading to a proper decay length of the order of 10 cm, which is exactly the scale which disappearing track searches are most sensitive to without requiring large Lorentz boosts. As a consequence, these analyses can be used to constrain wino-like charginos with masses up to 430 GeV. As to a pure higgsino system, with a proper decay length of the order of 1 cm a sizeable boost is needed for the LHC experiments to be able to measure any track at all, and the remaining signal events must obviously be numerous enough to overcome the backgrounds, mostly given by $t\bar{t}$ and $W+$ jet events. Despite these obstacles, by recasting the published disappearing track analyses one can obtain limits on the higgsino mass which are claimed to start being competitive with the LEP bounds quoted below, excluding higgsino masses of the order of 100 GeV [77, 78], although these analysis reinterpretations are subject to considerable theoretical uncertainty. Unfortunately, the excluded cross-sections found by the latest and presently most constraining analysis, using 36 fb^{-1} of 13 TeV data at ATLAS [75], have not been released to the public. This makes any attempt at recasting this analysis for the higgsino case doubtful at least, so we will refrain from doing so.

For the fine-tuned case of extremely degenerate higgsinos with $\Delta m_{\chi_1^\pm - \chi_1^0} < m_{\pi^\pm}$ the chargino is stable on collider scales, and bounds on massive stable charged particles apply. CMS has searched for stable massive charged particles [79] with a 13 fb^{-1} sample of 13 TeV data. When reinterpreting their limits for higgsinos, the 95% exclusion bound is between 800 and 900 GeV. An older ATLAS analysis [80] with 19 fb^{-1} of 8 TeV data puts an explicit bound of 620 GeV on long-lived charginos, which however correspond to mixed wino-higgsinos. The corresponding update of the charged R-hadron search with 3.2 fb^{-1} at 13 TeV [81] does not provide any limits on charginos, but given the smaller integrated luminosity their gluino limits are somewhat weaker than those of [79], which should therefore represent the strongest bound at present. In the transitional regime of mass splittings between about 145 MeV and 140 MeV, further bounds come from searches for energy loss from ionization [82], excluding mixed higgsino-wino chargino masses of up to 482 GeV for a decay length of 4.5 m. This translates into a conservative bound on the mass of a higgsino-like chargino with the same lifetime of around 350 GeV.

3.3.2 Higgsinos at LEP

Direct higgsino pair production at an electron-positron collider allows to study the higgsino system in a much cleaner environment than at a hadron collider. The backgrounds are considerably lower, and the energy of the colliding particles is fixed (as opposed to a proton-proton machine, where only the distributions of parton energies are known). There is relatively less interest in studying the decays of heavier superparticles into the higgsinos in order to learn about the higgsino system, since the cross-section of direct higgsino production is typically the largest among the superpartners for a higgsino LSP. Hence the focus will be on Drell-Yan production via Z boson or photon exchange. As opposed to hadron colliders, thanks to the clean environment even the case of small mass splittings around 1 GeV can be studied using the monophoton signature.

The most stringent lower phenomenological limits on the higgsino mass have long been given by the null result of chargino searches at the LEP-2 collider. The LEP SUSY working group has combined the $\sqrt{s} \leq 208$ GeV data of the four LEP experiments, searching for chargino pair production events with subsequent decay into the lightest neutralino. For sizeable mass splittings $\Delta m_{\chi_1^\pm - \chi_1^0} \gtrsim 10$ GeV, this leads to a rather robust combined limit of $\sqrt{s}/2 \approx 104$ GeV on the chargino mass [83]. For lower mass splittings, the decay products become too soft to be reliably reconstructed. The preferred analysis strategy in the case of small $\Delta m_{\chi_1^\pm - \chi_1^0}$ requires a hard ISR photon, and therefore the limits become somewhat weaker [84].

Specifically, the limits on the lightest higgsino-like chargino mass derived from ALEPH, DELPHI, L3 and OPAL data for low mass splittings [71] are summarized in Table 3.1. Note that for mass splittings below the charged pion mass $m_{\pi^\pm} = 140$ MeV, the chargino becomes effectively stable on collider timescales (see Fig. 3.8), and therefore the limits are given by heavy stable charged particle searches which are again more powerful. The global lower limit is quoted by the LEP SUSY working group as $m_{\chi_1^\pm} > 92.4$ GeV at 95% CL.

The LEP exclusion limit is shown in green in Fig. 3.9. For intermediate mass splittings

$\Delta m_{\chi_1^\pm-\chi_1^0}$ [GeV]	approximate limit on $m_{\chi_1^\pm}$ [GeV]
> 10	104
> 3	102
> 0.25	95
> 0.14	93
< 0.14	104

Table 3.1: Approximate LEP-2 limits [71] on the higgsino-like chargino mass as a function of the chargino-neutralino mass splitting $\Delta m_{\chi_1^\pm-\chi_1^0}$.

it continues to be the strongest bound.

3.4 Prospects

3.4.1 The LHC at high luminosity and future hadron colliders

The discovery prospects for isolated light and near-degenerate higgsinos at the high-energy and high-luminosity LHC, as well as at future possible hadron colliders up to a hypothetical centre-of-mass energy of 100 TeV, have been studied extensively in many recent works [77, 85–91]. We have already reported the principal signatures in Section 3.3.1, so we will briefly review some main results on the anticipated discovery potential.

Concerning the signature of an ISR jet, missing transverse energy and possible soft dileptons, which is already providing some constraints with LHC-13 data (see Section 3.3.1), Ref. [85] argues that even the near-degenerate case of $\Delta m_{\chi_1^\pm-\chi_1^0} = 3.5$ GeV can be probed with a combined analysis taking into account 0-, 1- and 2-lepton events at 14 TeV. For this mass splitting, with an integrated luminosity of 3000 fb^{-1}, the reach will extend to chargino masses of at least between 100 and 120 GeV with an assumed systematic uncertainty of 5% in the background, and up to between 200 and 250 GeV if this uncertainty can be reduced to 1%. In the latter case a 5σ discovery would be possible for $\chi_1^\pm \approx 120$ GeV [85]. Refs. [87] emphasize that the monojet and monophoton searches, taken on their own, are not sensitive enough to probe the higgsino sector even at high luminosity, but that for mass splittings of the order of 10 GeV a 5σ discovery should be possible in the dilepton channel for chargino masses $\gtrsim 200$ GeV at 14 TeV with 1000 fb^{-1}. A 95% CL exclusion should, for similar mass splittings, already be possible with 100 fb^{-1} of data for $m_{\chi_1^\pm}$ around 180 GeV. Ref. [88] studies various benchmark points with varying mass splittings between about 5 and 50 GeV, concluding that with a mere 100 fb^{-1} at 14 TeV a 140 GeV chargino may be discovered if the mass splitting is large (a conclusion now superseded by the CMS analysis cited above, which shows no excess in that parameter region), again using both leptons and an ISR jet as a signature. Ref. [91] claims a sensitivity at the 2.9σ level for 14 TeV, 3000 fb^{-1} to a spectrum with a 120 GeV χ_1^0 and a mass splitting $\Delta m_{\chi_2^0-\chi_1^0}$ as low as 4 GeV by selecting collinear muon pairs in the final state, without however specifying the assumption for the background systematic uncertainty used for this analysis.

Ref. [89] claims a possible exclusion of a pure higgsino up to about 200 GeV in the monojet channel alone at 14 TeV with 3000 fb^{-1}, albeit assuming a rather optimistic

1% systematic uncertainty on the background. When increasing this uncertainty to 2% the limits obtained are no stronger than the present ones from LEP. For a future 100 TeV collider, the claimed 95% CL exclusion reach increases to between 500 and 900 GeV depending on the assumed background uncertainty, again at 3000 fb^{-1} integrated luminosity and using monojets only. In particular, the interesting region of higgsino masses around 1.1 TeV, where a pure higgsino could be a thermal WIMP dark matter candidate (see Section 3.5), would be out of reach for the monojet search even at 100 TeV. Similar conclusions are reached in [90].

Searches with three leptons in the final state were argued in [92] to provide increased sensitivity for somewhat larger mass splittings, but are mainly applicable to mixed electroweakino systems.

Concerning the disappearing track signature which is applicable for very small mass splittings, corresponding to macroscopic decay lengths, Ref. [89] finds that the LHC-14 with 3000 fb^{-1} might be sensitive to pure higgsinos with radiatively induced mass splittings up to 140 GeV, with discovery possible for up to 95 GeV. By scaling the assumed background by a factor 5, upwards or downwards, the mass reach becomes smaller or larger by about 50 GeV. At a 100 TeV collider the sensitivity for exclusion is similarly estimated at 615^{+135}_{-130} GeV, and that for discovery at 485^{+110}_{-105} GeV. Ref. [86] argues that, using a modified forward tracking system, the sensitivity of ATLAS and CMS could be improved to obtain a sensitivity to up to 420 GeV pure higgsinos for this kind of signature, and a discovery might be possible up to 380 GeV. At a 100 TeV collider, using a similar tracker and analysis strategy is claimed to allow the exclusion or discovery of pure higgsino dark matter at $m_\chi = 1.1$ TeV.

Ref. [77] argues that a mere change of the analysis strategy at LHC may already improve the sensitivity to disappearing tracks significantly, without the need for a dedicated hardware upgrade. Specifically, the present ATLAS searches require at least four hits in their pixel detector to reconstruct a disappearing "tracklet". By changing this requirement to ask for only one hit in the Insertable B-Layer, the innermost tracker component, and a further hit in the pixel tracker, the minimum detectable track length could be reduced. According to [77] this would allow to reach a sensitivity to up to around $500 - 600$ GeV higgsino masses at LHC-14 with 3000 fb^{-1}, depending on the assumptions for the background. Moreover, it is claimed that a 33 TeV collider with suitable hardware should be sensitive to a thermal relic higgsino.

It is also of interest to study mass spectra where not only the higgsinos are light, but also the stop squarks are within collider reach. The motivation for this hypothesis is that, next to the higgsinos, the third-generation squarks have the most direct impact on the Higgs potential and hence on the electroweak scale: The higgsino mass parameter μ contributes to the Higgs mass at the tree level, see Eq. (2.8), while the top-stop system couples to the Higgs sector at one loop with an $\mathcal{O}(1)$ Yukawa coupling. Second- and third-generation squarks, by contrast, can be much heavier without paying a large fine-tuning price. While the gluino is often included in the list of particles which need to be light, the radiative corrections due to the gluino, being two-loop suppressed, can still be modest if the mediation scale is low. Moreover, the gluino mass is by now already constrained to be upwards of about 2 TeV by LHC-13 data, depending on the details of the spectrum. This motivates studying a spectrum where the only particles produced in the strong interactions are stop (and possibly sbottom) pairs.

For this kind of spectrum within the MSSM, we will briefly recapitulate the main findings of the analysis of [5], even though the benchmark points originally provided in that paper have now been made obsolete by LHC data. The conclusion, however, continues to hold nevertheless: Assuming that an excess in b-jets and missing energy will eventually be found at high luminosity at the LHC or at a future hadron collider, the specifics of this excess can be used to discriminate between a mass spectrum with light higgsinos on the one hand, and a generic MSSM with a bino LSP, and possibly an intermediate-mass wino-like chargino and neutralino on the other.

To see this, let us consider stop pair production and the subsequent cascade decays into the lightest supersymmetric particle. We distinguish three scenarios:

1. electroweak-scale near-degenerate higgsinos are the lightest MSSM superpartners, and the masses of all other superpartners are at least at the TeV scale, including the relatively light stops,
2. the LSP is an electroweak-scale bino-like neutralino, and intermediate wino-like neutralinos and charginos exist at about twice its mass, as suggested by gaugino mass unification,
3. the LSP is an electroweak-scale bino-like neutralino, while the winos and higgsinos are heavy along with the other superpartners.

In case 1., the decay chain for a $\tilde{t}_1$ is always $\tilde{t}_1 \to t\chi^0_{1,2}$ or $\tilde{t}_1 \to b\chi^\pm_1$, see Fig. 3.3. Since χ^0_2 and $\chi^\pm_1$ decay into χ^0_1 and very soft hadrons, photons or leptons, their decays are effectively invisible, hence the signature of a decaying $\tilde{t}_1$ is always a hard b-jet and missing transverse energy. In case 2., some fraction of stops will also decay into intermediate charginos and neutralinos, which can subsequently decay leptonically into χ^0_1. The signature may therefore involve a hard b-jet, missing transverse energy, and a hard lepton; dedicated cuts on the signal can be used to discriminate between this case and case 1. In case 3., the only possible decay mode is $\tilde{t}_1 \to t\chi^0_1$. The resulting b-jet spectrum is different from that of case 1., where some fraction of the stops decay directly into b-quarks without an intermediate top, and can again be distinguished from it with appropriate cuts.

For a quantitative analysis and a detailed search strategy using some specific MSSM spectra, we refer to [5] (noting however that, as already mentioned, the benchmark points presented in that paper are by now outdated — it would be a worthwhile line for future research to update this analysis for LHC-14).

3.4.2 Future e^+e^- colliders

A future linear collider has frequently been argued to be an ideal experiment for probing the higgsino sector, see for instance [8, 93–95]. Indeed, the clean environment of a lepton collider together with a well-defined (and tunable) initial-state energy, and the possibility exploiting beam polarization would give a linear collider unprecedented reach and precision to not only discover light higgsinos, but also to measure their properties in some detail. To corroborate this point, we will now summarize the case study of [8] where the ILC potential was analysed for two particular benchmark points of quasi-degenerate electroweak-scale higgsinos.

The analysis in [8] assumed chargino pair production or $\chi_2^0\chi_1^0$ production in combination with an ISR photon at an ILC-like linear collider with beam polarization and a centre-of-mass energy $\sqrt{s} = 500$ GeV. The higgsino masses for the two benchmark models providing the signal are listed in 3.2. The spectra also contain the complete set of MSSM

	$m_{\chi_1^0}$	$\Delta m_{\chi_1^\pm-\chi_1^0}$	$\Delta m_{\chi_2^0-\chi_1^0}$
benchmark point I	164.17 GeV	1.60 GeV	2.70 GeV
benchmark point II	166.59 GeV	0.77 GeV	1.04 GeV

Table 3.2: Benchmark points used in the analysis of [8].

particles, with the remaining electroweak gauginos, squarks, sleptons and non-SM Higgs bosons in the multi-TeV range, and a SM-like Higgs boson in the range 125 ± 3 GeV (assuming a 3 GeV theory uncertainty in the determination of its mass). Signal events were generated with `Whizard` [96], decayed with `Pythia` [97] and processed with the fast detector simulation `SDV` [98] to model the anticipated response of the ILD detector. Fig. 3.10 shows the expected cross section for chargino pair production and neutralino pair production as a function of $\sqrt{s}$ for benchmark point II with two different beam polarizations.

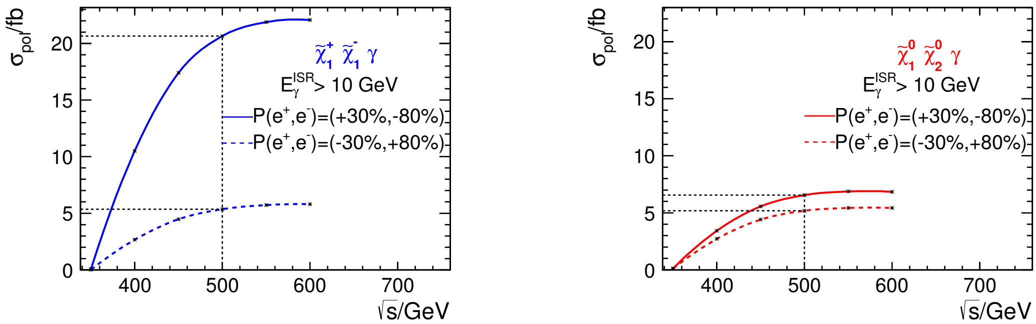

Figure 3.10: Higgsino cross sections for benchmark point II at the ILC as a function of the centre-of-mass energy $\sqrt{s}$. Left: chargino pair production, right: neutralino pair production. The solid and dashed lines represent two different beam polarization configurations.

Standard Model background events can arise from e^+e^- collisions or from photon scattering, since the ILC beam contains a considerable photon component due to the strong fields between the colliding bunches. The background events can be divided into three classes: 2-, 4- and 6−fermion events from e^+e^- collisions, 3- and 5-fermion events from $e\gamma$ scattering, and 2- and 4-fermion events from pure photon scattering. The missing energy signature characteristic of a higgsino pair recoiling from a photon can me mimicked either by Standard Model events involving either neutrinos, or by events with energy lost due to detector acceptance. The former arise mainly from τ pairs produced in $e^+e^- \to 2f$ or $e^+e^- \to 4f$ events, and can be brought under control with relatively mild cuts. A potentially severe background of the latter type are processes of the $\gamma\gamma \to 2f$ type where the colliding electrons leave through the outgoing beam pipe unseen, carrying away most of the energy, but having produced a soft fermion pair through virtual photons. This background is efficiently suppressed by demanding an ISR photon, which would deflect the corresponding beam electron into the acceptance region of the detector. A second important background is arising from $e\gamma \to 3f$ processes, specifically from t-channel electron exchange where the outgoing electron carries away a large fraction of the energy into the beam pipe, and the ISR photon can now recoil against an $f\bar{f}$ pair.

A number of preselection cuts are applied to reduce these backgrounds, the most important being

- requiring exactly one reconstructed photon with energy > 10 GeV (which reduces the number of $\gamma\gamma \to 2, 4\,f$ events by about a factor 30)
- requiring all other reconstructed particles to be central, at an angle of at least 20° with respect to the beam axis, and to have an energy less than 5 GeV (which reduces the number of $ee \to 2, 4, 6\,f$ events by two orders of magnitude)
- requiring missing energy of at least 300 GeV, with the missing momentum vector within the acceptance region of the detector.

To single out chargino pair events, one imposes the following selection cuts:

- the final state should contain a charged lepton and a pion, from one of the charginos decaying leptonically and the other hadronically, allowing also for a π^0 with the BP I events to take into account two-pion decays,
- $E^*_\pi < 3$ GeV, where the E^*_π variable is defined by

$$E^*_\pi = \frac{(\sqrt{s} - E_\gamma)E_\pi + \vec{p}_\pi \cdot \vec{p}_\gamma}{\sqrt{s'}} \tag{3.1}$$

  and where $\sqrt{s'}$ is the reduced centre-of-mass energy of the system recoiling against the ISR photon,

$$s' = s - 2\sqrt{s}E_\gamma \tag{3.2}$$

  (this provides an efficient discrimination against τ events),
- further dedicated cuts to remove $\gamma\gamma \to \tau^+\tau^-$ events at large $\sqrt{s'}$

The selection cuts to single out neutralino events are

- the final state should contain an additional "soft" photon, since the main decay mode for $\chi^0_2 \to \chi^0_1$ for these mass splittings is the radiative one, see Fig. 3.6,
- the scattering angle of the soft photon should satisfy $|\cos\theta_{\gamma_{\text{soft}}}| < 0.85$,
- and $E^*_{\gamma_{\text{soft}}} > 0.5$ GeV, where the $E^*_{\gamma_{\text{soft}}}$ variable is defined by

$$E^*_{\gamma_{\text{soft}}} = \frac{(\sqrt{s} - E_\gamma)E_{\gamma_{\text{soft}}} + \vec{p}_{\gamma_{\text{soft}}} \cdot \vec{p}_\gamma}{\sqrt{s'}}\,. \tag{3.3}$$

Table 3.3 shows the resulting cut flow for the two benchmark points in comparison with SM background events.

Importantly, the signal with an assumed integrated luminosity of 500 fb^{-1} would not only be large enough to allow for a higgsino discovery, but the signal event shape would allow for obtaining a rather accurate measurement of the $\chi^\pm_1$ and χ^0_2 masses, as well as the mass difference with the lightest neutralino. For example, for the chargino produced at the threshold, the reduced centre-of-mass energy $\sqrt{s'}$ defined in Eq. (3.2) is twice the

	BP I		BP II		Standard Model		
	$\chi_1^+\chi_1^-\gamma$	$\chi_1^0\chi_2^0\gamma$	$\chi_1^+\chi_1^-\gamma$	$\chi_1^0\chi_2^0\gamma$	$ee \to 2,4,6f$	$e\gamma \to 3,5f$	$\gamma\gamma \to 2,4f$
no cut	38672	24250	38130	23940	2.64×10^7	8.88×10^7	9.76×10^8
1 hard ISR photon	30058	9551	29675	9317	3.16×10^6	1.51×10^7	1.78×10^7
others central, soft	20611	6615	22156	7110	9092	5.97×10^5	1.24×10^6
missing energy	19872	6365	21558	6872	5731	1.18×10^5	3.31×10^5
$l^\pm\pi^\pm(\pi^0)$	5509	134			38	6197	13991
$E_\pi^* < 3\,\text{GeV}$	4435	103			0	2635	6162
$\gamma\gamma \to \tau\tau$ cuts	3813	97			0	2564	1452
$(E_{\text{miss}} > 350\,\text{GeV})$	3812	97			0	1016	511
$l^\pm\pi^\pm$			5489	38	19	2478	6754
$E_\pi^* < 3\,\text{GeV}$			5489	38	0	1465	4755
$\gamma\gamma \to \tau\tau$ cuts			4600	36	0	1417	782
$(E_{\text{miss}} > 350\,\text{GeV})$			4599	36	0	536	218
soft γ	53	1733	155	5224	399	1217	2254
$\lvert\cos\theta_{\gamma_{\text{soft}}}\rvert < 0.85$	38	1467	120	4538	233	800	1145
$E_{\gamma_{\text{soft}}}^* > 0.5\,\text{GeV}$	19	1395	22	4095	109	242	413
$(E_{\text{miss}} > 350\,\text{GeV})$	19	1395	22	4095	90	180	384

Table 3.3: Cut flow showing preselection cuts as well as selection cuts for chargino events with larger Δm and smaller Δm and neutralino events. The event rates correspond to an integrated luminosity of 500 fb^{-1} at the $\sqrt{s} = 500$ GeV ILC with -80% beam polarization for the electrons and 30% beam polarization for the positrons.

chargino mass, hence

$$m_{\chi_1^\pm} = \frac{1}{2}\sqrt{s - 2\sqrt{s}E_\gamma}\Big|_{\text{threshold}} . \tag{3.4}$$

Fitting the background part of the event distribution with a two-parameter exponential and the signal part near the threshold region with a straight line, the threshold value of $\sqrt{s'}$ can be determined from the simulated data, see Fig. 3.11. This gives a reconstructed chargino mass of $m_{\chi_1^\pm} = 168.0 \pm 1.4$ GeV for benchmark point I and of $m_{\chi_1^\pm} = 168.6 \pm 1.0$ GeV for benchmark point II, to be compared with the input masses of 165.77 GeV and 167.36 GeV respectively.

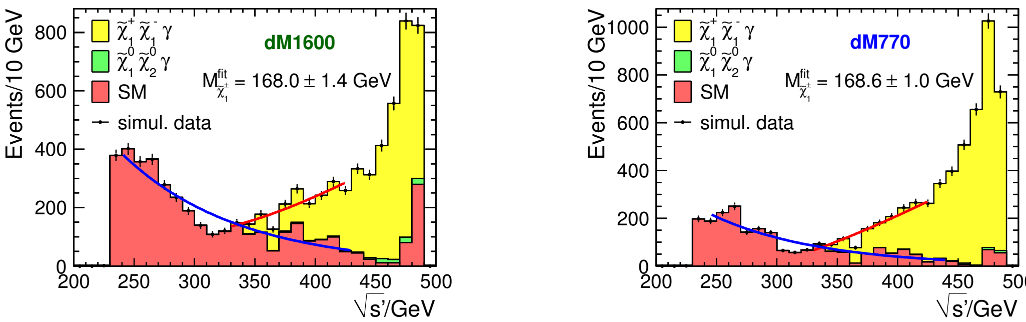

Figure 3.11: Chargino mass reconstruction from the $\sqrt{s'}$ distribution of signal and background events, taken from [8]. Left: benchmark point I, right: benchmark point II. The blue line is given by an exponential fitted to the background in the signal-free region, the red line by a linear fit to the signal near the threshold region. Their intersection, i.e. the onset of the signal, gives the threshold mass $2m_{\chi_1^\pm}$.

A similar strategy can be adopted to fit the heavier neutralino mass, see Fig. 3.12. For splittings as small as those considered here, the two neutralino masses can be taken to be approximately equal. The reconstructed neutralino masses are $m_{\chi_2^0} = 168.2 \pm 1.6$ GeV for benchmark point I and of $m_{\chi_1^\pm} = 166.3 \pm 0.8$ GeV for benchmark point II, to be compared with the input masses of 166.9 GeV and 167.6 GeV respectively.

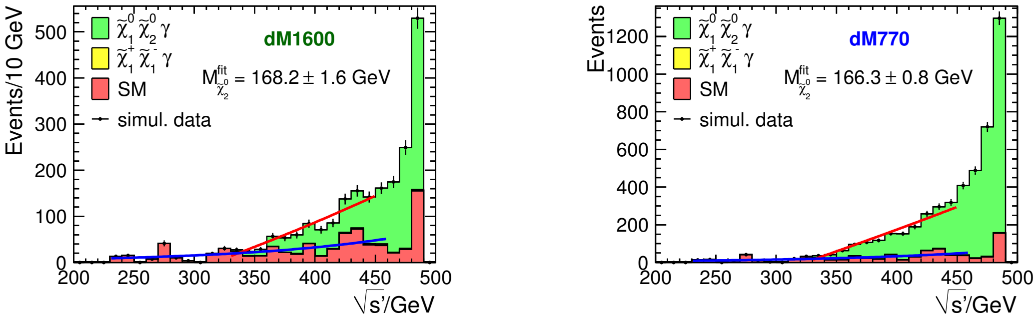

Figure 3.12: Neutralino mass reconstruction from the $\sqrt{s'}$ distribution of signal and background events, taken from [8], as in Fig. 3.11 for the chargino case. The intersection of the red line (signal fit) with the blue line (background fit) gives the threshold mass $m_{\chi_2^0} + m_{\chi_1^0}$.

A relatively loose cut on missing energy is required for this procedure, in order to obtain a good fit to the background events in a signal-free region. However, one may further reduce the background events by increasing the missing energy cut to 350 GeV, the result of which is also shown in Table 3.3. This allows to also reconstruct the mass difference $\Delta m_{\chi_1^\pm - \chi_1^0}$ with good precision. To this end, one uses the fact that $E_\pi^* \approx \Delta m_{\chi_1^\pm - \chi_1^0}$ at the threshold (where the quantity E_π^* is defined in Eq. (3.1)). Subtracting the SM background, and fitting the signal events near the threshold (at $\sqrt{s'} < 345$ GeV for this scenario) with a Gaussian, one obtains a reconstructed value of $\Delta m_{\chi_1^\pm - \chi_1^0} = 1.63 \pm 0.27$ GeV for benchmark point I and $\Delta m_{\chi_1^\pm - \chi_1^0} = 0.81 \pm 0.04$ GeV, to be compared with the input values of 1.60 GeV and 0.77 GeV respectively; see Fig. 3.13. A determination

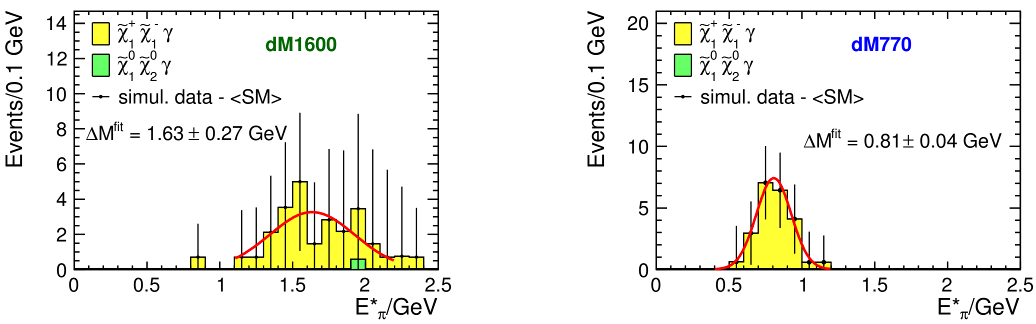

Figure 3.13: Reconstruction of the mass difference $\Delta m_{\chi_1^\pm - \chi_1^0}$ from the E_π^* distribution of signal events near the threshold, taken from [8]. Left: benchmark point I, right: benchmark point II. The red line shows a Gaussian fit to the distribution, with the mean value corresponding to the approximate mass difference.

of the neutralino-neutralino mass difference at a comparable precision would require a different strategy, and has not been studied.

Two other characteristics of the higgsino system which can be reconstructed precisely with this data sample are the chargino pair production cross section and the neutralino cross section, using event rates for different beam polarisations. The relative precision is estimated at a few percent. Finally, the fundamental parameters underlying the corresponding model, notably the μ parameter, the electroweak gaugino masses and $\tan\beta$, can also be constrained, although there is some considerable degeneracy in mapping the experimental observables to these parameters, μ being the only one among them which can be unambiguously fixed at the level of about 4%.

In summary, the analysis of [8] demonstrates that with a linear collider, it would be possible to rather precisely reconstruct the characteristics of a light higgsino system, even for sub-GeV mass differences. The same task would be extremely difficult if not impossible at a hadron collider.

3.5 Higgsino dark matter

In R-parity preserving supersymmetry, if the lightest supersymmetric particle is a neutralino, then it can be a good candidate for thermal dark matter.[1] This is because it is stable (by R-parity), and because its annihilation cross-section is generically of the right order of magnitude to produce an abundance of thermal relics which is compatible with observation, provided its mass is not too far from the electroweak scale.

This is true, in particular, for the lightest higgsino-like neutralino if the electroweak gaugino masses M_1 and M_2 satisfy $|M_{1,2}| \gg |\mu|$. However, in this almost-pure higgsino limit, the annihilation cross-section for an electroweak-scale μ is somewhat larger than the observed value $\Omega h^2 = 0.1199 \pm 0.0022$ [100]. More precise calculations yield that $|\mu| = 1.1$ TeV leads to the correct relic abundance in this limit. While such large values of the μ parameter are no longer well motivated from fine-tuning arguments, thermal dark matter provides an independent motivation to study them. Alternatively, one could consider lighter higgsinos as dark matter candidates if they were created through non-thermal processes.

In the standard WIMP freeze-out scenario, the thermal history of the universe in the presence of a neutralino as the lightest supersymmetric particle can be summarized as follows. Assuming a high reheating temperature $T_{\mathrm{RH}} \gg |\mu|$, the dark matter particle and the Standard Model particles are initially in chemical and thermal equilibrium. Once the temperature falls below $\sim |\mu|$, dark matter particle pairs can no longer be produced from SM particle collisions, while they can still annihilate. The number density decreases as a consequence of lightest neutralinos annihilating and, in particular in the higgsino-like case where they are almost degenerate with the other higgsinos, also coannihilating with similarly heavy states. As the universe expands, the annihilation rate eventually becomes of the order of the Hubble rate and the dark matter sector "freezes out": the higgsino number density remains approximately constant from that point on. Eventually even kinetic equilibrium with the Standard Model particles will be lost as the elastic scattering rate drops below the Hubble rate; all higgsino-like states will end up as lightest neutralinos.

A detailed computation of the thermal relic density of a TeV-scale higgsino-like neutralino

[1] See e.g. [99] for a general review.

χ_1^0 must take into account coannihilation [101] with χ_2^0 and $\chi_1^\pm$, which will play a large role in depleting the relic density, since these states are guaranteed to be almost mass-degenerate with χ_1^0 in the pure higgsino limit. Since the higgsino is non-relativistic at the time of freeze-out, and heavier than the electroweak gauge bosons, non-perturbative corrections due to multiple gauge boson exchange might also be considered in computing the annihilation cross-section [102]. However, in the higgsino case this effect turns out to be numerically negligible [103]. The thermal neutralino relic density as a function of the higgsino mass is shown in Fig. 3.14.

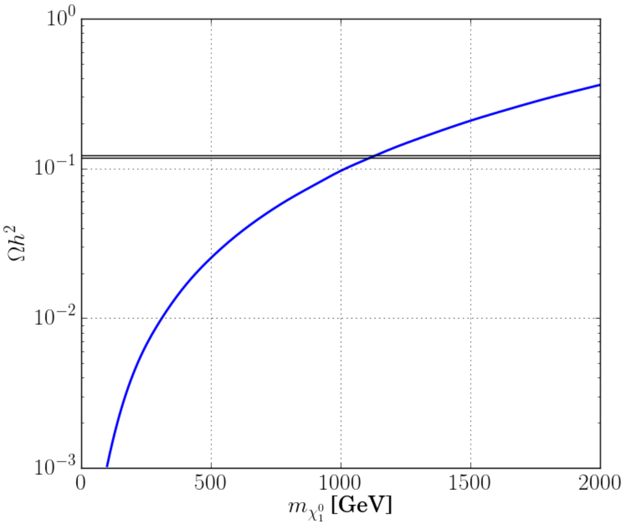

Figure 3.14: Thermal neutralino relic density for an almost-pure higgsino as a function of its mass (assuming that χ_1^0 and χ_2^0 are not completely mass-degenerate, so that χ_1^0 is the only cosmologically long-lived state). The horizontal line show the 1σ band of the Planck satellite measurement $\Omega h^2 = 0.1199 \pm 0.0022$ [100]. Generated with `micrOMEGAs` [104].

As reviewed in Section 3.4, it would take a 100 TeV hadron collider running at high luminosity to probe a pure higgsino thermal relic in the laboratory. However, almost-pure higgsino dark matter can also be constrained by direct detection experiments, even though the higgsino-nuclear cross sections are suppressed and mostly loop-induced in the near-degenerate case [105]. For example, the reach of the upcoming LZ experiment [106] is studied in [90], where it is shown that a near-degenerate higgsino could be excluded for mass splittings down to around 1 GeV.

Mass splittings below this value are more difficult to constrain. This is in part due to an accidental cancellation between the one-loop diagrams mediating the effective WIMP-nucleon interaction, which happens to be particularly effective around a Higgs mass around 125 GeV. For asymptotically small wino or bino admixtures, the higgsino-nucleon spin-independent cross section was calculated to be of the order of 10^{-49} cm^2 [107] and thus below the "neutrino floor" of coherent neutrino-nucleus scattering, which constitutes an intrinsic barrier for any direct detection experiment relying on WIMP-nucleon scattering.

However, in the pure higgsino limit the higgsino-like neutralinos χ_1^0 and χ_2^0 become effectively a Dirac fermion with a large coupling to the Z boson and a correspondingly large nuclear cross section. In fact, once the mass splitting $\Delta m_{\chi_2^0-\chi_1^0}$ is small enough to allow for inelastic scattering $N\chi_1^0 \to N'\chi_2^0$ through Z exchange, the WIMP-nucleon cross section becomes of the order of 10^{-39} cm^2, well above the present direct detection bounds. This happens for a mass splitting which is of the order of the present-day WIMP kinetic energy, or around 100 keV, corresponding to electroweak gaugino masses around $M_{1,2} \sim 10^{8-9}$ GeV. Thus, for probing higgsino dark matter in direct detection experiments, the situation is somewhat similar to that for higgsinos at hadron colliders: Sizeable mass splittings can be constrained, as well as the almost-degenerate case, while there remains a window of intermediate mass splittings which are difficult to constrain.

Indirect dark matter searches can be used to put further constraints on higgsino dark matter, independently of the mass splittings, although presently they are not yet sensitive to thermal higgsinos at 1.1 TeV. These constraints are obtained as follows: Neutralinos will annihilate with each other in regions of sizeable dark matter density, for example in galactic halos. The main annihilation channels for higgsinos are W and Z bosons, whose subsequent decays will produce charged cosmic rays, such as antiprotons, as well as neutral photons, which can be detected by satellite experiments. Unfortunately the constraints from such indirect dark matter searches tend to be subject to large astrophysical uncertainties, related to the difficulties in modelling cosmic ray propagation and the unknown dark matter distribution. Recent studies [108, 109] based on the limits of [110] (found by analysing AMS-02 antiproton data [111]) find that non-thermal higgsino dark matter is excluded up to 800^{+130}_{-320} GeV, with the large uncertainty reflecting said astrophysical uncertainties as well as uncertainties in the antiproton production cross section. Whether or not future γ ray telescopes such as CTA will eventually have sensitivity to thermal higgsinos is presently still unclear [108]. Finally, for completeness we mention that recently there have been proposals for constraining almost-pure higgsino dark matter through its capture in compact astrophysical objects [108, 112].

Chapter 4

Conclusions and outlook

Higgsino-like neutralinos and charginos continue to be promising candidates for supersymmetric particles with electroweak-scale masses. From the theoretical point of view, light higgsinos used to be well motivated before the LHC era by the argument that, at least in the simplest models, the higgsino mass directly enters in the Higgs boson potential. The most natural expectation for the higgsino mass is therefore the electroweak scale. One might now argue that similar predictions have unsuccessfully been brought forward for the stop and gluino masses, which strongly affect the electroweak scale through loop corrections and therefore should not be much heavier — an expectation which is increasingly in conflict with observation. Indeed, the recent bounds from LHC-13 clearly show that at least a little hierarchy must exist between the electroweak scale and any possible supersymmetric extension of the Standard Model, at least as far as the coloured superpartners are concerned. However, this little hierarchy need not necessarily concern the higgsino sector, which remains largely unconstrained so far. Given also that the origin of the higgsino mass might well be different from the masses of squarks and sleptons, as the μ parameter is singled out by being allowed by supersymmetry, it is certainly possible that the higgsinos could have electroweak-scale masses, and supersymmetry with light higgsinos remains a subject well worth being studied further.

In this Report we have reviewed a number of theoretical approaches to the μ problem, showing how a higgsino mass parametrically below the UV completion scale could be generated in various ways. Most of these approaches attempt to connect μ with the scale of supersymmetry breaking, by studying UV completions which forbid a tree-level μ term but generate an effective μ parameter of the order of the gravitino mass in either the Kähler potential or the superpotential. In some of them (such as in hybrid gauge-gravity mediation with large messenger numbers) the resulting μ parameter is actually predicted to be smaller than the scale of the typical soft masses. This in turn motivates studying the phenomenology of light and higgsino-like charginos and neutralinos with masses of the order of the electroweak scale. We have reviewed the current constraints and some promising search strategies at future collider experiments. It turns out that the discovery reach of the LHC and of a possible future hadron collider depends crucially on the mass splittings between charginos and neutralinos, with the arguably most interesting region of mass splittings around a GeV being the most difficult to explore. To properly probe this region, the precision of an electron-positron collider will be needed. We have also briefly sketched the present situation for a pure higgsino thermal dark matter candidate.

Experiment will ultimately decide whether there are TeV-scale squarks and gluinos (at this point it is worth noting that the LHC is still not running at its maximal energy of 14 TeV, and that the 13 TeV run has collected only about 1% of the total projected integrated luminosity to be recorded during the LHC's lifetime). Complementing these searches, chargino and neutralino searches will increasingly probe and constrain the electroweakino sector, which should lead to improving constraints on higgsinos with mass splittings of the order of 10 GeV, and on higgsinos whose masses are degenerate to the point that the chargino is effectively stable. The region in the parameter space, for which there is presently the least hope to reach enough sensitivity at the LHC to explore it is that of intermediate mass splittings between about 300 MeV and 8 GeV. Here neither searches for stable or long-lived particles nor dileptons searches are effective. It would be very interesting to study potential dedicated experimental strategies to access this region at hadron colliders, which could be a promising direction for continuing the line of research presented in Chapter 3 of the present Report. Such studies might also serve as a guideline for the design of a possible future hadron collider, or even more for a future lepton colliders. As we have shown, a linear electron-positron collider would be an excellent instrument to study light higgsinos, but our analysis was limited to a few single benchmark points. More complete studies of the parameter space of light higgsino models, with regards to the discovery potential of a future linear collider, may also be a promising avenue for future phenomenological studies.

On the theoretical side, an evident question to raise is that of the origin of the supersymmetric little hierarchy, and of its potential connection to light higgsino physics. We have presented some possible approaches in Chapter 2, for example relating μ to the breaking of a Peccei-Quinn symmetry rather than supersymmetry breaking, or proposing a focus point cancellation between the various radiative corrections to the Higgs potential. However, it seems fair to say that no complete and satisfactory mechanism is known as yet for generating an electroweak scale significantly below the soft mass scale.

Much recent research has focussed on exploring the possibility to protect the electroweak scale from large quantum corrections by introducing additional symmetries, in keeping with the usual effective field theory reasoning, in a manner that can be effectively hidden from the LHC. In this way one may attempt to bridge the discrepancy between the electroweak scale and a somewhat higher cutoff scale, where the theory could be UV-completed by a supersymmetric model. However, concrete models (see e.g. [113] for some recent examples) invariably need to pay a significant price in complicating the theory, assuming in particular a TeV-scale particle content which far exceeds that of the MSSM and typically some fairly intricate symmetry structures. One could speculate if a solution to the little hierarchy problem (and perhaps even the "big", electroweak one) might not rather be found by examining the structure of possible UV completions. This rests on the bold assumption that the said usual effective field theory reasoning fails somewhere, in a manner that is yet to be understood, when applied to the electroweak scale (or perhaps when applied to a UV completion at the Planck scale). Yet this conclusion may be the one we will be finally forced to draw in case that, in the most pessimistic of scenarios, the remaining years of the LHC running will provide no sign of new physics whatsoever. It remains to be seen whether, in any hypothetical approach that tries to predict large hierarchies of scales in the infrared from the ultraviolet structure of some UV completion of the Standard Model, there is a place for light higgsinos, or even for supersymmetry. However, as we have also shown, there is in principle no obstacle in extrapolating a model with light higgsinos to very large scales, and to match it with the MSSM, provided that the low-energy theory contains also a second Higgs doublet (and provided that the associated fine-tuning problem can indeed be relegated to the UV completion).

Interesting connections exist between the subject of light higgsinos presented here and between non-supersymmetric models of electroweakly interacting fermionic dark matter. In fact, none of the reasoning reviewed in Section 3.5 relies on the dark matter candidate being the higgsino of the MSSM; all that is needed is a (pseudo-)Dirac fermion doublet with the quantum numbers $(\mathbf{1},\mathbf{2})_{\frac{1}{2}} \oplus (\mathbf{1},\mathbf{2})_{-\frac{1}{2}}$, the only role of the electroweak gauginos of the MSSM being to provide a mass splitting large enough to avoid a conflict with direct detection. This model can be generalized to allow for different electroweak representations [114] or for several electroweak multiplets in different representations, even ones which do not appear in the MSSM [12], applying the lessons learned from supersymmetric dark matter about the importance of effects such as coannihilation. Collider constraints and the future discovery potential for such "non-supersymmetric electroweakinos" represent a related and interesting field of research, and one which again allows to draw from experience gained by studying the higgsino system.

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Annexe

- CV
- Rapport sur les stages et collaborations
- Contrats de recherche obtenus
- Direction d'étudiants
- Liste de publications
- Résumé du mémoire en français
- Tirés à part des publications pertinentes