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Breaking of R-parity and supersymmetry in supersymmetric models

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ACADEMIC DISSERTATION
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Preface

First, I would like to thank my parents Kerttu and Rainer for the last 28 years, my sister Pia for the last 25 years and Anna for being there.

Of my collaborators I would like to extend my special gratitude to Katri Huitu. Without her supervision, collaboration and support this work would not have been possible. I would also like to thank Jukka Maalampi for advice, collaboration and everything. Katri Huitu, Jukka Maalampi and Claus Montonen have read through this thesis and given valuable comments and corrections. It has been a great pleasure to work with Mariana Frank, Homayoun Hamidian, Yoshiharu Kawamura, Tatsuo Kobayashi, Pran Pandita and Da-Xin Zhang. I have learned a lot from all of them, for which I will be always in debt.

This work has been mainly done in the Helsinki Institute of Physics. My studies have mostly taken place in the Physics Department of the University of Helsinki. I have had many excellent teachers, like Masud Chaichian, whose advice, support and the experience gained in working with him has been invaluable. HIP and the Physics Department have been good places for study and research. The administrative branch of the Institute deserves special thanks for making everything run smoothly. Part of my work has been financed by the Magnus Ehrnrooth foundation and the Academy of Finland.

There are people who have made life worth living for during the past years (even though the presence of some of them may have inadvertently delayed my dissertation, but only slightly). It would be unfair to name them here, because the list would be long and necessarily incomplete. But you know who you are, and I hope we will always be friends.[1]

Lastly, I would like to extend my special gratitude to the taxpayers from Finland and all over the world for supporting my research.

[1] Of a specific request, I would like to mention a certain couple who has always helped me when I have been hungry.

Contents

Chapter 1

About this thesis

1.1 Composition of this work

This doctoral thesis consists of the introduction and the following research publications [1, 2, 3, 4, 5, 6, 7]:

1. **Radiative symmetry breaking and the $b \to s\gamma$ decay in the generalized GMSB models, by Homayoun Hamidian, Katri Huitu, Kai Puolamäki and Da-Xing Zhang. Published in Physics Letters B448 (1999) 234, hep-ph/9808341.**

   General class of models with gauge mediated supersymmetry breaking is studied. The radiative symmetry breaking mechanism and the particle spectra predicted by this class of models is studied and the $b \to s\gamma$ decay branching ratio in these models is also calculated.

2. **$SO(10)$ GUTs with gauge mediated supersymmetry breaking, by Mariana Frank, Homayoun Hamidian and Kai Puolamäki. Published in Physics Letters B456 (1999) 179, hep-ph/9903283.**

   A general class of $SU(10)$ grand unified theories is investigated within the framework of gauge mediated supersymmetry breaking. A most general messenger sector is assumed and the Standard Model gauge group is embedded into either $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ or $SU(2)_L \times U(1)_{I_{3_R}} \times U(1)_{B-L}$ left-right symmetry groups. It is found that the requiring of the perturbativity of the gauge couplings and the gauge unification leads to an almost unique messenger sector and testable predictions for the sparticle masses.

3. **The supersymmetric spectrum in $SO(10)$ GUTs with gauge mediated supersymmetry breaking, by Mariana Frank, Homayoun**

Hamidian and Kai Puolamäki. Published in Physical Review D60 (1999) 095011, hep-ph/9904458.

The analysis of Paper 2 is extended by studying particle spectrum and the properties of the models in more detail.

4. **Phenomenological constraints on $SU(5)$ GUTs with non-universal gaugino masses, by Katri Huitu, Yoshiharu Kawamura, Tatsuo Kobayashi and Kai Puolamäki. Published in Physical Review D61 (1999) 035001, hep-ph/9903528.**

   Phenomenological aspects of supersymmetric $SU(5)$ grand unified theories are studied with non-universal gaugino masses. For large $\tan\beta$, constraints arising from the requirement of a successful electroweak symmetry breaking and the positivity of stau mass squared as well as the $b \to s\gamma$ decay rate are investigated. The nature of the lightest supersymmetric particle, neutralino or stau, is determined in the allowed region. Examples of mass spectra are given. The study concentrates on the large $\tan\beta$ scenario, and care has been taken to take the relevant next-to leading order radiative corrections properly into account.

5. **Generic gravitational corrections to gauge couplings in $SU(5)$ GUTs, by Katri Huitu, Yoshiharu Kawamura, Tatsuo Kobayashi and Kai Puolamäki. Published in Physics Letters B468 (1999) 111, hep-ph/9909227.**

   Non-universal corrections to gauge couplings due to higher dimensional operators are studied in supersymmetric $SU(5)$ theories. It is found that the corrections can push up the unification scale, thus avoiding the relatively short nucleon lifetime prediction characteristic of the $SU(5)$ GUT models.

6. **Critical basis dependence in bounding R-parity breaking couplings from neutral meson mixing, by Katri Huitu, Kai Puolamäki and Da-Xing Zhang. Published in Physics Letters B446 (1999) 285, hep-ph/9808338.**

   The limits on the products of two λ'-type R-parity violating couplings are derived from neutral meson mixing. It is emphasized that the choice of the basis of quarks is important in quoting the limits.

7. **Higgs sector and R-parity breaking couplings in models with broken $U(1)_{B-L}$ gauge symmetry, by Kai Puolamäki. Published in Physical Review D62 (2000) 055010, hep-ph/0004239.**

Four different supersymmetric models based on $SU(2)_L \times U(1)_R \times U(1)_{B-L}$ and $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ left-right symmetry groups are studied. In these models the $U(1)_{B-L}$ symmetry is broken spontaneously by the vacuum expectation value (VEV) of a sneutrino field. Explicit formulas for masses and mixings in the physical lepton fields are found. The spontaneous symmetry breaking mechanism fixes the trilinear R-parity breaking couplings. In particular, a potentially large trilinear lepton number breaking coupling, which is unique to left-right models, is found.

1.2 Some abbreviations used in this thesis

CP A combined charge conjugation (C) and parity (P) transformation. The Standard Model is invariant under CP-transformation, except for the strong QCD phase related to the $SU(3)$ gauge symmetry and a phase in the Yukawa matrices. Supersymmetric model may have more sources of CP-violation.

FCNC Flavor-changing neutral currents, a common name for processes in which a quark or lepton is transformed into a quark or lepton of same charge, but of a different family. In the Standard Model the FCNC processes are strongly suppressed, making many of the FCNC processes sensitive to the radiative effects of non-Standard Model physics.

GeV A unit of energy or mass. The mass of the proton is about 1 GeV.

GMSB Gauge mediated supersymmetry breaking, a mechanism in which the supersymmetry breaking effects are transmitted to the visible sector by gauge interactions.

GUT Grand Unified Theory. A hypothetical gauge field theory, relevant at very high energies, $M_{GUT} \sim 10^{16}$ GeV, that unifies the electromagnetic, weak and color interactions.

LEP An electron-positron collider at CERN.

LSP Lightest supersymmetric particle, which is often a neutralino. In GMSB models the LSP is typically the gravitino. If R-parity is conserved the LSP is a stable particle.

MSSM Minimal Supersymmetric Standard Model, a supersymmetric extension of the Standard Model.

NLSP Next-to-lightest supersymmetric particle. In GMSB models the LSP is the gravitino, and the NLSP is typically a neutralino or the stau.

QCD Quantum chromodynamics, a gauge theory of the strong, or color, interactions.

QED Quantum electrodynamics, a gauge theory of electromagnetic interactions.

R-parity A quantum number, defined by $R = (-1)^{3(B-L)+2S}$, where B and L are baryon and lepton numbers and S is the spin of the field in question. Conservation of R-parity implies conservation of $B - L$ quantum number, which in turn implies that the nucleon is stable.

SM The Standard Model of particle physics, the gauge theory of electromagnetic, weak and strong interactions. The Standard Model describes all fundamental forces, except gravity.

SUGRA Supergravity, a non-renormalizable theory of gravity obtained by localizing supersymmetry transformations. Gravity interactions may transmit the effect of supersymmetry breaking to the visible sector of the model.

SUSY Supersymmetry, a symmetry relating bosonic and fermionic degrees of freedom in a quantum field theory.

SUSYLR The supersymmetric left-right model, an extension of MSSM, obeying gauge symmetry $SU(2)\times SU(2)\times SU(4)$ or some of its rank-five subgroups.

VEV The vacuum expectation value of a scalar field.

Chapter 2

Supersymmetric gauge field theories

2.1 Historical background

The fundamental interactions of nature, apart from gravity, are described successfully in terms of relativistic gauge field theories. The gauge theory of electromagnetic interactions, the quantum electrodynamics (QED), is probably *the* most accurately tested physical theory. It is based on $U(1)$ gauge symmetry, having the photon as a mass-less gauge boson to mediate the interaction between electrically charged particles. The electromagnetic interactions described by the QED involve vector like (V) fermion currents which conserve parity (P). There exists another class of interactions that breaks the parity symmetry, namely weak interactions. These were originally incorporated in the so-called Fermi theory or $V-A$ theory, where the weak interactions are described in terms of non-renormalizable four-fermion operators.

The QED and the Fermi theory were later unified to what is currently know as the electroweak theory, a gauge field theory based on the $SU(2)_L \times U(1)_Y$ [8, 9]. The fundamental spin-$\frac{1}{2}$ fermions of the electroweak theory are the three families of leptons and quarks. The theory predicts that there are three heavy weak gauge bosons, two electrically charged and one neutral. When the electroweak gauge bosons are integrated out of the theory there appears, apart from a charged current Fermi operator that had been observed in experiments, also a neutral current operator. The presence of this neutral current interaction was first confirmed in 1973 by the detection of the reaction $\nu_\mu/\overline{\nu}_\mu + N \rightarrow \nu_\mu/\overline{\nu}_\mu + (\text{hadrons})$, where $\nu_\mu/\overline{\nu}_m$ is the muon neutrino or its antiparticle and N is a nucleon [10]. The electroweak gauge bosons ($W^\pm$, Z)

have been later directly observed [11, 12] and their properties and couplings have been measured to a great precision.

A corner stone of the electroweak theory is that the masses of gauge bosons and fermions are generated by the so-called Higgs mechanism [13, 14, 15]. The Higgs mechanism is needed to break the $SU(2)_L \times U(1)_Y$ gauge symmetry to the residual electromagnetic $U(1)_{em}$ symmetry. The Higgs boson is the only particle in the electroweak model which has not been directly observed so far. The search of the Higgs boson has been, and will be, one of the top priorities in the present and planned future particle accelerator experiments.

The electroweak theory combined with the quantum chromodynamics (QCD), the $SU(3)_C$ gauge theory of strong interactions, is known as the Standard Model (SM) of particle physics. It describes all the basic interactions of nature, except gravity. The parameters of the Standard Model have been measured and verified thoroughly, and to date it is an accurate description of these interactions as far as any experiment can tell — at least when massive neutrinos are included in the model. Explaining the recent results of the Kamiokande underground experiment [16] on atmospheric neutrinos and the other indications of neutrino oscillations (the deficit of solar neutrinos [17] and the claimed observation of the $\nu_e \rightarrow \nu_\mu$ oscillation in a laboratory experiment [18]) is not possible within the framework of the original version of the Standard Model where neutrinos are strictly massless. The model should be modified to allow for non-vanishing neutrino masses by extending its field content or the underlying gauge symmetry in a way that does not spoil its well-tested predictions.

The Standard Model, while providing an accurate description of nature, is not yet fully satisfactory in the sense that it has a great number of *ad hoc* parameters. The masses and mixings of fermions, as well as the CP phase factor taking care of the observed CP violation of weak interactions, for example, have their origins in the Yukawa couplings for which the Standard Model provides no explanation. The Standard Model does not explain gravity. The so-called hierarchy problem is related to the radiative corrections to the Higgs mass: barring fine-tunings the Higgs mass term receives radiative corrections that are proportional to the scale of non-standard model physics ($\Lambda_{\mathrm{CUTOFF}}$).

A very important step forward was the discovery of supersymmetry in early 70's [19, 20, 21]. When implemented in the Standard Model, supersymmetry stabilizes the scalar mass terms of the scalar potential and, as it was later discovered, it leads to a unification of the gauge couplings near the Planck scale [22, 23, 24].

The original research papers included in this thesis concern the supersym-

metric version of the Standard Model and its extensions. The main emphasis is put on the analysis of supersymmetry breaking mechanisms and R-parity breaking and their experimental verification.

2.2 Hierarchy problem

There are at least two vastly different scales in nature. One is the scale of the electroweak physics, described by the Standard Model, around 100 GeV, and another is the scale of gravity. In four space-time dimensions the mass scale of gravity is given by Newton's constant G_N [25]

$$G_N = \frac{1}{8\pi M_{Pl}^2}, \tag{2.1}$$

where M_{Pl} is the reduced Planck scale $M_{Pl} = 2.4 \times 10^{18}$ GeV. There is a 16 orders of magnitude difference between these two fundamental scales.

The problem of having two vastly separate scales is most acute in computing the radiative corrections to the masses of scalar particles, such as the Higgs boson. The part of the potential of the fundamental scalar field of the Standard Model that is relevant for the Higgs mechanism reads

$$V = m_H^2 H H^* + \frac{1}{2} f \left(H H^*\right)^2, \tag{2.2}$$

where H is the Higgs scalar. The gauge symmetry is broken by a non-vanishing vacuum expectation value of the Higgs field H. The vacuum expectation value of the Higgs field contributes to the masses of the weak gauge bosons. The VEV is expected to be $\langle H H^* \rangle = (174 \text{ GeV})^2$ [25]. The minimization of the potential leads to the relation

$$\left\langle \frac{\partial V}{\partial H H^*} \right\rangle = m_H^2 + f \langle H H^* \rangle = 0. \tag{2.3}$$

The mass term $-m_H^2$ must be of the order of $(100 \text{ GeV})^2$ or less. The quantity m_H^2, however, receives enormous radiative corrections from any heavy particle that couples radiatively to the Higgs scalar. This can be seen explicitly by considering one-loop corrections in a toy model involving two complex scalar fields ϕ and φ and a Majorana fermion field Ψ. The relevant part of the Lagrangian density reads

$$-\mathcal{L} \;=\; m^2 \phi \phi^* + \mu^2 \varphi \varphi^* + \frac{1}{2} \mu' \overline{\Psi} \Psi$$

$$+\lambda^2\phi\phi^*\varphi\varphi^* + \frac{1}{2}\lambda'\overline{\Psi}\left(\phi P_L + \phi^* P_R\right)\Psi$$
$$+\left[\frac{1}{2}\lambda''\left(m+2\mu\right)\phi\varphi^2 + \mathrm{HC}\right]. \tag{2.4}$$

At tree level the mass of the scalar ϕ is

$$m_\phi^2 = \frac{1}{2}\left\langle\frac{\partial^2 V}{\partial \mathrm{Re}\phi^2}\right\rangle = m^2. \tag{2.5}$$

The masses of the scalar φ and the fermion Ψ are at tree level

$$m_\varphi^2 = \mu^2,\ m_\Psi = \mu', \tag{2.6}$$

respectively.

The radiative correction to the mass of ϕ can be found, for example, by using the effective potential approach [26]. The effective potential is at one-loop level in the $\overline{MS}$ scheme

$$V = V_0 + \Delta V_{1-\mathrm{loop}} + \ldots, \tag{2.7}$$

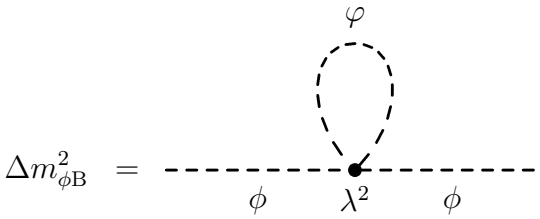
where V_0 is the tree level potential given by Equation (2.4) and $\Delta V_{1-\mathrm{loop}}$ is the one-loop radiative correction given by

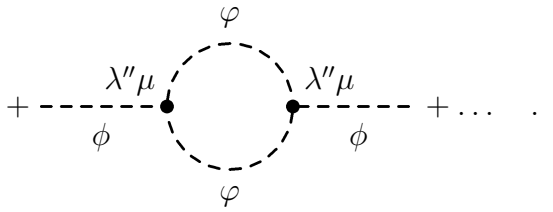
$$\Delta V_{1-\mathrm{loop}} = \frac{1}{64\pi^2}\sum_{k=\mathrm{all\ fields}}(-1)^{2J_k}\left(2J_k+1\right)m_k^4\left(\ln\frac{m_k^2}{Q^2} - \frac{3}{2}\right), \tag{2.8}$$

where m_k are the masses of the fields in the theory, J_k are their spin and Q is the renormalization scale. The one-loop correction Δm_ϕ^2 to the mass of the scalar ϕ is given by

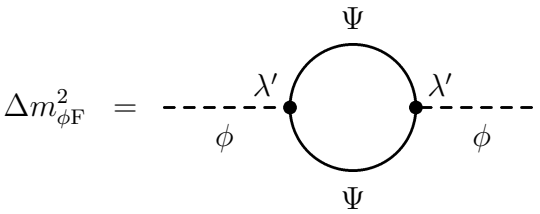
$$\Delta m_\phi^2 = \frac{1}{2}\left\langle\frac{\partial^2 \Delta V}{\partial \mathrm{Re}\phi^2}\right\rangle. \tag{2.9}$$

This can be divided into parts corresponding to the boson loops (2.10) and fermion loop (2.11), i.e., $\Delta m_\phi^2 = \Delta m_{\phi/\mathrm{B}}^2 + \Delta m_{\phi/\mathrm{F}}^2$:





$$+ \ldots \quad . \tag{2.10}$$



$$+ \ldots \quad . \tag{2.11}$$

The one-loop radiative correction due to boson loops is given by

$$\Delta m^2_{\phi/\mathrm{B}} = -\frac{1}{16\pi^2}\lambda^2\mu^2\left[1 - \left(1 + 2\frac{\lambda''^2}{\lambda^2}\right)\ln\frac{\mu^2}{Q^2}\right] + \ldots, \tag{2.12}$$

where only the terms proportional to μ^2 have been shown. The corrections due to the fermion loops are in turn given by

$$\Delta m^2_{\phi/\mathrm{F}} = \frac{1}{16\pi^2}\lambda'^2\mu'^2\left(1 - 3\ln\frac{\mu'^2}{Q^2}\right) + \ldots \quad . \tag{2.13}$$

The crucial point is that the corrections to the scalar masses are generally proportional to the mass of the heaviest particle in the theory. Even if the heavy particle had no direct coupling with the light degrees of freedom in this simple theory, higher loop effects would anyway introduce this proportionality in any realistic model.

The exact nature of the quadratic dependence on the radiative corrections is determined by the details of the renormalization scheme. In the cutoff regularization [26], where the momentum integrals are cut off by the factor $e^{-p^2/\Lambda^2_{UV}}$ (p^2 is the Euclidean momentum), the leading term in the corresponding radiative correction is proportional to Λ^2_{UV}. The cut-off scale Λ_{UV} can be understood physically as the scale at which new physics becomes relevant.

In gauge field theories the natural scale of new physics would be near the Planck scale M_{Pl}. Requiring the scalar potential of the electroweak model to

remain at the scale of $\langle HH^* \rangle = (7.3 \times 10^{-17} M_{Pl})^2$ requires thus an extreme fine-tuning of the counter terms that have a *quadratic* dependence on the scale of the beyond-the-Standard-Model physics.

2.2.1 Supersymmetry

Supersymmetry [19, 20, 21] solves the fine-tuning problem by requiring the existence of an equal number of bosonic and fermionic degrees of freedom and imposing constraints on their couplings. The dangerous radiative corrections to the masses of scalar fields generated by boson and fermion loops have an opposite sign and, if supersymmetry is strictly obeyed, they cancel each other. In the example above the supersymmetric limit is reached by choosing the dimensionless coupling constants to be equal ($\lambda = \lambda' = \lambda''$) and by requiring the boson field φ and Majorana fermion Ψ to have an equal mass ($\mu = \mu'$). This limit corresponds to a subset of an $N = 1$ supersymmetric theory defined by a superpotential

$$W = \frac{1}{2} m \phi_S^2 + \frac{1}{2} \mu \varphi_S^2 + \frac{1}{2} \lambda \phi_S \varphi_S^2, \tag{2.14}$$

where ϕ_S and φ_S denote chiral superfields. The superfield ϕ_S contains the complex scalar field ϕ and a Majorana fermion and φ_S contains the complex scalar φ and the Majorana fermion Ψ.

2.2.2 Extra dimensions

Another solution to the hierarchy problem, much discussed in literature recently, is to bring the Planck scale down to the vicinity of the electroweak scale. In this case the quadratic corrections to the masses of scalars are typically of the order of the electroweak masses, and the fine-tuning problem is avoided.

The Planck scale can be decreased by introducing extra spatial dimensions in which only gravity can propagate, in addition to the four ordinary space-time dimensions [27, 28, 29, 30, 31]. In this scenario the Standard Model gauge interactions, which have been tested up to the electroweak scale, are confined to a narrow region, a brane, in the direction of the extra dimensions, the width being of the order of $(1\ \mathrm{TeV})^{-1}$ or less. It is experimentally known that the gravity obeys the four-dimensional Gauss law $V(r) \propto 1/r$ at scales larger than about one millimeter [32]. The extra dimensions must therefore be compactified at a radius $R \lesssim 1$ mm. The Gauss law in $4 + n$ dimensions is

$$V(r) \propto -\frac{1}{M_{Pl(4+n)}^{2+n}} \frac{1}{r^{1+n}}, \ r < R, \tag{2.15}$$

where $M_{Pl(4+n)}$ is the fundamental scale of gravity in $4+n$ dimensional space-time. The potential obeys

$$V(r) \propto -\frac{1}{M_{Pl(4+n)}^{2+n} R^n} \frac{1}{r}, \; r > R \tag{2.16}$$

at the scales larger than the compactification radius R, where

$$M_{Pl(4+n)}^{2+n} R^n \sim M_{Pl}^2 \tag{2.17}$$

is the effective four-dimensional Planck mass observed at large scales.

The condition (2.17) can be satisfied with two extra dimensions ($n = 2$) so that the Planck scale $M_{Pl(4+n)}$ is close to the electroweak scale and R is consistent with the results of gravity measurements. If there are more extra dimensions, the compactification radius is correspondingly smaller; in the limit of an infinite number of extra dimensions the radius is $R \sim M_{Pl(4+n)}^{-1}$.

If the fundamental scale of gravity, in this case $M_{Pl(4+n)}$, is low enough, one does not need to introduce supersymmetry to avoid fine tuning in the effective quantum field theory. Nevertheless, supersymmetry may still be needed for the consistency of the underlying fundamental theory.

2.2.3 Composite scalars

Another suggested solution to the fine-tuning problem is that the Higgs scalar does not exist on the fundamental level. In the so-called technicolor models, the Higgs field is a composite object, namely a condensation of fermion field, which transform under some hypothetical gauge group that becomes strong at scale $\Lambda_{TC} \sim 1$ TeV; that is, the coupling constant of the gauge group satisfies $\alpha_{TC}(\Lambda_{TC}) = g_{TC}^2(\Lambda_{TC})/(4\pi) \sim 1$ [33, 34, 35, 36]. It has turned out to be difficult to construct a model that at the same time generates masses for all fermions, including the top and bottom quarks, and is not in contradiction with the experimental measurements involving electroweak gauge bosons and flavor changing neutral currents.

2.3 Supersymmetry algebra

The supersymmetry algebra is an extension of the Poincaré algebra, i.e. the algebra of proper Lorentz transformations of $SO(1,3)$ and space-time translations [19, 20, 21, 37]. The generators of the space-time translations P_μ correspond to four-momentum and the generators of the proper Lorentz transformations $M_{\mu\nu}$ of $SO(1,3)$ energy-momentum tensor.

The Coleman-Mandula theorem states that the most general Lie algebra of symmetries of S-matrix of a local relativistic quantum field theory in four dimensional space-time contains only the Poincaré generators and a finite number of internal symmetry generators that belong to a Lie algebra of a compact Lie group [38].[1] The generators of the internal symmetry are necessarily scalars under the Poincaré group, i.e. they commute with the Poincaré generators.

Within the assumptions of the Coleman-Mandula theorem there is no room for supersymmetry, the symmetry between the bosonic and fermionic degrees of freedom. This restriction can be avoided by allowing also anti-commutators within the defining relations of the algebra. The resulting algebra is called a graded Lie algebra [39].

The commuting part of the supersymmetry algebra (denoted generically by the generator M) is fixed by the Coleman-Mandula theorem. The anti-commuting part (denoted by Q) obeys the following schematic relations:

$$\{Q, Q\} = M, \qquad [M, Q] = Q. \tag{2.18}$$

A minimal non-trivial supersymmetry algebra is given by the momentum generators P_μ and the supersymmetry generators Q_α and $\overline{Q}_{\dot{\alpha}}$ ($\alpha = 1, 2$) behaving like Weyl spinors under the Poincaré transformations:

$$[Q_\alpha, P_\mu] = \left[\overline{Q}_{\dot{\alpha}}, P_\mu\right] = [P_\mu, P_\nu] = 0, \tag{2.19}$$

$$\{Q_\alpha, Q_\beta\} = \left\{\overline{Q}_{\dot{\alpha}}, \overline{Q}_{\dot{\beta}}\right\} = 0, \tag{2.20}$$

$$\left\{Q_\alpha, \overline{Q}_{\dot{\beta}}\right\} = 2\sigma^\mu_{\alpha\dot{\beta}} P_\mu. \tag{2.21}$$

The supersymmetry generators commute with the energy-momentum tensor, $[Q_\alpha, M_{\mu\nu}] = \left[\overline{Q}_{\dot{\alpha}}, M_{\mu\nu}\right] = 0$. These relations describe the simplest supersymmetric algebra in which there is only one set of supersymmetry generators Q_α and $\overline{Q}_{\dot{\alpha}}$ ($N = 1$). In a more general case one can have an arbitrary number N of such generators. A representation of the supersymmetry algebra, a supermultiplet, describes fields with $N + 1$ different kinds of helicities, separated by a half unit of spin [40]. If $N \geq 2$ the multiplet of left-handed leptons or quarks would also contain their right-handed partners (plus complex scalars). Due to the supersymmetry the left- and right-handed components of the matter fermions should transform similarly under the gauge group. However, it is known from experiments that the left- and right-handed fermions have different gauge quantum numbers in the $SU(2)_L \times U(1)_Y$ gauge symmetry.

[1] In the derivation of the Coleman-Mandula theorem it is further assumed that there are finite number of particles of given mass and that there is an energy gap between the vacuum and one-particle states.

For this reason it is thought that *only* $N = 1$ supersymmetry has relevance for the low-energy phenomenology.

A finite element of the supersymmetry group is [37]

$$G\left(x^{\mu}, \xi, \overline{\xi}\right) = e^{i\left(\xi Q + \overline{\xi Q} - x^{\mu} P_{\mu}\right)}, \tag{2.22}$$

where ξ_{α} and $\overline{\xi}_{\dot{\alpha}}$ are two-component anti-commuting Grassmann variables. A linear representation of the supersymmetry group is given by

$$\begin{aligned} P_{\mu} &= i\partial_{\mu}, \\ iQ_{\alpha} &= \frac{\partial}{\partial\theta^{\alpha}} - i\sigma^{\mu}_{\alpha\dot{\alpha}}\overline{\theta}^{\dot{\alpha}}\partial_{\mu}, \\ i\overline{Q}_{\dot{\alpha}} &= -\frac{\partial}{\partial\overline{\theta}^{\dot{\alpha}}} + i\theta^{\alpha}\sigma^{\mu}_{\alpha\dot{\alpha}}\partial_{\mu}. \end{aligned} \tag{2.23}$$

In group theory the covariant derivative D is an object that (anti)commutes with the generators of the group. In the case of supersymmetry algebra one has schematically $\{D, Q\} = \{D, \overline{Q}\} = \{\overline{D}, Q\} = \{\overline{D}, \overline{Q}\} = 0$, or equivalently, $[D, G] = 0$, where G is an element of the supersymmetry algebra given in Equation (2.22). The covariant derivative can thus be used to impose constraints that are invariant under supersymmetry transformations. In the representation (2.23) the covariant derivatives are

$$\begin{aligned} D_{\alpha} &= \frac{\partial}{\partial\theta^{\alpha}} + i\sigma^{\mu}_{\alpha\dot{\alpha}}\overline{\theta}^{\dot{\alpha}}\partial_{\mu}, \\ \overline{D}_{\dot{\alpha}} &= -\frac{\partial}{\partial\overline{\theta}^{\dot{\alpha}}} - i\theta^{\alpha}\sigma^{\mu}_{\alpha\dot{\alpha}}\partial_{\mu}. \end{aligned} \tag{2.24}$$

A *superfield* is a representation of the supersymmetry algebra. A general superfield can be expressed as a finite series in terms of the Grassmann variables θ and $\overline{\theta}$. A superfield Φ satisfying a covariant constraint

$$\overline{D}_{\dot{\alpha}}\Phi = 0 \tag{2.25}$$

is called a (left) chiral superfield. (A right chiral superfield $\Phi^{\dagger}$ would satisfy the constraint $D_{\alpha}\Phi^{\dagger} = 0$.) The chiral superfield can be expressed in terms of the component fields as

$$\Phi(y^{\mu}, \theta) = \phi(y) + \sqrt{2}\theta\Psi(y) + \theta\theta F(y), \tag{2.26}$$

where $y^{\mu} = x^{\mu} + i\theta\sigma^{\mu}\overline{\theta}$. Here ϕ is a complex scalar field, Ψ is a Majorana fermion and F is an auxiliary field that can be eliminated using the equations

of motion. Since $\overline{D}$ is a linear differential operator, holomorphic functions of chiral superfields are also chiral superfields.

In a supersymmetric theory the total action is invariant under supersymmetry transformations. The F-component of a chiral superfield W (the coefficient of θ^2 in the expansion with the respect of the Grassmann variables θ_α and $\overline{\theta}_{\dot{\alpha}}$) transforms as a total derivative in infinitesimal supersymmetry transformations, i.e., $\delta F = \int d^2\theta \delta W = i\sqrt{2}\partial_\mu \Psi \sigma^\mu \xi$. This expression can thus be used in constructing invariant Lagrangians, because the variation of the total action vanishes, $\delta S = \int d^4x d^2\theta \delta W = 0$.

Another total derivative can be constructed from the D-term (the coefficient of $\theta^2\overline{\theta}^2$) of a product of two chiral superfields: $\int d^2\theta d^2\overline{\theta} \Phi^\dagger \Phi$.

A chiral superfield contains a complex scalar (spin-0) and a Majorana fermion (spin-$\frac{1}{2}$). A full model that has the Standard Model as low energy limit must, of course, contain spin-1 gauge bosons as well. These are incorporated in the so-called vector superfields V which consist of a spin-$\frac{1}{2}$ Majorana fermion (gaugino) and a spin-1 gauge field. The vector superfields, and thus the spin-1 vector field, can be required to be real, that is, they satisfy the covariant constraint

$$V(x, \theta, \overline{\theta}) = V(x, \theta, \overline{\theta})^\dagger, \tag{2.27}$$

which defines an irreducible representation of the supersymmetry algebra.

In a non-supersymmetric theory a scalar or fermion field Φ transforms in a gauge transformation as

$$\Phi \to e^{-2igT_a\Lambda^a}\Phi, \tag{2.28}$$

where g is the gauge coupling, T_a are the generators of the gauge group, and Λ^a are some arbitrary real numbers. In the supersymmetric generalization of the gauge transformation the field Φ and the numbers Λ^a are replaced by chiral superfields. Then the gauge superfield V transforms as

$$e^{2gT_aV^a} \to e^{-2igT_a\Lambda^a} e^{2gT_aV^a} e^{2igT_a\Lambda^a}. \tag{2.29}$$

The gauge superfields V^a appear in the Lagrangian in the form of field strength superfield W defined as

$$W_\alpha^a \equiv \overline{D}^2 \left[D_\alpha V^a + igf^{abc} \left(D_\alpha V_b\right) V_c \right] + \mathcal{O}\left(V^3\right). \tag{2.30}$$

For an abelian group the terms proportional to $\mathcal{O}(V^2)$ vanish. For a non-abelian group the the terms $\mathcal{O}(V^3)$ vanish identically in the so-called Wess-Zumino gauge by the defining condition of the gauge [37]. The W_α^a is a chiral superfield, as defined in (2.25), because of the property $\overline{D}^3 = 0$ of the

covariant derivative. The most general renormalizable Lagrangian involving chiral and gauge superfields invariant under the gauge and global $N = 1$ supersymmetry transformations can then be written in the Wess-Zumino gauge as

$$\mathcal{L} = \int d^2\theta d^2\overline{\theta}\Phi_i^\dagger e^{2gT_aV^a}\Phi_i + \left[\int d^2\theta \left(\frac{1}{64}W^{a\alpha}W_{a\alpha} + W\left(\Phi_i\right)\right) + \mathrm{HC}\right], \tag{2.31}$$

where the superpotential $W(\Phi_i)$ is a holomorphic function of the chiral superfields Φ_i.

An important consequence of the supersymmetry is that the vacuum energy vanishes [41]. The Hamiltonian of a globally supersymmetric theory can be written with the help of the anti-commutation relation (2.21) as

$$H = P_0 = \frac{1}{4}\left(Q_1\overline{Q}_{\dot{1}} + \overline{Q}_{\dot{1}}Q_1 + Q_2\overline{Q}_{\dot{2}} + \overline{Q}_{\dot{2}}Q_2\right). \tag{2.32}$$

The vacuum energy is always non-negative, since it is a sum of squares of hermitean operators. The vacuum is invariant under supersymmetry transformations ($Q_\alpha\left|0\right\rangle = 0$ and $\overline{Q}_{\dot{\alpha}}\left|0\right\rangle = 0$) if and only if the vacuum energy is zero:

$$\left\langle H\right\rangle = \left\langle 0\left|H\right|0\right\rangle = \frac{1}{4}\left(\left|Q_1\left|0\right\rangle\right|^2 + \left|\overline{Q}_{\dot{1}}\left|0\right\rangle\right|^2 + \left|Q_2\left|0\right\rangle\right|^2 + \left|\overline{Q}_{\dot{2}}\left|0\right\rangle\right|^2\right) = 0. \tag{2.33}$$

Thus in all supersymmetry preserving ground states one has a vanishing potential $\left\langle V\right\rangle = 0$ to all orders of perturbation theory. On the other hand, a potential with a positive vacuum energy $\left\langle V\right\rangle > 0$ implies that the supersymmetry is broken dynamically.

In locally supersymmetric supergravity models the situation is somewhat different: the ground state of the potential may have a *negative* energy. The positive contribution of the supersymmetry breaking effects to the scalar potential can cancel the negative contribution, resulting in a vanishing vacuum energy. A vanishingly small vacuum energy is desirable, since it would imply a very small cosmological constant. It would be consistent with observational data suggesting that the cosmological constant Λ is no larger than $|\Lambda| \lesssim (10^{-12}\ \mathrm{GeV})^4$ [42, 43, 44, 45].

2.4 Realistic models

2.4.1 Minimal Supersymmetric Standard Model

The simplest phenomenologically viable supersymmetric model, the Minimal Supersymmetric Standard Model (MSSM) [46, 47], can be constructed

starting from the Standard Model. The lepton and quark fields are replaced by the respective chiral superfields containing a spin-0 scalar and a spin-$\frac{1}{2}$ fermion field and the gauge fields are replaced by gauge superfields containing the spin-1 gauge boson and spin-$\frac{1}{2}$ gaugino field. The MSSM obeys the $SU(3)_C \times SU(2)_L \times U(1)_Y$ gauge symmetry.

In the Standard Model there is only one Higgs doublet, H, which generates masses of quarks through the couplings $\overline{u}_R u_L H$ and $\overline{d}_R d_L H^\dagger$. In a supersymmetric theory the latter term is, however, forbidden as it is not invariant in a supersymmetry transformation. Therefore two Higgs doublet superfields with opposite hypercharges are needed.

The field content of the MSSM is thus the following. The gauge fields are arranged in the gauge supermultiplets

$$\begin{aligned} G_a &\quad (\mathbf{8},\mathbf{1},0), \\ W_a &\quad (\mathbf{1},\mathbf{3},0), \\ B_a &\quad (\mathbf{1},\mathbf{1},0), \end{aligned} \tag{2.34}$$

where the numbers in parenthesis denote the $SU(3)_C \times SU(2)_L \times U(1)_Y$ representation. The Higgs fields as well as lepton and quark fields are described together with their superpartners in terms of the following chiral superfields:

$$\begin{aligned} H_u &= \begin{pmatrix} H_u^0 \\ H_u^- \end{pmatrix} \quad \left(\mathbf{1},\mathbf{2},-\frac{1}{2}\right), \\ H_d &= \begin{pmatrix} H_d^+ \\ H_d^0 \end{pmatrix} \quad \left(\mathbf{1},\mathbf{2},\frac{1}{2}\right), \\ L_i &= \begin{pmatrix} \nu_{Li} \\ e_{Li}^- \end{pmatrix} \quad \left(\mathbf{1},\mathbf{2},-\frac{1}{2}\right), \\ E_i &= e_{Ri}^+ \quad (\mathbf{1},\mathbf{1},1), \\ Q_i &= \begin{pmatrix} u_{Li} \\ d_{Li} \end{pmatrix} \quad \left(\mathbf{3},\mathbf{2},-\frac{1}{6}\right), \\ D_i &= d_{Ri} \quad \left(\mathbf{3}^*,\mathbf{1},\frac{2}{3}\right), \\ U_i &= u_{Ri} \quad \left(\mathbf{3}^*,\mathbf{1},-\frac{1}{3}\right), \end{aligned} \tag{2.35}$$

where $i = 1,2,3$ is the flavor index and the color indices have been suppressed. The most general renormalizable superpotential can then be written as

$$\begin{aligned} W_{\text{MSSM}} &= \lambda_{eij} E_i L_j^T i\tau_2 H_d + \lambda_{dij} D_i Q_j^T i\tau_2 H_d + \lambda_{uij} U_i Q_j^T i\tau_2 H_u \\ &\quad + \mu H_u^T i\tau_2 H_d + W_{\not{L}} + W_{\not{B}}, \end{aligned} \tag{2.36}$$

where λ_e, λ_d and λ_u are the Yukawa coupling matrices of leptons and quarks and $W_{\not L}$ and $W_{\not B}$ are the lepton and baryon number violating contributions to the superpotential.

One shortcoming of the MSSM is that it predicts, in contrast to the SM, the existence of baryon and lepton number violating interactions, which are strongly disfavored by the observed stability of the nucleons. This problem is usually solved by introducing an extra symmetry, the so-called R-parity, that forbids these interactions, i.e. $W_{\not L} = W_{\not B} = 0$. The R-parity will be discussed later in Chapter 4.

In addition to the hierarchy problem, the Standard Model has a unification problem: while the $SU(2)_L$ gauge coupling α_2 is (in the $\overline{MS}$ scheme) unified to the $U(1)_Y$ gauge coupling α_1 at $M_X \simeq 10^{13}$ GeV, $\alpha_1(M_X) = \alpha_2(M_X)$, the $SU(3)_C$ color gauge coupling α_3 misses this unification. The color and isospin gauge couplings meet, $\alpha_3(M'_X) = \alpha_2(M'_X)$, at scale $M'_X \simeq 10^{16}$ GeV. This invalidates the grand unification of the Standard Model forces, which is required by many candidates for an underlying theory. With the supersymmetric extension of the Standard Model all three gauge couplings meet at the scale $M_{GUT} \simeq 2 \times 10^{16}$ GeV, providing the masses of the supersymmetric partners are around 1 TeV. The apparent unification of the gauge couplings, $\alpha_1(M_{GUT}) = \alpha_2(M_{GUT}) = \alpha_3(M_{GUT})$, is one of the great advantages of the MSSM as compared with the SM [22, 23, 24].

The lightest supersymmetric particle (LSP) is stable, if the baryon and lepton numbers are conserved. If the lightest supersymmetric particle is the neutralino, it may contribute to non-baryonic cold dark matter density of the universe [48, 49]. The recent observations seem to favor an accelerating flat universe, with matter density of about one-third of the critical density [50, 51].

Consistency with the recent evidence of non-vanishing neutrino masses [16] may require the addition of a gauge-singlet sterile neutrino $N_i = \nu_{Ri}$ and term

$$W_{\text{Dirac}} = \lambda^D_{\nu ij} N_i L_j i\tau_2 H_u, \tag{2.37}$$

into the theory. The term (2.37) would give origin for a Dirac mass term. A violation of the lepton number, and thus R-parity, may generate an effective Majorana mass for neutrinos, even without the sterile neutrino and Dirac mass term.

2.4.2 Supersymmetric left-right model

As discussed above, the MSSM does not explain, in contrast to the non-supersymmetric $SU(2) \times U(1)$ model, why the baryon and lepton numbers

should be conserved. Furthermore, the MSSM gives no explanation for the apparent asymmetry of the interaction of the left- and right-handed fermions. Let us now discuss an extended model, the supersymmetric left-right model, that does not have these shortcomings.

The so-called E-chain ($\ldots \supset E_7 \supset E_6 \supset SO(10) \supset \ldots$) of grand unification has two physically viable symmetry breaking patterns [52]:

$$SO(10) \supset SU(5) \times U(1) \supset SU(3) \times SU(2) \times U(1) \quad \text{and} \tag{2.38}$$

$$SO(10) \supset SU(2) \times SU(2) \times SU(4). \tag{2.39}$$

The breaking chain (2.38) leads to the $SU(5)$ grand unified symmetry (GUT) as an intermediate step, which breaks further to the MSSM symmetry $SU(3) \times SU(2) \times U(1)$ at the scale of around 10^{16} GeV [53, 54, 55, 56]. In the second breaking chain (2.39) the $SO(10)$ grand unified symmetry breaks, instead to the SM gauge symmetry, to the symmetry $SU(2) \times SU(2) \times SU(4)$, a maximal subgroup of $SO(10)$, where the left- and right-handed fermions are treated on the same footing, i.e. the resulting theory is left-right symmetric. The $SU(4)$ part is further broken to $SU(3)_C \times U(1)_{B-L}$, where $SU(3)_C$ is the color gauge group. The $U(1)_{B-L}$ implies the conservation of baryon number minus lepton number, $B - L$, which is thus an automatic consequence of gauge symmetry in this model. One of the $SU(2)$ symmetry groups, denoted by $SU(2)_L$, is associated with the ordinary weak isospin, while the other one, denoted by $SU(2)_R$, is its counterpart for right-handed fermions [2, 3, 56].

The gauge quantum numbers are defined so that the electric charge Q, the generator of the gauge symmetry $U(1)_{em}$ of the electromagnetic interaction, is given by

$$Q = I_{3L} + I_{3R} + (B - L), \tag{2.40}$$

where I_{3L} is the $SU(2)_L$ isospin quantum number, I_{3R} is the respective $SU(2)_R$ quantum number and $(B - L)$ is the $U(1)_{B-L}$ charge. The weak hypercharge Y is defined as $Y = I_{3R} + (B - L)$.

A typical phenomenologically viable particle content of a left-right model obeying $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$ gauge symmetry is given by

$$L_L^i = \begin{pmatrix} \nu_L^i \\ e_L^i \end{pmatrix} \qquad \left(\mathbf{2}, \mathbf{1}, -\tfrac{1}{2}, \mathbf{1}\right),$$

$$L_R^i = \begin{pmatrix} e_R^i \\ \nu_R^i \end{pmatrix} \qquad \left(\mathbf{1}, \mathbf{2}, \tfrac{1}{2}, \mathbf{1}\right),$$

$$Q_L^i = \begin{pmatrix} u_L^i \\ d_L^i \end{pmatrix} \qquad \left(\mathbf{2}, \mathbf{1}, -\tfrac{1}{6}, \mathbf{3}\right),$$

$$Q_R^i = \begin{pmatrix} d_R^i \\ u_R^i \end{pmatrix} \qquad \left(\mathbf{1}, \mathbf{2}, \tfrac{1}{6}, \mathbf{3}^*\right),$$

$$\phi_k = \begin{pmatrix} \phi^0_{1k} & \phi^+_{2k} \\ \phi^-_{1k} & \phi^0_{2k} \end{pmatrix} \quad (\mathbf{2},\mathbf{2},0,\mathbf{1})\,(k=1,2),$$

$$\Delta_R = \begin{pmatrix} \frac{1}{\sqrt{2}}\Delta^-_R & \Delta^0_R \\ \Delta^{--}_R & -\frac{1}{\sqrt{2}}\Delta^-_R \end{pmatrix} \qquad (\mathbf{1},\mathbf{3},-1,\mathbf{1})\,,$$

$$\delta_R = \begin{pmatrix} \frac{1}{\sqrt{2}}\delta^+_R & \delta^{++}_R \\ \delta^0_R & -\frac{1}{\sqrt{2}}\delta^+_R \end{pmatrix} \qquad (\mathbf{1},\mathbf{3},1,\mathbf{1})\,,$$

$$\Delta_L = \begin{pmatrix} \frac{1}{\sqrt{2}}\Delta^-_L & \Delta^0_L \\ \Delta^{--}_L & -\frac{1}{\sqrt{2}}\Delta^-_L \end{pmatrix} \qquad (\mathbf{3},\mathbf{1},-1,\mathbf{1})\,,$$

$$\delta_L = \begin{pmatrix} \frac{1}{\sqrt{2}}\delta^+_L & \delta^{++}_L \\ \delta^0_L & -\frac{1}{\sqrt{2}}\delta^+_L \end{pmatrix} \qquad (\mathbf{3},\mathbf{1},1,\mathbf{1})\,. \tag{2.41}$$

The numbers in parenthesis denote the gauge quantum numbers of the representations under the $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$ gauge group.

The left- and right-handed lepton and quark superfields are accommodated in the $SU(2)_L$ and $SU(2)_R$ doublets $L^i_{L,R}$ and $Q^i_{L,R}$. The two Higgs bidoublets, ϕ_1 and ϕ_2, contribute to the electroweak symmetry breaking by obtaining non-vanishing vacuum expectation values $\langle\phi^0_{11}\rangle = v_d$ and $\langle\phi^0_{22}\rangle = v_u$. These VEVs are controlled by the mass of the the $SU(2)_L$ gauge boson $W^\pm_L$:

$$m^2_{W^\pm_L} = \frac{1}{2}g^2_L\left(v^2_u + v^2_d\right) = (81.2\text{ GeV})^2\,. \tag{2.42}$$

The ratio of the VEVs is denoted by $\tan\beta = v_u/v_d$ like in the MSSM.

The breaking of the intermediate $SU(2)_R \times U(1)_{B-L}$ symmetry into the hypercharge symmetry $U(1)_Y$ is due to the vacuum expectation value of the triplet Higgses, $\langle\delta^0_R\rangle = v_{\delta R}$ and $\langle\Delta^0_R\rangle = v_{\Delta R}$, or to the VEV of the right-handed sneutrino, $\langle\tilde{\nu}_R\rangle = \sigma_R$. The mass of the right-handed gauge boson $W^\pm_R$ is given by

$$m^2_{W^\pm_R} = \frac{1}{2}g^2_R\left(2v^2_{\Delta R} + 2v^2_{\Delta R} + \sigma^2_R\right), \tag{2.43}$$

where the contribution of the bidoublet VEVs v_u and v_d has been ignored. The experimental lower limit on the mass of the $W^\pm_R$ boson, obtained at the Tevatron, is 715 GeV [57, 25]. The magnitude of the right-handed VEVs, $v_{\Delta R}$, $v_{\delta R}$ and σ_R, must therefore be of the order of TeV or more. Limits arising from the measurements of flavor changing neutral currents may constrain the right-handed scale to be of the order of 20 TeV or higher [58, 7]. In GUT models the right-handed scale can be anything between the Planck scale and TeV-scale: in a typical model it is of the order of 10^{11} GeV, fixed by requiring the unification of the gauge coupling constants at some higher energy scale [56].

The particle content of the model may be expanded from what was presented in (2.41) by singlet or triplet superfields that have a vanishing $B-L$ quantum number [59, 60, 61]. The $SU(2)_R$ symmetry may in the low-energy theory be replaced by an $U(1)_{I_{3R}}$ symmetry [62]. In the $SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L}$ theory the matter field content can in fact be similar to that of the MSSM apart from additional right-handed neutrinos [7].

Although the supersymmetric left-right (SUSYLR) model may at low energies be phenomenologically indistinguishable from the MSSM, it differs from that in many essential respects. In this model the parity symmetry, while valid at high energies, is broken dynamically in low-energy phenomena, in contrast to the SM and MSSM, where the parity violation is put in quite arbitrarily by hand. This model also involves naturally right-handed neutrinos whose existence may be necessary for explaining the observed small non-vanishing neutrino mass. Furthermore, in the SUSYLR model the R-parity symmetry is automatically present due to the $U(1)_{B-L}$ gauge symmetry while in the MSSM it must be separately introduced as an extra symmetry. On the other hand, in the SUSYLR model the R-parity can be spontaneously broken, resulting in a predictive pattern of R-parity violating couplings, as will be discussed in Section 4. The left-right symmetries can be used to make the supersymmetric contribution to the CP violating phases vanish [63, 64, 65, 66]. The doubly charged Higgs bosons and Higgsinos are phenomenologically particularly interesting, since they can be arbitrarily light, are relatively easy to observe and have no MSSM counterparts.

Chapter 3

Breaking of supersymmetry

If supersymmetry were unbroken, the masses of the component fields of a chiral superfield, like selectrons $\tilde{e}_1$ and $\tilde{e}_2$ and the electron e, should obey the relation $2m_e^2 = m_{\tilde{e}_1}^2 + m_{\tilde{e}_2}^2$. This would imply the existence of some supersymmetric particles with masses of the same order of magnitude as those of their Standard Model counterparts. This is a special case of a general tree-level result of the theories of chiral superfields [67] stating

$$\mathrm{STr}M^2 = \sum_J (-1)^{2J} (2J+1)\, m_J^2 = 0, \tag{3.1}$$

where STr denotes so-called supertrace over the real fields of spin J and mass m_J.

No such light superpartners of any Standard Model particles has been found, however. The current limits from the direct searches for supersymmetric particles at the LEP and Tevatron experiments on the masses of the supersymmetric particles in the MSSM are listed in Table 3.1.

The exact mass limits depend on the experimental setup and the composition of the fields. The limits obtained from indirect measurements, such as the width of the $b \to s\gamma$ decay, can sometimes be used to set better, but more model-dependent, limits than those obtained from the direct searches, as is pointed out in Papers 1 and 4.

The effective supersymmetry breaking mass terms, which give the spartners a mass consistent with the experimental limits, must thus be at least of the order of the electroweak scale if the supersymmetric model is to remain a viable alternative. The tree-level scalar potential of a supersymmetric theory can be written in terms of the D and F terms as

$$V = \frac{1}{2} D_a D^a + F_i F^i, \tag{3.2}$$

particle	lower limit on mass
chargino (mixtures gaugino and Higgsino)	45–150 GeV
neutralino (mixtures gauginos and Higgsinos)	23–83 GeV
selectron or smuon	84–89 GeV
stau	71 GeV
sneutrino	43 GeV
stop	87–90 GeV
gluino	180–190 GeV
squarks	230–260 GeV

Table 3.1: Experimental limits on the masses of supersymmetric particles [25].

where the indices a denote the gauge indices and i the chiral superfields of the model. It can be shown from general arguments that if supersymmetry is broken dynamically, the vacuum expectation value of the potential is lifted, $\langle V \rangle > 0$, as was discussed in Section 2.3. Therefore, one must have either a non-vanishing F-term or D-term in the minimum of the scalar potential. The F-term supersymmetry breaking (i.e., $\langle F_i \rangle \neq 0$) does not alter the supertrace sum rule (3.1) at tree level [68, 69]. The supertrace sum rule is not violated by the D-term supersymmetry breaking [70] ($\langle D_a \rangle \neq 0$), either, if the related gauge symmetry is anomaly-free [67, 71]. Consequently, if the supersymmetry is broken spontaneously at tree level, then in any reasonable model there should exist light supersymmetric particles, in contradiction with the experimental results.

The tree level supersymmetry breaking is thus unacceptable from the phenomenological point of view. However, the supertrace rule (3.1) is in general modified by radiative corrections. In a typical scenario the model can be divided into a hidden and a visible sector. In a typical scenario the visible sector of the model is the MSSM or some GUT model. The supersymmetry is broken in the hidden sector. By hidden sector one means the part of the model that has not been directly observed by the current experiments, because it contains fields that are either too heavy or too weakly coupled. When the hidden sector is integrated out, it is assumed to create, however, supersymmetry breaking mass terms of the order of 1 TeV to the Lagrangian.

3.1 Effective theory: boundary conditions

As discussed in Section 2.2, one of the main motivations for introducing the supersymmetry is that in supersymmetric theories the quadratic divergences of scalar masses are canceled. In a non-supersymmetric theory the quadratic divergences would push all scalar masses naturally close to the Planck scale, M_{Pl}. The cancellation of quadratic divergences is one example of the so-called non-renormalization theorems, which state that the superpotential of an $N = 1$ supersymmetric theory is renormalized, except by finite amounts, only by wave function renormalization [72].

The mass terms of scalars are included in the superpotential part of the Lagrangian, which can be written as

$$\mathcal{L}_W = \int d^2\theta W(\Phi) + \text{HC}. \tag{3.3}$$

The non-renormalization theorem can be proven heuristically by using superspace formalism [37]: The radiative corrections to the effective action can always be written as an integral over $d^4\theta$ with no superspace delta-functions. On the other hand, $\mathcal{L}_W$ can be written as a $d^4\theta$ integral only with superspace delta-functions. Consequently, the scalar mass terms appearing in the superpotential are not renormalized. If this property is not spoiled, the breaking of supersymmetry preserves naturalness. Terms that break the supersymmetry while preserving the naturalness of the theory are called soft supersymmetry breaking terms.

One can add to the Lagrangian (2.31) of a supersymmetric model involving chiral superfields Φ_k and general gauge superfields V the following super-renormalizable terms [73]:

$$\begin{aligned}\mathcal{L}_{\cancel{SUSY}} &= \int d^4 U_{kl}\theta\Phi_k^\dagger e^{2gt_k^i V_i}\Phi_l + \left[\int d^2\theta\left(\frac{1}{64}NW^\alpha W_\alpha + S_j W_j(\Phi_k)\right) + \text{HC}\right] \\ &= -\frac{1}{2}\tilde{m}_{kl}^2\phi_k^\dagger\phi_l + \frac{1}{2}\tilde{M}\lambda\lambda + A_j W_j(\phi_k) + \text{HC},\end{aligned} \tag{3.4}$$

where the following set of dimension-zero spurions have been introduced:

$$U_{kl} = \tilde{m}_{kl}^2\theta^2\overline{\theta}^2, \qquad N = \tilde{M}\theta^2, \qquad S_j = A_j\theta^2. \tag{3.5}$$

The W_j denote the various terms contained in the superpotential of the model, $W = \sum_j W_j$.

The mass terms arising from the supersymmetry breaking and generated by spurions (3.5) can be thought as constant background fields [74]. They preserve the naturalness of the theory.

If there are no gauge singlet matter fields, then also the following dimension-three terms preserve naturalness [75]:

$$\mathcal{L}'_{\not SUSY} = R_{klm}\phi_k^\dagger\phi_l\phi_m + \frac{1}{2}m_{Fij}\Psi_i\Psi_j + m_{Aia}\Psi_i\lambda_a + \text{HC}. \qquad (3.6)$$

The term proportional to m_A is possible only if there exist chiral superfields in the adjoint representation of the gauge group, and the term proportional to m_F can always be rotated away by a re-definition of the superpotential and the soft supersymmetry breaking terms in $\mathcal{L}_{\not SUSY}$. Most supersymmetry breaking scenarios create an insignificantly small trilinear term R_{klm}. Therefore, in phenomenological studies the soft supersymmetry breaking terms contained in $\mathcal{L}'_{\not SUSY}$ are usually ignored. This is the case also in the present work.

The outline of a typical phenomenological study of supersymmetry breaking is as follows. The low-energy effective theory is an $N = 1$ supersymmetric model, like the MSSM or SUSYLR model. The soft supersymmetry breaking terms are generated by some mechanism, for example via some supergravity model or a model with gauge mediated supersymmetry breaking, at some pre-defined scale. The effects of the hidden sector (where supersymmetry is broken) are integrated out at a chosen scale, leading to some pattern of the values of the soft supersymmetry breaking couplings. They result in a set of parameters at the electroweak scale. One can then investigate the low-energy phenomena of electroweak interactions, set limits on the model parameters and study the manifestations of the model in the current and planned particle physics experiments.

3.2 Gravity mediated supersymmetry breaking

In the minimal supergravity scenario supersymmetry is broken at the scale $M_S \sim 10^{11}$ GeV in a hidden sector that has *only* gravitational interactions with the visible sector, i.e. with the particles of the MSSM. The breaking of supersymmetry is transmitted by gravity to the visible sector resulting in supersymmetry breaking effects of order of $M_{SUSY} \sim g_{gravity}M_S^2 \sim 1$ TeV, where $g_{gravity} = M_{Pl}^{-1}$.

When supersymmetry transformations are made local, one ends up with non-renormalizable theory of quantum gravity, the so-called supergravity theory. The gravitational interactions are transmitted by a supergravity multiplet, which contains a spin-2 graviton and a spin-$\frac{3}{2}$ gravitino. The supergravity Lagrangian contains a unique combination of the scalar fields, called

a Kähler potential G [76]:

$$G\left(\phi^{\dagger},\phi\right)=J\left(\phi^{\dagger},\phi\right)+\ln\frac{|W|^{2}}{M_{Pl}^{6}}, \tag{3.7}$$

where J is in the minimal case just a quadratic sum of fields [1],

$$J\left(\phi_{i}^{\dagger},\phi\right)=M_{Pl}^{-2}\phi_{i}\phi^{i*}. \tag{3.8}$$

The most general globally supersymmetric Lagrangian for chiral and vector superfields is in the limit of a flat space and small gravitational coupling given by

$$\mathcal{L}=\int d^{4}\theta\varphi^{\dagger}\varphi K\left(\Phi^{\dagger}e^{2gV},\Phi\right)+\left[\int d^{2}\theta\left(\varphi^{3}W(\Phi)+\varphi^{0}f_{ab}\left(\Phi\right)W_{a}^{\alpha}W_{\alpha b}\right)+\mathrm{HC}\right], \tag{3.9}$$

where φ is a so-called Weyl compensator to be discussed below, and the function K contains the contribution of the Kähler potential:

$$K=-3M_{Pl}^{2}e^{-\frac{1}{3}J}. \tag{3.10}$$

Since gravitational theory is not renormalizable, the function K and the superpotential W do generally contain non-renormalizable terms suppressed by factors of M_{Pl}^{-1}.

In the simplest case the hidden sector, where the supersymmetry is broken spontaneously, and the visible sector have mutually only gravitational interactions. The simplest way to accomplish this is to take the superpotential to be a sum of hidden and visible sector superfields [78, 79, 47], i.e.

$$W=h(z_{i})+g(y_{i}), \tag{3.11}$$

where z_i and y_i denote the chiral superfields of the hidden sector and the visible sector, respectively. In our discussion it is sufficient that just one hidden sector field, say $z_1\equiv z$, achieves a non-vanishing VEV and F-term:

$$\begin{aligned}\langle z\rangle &= aM_{Pl},\\ \left\langle\frac{\partial h}{\partial z}\right\rangle &= bM_{SUSY}M_{Pl},\\ \langle h\rangle &= M_{SUSY}M_{Pl}^{2}.\end{aligned} \tag{3.12}$$

[1] J can be always shifted to the minimal form by a suitable super-Weyl transformation [77].

Here a and b are dimensionless parameters of the order of unity, and $M_{SUSY} \sim 1$ TeV will eventually be the scale characteristic of supersymmetry breaking effects in the visible sector.

As a result of supersymmetry breaking the spin- $\frac{3}{2}$ gravitino obtains the mass

$$m_{3/2} = e^{\frac{1}{2}\langle G\rangle} M_{Pl} \simeq e^{\frac{1}{2}|a|^2} M_{SUSY}. \tag{3.13}$$

The graviton remains massless.

The Weyl compensator φ appearing in the global supergravity Lagrangian (3.9) is a non-dynamical chiral superfield [80, 37, 81, 77]. It is introduced to make the action manifestly invariant under Weyl rescaling. The Weyl rescaling of the metric is applied in order to (re-)normalize the gravitational coupling to $G_N = (8\pi M_{Pl}^2)^{-1}$ and to separate out the scalar auxiliary field of the supergravity multiplet. The gravitational effects lead also to non-standard kinetic terms of the type $(1 + \mathcal{O}(1)\langle J\rangle)\,\partial^\mu\phi\partial_\mu\phi^*$ for the scalar components of chiral superfields in the action (3.9). The Weyl rescaling is used to transform the gravitational coupling to the measured value and the the vacuum expectation value of the function J to zero, restoring the correct form of the kinetic terms.

The action (3.9) and the Kähler potential G given in Equation (3.7) are invariant under Weyl rescaling. The other parameters fields transform as

$$\begin{aligned} J &\rightarrow J - \tau - \tau^\dagger, \\ W &\rightarrow e^{\tau} W, \\ \varphi &\rightarrow e^{-\tau/3}\varphi, \end{aligned} \tag{3.14}$$

where τ is an arbitrary chiral superfield. All explicit mass scales M appearing in the Lagrangian are rescaled by $M \rightarrow e^{-\tau/3}M$. The super-Weyl symmetry is broken explicitly if the Weyl compensator has a non-vanishing vacuum expectation value, $\langle\varphi\rangle \neq 0$. The actual value of the VEV is irrelevant, since any non-vanishing VEV can be transformed to any non-vanishing chiral function by a properly chosen super-Weyl transformation. In particular, we can fix $\varphi = 1$ and then perform the super-Weyl transformation defined by $\tau = \frac{1}{2}|a|^2$, which disentangles the effect of supergravity from the matter Lagrangian.

The tree-level effective potential can be written as [47]

$$V = e^{\phi^{i*}\phi_i/M_{Pl}^2}\left(\left|\frac{\partial W}{\partial\phi_i} + M_{Pl}^{-2}\phi^{i*}W\right|^2 - 3M_{Pl}^{-2}|W|^2\right) + (\mathrm{D}-\text{terms}). \tag{3.15}$$

Using the Weyl-rescaled superpotential $\hat{W} \simeq e^{|a|^2/2}W = m_i^2 y_i + \frac{1}{2}\mu_{ij}y_i y_j + \frac{1}{3}\lambda_{ijk}y_i y_j y_k + \ldots$ the potential of the visible sector is in the leading order in

M_{SUSY}/M_{Pl} given by

$$V = \left|\frac{\partial \hat{W}}{\partial y_i}\right|^2 + V_{\not{SUSY}} + (\mathrm{D}-\mathrm{terms}) \tag{3.16}$$

where

$$\begin{aligned} V_{\not{SUSY}} &= m_{3/2}^2 \phi^{i*}\phi_i \\ &\quad + \left(\frac{1}{3} A\lambda_{ijk} m_{3/2}\phi_i\phi_j\phi_k + \mathrm{HC}\right) \\ &\quad + \left(\frac{1}{2} B\mu_{ij} m_{3/2}\phi_i\phi_j + \mathrm{HC}\right) \\ &\quad + \left(C m_i^2 m_{3/2}\phi_i + \mathrm{HC}\right) + \ldots \quad . \end{aligned} \tag{3.17}$$

The non-renormalizable terms involving the superfields of the visible sector have been left out from Equation (3.17), since they are suppressed by factors of M_{Pl}^{-1}. The values of the parameters A, B and C depend on the details of the hidden sector. For the model described above they are given by $A = (a + b^*)a^*$, $B = A-1$ and $C = A-2$, where a and b are defined in Equation (3.12). They may contain complex phases originating from the hidden sector. These phases can lead to observable CP-violating effects in the visible sector [82].

The soft supersymmetry breaking terms due to gravity are universal, as in the Equation (3.17), *if* the hidden sector fields z_i are singlets under the gauge group of the visible sector. Non-singlet contributions to the soft supersymmetry breaking terms would lead to mass patterns similar to the non-universal gaugino masses discussed in the next section.

Another visible sector effect generated by the minimal supergravity model is the mass term $\tilde{M}$ for the gauginos:

$$\tilde{M} = \frac{1}{2} e^{\frac{1}{2}\langle G\rangle} \left\langle \frac{\partial f_{ab}^*}{\partial z_i} \right\rangle \left\langle \frac{\partial G}{\partial z_i} \right\rangle M_{Pl}^3. \tag{3.18}$$

Gauginos will have mass terms if the gauge kinetic term is non-minimal ($f_{ab} \sim (\mathcal{O}(1) z_i/M_{Pl} + \ldots)\delta_{ab}$) and one of the quantities $\langle \partial G/\partial z_i\rangle \sim M_{Pl}^{-1}$ is non-vanishing. The gaugino mass term is then naturally of the order of the gravitino mass $\tilde{M} \sim m_{3/2}$.

The minimal supergravity model predicts universal soft mass-squared terms, gaugino mass terms and trilinear supersymmetry breaking terms. They are all expected to be of the order of gravitino mass $m_{3/2} \sim 1$ TeV. The universality is broken by radiative corrections as the supersymmetry breaking parameters are run down from the Planck scale or the GUT scale down to the electroweak scale.

The minimal supergravity model provides a simple and working scenario for the breaking of supersymmetry. It is a standard framework for analyzing the phenomenology of supersymmetry.

There is also a quantum contribution to the supersymmetry breaking masses, namely the conformal anomaly, which generates gaugino masses at one-loop level and mass-squared terms of scalars at two-loop level [83, 84]. This contribution is present in all hidden sector models, but in most models the tree level contributions are dominant, at least as far as the scalar mass terms are concerned. The conformal anomaly arises if the Weyl compensator needed to disentangle the supergravity from the Lagrangian has an auxiliary component H, i.e.

$$\varphi = 1 - H\theta^2. \tag{3.19}$$

In a classical treatment the auxiliary component can be rotated away, at least for the terms involving no explicit mass terms in the supersymmetric Lagrangian. However, when the Lagrangian (3.9) is properly regularized, it turns out that the anomaly-mediated contributions to the gaugino mass $\tilde{M}$, scalar mass $\tilde{m}^2$ and trilinear scalar coupling A are, schematically,

$$\begin{aligned} \tilde{M} &= \frac{\alpha}{4\pi}\beta H, \\ \tilde{m}^2 &= \frac{1}{2}\frac{d\gamma}{d\ln\mu}H^2, \\ A &= -\sum \lambda\gamma H. \end{aligned} \tag{3.20}$$

Here β is the beta function related to the gauge coupling α, γ is the anomalous dimension, defined as a derivative of the wave function renormalization Z, $\gamma = -\frac{1}{2}d\ln Z/d\ln\mu$, and λ is the Yukawa coupling.

The most obvious problem with the pure anomaly-mediated supersymmetry breaking is that the mass-squared term for a scalar transforming under an infrared-free gauge group is negative, at least if the Yukawa couplings are small. Examples of such scalars are the sleptons of the MSSM. One way to solve this problem is to add a positive universal contribution to the squared scalar masses.

The most distinctive signature of the models with anomaly mediated supersymmetry breaking in collider experiments would be light mass-degenerate winos and almost mass-degenerate same-flavor sleptons [85]. The gravitinos are typically quite heavy, since the supersymmetry breaking masses are suppressed by loop effects relative to the gravitino mass.

3.2.1 Non-universal gaugino masses and gauge couplings in $SU(5)$ GUT models

In the minimal supergravity model gaugino masses and gauge couplings unify at the grand unification scale, as was discussed in the previous section. There are, however, several viable supersymmetry breaking mechanisms that would lead to non-universal gaugino masses and gauge couplings at the GUT scale. One possible mechanism of that kind, studied in Papers 4 and 5 of this thesis, is the one where gaugino masses arise from a condensation of the F-component of a chiral superfield in one of the representations **24**, **75** or **200** in an $SU(5)$ grand unified theory (GUT) [86, 87, 88].

The gauge kinetic function $f_{\alpha\beta}$ is generally given by

$$\begin{aligned}\mathcal{L} &= \frac{1}{64}\int d^2\theta f_{\alpha\beta}(X)W^\alpha W^\beta + \mathrm{HC} \\ &= -\frac{1}{4}\sum_{\alpha\beta}\mathrm{Re}f_{\alpha\beta}(S)F^\alpha_{\mu\nu}F^{\beta\mu\nu} \\ &\quad -\left(\sum_{\alpha\beta\alpha'\beta'}\frac{1}{4}F^I_{\alpha'\beta'}\frac{\partial f_{\alpha\beta}(S)}{\partial S^I_{\alpha'\beta'}}\lambda^\alpha\lambda^\beta + \mathrm{HC}\right) + \ldots,\end{aligned} \tag{3.21}$$

where $\langle X\rangle = S + F\theta^2$. The superfield X (and its VEV) has been decomposed into a singlet part X^S (S^S and F^S) and non-singlet parts $X^{\mathbf{N}}$ ($S^{\mathbf{N}}$ and $F^{\mathbf{N}}$), so that $X = \sum_I X^I$. The gauge kinetic function can then be decomposed as

$$f_{\alpha\beta}(X) = f_0(X^S)\delta_{\alpha\beta} + \sum_N \xi_N(X^S)\frac{X^N_{\alpha\beta}}{M_{Pl}} + \mathcal{O}\left[\left(\frac{X^N_{\alpha\beta}}{M_{Pl}}\right)^2\right], \tag{3.22}$$

where f_0 and ξ_N are functions of the gauge singlet superfield X^S. The symmetric product of the product of two adjoint representations of $SU(5)$, like the expression $W^\alpha W^\beta$ in Equation (3.21), is decomposed as follows:

$$(\mathbf{24}\times\mathbf{24})_s = \mathbf{1} + \mathbf{24} + \mathbf{75} + \mathbf{200}. \tag{3.23}$$

Thus, the non-singlet representations X^N of the chiral superfield X allowed as a linear term in the gauge kinetic function $f_{\alpha\beta}$ in Equation (3.21) are **24**, **75** and **200**.

In the Papers 4 and 5, where the phenomenology of non-universal gaugino masses is investigated, two basic assumptions have been made. First, the supersymmetry is assumed to be broken by the F-components of X, i.e. $F^I = \mathcal{O}(m_{3/2}M_{Pl})$, where $m_{3/2}$ is the gravitino mass. The second assumption

is that the GUT symmetry is broken at the GUT scale down to the Standard Model gauge symmetry by nonzero VEVs $S^{\mathbf{N}}$ of the non-singlet scalar field.

The unification condition for the gauge couplings at the GUT scale M_{GUT} is given by

$$\frac{\alpha_a^{-1}(M_{GUT})}{4\pi}\delta_{ab} = \langle \mathrm{Re} f_{ab} \rangle. \tag{3.24}$$

The correction to the universality of the gauge couplings ($\delta\alpha_a^{-1} = \alpha_a^{-1} - \alpha_{GUT}^{-1}$) caused by the VEVs $S^{\mathbf{N}}$ of the non-singlet fields can be expressed as

$$\begin{pmatrix} \delta\alpha_1^{-1}(M_{GUT}) \\ \delta\alpha_2^{-1}(M_{GUT}) \\ \delta\alpha_3^{-1}(M_{GUT}) \end{pmatrix} = \alpha_{GUT}^{-1}(M_{GUT}) \begin{pmatrix} -\frac{1}{2\sqrt{15}} & -\frac{3}{2\sqrt{15}} & \frac{2}{2\sqrt{15}} \\ -\frac{5}{6} & \frac{1}{2} & \frac{1}{6} \\ \frac{10}{2\sqrt{21}} & \frac{2}{2\sqrt{21}} & \frac{1}{2\sqrt{21}} \end{pmatrix} \begin{pmatrix} x_{\mathbf{24}} \\ x_{\mathbf{75}} \\ x_{\mathbf{200}} \end{pmatrix}. \tag{3.25}$$

The $x_{\mathbf{N}}$'s parametrize the VEVs of the Higgs fields in specific directions,

$$x_{\mathbf{N}} \simeq \frac{\sqrt{2}\xi_N\left(S^S\right)}{f_0\left(S^S\right)}\frac{S^{\mathbf{N}}}{M_{Pl}}. \tag{3.26}$$

Their magnitude is supposed to be of the order $S^{\mathbf{N}}/M_{Pl} \lesssim M_{GUT}/M_{Pl} \sim 1/100$ or less.

The corrections can be large enough to push the $SU(5)$ unification close to the Planck scale, as shown in Figure 3.1 taken from Paper 5. The $SU(5)$ GUT models are known to suffer from the fact that they typically predict the nucleon to decay faster than is allowed by the experimental data [25]. With the increased unification scale the nucleon decay width falls, however, below the experimental bounds [89, 90, 91, 92, 93, 94].

The gaugino mass $\tilde{M}_a$ at the unification scale, derived from Equation (3.21), is given by

$$\tilde{M}_a(M_{GUT})\delta_{ab} = \sum_{Ia'b'} \frac{F^I_{a'b'}}{2\mathrm{Re} f_{ab}\left(S\right)}\frac{\partial f_{ab}\left(S\right)}{\partial S^I_{a'b'}}. \tag{3.27}$$

If the supersymmetry breaking is dominated by one of the F-components the relative masses of gauginos obey relations presented in Table 3.2.

The particle spectrum of models with non-universal gaugino masses and large $\tan\beta$ are analyzed in Paper 4 of this thesis. The experimental bounds on the branching ratio of the decay $b \to s\gamma$ and the requirement that there is a successful electroweak symmetry breaking and the unification of the Yukawa couplings of the b quark and τ lepton were used to constrain the allowed range of parameters. The model where $F_{\mathbf{24}}$ dominates the gaugino mass generation favors the bino as the LSP. In models where the $F_{\mathbf{75}}$ and $F_{\mathbf{200}}$ are dominant the lightest neutralino and the lightest chargino are almost degenerate in

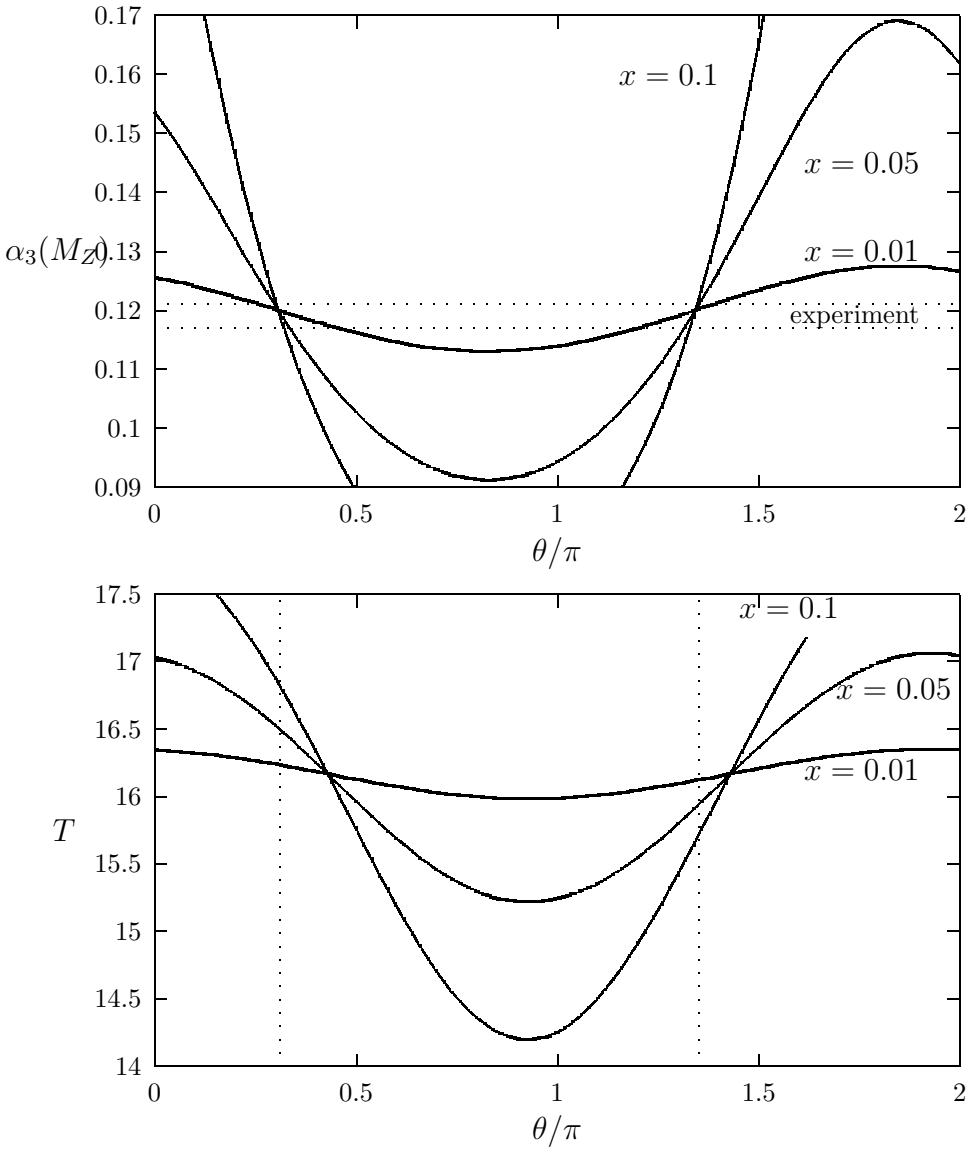


Figure 3.1: The color gauge coupling $\alpha_3(M_Z)$ and the unification scale $T = \log_{10} M_{GUT}$ [GeV] of the gauge couplings of the MSSM as functions of θ/π, where the non-universal contribution to the gauge couplings at the GUT scale is assumed to originate from a linear combination to the direction $\langle \mathbf{24} \rangle \sin\theta + \langle \mathbf{75} \rangle \cos\theta$. The plots are given for three different values of $x^2 = x_{\mathbf{24}}^2 + x_{\mathbf{75}}^2$, where $x_{\mathbf{24}}$ and $x_{\mathbf{75}}$ are defined in Equation (3.25). There exists a solution at $\tan\theta \simeq 1.4$, denoted by dotted vertical lines, that allows a large unification scale M_{GUT}, while at the same time the correction to $\alpha_3(M_Z)$ is very small due to cancellations between **24** and **75** contributions.

F_X	M_1^{GUT}	M_2^{GUT}	M_3^{GUT}	M_1^{mz}	M_2^{mz}	M_3^{mz}
1	1	1	1	0.4	0.8	2.9
24	-0.5	-1.5	1	-0.2	-1.2	2.9
75	-5	3	1	-2.1	2.5	2.9
200	10	2	1	4.1	1.6	2.9

Table 3.2: Gaugino masses in the case where one of the F-components dominates the supersymmetry breaking. The gluino mass at the GUT scale is normalized to unity.

mass. This would be a challenging situation from the experimental point of view, since a chargino decaying into soft pions and an invisible neutralino would be quite difficult to detect [95, 96, 97].

3.3 Gauge mediated supersymmetry breaking

If the effects of supersymmetry breaking are transmitted from the hidden sector to the visible sector by gravitational interactions, the supersymmetry breaking scale M_S must be, by dimensional arguments, of the order of $M_S \sim 10^{11}$ GeV. If the supersymmetry breaking interactions are transmitted by gauge interactions, with couplings of order $\alpha_{\text{gauge}} \sim 0.01$, then the supersymmetry breaking scale can be considerably lower, $M_{SUSY} \sim \alpha_{gauge} M_S$, i.e. $M_S \simeq 100$ TeV.

In the so-called gauge mediated supersymmetry breaking (GMSB) models [34, 98, 99, 100, 101, 102, 103, 104, 105, 106] the supersymmetry is broken in a hidden sector, resulting in a non-vanishing auxiliary term of some MSSM singlet chiral superfield X:

$$\langle X \rangle = S + F\theta^2. \tag{3.28}$$

This singlet field X is assumed to couple to the messenger fields Φ_i and their conjugate fields $\overline{\Phi_i}$ (or to an adjoint messenger field Q) via a Yukawa interaction:

$$\mathcal{L} = \int d^2\theta X \left(\lambda_i \Phi_i \overline{\Phi_i} + \frac{1}{2} \lambda_j Q_j^2 \right) + \text{HC}. \tag{3.29}$$

The messenger fields have, by definition, non-vanishing quantum numbers under the gauge symmetry of the visible sector.

The spinor components of Φ_i and $\overline{\Phi_i}$ form Dirac fermions with masses $m_{Di} = \lambda_i S$, while the non-vanishing F-term splits the mass of the scalar

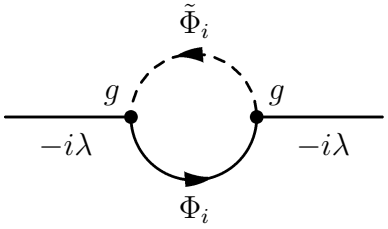
components into two non-degenerate values. The scalar components have the following mass-squared matrix

$$V_{mass} = \begin{pmatrix} \Phi_i^\dagger & \overline{\Phi}_i \end{pmatrix} \begin{pmatrix} |\lambda_i S|^2 & \lambda_i^* F^* \\ \lambda_i F & |\lambda_i S|^2 \end{pmatrix} \begin{pmatrix} \Phi_i \\ \overline{\Phi}_i{}^\dagger \end{pmatrix}, \tag{3.30}$$

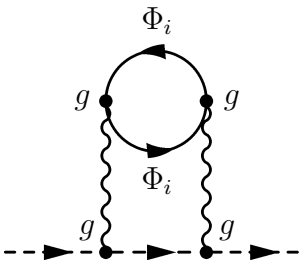
with mass eigenvalues

$$m_{Si}^2 = m_{Di}^2 \left(1 \pm \left|\frac{F}{\lambda_i S^2}\right|\right). \tag{3.31}$$

The splitting of the scalar masses breaks the supersymmetry explicitly. The breaking of supersymmetry is transmitted to the visible sector by gauge interactions. The gaugino masses are induced at one-loop level and the scalar mass-squared terms at two-loop level through the following diagrams [107, 108]:

$\tilde{M} =$ 

(3.32)

$\tilde{m}_k^2 =$ 

$+ \ldots$. (3.33)

The loop contributions (3.32) and (3.33) can be calculated, most easily, from the wave-function renormalization [109].

The renormalization group equation of the gauge coupling α is

$$\frac{d}{d\ln\mu}\alpha^{-1} = \frac{1}{2\pi}\beta, \tag{3.34}$$

where the one-loop beta-function coefficient β is given by $\beta = 3C - S_G$, C is the Casimir index for the gauge group, and S_G is the sum of the Dynkin indices of all chiral superfields.

The particle content of the model changes at the messenger mass scale $\Lambda_M = \lambda S$, where the messenger fields decouple. The beta-function β can thus be expressed as:

$$\beta(\mu) = \begin{cases} \beta_{MSSM} - N & , \quad \mu > \Lambda_M \\ \beta_{MSSM} & , \quad \mu < \Lambda_M \end{cases} , \tag{3.35}$$

where β_{MSSM} is the beta-function of the model, without the contribution of the messenger fields.

Using Equations (3.34) and (3.35) and the expression (3.24) the gauge kinetic function can be written as

$$\mathrm{Re}f(X,\mu) = \frac{\alpha^{-1}(\Lambda_{UV})}{4\pi} + \frac{\beta_{MSSM} - N}{16\pi^2} \ln\frac{\lambda X}{\Lambda_{UV}} + \frac{\beta_{MSSM}}{16\pi^2} \ln\frac{\mu}{\lambda X}, \tag{3.36}$$

where $\Lambda_{UV} > \Lambda_M$ is some constant scale. The gaugino mass at the messenger mass scale is then given by Equation (3.27):[2]

$$\tilde{M} = \frac{F}{2\mathrm{Re}f(S)} \frac{\partial f(S)}{\partial S} = \frac{\alpha}{4\pi} N \frac{F}{S}. \tag{3.37}$$

The contribution to scalar mass terms can be calculated analogously starting from the wave function renormalization term of the chiral superfield Q:

$$\mathcal{L} = \int d^4\theta Z_Q\left(X, X^\dagger\right) Q^\dagger Q. \tag{3.38}$$

The mass-squared terms $\tilde{m}_k^2$ of the scalar fields are

$$\tilde{m}_k^2 = 2C_k \frac{\alpha^2}{16\pi^2} N \left|\frac{F}{S}\right|^2 = 2C_k \frac{\tilde{M}^2}{N}, \tag{3.39}$$

where C_k is the quadratic Casimir invariant of the representation of the chiral superfield in question. The supersymmetry breaking contributions to the gaugino and scalar mass terms are thus of the same order of magnitude.

There is no contribution to the trilinear A-terms at the messenger mass scale. Nevertheless, these terms obtain a non-vanishing values between the messenger mass scale and the electroweak scale due to radiative corrections.

The gravitino mass is proportional to the product of the gravitational coupling constant and the F-term that breaks the supersymmetry. In the GMSB model the gravitino mass is thus typically of the order of 1 eV, making the gravitino the lightest supersymmetric particle in that model. The

[2] It has been assumed that $x_i = \left|F/\lambda_i S^2\right| \ll 1$. The mass formulas (3.37) and (3.39) are accurate to the level of one per cent if $x_i \lesssim 0.85$. [108]

gravitino cannot be heavier that a few keV in order for its relic density not to over-close the universe [110]. A light gravitino would couple extremely weakly to the gauge or matter fields, the coupling being suppressed by factors of M_{Pl}^{-1}. If a pair of gravitinos is produced in a collider as a result of a LSP neutralino decay into a gravitino and photon, it would show up as missing energy.

The mass of a gaugino is proportional to the corresponding gauge coupling constant and to the sum of the Dynkin indices of the messenger fields. The exact mass spectrum of gauginos depends on the choice of the messenger fields. The bino, the supersymmetric partner of the $U(1)_Y$ gauge boson, is usually the lightest gaugino, and it is never heavier than gluinos due to the relatively small hypercharge gauge coupling and suppressing group theoretical factors.

The soft mass-squared terms of the scalars are of the order of the squared masses of gauginos multiplied by some group theoretical factors. Roughly speaking, the more gauge quantum numbers a particular chiral superfield has, the heavier is its scalar component. The left-handed squarks are the heaviest supersymmetric particles. The masses of the right-handed squarks are always more or less equal to the masses of gluinos. The left-handed sleptons and the right-handed slepton are typically lighter than squarks. The lightest of the sleptons is often the supersymmetric partner of the right-handed tau lepton, the stau. This is because the non-diagonal element of the 2×2 mass matrix of staus has a contribution which is proportional to the Yukawa coupling of the corresponding lepton of the family in question. Since the Yukawa couplings are the largest in the third family, the mass of the lighter of the two stau mass eigenstates is smaller than the mass of the light selectron or smuon. Consequently, the lightest supersymmetric particle is in most scenarios either stau or possibly some of the neutralinos.

All mass-squared terms evaluated at the messenger scale are according to Equation (3.39) non-negative. However, the heavy top quarks contribute to the renormalization group evolution of the mass-squared terms $m_{H_k}^2$ of the Higgs scalars. In physically viable scenarios making it negative near the electroweak scale as it should be in any physically viable scenario.

The so-called SUSY flavor problem arises when the flavor-non-universal masses of squark and slepton masses cause unacceptably large flavor violations for example in the $\mu \to e\gamma$ decay or in the $K - \overline{K}$ system. The flavor violations would be in control if the relative non-universality of the masses of squarks and sleptons is at most of the order of one per mille [111].

The GMSB contributions to the squark and slepton masses are the same for all families, because the gauge interactions are flavor-universal. Due to the relatively low value of supersymmetry breaking in the GMSB theories

(as compared for example with the SUGRA GUT models) the universality is maintained as mass parameters evolve from the messenger scale to the electroweak scale. The only non-universality originates from the Yukawa couplings. The contribution related to gravity is of the order of M_S^2/M_P, and it is less than one per mille, as required, if the supersymmetry breaking scale in the hidden sector is of the order of $M_S \sim 10^{15}$ GeV or less.

Another attractive feature of the GMSB theories is that they do not introduce any extra complex phases in addition to the one present in the non-supersymmetric Standard Model. There can be a phase in the VEV of the singlet field $\langle X \rangle$ that contributes directly to the gaugino mass terms. Nevertheless, *if* there is only one such field — as is the case in the simplest models — this extra phase can be rotated away and it is thus not observable in the visible sector.

The set of messengers is *a priori* unrestricted. In the so-called minimal GMSB model [104, 105, 106] the messenger fields transform according to the representation $\mathbf{5} + \overline{\mathbf{5}}$ of $SU(5)$. The representation $\mathbf{5}$ decomposes under the MSSM gauge group as $\mathbf{5} = L + D$, where the L denotes a field that has the same gauge quantum numbers as the lepton doublet and D the field with gauge quantum numbers common with the right-handed down-quarks. In the minimal model the Yukawa couplings related to the messenger multiplet $\mathbf{5}$ are required to unify at the GUT scale, i.e., $\lambda_D(M_{GUT}) = \lambda_L(M_{GUT})$.

In Paper 1 and in [112] the following general set of messenger fields is considered:

$$\begin{aligned}
n_Q : &\qquad Q + \overline{Q} = (\mathbf{3}, \mathbf{2}, \frac{1}{6}) + \text{conj.}, \\
n_U : &\qquad U + \overline{U} = (\overline{\mathbf{3}}, \mathbf{1}, -\frac{2}{3}) + \text{conj.}, \\
n_D : &\qquad D + \overline{D} = (\overline{\mathbf{3}}, \mathbf{1}, \frac{1}{3}) + \text{conj.}, \\
n_L : &\qquad L + \overline{L} = (\mathbf{1}, \mathbf{2}, -\frac{1}{2}) + \text{conj.}, \\
n_E : &\qquad E + \overline{E} = (\mathbf{1}, \mathbf{1}, 1) + \text{conj.},
\end{aligned} \tag{3.40}$$

where the multiplicities of the messenger fields are denoted by n_Q, n_U, n_D, n_L and n_E and the numbers in parenthesis denote their $SU(3)_C \times SU(2)_L \times U(1)_Y$ quantum numbers.

One can set *lower limits* on the multiplicities of various messenger fields by requiring that all gauginos obtain a non-vanishing mass. Given fixed gauge couplings at the electroweak scale, the messenger fields increase the value of gauge couplings at the GUT scale. One can thus set *upper limits* on the number of messengers by requiring that the gauge couplings remain

perturbative up to some high scale, like the Planck scale. There are 53 sets of messenger multiplicities that satisfy these requirements. Some of these sets are redundant, since they result in a similar pattern of supersymmetry breaking mass terms. Taking this redundancy into account one is left with 32 different possibilities.[3]

In Paper 1 the radiative symmetry breaking and the width of the radiative decay $b \to s\gamma$ are analyzed. The GMSB mechanism creates neither a non-vanishing bilinear supersymmetry breaking term, nor a bilinear term in the superpotential. If the bilinear supersymmetry breaking term is assumed to vanish at the messenger mass scale, the particle spectrum and $\tan\beta$ parameter of the model can be fixed for given messengers multiplicities, supersymmetry breaking scale and sign of the mu-term.

Papers 3 and 4 of this thesis analyze the gauge mediated supersymmetry breaking mechanism in the framework of the $SO(10)$ GUT with an intermediate SUSYLR symmetry. Two fundamental scales were assumed to exist in the theory, the scale of Grand Unification near the Planck scale and the messenger scale near the electroweak scale. At the intermediate energies the theory is assumed to be the SUSYLR model with the messenger fields added. Below the messenger scale the model is effectively the MSSM.

The requirement of radiative symmetry breaking, perturbativity up to the GUT scale, unification of gauge couplings scale and the experimental lower limit on the mass of the stau scalar leads to a very restricted set of messenger multiplicities. This makes the scenario a predictive alternative to the models motivated by the MSSM. All models have a NLSP that is either stau or neutralino, the charged slepton being also light (the LSP is the gravitino). Squarks and gluinos have mass of 600 GeV or more, and the rest of the supersymmetric and Higgs particles have a mass between the mass of the electroweak gauge bosons and that of squarks.

[3]The messenger multiplicities could be further constrained by requiring that they form a complete representation of some GUT group, like $SU(5)$ or $SO(10)$, or that the gauge couplings unify at some scale.

Chapter 4

R-parity symmetry

4.1 Baryon and lepton numbers

In the Standard Model the conservation of baryon (B) and lepton (L) numbers are accidental symmetries. Given the gauge symmetry and the particle content, the model does not allow renormalizable interactions that violated either the baryon or lepton number. In the MSSM the additional squark, slepton and Higgsino fields make the renormalizable baryon and lepton number violating interactions possible. It is assumed that the baryon and lepton number violations are suppressed by the conservation of the so-called R-parity [113, 90, 114]. The R-parity is a multiplicative quantum number defined by $R = (-1)^{3(B-L)+2S}$ where B and L are baryon and lepton numbers of respective fields and S is their spin.

Nevertheless, there is no *a priori* reason for R-parity to be conserved. The R-parity can in fact be violated, as long as the R-parity breaking couplings are sufficiently small, so that their effects have not yet been seen in experiments. There are two kinds of breaking mechanisms of the R-parity, explicit and spontaneous.

In the *spontaneous* R-parity violation the lepton number is broken due to a non-vanishing vacuum expectation value of a neutral scalar field, such as sneutrino, that carries a lepton number. The spontaneous R-parity violation conserves the baryon number, since a non-vanishing VEV of a field carrying color quantum number would violate the $SU(3)_C$ gauge symmetry. In MSSM all scalar fields carrying color charge are electrically charged: spontaneous violation of baryon number at tree level would thus also violate the conservation of $U(1)_{em}$ electromagnetic gauge symmetry.

If the full Lagrangian of the model contains R-parity violating terms that cannot be expressed in terms of VEVs by a suitable rotation of superfields,

then the R-parity is said to be broken *explicitly*.

The gauge symmetry allows the following renormalizable R-parity violating terms in the superpotential of MSSM:

$$W_{\not R} = W_{\not L} + W_{\not B}, \tag{4.1}$$

where

$$W_{\not L} = \lambda_{ijk} E_i L_j^T i\tau_2 L_k + \lambda'_{ijk} D_i Q_j^T i\tau_2 L_k + \epsilon_k L_k H_u \tag{4.2}$$

and

$$W_{\not B} = \lambda''_{ijk} U_i D_j D_k. \tag{4.3}$$

The terms appearing in $W_{\not L}$ violate lepton number, while $W_{\not B}$ violates the baryon number. Since existence of both baryon and lepton number violating couplings would lead to a rapid nucleon decay [90, 89, 115], one must assume that either baryon or lepton number violating terms vanish to a high accuracy ($W_{\not B} = 0$ or $W_{\not L} = 0$).

The trilinear R-parity violating Yukawa couplings can be constrained using the results of low energy experiments, for example in the Paper 6 of this thesis bounds were derived on the products of the λ'-type couplings using experimental results on the mass difference in the $K - \overline{K}$ and $B_d - \overline{B}_d$ systems.

There are also the corresponding R-parity violating soft supersymmetry breaking contributions to the scalar potential:

$$V_{\not L} = A_{ijk} E_i L_j^T i\tau_2 L_k + A'_{ijk} D_i Q_j^T i\tau_2 L_k + B_k L_k H_u \tag{4.4}$$

and

$$V_{\not B} = A''_{ijk} U_i D_j D_k. \tag{4.5}$$

The lepton doublets L_i and the Higgs doublet H_d have the same quantum numbers. If the lepton number is broken ($W_{\not L}, V_{\not L} \neq 0$) the Higgs and lepton fields are no longer unambiguous, and they can be re-defined by unitary rotation of $\overline{L} = (H_d L_1 L_2 L_3)^T$.

4.2 Gauged R parity symmetry

The R-parity symmetry is a continuous global $U(1)$ symmetry, where each field has a $U(1)$ quantum number proportional to the $B-L$ quantum number. If the global symmetry is broken spontaneously, a massless Goldstone scalar, a majoron (J), should exist [116, 117, 118]. This implies a novel decay channel $h \to JJ$ for the Higgs scalar h. The sneutrino acquiring a VEV must be a $SU(2)_L \times U(1)_Y$ singlet, i. e. a right-handed sneutrino, since

otherwise the majoron would contribute to the invisible decay width of the Z boson.

A more elegant solution to the proton decay problem is to have the R-parity as a local, instead of a global, symmetry. The R-parity is then a discrete subgroup of the gauge symmetry. The massless Goldstone mode is absorbed into the longitudinal polarization mode of some extra neutral gauge boson. This is the case for example in the SUSYLR models, where the R-parity symmetry is a subgroup of $U(1)_{B-L}$ gauge symmetry. The gauge symmetries are local, and thus protected from gravitational corrections, making them an attractive alternative for *ad hoc* global symmetries.

4.2.1 R-parity breaking SUSYLR models

The supersymmetric left-right (SUSYLR) model contains three left-handed neutrinos and three right-handed neutrinos and their supersymmetric partners, sneutrinos. The sneutrinos can obtain a non-vanishing vacuum expectation value in the process of spontaneous symmetry breaking. In some versions of the SUSYLR it is in fact inevitable that at least one of the right-handed sneutrinos obtains a non-vanishing VEV [59, 119, 120, 121]. The R-parity breaking in the SUSYLR models has been investigated in Paper 7 of this thesis.

The spontaneous R-parity breaking in the SUSYLR model takes effect solely through the vacuum expectation value of the sneutrino field. The R-parity violation is manifested in the fact that the mass matrix of neutralinos is mixed with the neutrino mass matrix, and the chargino mass matrix is mixed with the mass matrix of charged lepton. Similarly, the scalar mass eigenstates are mixtures of the Higgs and slepton fields. The R-parity violating mixing terms in the fermion and scalar matrices are proportional to the VEV's of the sneutrino field.

The VEV's of the left-handed sneutrinos, $\langle \tilde{\nu}_L \rangle = \sigma_{Lk}$, contribute to the masses of the electroweak gauge bosons and light neutrinos. Taking the gauge boson masses into account and requiring that the Yukawa coupling of the top quark remains perturbative sets an upper limit on the VEVs: $\sum_k |\sigma_{Lk}|^2 \lesssim (168 \text{ GeV})^2$. Requiring the neutrino masses to remain below their current experimental limits constrains the VEV to be at most of the order of $|\sigma_{Le}| \lesssim$ MeV for the electron sneutrino, of the order of $|\sigma_{L\mu}| \lesssim$ GeV for the muon sneutrino and of the order of $|\sigma_{L\tau}| \lesssim 10$ GeV for the tau sneutrino, barring cancellations in the neutrino/neutralino mass matrix.

From the phenomenological point of view the breaking of the R-parity via a non-vanishing VEV of a left-handed sneutrino is very similar to the bilinear R-parity breaking in the MSSM. In fact, the VEV of a left-handed

sneutrino can always be rotated into a bilinear R-parity breaking parameter and trilinear R-parity breaking terms in Equation (4.2) that are proportional to the corresponding Yukawa couplings [122, 123].

The VEV's of the right-handed sneutrinos, $\langle\tilde{\nu}_R\rangle = \sigma_R$, can be of the order of the mass of the right-handed gauge bosons. As a result of the non-vanishing VEVs of the right-handed sneutrino, $\sigma_R \neq 0$, the lepton mass eigenstates can be mixtures of gauginos, Higgsinos and lepton interaction eigenstates. The couplings of physical leptons, for example, to the gauge bosons of the Standard Model remain however to leading order unmodified, if the mixing is due to VEVs that are singlets under the SM gauge group. Furthermore, in the decoupling limit the couplings of charged leptons to the physical Higgs boson approach their Standard Model values.

In the SUSYLR model, in the limit where $\sigma_L = 0$ and the contribution of the gaugino mass terms to the neutrino masses is dominant, the neutrino obtains via a see-saw mixing with gauginos a Majorana mass

$$m_\nu \simeq \frac{m_D^2}{\tilde{M}}, \tag{4.6}$$

where $\tilde{M}$ is an effective gaugino mass term. The quantity m_D^2 is given by

$$m_D^2 = \frac{\lambda_\nu^2 \sigma_R^2 m_{Z_L}^2 \cos^2\beta}{\lambda_\nu^2 \sigma_R^2 + \mu^2}, \tag{4.7}$$

where λ_ν is the Yukawa coupling of the neutrino and μ is the mass parameter appearing in the superpotential.

The R-parity violating couplings are determined by the spontaneous symmetry breaking. In $SU(2)_R \times U(1)_{I_{3R}} \times U(1)_{B-L}$ model the effective R-parity breaking tri-linear couplings are proportional to the Yukawa matrices of down-type quarks and leptons. They are naturally suppressed below the current experimental limit. However, the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ SUSYLR-model has more interaction eigenstates sharing the hypercharge quantum number with the right-handed charged lepton. In Paper 7 it was found that it is possible to have a large R-parity breaking coupling that is not suppressed either by neutrino mass constraints or small Yukawa couplings. A potentially large R-parity violating coupling between an up- and down type quark, squark, and a charged lepton, proportional to the $SU(2)_R$ gauge coupling constant g_R, is due to the $SU(2)_R$ wino component in the right-handed part of the charged lepton mass eigenstate. It is given by

$$\mathcal{L} = -g_R x_k \tilde{d}_{Ri} \overline{u_i^c} P_R e_k + HC, \tag{4.8}$$

where $|x_k| \leq 1$, $k = e, \mu, \tau$, are dimensionless parameters and e_k is the lepton mass eigenstate. The parameter x_e is bound by constraints from neutrinoless

double beta decay, but in the case of heavy lepton families ($k = \mu, \tau$) the value of the parameter x_k can be of the order of unity. The coupling is the same for all quark families ($i = 1, 2, 3$) due to the universality of gauge couplings.

An advantage of the spontaneous R-parity violation realized in the SUSYLR model is that the baryon number is automatically protected by gauge symmetry, unlike in the MSSM. The R-parity breaking couplings can be parametrized in terms of few parameters, and the pattern of lepton-number violating couplings, if observed, will be distinctive signatures of the SUSYLR models.

Chapter 5

Concluding remarks

During the past ten years our understanding of the fundamental theories of nature has taken great leaps forward. The Standard Model of particle physics has become the most accurately measured physical theory. Nevertheless, the quest is far from complete. The Higgs sector of the Standard Model remains experimentally untested. Furthermore, there are theoretical arguments, like the hierarchy problem, that suggest that the Standard Model will not by far be the last word, the Theory of Everything.

The supersymmetry provides a logical solution to the hierarchy problem. Unfortunately, the supersymmetry brings with it over a hundred new parameters, reducing the predictability of the model significantly. The supersymmetry also brings problems: how is the supersymmetry broken at the electroweak scale and are baryon and lepton numbers conserved?

Most of the new parameters introduced by supersymmetric models are associated with the supersymmetry breaking. By assuming a particular supersymmetry breaking mechanism one can reduce the number of free parameters from a hundred to just a few. The model becomes easier to deal with and more predictive, and it can be used as a test case in designing new experiments and predicting what the new physics would look like, and ultimately give hints of the Theory of Everything, the theory that explains all particle interactions, including gravity. In fact, the consequences of supersymmetry breaking might be the *only* direct consequence of quantum gravity theory that we are able to measure in a foreseeable future.

Another mystery relates to the apparent conservation of baryon and lepton numbers in nature, which is most manifestly apparent in the stability of the proton. The natural situation in many models is, however, that lepton and baryon number are broken, not conserved. In particular, in the minimal supersymmetric model the lepton and baryon number (or R-parity) violating interactions are *a priori* unrestricted. Therefore one has to introduce some

mechanism that makes these interactions small in order to comply with the experimental data. One such mechanism is given by the supersymmetric left-right (SUSYLR) model, where the R-parity is conserved by the gauge symmetry and the R-parity violation, if it takes place, is spontaneous. The pattern of R-parity breaking couplings is unique to the spontaneous symmetry breaking mechanism, as the resulting spectrum of particle masses is a distinctive feature of the supersymmetry breaking mechanism. Observing the consequences of these couplings at relatively low energies can thus give information on the physics at the energies beyond the direct reach of the experiments.

The nature may be more complex (or simple) than we have thought. But if something along the lines speculated in these studies is observed in the current or planned accelerator or other particle physics experiments one has here a well thought-out framework and tools for studying nature.

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