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Higgs Bosons as Probes of Nonminimal Supersymmetric Models

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ACADEMIC DISSERTATION

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Abstract

The experiments at the Large Hadron Collider (LHC) have confirmed that the Standard Model (SM) is a good description of particle physics at the electroweak scale. The Standard Model is still incomplete, since it does not explain *e.g.* neutrino masses, dark matter, dark energy or gravity.

Supersymmetry is a well motivated way to extend the Standard Model. The minimal supersymmetric Standard Model (MSSM) has become somewhat fine-tuned after the first run of the LHC and therefore the detailed study of non-minimal supersymmetric models is highly motivated.

We studied phenomenological implications of some nonminimal supersymmetric models especially in the light of the recent discovery of a Higgs boson. Our studies focused on supersymmetry without R-parity and left-right symmetric supersymmetric models.

In the MSSM a 125 GeV Higgs requires rather heavy superpartners. In nonminimal models the Higgs mass can be lifted by contributions from new particles at tree-level, loop-level or by mixing effects. We found that if we introduce spontaneous R-parity violation, the mixing between the SM-like Higgs and a right-handed sneutrino can increase the mass of the SM-like Higgs if the sneutrino-like state is lighter than 125 GeV. One does not need as heavy superpartners as in the MSSM and thus fine-tuning is not as severe.

If the Higgs mass gets additional contributions, the squarks of the third generation can be more easily within the reach of the LHC. If R-parity is not imposed, the squarks can have new decay modes, which can have a large branching fraction. As an example we studied a model, where R-charges are identified with the lepton number and found that the discovery potential for the $\tilde{t} \to be^+$ mode is well beyond 1 TeV squark masses.

In left-right supersymmetry we studied the Higgs decay modes and the option of having a right-handed sneutrino as a dark matter candidate. We found that a loop-induced mixing of the bidoublets can either enhance or suppress the Higgs coupling to bottom quarks and thus change the signal strengths considerably. However, in the scan there was also a large number of points, where the couplings behaved close to those of the SM.

The right-handed sneutrino is a part of a doublet in left-right symmetric models. We found that sneutrinos may annihilate via a D-term coupling to the Higgs and produce the observed relic density. If we assume the gauge coupling of the right-handed gauge interactions to be the same as for the left-handed ones, we were able to predict a range of masses for the sneutrino. We also showed that with 100 fb^{-1} we could get a signal of superpartners from the sleptonic decays of the right-handed W_R-boson, if its mass is below 3 TeV.

Tiivistelmä

Large Hadron Collider -törmäyttimellä (LHC) suoritetut kokeet ovat osoittaneet, että hiukkasfysiikan standardimalli on hyvä kuvaus aineen rakenteesta ja vuorovaikutuksista nykyisillä kiihdytinenergioilla. Standardimalli ei kuitenkaan voi olla lopullinen hiukkasfysiikan teoria, sillä se ei selitä mm. neutriinojen massoja, pimeää ainetta, pimeää energiaa tai painovoimaa.

Standardimallissa alkeishiukkasten massat syntyvät ns. Higgsin mekanismin avulla. Siinä Higgsin kentän arvo on nollasta eroava, kun systeemin energia on pienin. Tätä kentän arvoa kutsutaan tyhjiöodotusarvoksi. Higgsin kentällä on myös vastaava hiukkanen, ns. Higgsin bosoni. Higgsin kentän tyhjiöodotusarvo antaa muille hiukkasille massan, joka on verrannollinen kyseisen hiukkasen ja Higgsin bosonin välisen vuorovaikutuksen suuruuteen. Kun Higgsin kentällä on tyhjiöodotusarvo, standardimallin liikeyhtälöiden ratkaisut eivät ole teorian mittasymmetrian mukaisia, vaikka itse yhtälöt ovat symmetriset — tätä kutsutaan mittasymmetrian spontaaniksi rikkoutumiseksi.

Supersymmetria on yksi eniten tutkituista tavoista laajentaa standardimallia. Supersymmetria liittää jokaiseen hiukkaseen superpartnerin, jonka ominaisuudet ovat muuten samanlaiset kuin alkuperäisellä hiukkasella, mutta spin eroaa puolella yksiköllä. Supersymmetrisissä malleissa on aina useampi Higgsin bosoni ja osalla uusista Higgsin bosoneista on sähkövaraus. LHC:n ensimmäisten vuosien tulokset ovat poissulkeneet suuren osan yksinkertaisimman supersymmetrisen mallin (minimaalinen supersymmetrinen standardimalli, MSSM) parametriavaruudesta. Tämän johdosta ei-minimaalisten supersymmetristen mallien tutkimus on perusteltua.

Tutkimme tässä väitöskirjassa supersymmetristen mallien fenomenologiaa erityisesti vuonna 2012 löydetyn Higgsin bosonin ominaisuuksien pohjalta. Erityisesti tarkastelimme malleja, joissa ei ole R–pariteettia sekä vasen-oikea-symmetristä mallia.

MSSM:ssa kevyimmän Higgsin bosonin massa ilman kvanttikorjauksia voi olla korkeintaan Z-bosonin massan verran. Löydetyn Higgsin bosonin massa on tätä suurempi ja niin suuri, että tarvittavat kvanttikorjaukset ovat varsin suuria ja erityisesti top-kvarkin superpartnerin tulisi olla noin $1,5$ TeV:n painoinen tai raskaampi. Ei-minimaalisissa supersymmetrisissä malleissa Higgsin bosonin massaraja voi olla korkeampi kuin MSSM:ssä tai uusien hiukkasten aiheuttamat kvanttikorjaukset voivat nostaa Higgsin massaa, jolloin superpartnereiden ei tarvitse olla yhtä raskaita kuin MSSM:ssa.

R-pariteetti on MSSM:ssa lisäoletus, joka kieltää baryoni- tai leptonilukua muuttavat vuorovaikutukset. Jos R-pariteettia rikkovat vuorovaikutukset olisivat sallittuja, protonit hajoaisivat. R-pariteetin säilymislaki myös takaa, että kevyin superpartneri ei hajoa ja voisi muodostaa pimeän aineen. R-pariteetti voi rikkoutua spontaanisti, jos jokin sneutriinoista saa tyhjiöodotusarvon. Spontaani R-pariteetin rikko synnyttää vain leptonilukua muuttavia vuorovaikutuksia, joten protonit eivät hajoa.

Kun R-pariteetti rikkoutuu spontaanisti, sneutriinon säilyvät kvanttiluvut ovat samat kuin Higgsin bosonilla, joten ne voivat sekoittua. Havaitsimme, että

jos sneutriino on kevyempi kuin Higgsin bosoni, tämä sekoittuminen voi nostaa Higgsin bosonin massaa ja näin ei tarvita yhtä suuria kvanttikorjauksia kuin MSSM:ssä. Higgsin bosonin ja sneutriinon sekoittuminen johtaisi myös Higgsin bosonin tuottotodennäköisyyden pienenemiseen.

Jos Higgsin bosonin havaittu massa saadaan pienemmillä kvanttikorjauksilla kuin MSSM:ssä, voivat stop- ja sbottom-skvarkit olla kevyempiä. Jos ei oleteta R-pariteettia, voivat skvarkit hajota eri tavalla kuin yleensä supersymmetrisissä malleissa oletetaan ja tälläisten hajoamisten osuus voi olla suuri. Tästä esimerkkinä tutkimme mallia, jossa R-symmetrian varaus samaistetaan leptoniluvun kanssa. Mallissa kvarkkien superpartnereilla on leptoniluku ja ne voivat hajota kvarkiksi ja leptoniksi. Stop-skvarkin pääasiallinen hajoamiskanava voi olla $\tilde{t} \to be^+$. Tutkimuksemme perusteella LHC pystyy löytämään stop-skvarkin tässä mallissa, vaikka sen massa olisi yli 1 TeV:n.

Heikoissa vuorovaikutuksissa pariteettisymmetria rikkoutuu: Beetahajoamisessa syntyvien hiukkasten spinit ovat vasenkätisiä ja antihiukkasten oikeakätisiä. Vasen-oikea-symmetrisissä malleissa tämä selitetään siten, että on olemassa myös toinen (oikeakätinen) heikko vuorovaikutus, jossa hiukkasten ja antihiukkasten spinit ovat päinvastaiset, mutta tämä vuorovaikutus on spontaanin symmetriarikon seurauksena paljon tunnettua (vasenkätistä) heikkoa vuorovaikutusta heikompi.

Vasen-oikea symmetrisen mallin osalta tutkimme Higgsin bosonin hajoamissuhteita sekä oikeakätistä sneutriinoa pimeän aineen kandidaattina. Kvanttikorjaukset aiheuttavat mallin Higgsin bosoneille sekoittumisen, jossa kevyimmän Higgsin bosonin hajoaminen b-kvarkkipariksi voi poiketa huomattavasti standardimallin ennusteesta. Toisaalta osassa datapisteistä hajoamissuhteet ovat lähellä standardimallin ennustetta.

Vasen-oikea symmetrisessä mallissa oikeakätinen neutriino ja sen superpartneri vuorovaikuttavat oikeakätisten heikkojen vuorovaikutusten kautta toisin kuin malleissa, joissa oikeakätisiä heikkoja vuorovaikutuksia ei ole. Tämä saa aikaan sen, että jos sneutriino on kevyin supersymmetrinen hiukkanen, niiden annihiloituminen varhaisessa maailmankaikkeudessa voi olla niin voimakasta, että jäljelle jää havaittu määrä pimeää ainetta. Jos oikeakätisten heikkojen vuorovaikutusten voimakkuus tunnetaan, havaitusta pimeän aineen määrästä voidaan ennustaa kevyimman sneutriinon massa. Osoitimme myös, että jos oikeakätisten vuorovaikutusten W-bosoni on kevyempi kuin 3 TeV, sen hajoamiset leptonien superpartnereiksi voivat antaa signaalin supersymmetriasta.

Acknowledgements

The completion of this thesis is a certain kind of an endpoint, though this will not mean that I would quit doing particle physics. It has been a privilege to work in the field during the time when new territories have been explored with the Large Hadron Collider.

First I wish to thank my supervisor, professor Katri Huitu for the guidance throughout the whole work. We have had long discussions on physics and other matters. I have learned a lot about the importance of collaborating, being cautious on interpreting experimental data, presenting your thoughts clearly and finding the physics behind the results.

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Chapter 1

Introduction

1.1 The era of LHC and the need for new physics

The previous few years have been exceptional in particle physics since we have witnessed the exploration of a new energy scale at the Large Hadron Collider (LHC). The proton-proton experiments at center-of-mass energies of 7 TeV and 8 TeV led to the discovery of a new boson [1,2], whose properties are compatible with the Standard Model Higgs boson[1]. In addition, the experiments have given new, more stringent bounds on many experimental observables. Essentially all experimental data from the LHC is consistent with the Standard Model.

Although the Standard Model (SM) has been extremely successful, there is still room for new physics and a need for it. The LHC has been able to exclude only a couple of models completely. The Higgs data exclude the extension of the Standard Model by a fourth fermion generation and the fermiophobic Higgs model [3,4]. Also a few minimal supersymmetric models can be considered to be excluded since they do not allow a heavy enough Higgs boson without extremely heavy supersymmetric partners [5].

The most direct evidence for the incompleteness of the Standard Model comes from neutrino experiments, where the existence of non-zero neutrino masses has been confirmed during the last two decades [6–10]. Majorana neutrino masses can be included into the Standard Model by introducing a $d = 5$ operator [11]. This operator can be a remnant of a seesaw mechanism [12–17] after the heavy particles have been integrated out. The absolute mass scale, the mass hierarchy and the Dirac or Majorana nature of neutrinos still remain to be solved experimentally.

There is also indirect evidence that requires extensions of the Standard

[1]Various authors can be attributed to the symmetry breaking mechanism and the associated scalar particle, both commonly carrying the name of Peter Higgs. In addition to Higgs, at least Philip Anderson, Robert Brout, Francois Englert, Jeffrey Goldstone, Gerald Guralnik, Carl Hagen, Tom Kibble, Yoichiro Nambu, Abdus Salam and Steven Weinberg were involved in building the mechanism of electroweak symmetry breaking. For brevity, the particle will be called Higgs boson in this thesis.

Model. The power energy spectrum of the cosmic microwave background [18] can be explained with the so called ΛCDM-model [19–22], where in addition to baryonic matter there is cold dark matter and dark energy. The existence of dark matter can also be inferred from the measurements of galactic velocities [23, 24] and the existence of dark energy is supported by the accelerated expansion of the Universe [25–28]. In the Standard Model there are no candidates for dark matter nor is there an explanation for dark energy.

The Sakharov conditions [29] for matter-antimatter asymmetry require the violation of C- and CP-symmetries. In the Standard Model the charge conjugation symmetry is broken by weak interactions [30] but the only CP-violating effects come from the single phase of the CKM matrix [31, 32]. So far all observed CP-violating reactions [33–37] have been explained with this single parameter [38, 39] but the matter-antimatter asymmetry of the Universe is harder to explain [40, 41].

The value of the Higgs boson mass also gives reasons to expect new physics. First one of the Sakharov conditions requires a departure from thermal equilibrium. However in the SM there is no first order phase transition with a Higgs mass of 125 GeV [42, 43] and hence electroweak baryogenesis is not possible without extending the SM. Second the stability of the Standard Model scalar potential depends on the Higgs boson mass. The running of the quartic coupling turns it negative at an energy scale of the order of 10^{10} GeV [44–46]. Hence we live in a metastable vacuum close to the border where the vacuum becomes too short-lived or unstable. Hence any extension should improve the stability of the scalar potential. Since we are not far from the stable region, even some of the simplest extensions can make the vacuum stable [47–49].

The Standard Model has only one dimensionful parameter, the Higgs boson mass term. Since scalar masses are not protected by any symmetry[2], they get quantum corrections from several terms of the Lagrangian [50]. In cutoff regularization the correction diverges quadratically, *i.e.* the quantum correction to the mass is proportional to the cutoff scale. In the renormalization procedure one fixes the mass to the observed value by adding a counterterm. If the cutoff scale is taken from the other known dimensionful quantity, the gravitational constant, we end up to the Planck scale, which is 16 orders of magnitude larger than the Higgs boson mass. Hence the counterterm and the quantum corrections need to match with an enormous precision and still leave a nonzero result. In addition, the procedure is not stable against radiative corrections [51, 52].

In the Standard Model this fine-tuning can be thought of an artifact of the regularization procedure (see, *e.g.* [53]) since there are no other explicit mass scales than the electroweak scale. In dimensional regularization there are no quadratic divergences. The problem arises when one uses any model with new mass scales to explain *e.g.* neutrino masses or grand unification. In such a case the question arises how can a large hierarchy between mass scales be maintained

[2] Fermion masses are protected by chiral symmetry, gauge boson masses by gauge symmetry. This means that the quantum corrections can come only from the terms that break this symmetry and hence the correction is always proportional to the mass itself. By simple dimensional analysis this means that the corrections depend logarithmically on the scale.

against quantum corrections.

There are also some hints from the precision frontier. The measured value of the anomalous magnetic moment of the muon [54] is in slight tension with the best theoretical computations [55–58]. The difference cannot yet be considered decisive. The next generation of experiments should reduce the experimental uncertainty and if the theoretical computation will improve as much, the difference could become significant.

Flavor physics has provided some small hints of new physics and on the other hand very tight constraints on it. The largest deviations from theoretical expectations come in the decays of B-mesons to $D^{(*)}\tau\nu$ [59–62]. The difference from SM predictions is at most at the level of 3σ. In any case, if the errors are estimated correctly, measuring a large number of observables should produce deviations larger than 3σ in about 0.3% of the measurements and hence definite conclusions cannot be made yet.

The most constraining are the results on the branching ratios of $B_{s,d} \to \mu^+\mu^-$. The LHCb collaboration has found a signal for the former with a significance of 4.0σ and the best fit of the branching ratio close to SM expectations, whereas there is an upper limit for the latter, not too much above the SM prediction [63–65]. Many scenarios beyond the SM can easily enhance these branching ratios by a factor of ten or more [66–68]. Hence these branching ratios constrain the parameter spaces of many models outside the reach of direct searches (*e.g.* for minimal supergravity, see [69]).

The LHC is expected to function for at least two decades with the center-of-mass energy of 14 TeV. The higher collision energy will make it possible to probe even higher mass scales, up to a few TeV for strongly interacting particles.

The expectations of finding new physics beyond a SM-like Higgs were based on the Standard Model not to be fine-tuned in the context of the larger theory. Hence there should be a mass scale not too far from the electroweak scale. The first run of LHC already makes some fine-tuning inevitable.

1.2 The Higgs boson as a portal to beyond the Standard Model

There are several ways to go beyond the Standard Model (BSM) without getting into conflict with existing experimental data. Some of them will be reviewed in the following chapters.

The energy scale below 100 GeV was explored thoroughly by the precision experiments at the Large Electron-Positron collider (LEP). The LEP data is in essential agreement with the Standard Model. The only known particles not studied by the LEP are the top quark and the Higgs boson. The top quark was discovered at the Tevatron [70, 71] but due to the smaller production cross section at 2 TeV collision energy and smaller luminosity the number of produced top quarks was rather limited compared to the first run of the LHC. Both top and Higgs can be produced copiously at the 14 TeV phase of the LHC and uncertainties related to their properties will be reduced during the next run of the LHC.

The anomalous decays of the top quark could give hints of new physics but so far nothing exceptional has been observed. For instance if we had a light charged Higgs, the top quark could decay via $t \to H^+ b$ but the upper limit for the branching ratio for such a decay is constrained to be at a level of one percent [72–76]. This is already enough to rule out a light charged Higgs in the Minimal Supersymmetric Standard Model (MSSM) for some values of the charged Higgs boson mass. There is no evidence of deviations from the SM in other decays either [77, 78].

The Higgs boson is a potential channel for studying new physics. It couples most strongly to heavy particles and hence the indirect effects of unknown particles may be seen in Higgs physics. In addition, as a scalar field, the Higgs field has the dimension of mass in natural units, whereas fermions have dimension mass$^{3/2}$. Since renormalizable terms have a dimension of mass to a power four or less, the Higgs scalar may have renormalizable couplings to unknown particles that a fermionic field cannot have.

The leading Higgs production mechanism at the LHC is gluon fusion, where the leading order is at one-loop level [79,80] and hence it is a potential probe for new strongly interacting particles. The rough agreement of the Higgs production with SM expectations rules out the fourth generation of quarks [3,4]. The major problems in this production mechanism are that it is difficult to tag so that a limited number of final state are possible to study[3] and the theoretical errors on the Standard Model prediction are quite large and somewhat uncertain [81, 82].

Among the various decay channels of the Higgs boson the decays to $\gamma\gamma$ and $Z\gamma$ are also one-loop processes at leading order [83, 84]. They may be mediated by any charged particle coupling to the Higgs. The interpretation of any deviation is not straightforward, since the contribution of a given particle depends on its mass, spin and charge — in addition to possible deviations coming from the theoretical uncertainties in the production cross section and other decay widths. In any case the one-loop mediated production and decay channels complement the flavor physics tests.

If there are new particles that share the good quantum numbers with the Higgs boson they may mix. This mixing will lead to altered couplings compared to SM predictions. The LHC can improve the accuracy of coupling measurements to below ten percent errors in the best channels [85] and that could be sensitive enough to imply a deviation from the SM.

There are two important Higgs couplings that have not been measured so far directly. The Yukawa coupling to top quarks is mainly responsible for Higgs production in gluon fusion, but its direct measurement requires observing the associated $pp \to t\bar{t}h + X$ production. This is an important cross-check of the consistency of the standard picture, since there could be unknown colored particles coupling to the Higgs. The other important thing to measure is the Higgs

[3] All final states consisting of only jets are hidden in QCD backgrounds, which are larger by roughly six orders of magnitude. Final states with leptons or photons are accessible.

self coupling. In the SM it is determined once the Higgs mass is known but the coupling could deviate from its standard value in BSM models. The four-point coupling does not seem to be experimentally accessible at the LHC [86], but there is a chance of measuring the three-point coupling [87–89].

1.3 Structure of the thesis

After this introductory chapter motivating the need for physics beyond the Standard Model we study in more detail the physics of the Higgs mechanism and extended scalar sectors in chapter 2. Chapter 3 is devoted to supersymmetry and its implications on the Higgs sector. In chapter 4 we discuss supersymmetry without R-parity and and especially R-symmetric models and spontaneous R-parity violation. Chapter 5 discusses left-right symmetric models and their supersymmetric extensions. In chapter 6 we summarize the results.

The idea in this thesis is to use the SM-like Higgs particle as a portal to study supersymmetric models. Supersymmetry remains as a well-motivated scenario beyond the Standard Model and can solve many of the problems of contemporary physics described in section 1.1. Minimal supersymmetric models are somewhat fine-tuned but nonminimal supersymmetric models can reduce the amount of fine-tuning, which motivates the study of these extensions in the view of the Higgs discovery and the LHC bounds on supersymmetry searches. Supersymmetry, even in its minimal realization, has an extended Higgs sector. Nonminimal supersymmetric models usually have a rather large Higgs sector. In each case the other particles of the model can leave fingerprints in the SM-like Higgs, which can then be observed in collider experiments.

1.3.1 List of publications and author's contribution

The publications included in this thesis are

I **Higgs sector of NMSSM with right-handed neutrinos and spontaneous R-parity violation**,
Katri Huitu and Harri Waltari, JHEP **1411** (2014) 053.

II **Left-right supersymmetry after the Higgs boson discovery**,
Mariana Frank, Dilip Kumar Ghosh, Katri Huitu, Santosh Kumar Rai, Ipsita Saha and Harri Waltari, Phys. Rev. **D90** (2014) 115021.

III **Light top squarks in $\mathbf{U(1)_R}$-lepton number model with a right-handed neutrino and the LHC**,
Sabyasachi Chakraborty, AseshKrishna Datta, Katri Huitu, Sourov Roy and Harri Waltari, Phys. Rev. **D93** (2016) 075005.

IV **Resonant slepton production and right sneutrino dark matter in left-right supersymmetry**,
Mariana Frank, Benjamin Fuks, Katri Huitu, Santosh Kumar Rai and Harri Waltari, accepted for publication in JHEP, arXiv 1702.02112.

Article **I** discusses the effects of spontaneous R-parity violation to the Higgs sector. The original idea was to look at bounds from lepton flavor violation, but after the Higgs discovery the focus turned more to the implications of spontaneous R-parity violation on the Higgs sector. The current author made all of the computations and figures and wrote the first draft of the manuscript, which was then jointly edited.

Article **II** discusses the implications of the Higgs bounds on left-right supersymmetry. The original idea was that the model has a light doubly charged Higgs, which could have an effect on the one-loop decays $h \to \gamma\gamma$ and $h \to Z\gamma$. We performed a thorough parameter scan of the model and discussed the implications. The current author was responsible of making the spectrum and coupling generator and participated in making the modifications to HIGLU. The text was written jointly, the current author was responsible for sections 3, 4.3, figures 1–2 and minor parts of sections 2 and 5.

Article **III** studies the top squarks in an R-symmetric model. The current author made some of the background analyses and wrote some parts of the text as well as participated in the editing of the text.

Article **IV** looks at the possibility of having right-handed sneutrino dark matter in the left-right supersymmetric model. We studied the possibility of producing right-handed sleptons through the decay of the W_R boson. The current author was responsible of producing the benchmarks and computing the constraints from the relic density. The text was written jointly, the current author being responsible for sections 3 and 4 and parts of section 2.

1.3.2 Conventions

Throughout this text the natural system of units is used, where $\hbar = c = 1$. Masses, momenta and energies are expressed in electronvolts. The zeroth component of four-vectors is timelike and the metric tensor is $g_{\mu\nu} = \mathrm{diag}(1, -1, -1, -1)$. Repeated indices are summed over. The Feynman slash means $\not{p} = p_\mu \gamma^\mu$. The vacuum expectation value of the SM Higgs field is 246 GeV and, unless otherwise stated, complex neutral scalar fields are of the form $\varphi^0(x) = \frac{1}{\sqrt{2}}(H(x) + iA(x) + v)$, where $H(x)$ and $A(x)$ are real scalar fields.

The SM Higgs boson or the SM-like Higgs boson[4] in extended models is denoted by h, other CP-even Higgs bosons by H_i and CP-odd Higgs bosons by A_i.

From chapter 3 onwards the spinors will be written by using the van der Waerden notation of dotted and undotted two-component spinors. We use $\sigma^\mu = (\mathbf{1}, \sigma^i)$ and $\bar{\sigma}^\mu = (\mathbf{1}, -\sigma^i)$ where σ^i are the Pauli matrices.

Missing transverse energy is denoted by $\not{E}_T$.

Unless otherwise stated, values of experimental quantities are from [90].

[4] The neutral CP-even SU(2) doublet scalar with the VEV closest to 246 GeV.

Chapter 2

Fundamentals of Higgs physics

2.1 Spontaneous symmetry breaking

In physics the concept of symmetry is somewhat different from its everyday use. In addition to everyday symmetries (of *e.g.* geometrical shape) we use the concept to describe the fact that equations of motion remain unchanged after some transformations. For instance, we might move a system to another place. If the equations of motion depend only on mutual distances the system will behave identically in the new place. This is called translational symmetry. Other space-time symmetries include invariance under rotations, reflections and time reversal.

In particle physics the most utilized symmetries are internal symmetries. The simplest example of such a symmetry is the fact that multiplying all the wave functions by a common phase does not produce any experimental consequences. When we allow the phase to be position-dependent and require invariance in local phase transformations, we need an additional field to compensate for additional terms from derivatives of the wave function. This field can be identified with the electromagnetic field.

A symmetry may be broken explicitly. For instance the flavor symmetry between u, d and s quarks is broken by their different masses or the isospin symmetry between the proton and the neutron is broken by their electric charges (and small mass difference). Even in this case there are some remnants of the symmetry left.

Spontaneous symmetry breaking means that the theory (*i.e.* the action) is symmetric but the solutions to the equations of motion are not. A common example is a ferromagnet. The equations of electromagnetism are rotationally invariant. When the magnet is in a magnetized phase, the direction of the magnetic field picks randomly one direction and hence the rotational invariance is broken. There is still left a symmetry in rotations around the axis pointing along the magnetic field.

The phenomenological description of spontaneous symmetry breaking is due to Ginzburg and Landau [91]. They constructed the free energy of a superconductor using a complex effective wave function $\Psi(x,t)$ to describe the superconducting electrons. The expression for free energy they used is of the form

$$F[\Psi,\vec{A}] = F_0 + \int dV \left[\alpha|\Psi|^2 + \beta|\Psi|^4 + \gamma|\vec{D}\Psi|^2 + \kappa(\nabla\times\vec{A})^2\right], \tag{2.1}$$

where $\vec{D} = \nabla - iq\vec{A}(x,t)$ is the covariant derivative. The functional is invariant under the following transformations: $\Psi(x,t) \to \Psi(x,t)e^{i\theta(x,t)}$, $\vec{A}(x,t) \to \vec{A}(x,t) + q^{-1}\nabla\theta(x,t)$.

The system is not stable unless β, γ and κ are positive. The minimum is obtained by varying the free energy with respect to Ψ and $\vec{A}(x)$. This gives the equations

$$\frac{\delta F}{\delta\Psi^*} = -\gamma(\nabla - iq\vec{A})^2\Psi + (\alpha + 2\beta|\Psi|^2)\Psi = 0, \tag{2.2}$$

$$\frac{\delta F}{\delta\vec{A}} = i\gamma(\Psi^*\nabla\Psi - \Psi\nabla\Psi^*) + 2\gamma q^2|\Psi|^2\vec{A} + 2\kappa\nabla\times\vec{B} = 0, \tag{2.3}$$

where $\vec{B} = \nabla\times\vec{A}$.

The coefficients depend on temperature. Spontaneous symmetry breaking occurs when at some temperature α turns negative. In that case the minimum of the free energy occurs with $|\Psi| \neq 0$, which means that the vacuum state is not invariant under the gauge transformations. This leads to a phase transition. From equation (2.3) we find that in the absence of any electric field the current density is $\vec{J} \propto \nabla\times\vec{B} = -\frac{\gamma}{2\kappa}[i(\Psi^*\nabla\Psi - \Psi\nabla\Psi^*) + 2q^2|\Psi|^2\vec{A}]$. When $|\Psi| \neq 0$ at the ground state there is a current without any voltage applied, *i.e.* the system is in a superconducting state.

Spontaneous symmetry breaking is always related to degenerate vacua. If the Hamiltonian is invariant under a transformation U, *i.e.* $[U,H] = 0$, but the ground state $|\Psi_0\rangle$ is not, also the state $U|\Psi_0\rangle$ is degenerate to the ground state, since $H(U|\Psi_0\rangle) = UH|\Psi_0\rangle = E_0(U|\Psi_0\rangle)$, where E_0 is the energy of the ground state. This has some consequences. If the broken symmetry is discrete, the breaking leads to domain walls [92,93], since the vacuum configuration may be different in different parts of space. If the symmetry is continuous, there is a long-range mode[1], the so called Goldstone mode, which consists of excitations along the continuous set of degenerate vacuum states. In particle physics this corresponds to a massless particle.

2.2 The Higgs mechanism in gauge theories

The idea of using spontaneous symmetry breaking in particle physics models was introduced by Nambu and Jona-Lasinio [94–96].

[1] Correlations between field variables are not damped exponentially as the distance increases.

In relativistic quantum field theories spontaneous breaking of a continuous symmetry leads to a massless scalar, the Goldstone boson. The theorem was conjectured by Goldstone [97] and proved by him, Salam and Weinberg [98]. The first one to conjecture a loophole in this theorem was Anderson [99] and the discussion was continued by Klein and Lee [100]. They studied nonrelativistic models where spontaneous symmetry breaking does not lead to massless bosons and argued that a relativistic analogue should exist. Thereafter Gilbert showed that the exception found by Klein and Lee requires a unit vector in the time direction [101]. In relativistic theories you do not have such a preferred Lorentz frame. Hence the Goldstone theorem would hold in the relativistic case.

Higgs noted in [102] that in gauge theories the gauge condition is usually not Lorentz-invariant. For instance in using the radiation gauge ($\nabla \cdot \vec{A} = 0$) or temporal gauge ($A^0 = 0$) a preferred vector in the time direction exists. This allows to use the loophole in the Goldstone theorem found by Klein and Lee for relativistic gauge field theories.

Englert and Brout [103] and, independently, Higgs [104] noted that the interaction of gauge fields with a scalar with a vacuum expectation value (VEV) gives a mass term for the gauge bosons. Guralnik, Hagen and Kibble [105] then continued the work of Brout, Englert and Higgs and noted that spontaneous symmetry breaking indeed generates masses for gauge bosons but introducing explicit mass terms makes the original theory manifestly covariant (since the freedom of choosing a gauge is gone) and hence the Goldstone bosons reappear.

The Higgs mechanism can be stated in its simplest (and original) form by considering a complex scalar field in an U(1) gauge theory. We assume the Lagrangian to be of the form

$$\mathcal{L} = (D^\mu \varphi(x))^* (D_\mu \varphi(x)) - \frac{1}{4} F^{\mu\nu} F_{\mu\nu} + m^2 |\varphi(x)|^2 - \frac{\lambda}{4} |\varphi(x)|^4, \tag{2.4}$$

where $D_\mu = \partial_\mu - ieA_\mu(x)$ is the covariant derivative and $F_{\mu\nu} = \partial_\mu A_\nu(x) - \partial_\nu A_\mu(x)$ is the field-strength tensor. The last two terms are (minus) the scalar potential and the signs have been chosen so that with real parameters there will be spontaneous symmetry breaking. The Lagrangian has a U(1) symmetry, where the fields transform as $\varphi(x) \to \varphi(x) e^{i\theta(x)}$ and $A_\mu(x) \to A_\mu(x) + \frac{1}{e}\partial_\mu \theta(x)$. A mass term for the gauge field would not be gauge invariant.

The kinetic term is minimized by choosing $\varphi(x)$ to be a constant. The minimum of the scalar potential is at $|\varphi(x)|^2 = 2m^2/\lambda$, whenever $m^2 > 0$. Next we shall expand the Lagrangian around the minimum of the potential, *i.e.* write $\varphi(x) = \frac{1}{\sqrt{2}}(v + h(x) + ia(x))$, where $v = 2m/\sqrt{\lambda}$ is the VEV of the scalar field. The vacuum configuration picks a direction and thus is not gauge invariant.

The new Lagrangian is (omitting constants and linear terms)

$$\begin{aligned}\mathcal{L} = &-\frac{1}{4} F^{\mu\nu} F_{\mu\nu} + \frac{1}{2}\partial^\mu h \partial_\mu h + \frac{1}{2}\partial^\mu a \partial_\mu a - eA^\mu (v\partial_\mu a + h\partial_\mu a - a\partial_\mu h) + \frac{e^2 v^2}{2} A^\mu A_\mu \\ &+ \frac{e^2}{2}(h^2 + 2vh + a^2) A^\mu A_\mu - \frac{\lambda}{2} v^2 h^2 - \frac{\lambda}{8} h^2 a^2 - \frac{\lambda v}{4}(ha^2 + h^3) - \frac{\lambda}{16}(h^4 + a^4).\end{aligned} \tag{2.5}$$

Particle	SU(3)	SU(2)	U(1)
Left-handed quark doublets, $Q_L = (u_L \quad d_L)^T$	$\mathbf{3}$	$\mathbf{2}$	1/3
Right-handed up-type quarks, u_R	$\mathbf{3}$	1	4/3
Right-handed down-type quarks, d_R	$\mathbf{3}$	1	−2/3
Left-handed lepton doublets, $L_L = (\nu_L \quad \ell_L^-)^T$	1	$\mathbf{2}$	−1
Right-handed charged leptons, ℓ_R^-	1	1	−2
Gluons	$\mathbf{8}$	1	0
$W^\pm$, W^0	1	$\mathbf{3}$	0
B (hypercharge gauge boson)	1	1	0
Higgs boson	1	$\mathbf{2}$	1

Table 2.1: The fields of the Standard Model in the gauge basis and the corresponding representations of the gauge groups.

There are several important things to notice. First the gauge field aquires a mass ev after symmetry breaking. On the other hand the imaginary part of the scalar field becomes massless, *i.e.* it is the Goldstone boson. The original Lagrangian conserves the U(1) charge but after spontaneous symmetry breaking the terms of the form $hA^\mu A_\mu$ and h^3 break charge conservation. All of these are typical consequences of spontaneous symmetry breaking and they vanish in the limit $v \to 0$. One may also note that after symmetry breaking the mass term for the field $h(x)$ has the correct sign.

The remarkable thing about gauge theories is that one can choose a gauge, by setting $\theta(x)$ equal to the negative of the phase of $\varphi(x)$, such that the Goldstone boson vanishes completely from the Lagrangian. This choice of gauge is called unitary gauge. It leaves only the physical degrees of freedom in the Lagrangian. The Goldstone boson then becomes the longitudinal polarization state of the massive gauge boson [105].

The generalization of these results to the non-Abelian case were first derived by Kibble [106].

2.3 Properties of the Standard Model Higgs

The Standard Model is based on a $\mathrm{SU(3)_C \times SU(2)_L \times U(1)_Y}$ gauge symmetry. In addition to the gauge fields there are three generations of elementary fermions and the Higgs boson. The particle content of the SM is summarized in Tables 2.1 and 2.2.

The Standard Model Lagrangian is the most general gauge and Lorentz invariant expression with at most dimension four terms. It can schematically be written in the form

$$\mathcal{L} = -\frac{1}{4} F_a^{\mu\nu} F_{\mu\nu a} + i\bar{\Psi}_b \not{D} \Psi_b + (Y_{cd}(\bar{\Psi}_{Lc} \cdot \Phi)\Psi_{Rd} + h.c.) + |D^\mu \Phi|^2 + \mu^2 |\Phi|^2 - \frac{\lambda}{4}|\Phi|^4. \tag{2.6}$$

Index a runs over the various generators of the SM gauge group, indices b, c and d over the particle species, L and R refer to left- and right-chiral fermion

Particle	Spin	Mass (GeV)	Charge
Up quark	1/2	0.0023(7)	+2/3
Charm quark	1/2	1.28(3)	+2/3
Top quark	1/2	160(5)	+2/3
Down quark	1/2	0.0048(5)	−1/3
Strange quark	1/2	0.095(5)	−1/3
Bottom quark	1/2	4.18(3)	−1/3
Electron	1/2	5.110×10^{-4}	−1
Muon	1/2	0.1057	−1
Tau	1/2	1.7768(2)	−1
Neutrinos (3 generations)	1/2	Small	0
$W^{\pm}$	1	80.39(2)	±1
Z^0	1	91.188(2)	0
Photon	1	0	0
Gluon	1	0	0
Higgs boson	0	125.1(3)	0

Table 2.2: The mass eigenstates of the SM particles. All fermions except neutrinos are Dirac fermions, *i.e.* consist of both left- and right-handed spinors. Quark masses are running masses in the $\overline{\text{MS}}$ scheme. The mass of the top quark from the kinematical reconstruction of the decay products is 173.2(9) GeV. Neutrinos are linear combinations of flavor eigenstates [107]. Neutrino mass differences have been measured to be $\mathcal{O}(0.1)$ eV and the sum of their masses is constrained by cosmological observations to be of the same order of magnitude. In the SM they are assumed massless. The numbers in parentheses are the errors of the last digit, the relative errors in electron and muon masses are less than 10^{-7}.

fields, $F^{\mu\nu}$ is the field strength tensor, D_μ is the gauge covariant derivative, Ψ denotes a fermion field and Φ the Higgs doublet.

The Higgs boson is the hero and the villain of the Standard Model. The Higgs mechanism allows to generate masses for particles and allows to extend the model without violations of unitarity to very high energy scales[2]. On the other hand it creates the naturalness problem discussed in section 1.1 and also there is no explanation for the hierarchies in the Higgs couplings to fermions, which break flavor symmetries.

2.3.1 Electroweak symmetry breaking

The idea of the spontaneous breaking of a unified gauge theory of electromagnetic and weak interactions was developed independently by Weinberg [109] and Salam [110]. They applied the work of Kibble [106] to the proposition of Glashow [111] of using SU(2)×U(1) as the gauge group for electromagnetic and weak interactions of leptons.

The Higgs mechanism must break SU(2)×U(1) down to $\mathrm{U}(1)_{em}$ and hence the Higgs field must transform non-trivially under the electroweak gauge group. The minimal solution is to use a SU(2) doublet with hypercharge[3] $Y = \pm 1$. In the Standard Model the option $Y = +1$ is chosen. Hence the Higgs field is a doublet

$$\Phi = \begin{pmatrix} \varphi^+ \\ \varphi^0 \end{pmatrix}$$

with a charged and a neutral component after symmetry breaking.

The part of the Lagrangian relevant for electroweak symmetry breaking (EWSB) is

$$\mathcal{L} = (D^\mu\Phi)^\dagger D_\mu\Phi - \sum_{SU(2),U(1)} \frac{1}{4} F_a^{\mu\nu} F_{\mu\nu a} - \mu^2 \Phi^\dagger\Phi - \frac{\lambda}{4}(\Phi^\dagger\Phi)^2, \tag{2.7}$$

where $D_\mu = \partial_\mu - ig\frac{\tau^a}{2}W_\mu^a - ig'\frac{Y}{2}B_\mu$ is the covariant derivative and τ^a are the Pauli matrices. In the case of SU(2) the field strength tensor has an additional term due to the non-Abelian nature of the group: $F_a^{\mu\nu} = \partial^\mu W_a^\nu - \partial^\nu W_a^\mu + g\epsilon_{abc}W_b^\mu W_c^\nu$, where ϵ_{abc} is the fully antisymmetric Levi-Civita tensor.

Spontaneous symmetry breaking occurs when $\mu^2 < 0$. The minimum of the scalar potential is then at $\Phi^\dagger\Phi = -2\mu^2/\lambda$. There is a continuum of degenerate vacua. Of these we assign the VEV to the real part of the neutral component so that electric charge is conserved.

There is one combination of generators, the sum of the third component of SU(2) and hypercharge, that is left unbroken. We assign it to the electric charge. With three broken generators we get three Goldstone bosons, which can

[2] There will be a Landau pole for the U(1) gauge coupling at around 10^{34} GeV [108]. Below that energy scale the Standard Model is self-consistent.

[3] Other values of hypercharge for a SU(2) doublet violate charge conservation once the symmetry is spontaneously broken, since there are no neutral components.

be gauged away so that they become the longitudinal polarization states of the weak gauge bosons.

In the unitary gauge the Higgs field can be represented as

$$\Phi = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}$$

with $v^2 = -4\mu^2/\lambda$. The Lagrangian (2.7) now becomes

$$\mathcal{L} = \frac{1}{2}\partial^\mu h \partial_\mu h + \frac{g^2}{4}(v^2 + 2vh + h^2) W^{+\mu} W^-_\mu + \frac{1}{8}(-gW^0_\mu + g' B_\mu)^2 (v^2 + 2vh + h^2)$$
$$- \sum_{SU(2),U(1)} \frac{1}{4} F^{\mu\nu}_a F_{\mu\nu a} - (3\lambda v^2/8 + \mu^2/2) h^2 - \frac{\lambda v}{4} h^3 - \frac{\lambda}{16} h^4. \quad (2.8)$$

The charged gauge bosons $W^\pm$ acquire a mass $gv/2$. In the neutral sector the combination proportional to $-gW^0 + g'B$ gets a mass, but the orthogonal combination proportional to $g'W^0 + gB$ remains massless. The latter is the gauge boson of the unbroken $\mathrm{U}(1)_{em}$ symmetry, the photon. The normalized mass eigenstates are

$$Z^\mu = \frac{-gW^{0\mu} + g'B^\mu}{\sqrt{g^2 + g'^2}}, \quad m_Z = \frac{\sqrt{g^2 + g'^2}\, v}{2}, \quad (2.9)$$

$$A^\mu = \frac{g'W^{0\mu} + gB^\mu}{\sqrt{g^2 + g'^2}}, \quad m_A = 0. \quad (2.10)$$

One usually defines the so called weak mixing angle (or Weinberg angle) by setting $g/\sqrt{g^2 + g'^2} = \cos\theta_W$ and $g'/\sqrt{g^2 + g'^2} = \sin\theta_W$. One can measure $\sin^2\theta_W$ by several methods, including the comparison of muon neutrino and antineutrino scattering with electrons [112], the ratio of charged and neutral currents in neutrino-nucleus scattering [113], parity violation in Moller scattering [114] and asymmetries of Z-boson decays [115]. All of the measurements show that the so called ρ-parameter, defined $\rho = \frac{m_W^2}{m_Z^2 \cos^2\theta_W}$, is close to one as predicted by the SM[4].

We also find that there are couplings between a single Higgs boson and two W- or Z-bosons, which are absent in the symmetric phase. Hence the observation of the decays $h \to WW^*$ [116, 117] and $h \to ZZ^*$ [118, 119] already is a strong indication of electroweak symmetry breaking via the Higgs mechanism.

After inserting $v^2 = -4\mu^2/\lambda$ we find that the Higgs mass term gets the correct sign and the mass is $m_h = \sqrt{-2\mu^2}$. If one identifies the scalar resonance found by the LHC with the SM Higgs, one may determine all of the SM input parameters.

[4] The value of $\sin^2\theta_W$ and hence the ρ-parameter depends on the chosen renormalization scheme when loop corrections are taken into account. Hence there is no unique value for the ρ-parameter, but in any scheme the value is close to one.

2.3.2 Fermion masses

Fermion masses can be generated via so called Yukawa couplings. To be able to generate masses for all charged fermions one needs in addition to the usual Higgs representation also the charge-conjugated representation $\Phi^c = i\sigma_2\Phi^*$. The Yukawa terms of the Standard Model are

$$\mathcal{L}_Y = Y_{ij}^{(U)}(Q_L^i)^T\Phi^c\bar{u}_R^j + Y_{ij}^{(D)}(Q_L^i)^T\Phi\bar{d}_R^j + Y_{ij}^{(L)}(L_L^i)^T\Phi\ell_R^{+j} + \text{h.c.}, \quad (2.11)$$

where $Y^{(U,D,L)}$ are 3×3 matrices. To find the mass eigenstates one must diagonalize the Yukawa matrices. The diagonalizing transformations will not be the same for different matrices. This has its implications on charged weak interactions. Namely, the charged current interactions in the gauge basis can be written as

$$\mathcal{L}_{cc} = \frac{ig}{2}\bar{u}_j\gamma^\mu W_\mu(1-\gamma_5)d_j + \text{h.c.}. \quad (2.12)$$

The transformation to the mass basis can be done by unitary matrices, say $V_{(u)}$ for the up quark sector and $V_{(d)}$ for the down quark sector. Hence the charged current interactions in the mass basis are

$$\mathcal{L}_{cc} = \frac{ig}{2}\bar{u}_j V^*_{(u)jk}\gamma^\mu W_\mu(1-\gamma_5)V_{(d)kl}d_l + \text{h.c.}, \quad (2.13)$$

where $V^\dagger_{(u)}V_{(d)} \equiv V_{\text{CKM}}$ is the Cabibbo–Kobayashi–Maskawa (CKM) matrix [31, 32]. The CKM matrix is nearly diagonal so generation changing processes have smaller rates than processes involving only one generation.

One may notice that the coupling between the Higgs and fermions is $\lambda_{hf\bar{f}} = \sqrt{2}m_f/v$, *i.e.* the Higgs couples most strongly to heavy fermions. However, the fermion masses are free parameters so they have to be taken as inputs of the SM. There is no explanation for the nearly diagonal form of the CKM matrix either.

2.3.3 SM Higgs production at the LHC

The LHC is a circular proton-proton collider with a designed center-of-mass energy $\sqrt{s} = 14$ TeV. In 2010–2012 the LHC operated at energies of 7 TeV and 8 TeV and after the shutdown in 2015 the collision energy was 13 TeV.

The dominant production mechanism of the SM Higgs at the LHC is gluon fusion. At leading order it is a one-loop process [79, 80], mostly proceeding via a top quark loop, with a subleading contribution from bottom loops as shown in Figure 2.1. The next-to-leading order (NLO) QCD corrections were computed almost 25 years ago [120–122] and they almost double the production cross section from the leading order estimate. Even the next-to-next-to-leading order (NNLO) corrections are rather large [123–128], although smaller than the LO and NLO contributions. The NLO electroweak corrections are also known [129–132] and they are an order of magnitude smaller than the NLO QCD corrections. Also the mixed QCD-electroweak corrections $\mathcal{O}(\alpha\alpha_s)$ have

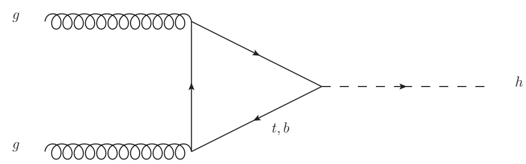

Figure 2.1: The leading order Feynman diagram of Higgs production via gluon fusion.

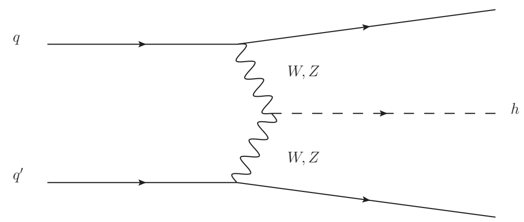

Figure 2.2: The leading order Feynman diagram of Higgs production via vector boson fusion. The recoiling quarks lead to jets with transverse momentum, which can be tagged.

been evaluated [133]. Recently the full N^{3}LO QCD corrections to the cross section were computed [134] after a number of partial results obtained over several years [135–141].

The predicted production cross section for gluon fusion with $m_h = 125$ GeV computed in [142] is 49 pb (55 pb) for $\sqrt{s} = 13$ TeV (14 TeV) with errors around 7%. Now that the N^{3}LO computation has been implemented the errors related to the QCD computation (scale variation, missing higher orders) are now at a level comparable to other uncertainties. The problem with gluon fusion is that the events are hard to distinguish from the background. Quarks are produced copiously via standard QCD processes and they outnumber any contribution coming from the Higgs decays. Hence only final states with leptons, photons or missing transverse energy can be identified.

There are two production modes based on electroweak production. They are subleading modes at hadron colliders but nevertheless important since they are easier to tag and allow to study the decays to quarks also. The first one is vector boson fusion (VBF) [143, 144]. There W- or Z-bosons are radiated off quarks and they thereafter collide forming a Higgs boson as shown in Figure 2.2. The quarks which radiate the gauge bosons get a kick and are seen as jets with transverse momentum, which can be tagged. The NLO QCD corrections to the total cross section were computed more than two decades ago [145] but since cuts on the jet momenta are made to reduce the background, the NLO differential cross section, computed in [146], is essential. These corrections are

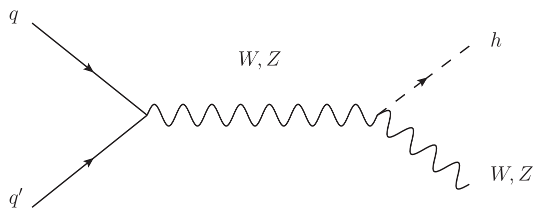

Figure 2.3: The leading order Feynman diagram of associated Higgs production with a vector boson. The additional boson makes the process easier to tag.

not as large as in gluon fusion but may reach 30%.

The NNLO corrections to VBF have been computed a few years ago [147, 148] and very recently also the differential cross section has been computed at NNLO in QCD [149]. The NNLO correction can be as large as 5%. The SM prediction for the total VBF cross section computed by the LHC Higgs cross section working group is 3.7 pb (4.2 pb) for $\sqrt{s} = 13$ TeV (14 TeV) [150]. The errors are below 5%.

The second electroweak production mode is the so called Higgs-strahlung or associated production, shown in Figure 2.3, where a Higgs boson is radiated from an off-shell W- or Z-boson [151, 152]. The decay products of the gauge boson can be identified. This mode is the leading production mode at electron-positron colliders, at hadron colliders it is comparable to VBF.

The NLO QCD corrections to this process have been computed nearly 25 years ago [153] and they increase the cross section at the LHC by roughly 30%. Most of the NNLO QCD corrections have also been computed 10 years ago [154] and corrections from top quarks at NNLO quite recently [155]. They increase the production cross section at the LHC by a few per cent compared to the NLO prediction. The NLO electroweak corrections have also been computed and they decrease the cross section by 5% [156]. The SM prediction for the cross section of WH production is 1.36 pb (1.50 pb) at $\sqrt{s} = 13$ TeV (14 TeV) and for ZH production 0.86 pb (0.96 pb). The errors are below 5%.

In the second run there is still one production mechanism that could be measured. The Higgs coupling to top quarks is large and hence it may be produced with a top-quark pair [157–162] as shown in Figure 2.4. The NLO QCD corrections have been computed [163–165] and at the LHC this enhances the cross section by about 20% compared to the LO result. Also the NLO EW correction has been computed recently [166]. The SM prediction for the cross section is 0.50 pb (0.60 pb) for $\sqrt{s} = 13$ TeV (14 TeV). The errors are below 10%.

The associated production with top quarks seems to be the simplest way to probe the Higgs coupling to top quarks. The cross section of gluon fusion depends on the top Yukawa coupling but it is not too sensitive to it. When the Yukawa coupling gets larger, the particle is heavier and the propagators in the

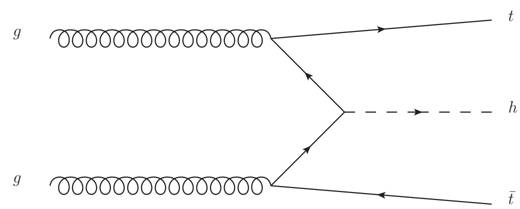

Figure 2.4: The leading contribution to associated Higgs production with a top quark pair.

loop are suppressed so that in the limit $m_t \gg m_h$ the amplitude approaches a constant value. The Higgs-top coupling also gives a subleading contribution to the decay $h \to \gamma\gamma$ but it can constrain the coupling only in a model-dependent way.

2.3.4 SM Higgs decay channels

The Higgs couples most strongly to heavy particles so it decays mainly to the heaviest particles kinematically available. With a mass of 125 GeV the top quark channel and on-shell gauge boson channels are not open and hence the dominant decay mode is $h \to b\bar{b}$. The subleading fermionic decay channels are $h \to \tau^+\tau^-$ and $h \to c\bar{c}$. Even though the decay $h \to W^+W^-$ is not allowed on-shell, the branching ratio for the off-shell decay $h \to WW^*$ is still rather large.

There are two rare decay modes which have become extremely important. The decay to ZZ^* which subsequently decay to four charged leptons has a small branching ratio. This channel has almost no background and provides a good mass resolution and a possibility to study the spin and parity of the particle. The one-loop mediated decay $h \to \gamma\gamma$ has also a small branching ratio but it provides a good mass resolution and a reasonable signal compared to the deviation of the background. These channels were the ones that made it possible to claim the discovery in July 2012 [1, 2].

The fermionic tree-level decay width is [167, 168]

$$\Gamma(h \to f\bar{f}) = \frac{N_c g^2 m_f^2 m_h}{32\pi m_W^2}\beta^3, \tag{2.14}$$

where N_c is the color factor (3 for quarks and 1 for leptons) and $\beta = (1 - 4m_f^2/m_h^2)^{1/2}$. For quarks the leading-log (LL) QCD correction to the decay width can be obtained by replacing the mass with the running quark mass evaluated at the Higgs boson mass [169]. For bottom quarks this correction is rather large bringing the decay width down by 50% hence effecting essentially to the total decay width. The $\mathcal{O}(\alpha_s)$ QCD correction [169] is positive but smaller than the LL contribution. These have also been computed in the limit $m_h \gg m_f$ in [170, 171].

The $\mathcal{O}(\alpha_s^2)$ QCD corrections in the limit of massless quarks were computed in [172, 173] and the correction $\mathcal{O}(m_f^2/m_h^2)$ to this in [174]. This correction is about one third of the $\mathcal{O}(\alpha_s)$ correction. Even the corrections $\mathcal{O}(\alpha_s^3)$ [175] and $\mathcal{O}(\alpha_s^4)$ [176] have been computed in the massless limit. Since $m_b^2/m_h^2 \simeq 10^{-3}$ the corrections from the finite quark masses are small. The differential cross section has been recently computed to $\mathcal{O}(\alpha_s^2)$ [177].

The NLO electroweak corrections have also been evaluated long time ago [178]. They are of the order of a few percent.

QCD backgrounds for $b\bar{b}$ production are large and hence one can search for $h \to b\bar{b}$ only using VBF or VH production so that either the recoiling jets or gauge boson can be tagged. The same applies to $h \to c\bar{c}$. The $h \to \tau^+\tau^-$ channel is cleaner so that also gluon fusion production can be considered.

The Higgs cannot decay to two on-shell gauge bosons but still the off-shell decays to WW^* and ZZ^* are relevant. The decay widths to these modes are [179]

$$\Gamma(h \to WW^*) = \frac{3g^4 m_h}{512\pi^3} F(x), \tag{2.15}$$

where

$$F(x) = \frac{3(1-8x^2+20x^4)}{(4x^2-1)^{1/2}} \arccos\left(\frac{3x^2-1}{2x^3}\right) - (1-x^2)\left(\frac{47}{2}x^2 - \frac{13}{2} + \frac{1}{x^2}\right) \\ - 3(1-6x^2+4x^4)\ln x, \; x = m_W/m_h, \tag{2.16}$$

and

$$\Gamma(h \to ZZ^*) = \frac{g^4 m_h}{2048\pi^3 \cos^4\theta_W}\left(7 - \frac{40}{3}\sin^2\theta_W + \frac{160}{9}\sin^4\theta_W\right) F(y), \tag{2.17}$$

where $y = m_Z/m_h$ and F is given in equation (2.16). (These results were derived already in [180] but the integral that gives $F(x)$ was evaluated numerically.)

The one-loop electroweak correction to these decays can be evaluated from the on-shell four-point amplitudes $f\bar{f}Vh$, which were first computed in the context of $e^+e^- \to Zh$ production [181–183]. For the SM Higgs with a mass of 125 GeV this gives an enhancement of a few percent [184].

The leptonic decays of gauge bosons do not have a too large background but channels with neutrinos produce a broad excess and hence make it more difficult to distinguish the resonance from background. The "golden channel" of four charged leptons from $h \to ZZ^*$ has almost no background and despite the small branching ratio it has been one of the most important channels so far since one can make a full kinematical reconstruction of the final state.

The Higgs decay to two photons is a one-loop process at leading order and it can be mediated by any charged particle that interacts with the Higgs. The partial width is [83, 168]

$$\Gamma(h \to \gamma\gamma) = \frac{\alpha^2 g^2 m_h^3}{1024\pi^3 m_W^2} \left| \sum_i N_i q_i^2 F_i \right|^2 , \tag{2.18}$$

where N_i is the number of colors and q_i the charge of particle i. F_i is a loop function that depends on the mass and spin of the particle. It is larger for heavy particles and in the limit where the particle in the loop is very heavy (light) compared to the Higgs[5], the functions have the limits $F_0 = -1/3$ (0), $F_{1/2} = -4/3$ (0) and $F_1 = 7$ (2), where the subscript indicates the spin of the particle. Hence vector bosons give the largest contribution to this decay mode. In the SM, the dominant contribution comes from the W boson with a subleading distructive contribution from the top quark loop. Due to the intermediate loop this decay mode is sensitive to new charged particles.

The NLO QCD corrections to this mode have been evaluated in [185]. Since quarks form only a subleading contribution to the amplitude the efect to the partial width is only an enhancement of about two percent.

In addition to the decay to quarks the Higgs can decay to jets also via the loop-induced decay to gluons [80]. This is the inverse of gluon fusion production and will be mediated mostly by the top quark loop. Due to large QCD backgrounds this mode will be hard to detect at hadron colliders. This decay mode will dominate over any contribution from light quarks, which would also be seen as dijets and hence one cannot determine Yukawa couplings for light quarks by collider experiments. Those Yukawa couplings can be mildly constrained from the total width but better constraints can be put by accurate atomic measurements [186].

Yet another loop-level decay is $h \to Z\gamma$. Like the diphoton decay the mediating particle can be any charged particle that couples also to the Z boson. The partial width [84] has a similar structure than in the case of the diphoton decay and again the dominant contribution comes from spin-1 mediators. The NLO QCD corrections to this mode have been evaluated but similarly to the diphoton channel the corrections are small, less than a percent [187].

If $h \to \gamma\gamma$ and $h \to Z\gamma$ can both be measured with a fairly good accuracy they will place very tight constraints on the charged particle content.

The predicted SM Higgs branching ratios are given in Table 2.3. The uncertainties in the SM predictions of decay modes are smaller than in the production modes since the uncertainties of parton distribution functions do not affect the decay processes. The predictions related to Higgs production and decay are among the most impressive results of perturbative QCD.

[5] Here it is also assumed that the particle gets its mass solely from the VEV of the Higgs field. If there are other scalars whose VEVs contribute to the mass, the amplitude will be suppressed.

Mode	Submode	Branching ratio
$b\bar{b}$		59.2%
WW^*		21.5%
	$\to 2\ell 2\nu$	2.3%
	$\to \ell\nu qq'$	9.4%
	$\to$hadronic	9.7%
gg (jets)		6.9%
$\tau^+\tau^-$		6.2%
$c\bar{c}$		3.0%
ZZ^*		2.7%
	$\to 4\ell$	0.027%
	$\to 4\nu$ (inv.)	0.11%
$\gamma\gamma$		0.23%
$Z\gamma$		0.16%

Table 2.3: The branching ratios of the Standard Model Higgs boson with a mass $m_h = 125.5$ GeV.The branching ratios are computed with HDECAY [188]. The submodes are computed with the experimental branching ratios from PDG. Here ℓ denotes any charged lepton.

2.3.5 Comparison of the 125 GeV boson and the SM Higgs

We can be certain that the SM needs to be extended. The SM offers a minimal model of EWSB and many extensions do have alternative ways to introduce masses. Although one typically needs a particle which somewhat resembles the SM Higgs boson, its properties will be modified from the SM predictions. Hence one needs to compare the new particle and the SM Higgs to see, if are there any hints of new physics and if not, how large contributions from new physics are still allowed.

Already from the decay products one could infer that the new particle is neutral and it is a boson. The spin and parity were studied from the angular correlations of the Higgs decaying to four leptons via weak gauge bosons. The SM quantum numbers $J^P = 0^+$ are favored over all other possibilities [189, 190]. If parity is not a good quantum number the positive parity component is dominant. The upper limit (at 95% CL) for the negative parity component is 43% measured from the ZZ^* channel [190].

The overall Higgs production rate, both untagged and with various tags related to the production mechanism is compatible with the SM prediction. The typical way to parametrize the experimental results is to define μ as the ratio of the experimentally fitted production cross section and the prediction for the SM Higgs. CMS finds $\mu = 1.00 \pm 0.14$ [191] and untagged, VBF tagged and VH tagged Higgs production to be within one standard deviation from the SM prediction. The $t\bar{t}h$ production seems to be larger than the SM prediction, $\mu_{tth} = 2.75 \pm 0.99$, but the result is not even strong enough to claim the observation of $t\bar{t}h$ production. ATLAS finds an overall rate $\mu = 1.18^{+0.15}_{-0.14}$ compared to the SM prediction [192]. ATLAS has all production modes roughly within one standard deviation from the SM prediction. The $t\bar{t}h$ production is also above the SM prediction, $\mu_{tth} = 1.7 \pm 0.8$ [193], but not enough to make any conclusions of deviations from the SM.

The consistency of the gluon fusion production rate with the SM can exclude the direct extension of the SM with a fourth generation of fermions [3, 4]. A new quark doublet should increase the gluon fusion production cross section by roughly a factor of nine [79] and that is not seen in the data.

The essential feature of SU(2)×U(1) breaking with a doublet field is the relationship between W and Z masses and the couplings $\lambda_{hWW}/\lambda_{hZZ} = m_W^2/m_Z^2$. This can be measured both from VH Higgs production and Higgs decays to WW and ZZ. Assuming the ratio to be positive ATLAS finds it to be $0.92^{+0.14}_{-0.12}$ times the SM value [192] and CMS finds it to be $0.94^{+0.22}_{-0.18}$ times the SM value [191], both consistent with the SM prediction.

Out of the individual decay modes, the decays to $\gamma\gamma$, ZZ^* and WW^* have been measured to an accuracy better than 20%. The decay modes to $b\bar{b}$ and $\tau^+\tau^-$ have a nonzero width but the errors in the signal strengths are still large. In all of these decay channels the signal strengths are reasonably close to the SM prediction [191, 192], at most slightly more than one standard deviation off. Other decay modes remain unobserved with upper limits clearly above the SM predictions [194–196]. So far there is no evidence of lepton flavor violating Higgs decays, although both ATLAS and CMS see a small excess in the channel $h \to \mu\tau$ [197–199].

In the SM there is only the rare decay $h \to ZZ^* \to 4\nu$, in which the decay products of the Higgs remain completely unobserved. In BSM scenarios with light dark matter coupling to the Higgs there could be new invisible decay modes. The $h \to$ invisible branching ratio can be constrained by looking at VH production in the cases where the vector boson can be fully reconstructed. The SM background comes from events where there is a Z boson decaying to neutrinos instead of the Higgs boson. Using this method and assuming SM production rates, ATLAS has placed an upper limit (at 95% CL) of 75% for the $h \to$ invisible branching ratio [200, 201] and CMS finds an upper limit of 24% [202]. Both experiments are consistent with the SM expectation.

The SM predicts the total decay width of the Higgs boson to be about 4 MeV. The resolution of the calorimeters will not allow to measure such a small width directly. Currently CMS finds an upper limit of 1.7 GeV from the combination of $h \to \gamma\gamma$ and $h \to ZZ^* \to 4\ell$ channels [191]. However one may indirectly constrain the decay width by comparing the on-shell and off-shell decays to ZZ^* [203]. With this method CMS obtains an upper limit of 13 MeV for the decay width [204], not too much above the SM prediction.

Altogether there are currently no clear indications of deviations from the SM Higgs predictions. However, the uncertainties are still large in many observables and they leave room for the fingerprints of new physics.

2.4 Extended Higgs sectors

One way to extend the Standard Model is to extend the sector of fundamental scalars. Such extensions can help to solve some of the problems of the Standard Model. In principle there are no limitations to the dimensionalities of the Higgs representations. However, from the known physics, models with doublets, singlets or triplets seem to be more motivated than larger representations. Next we briefly review the reasons for this and some of their basic features.

2.4.1 Additional doublets

The most studied extension of the scalar sector is the so called two Higgs doublet model (2HDM), originally proposed by T.D. Lee [205]. The original motivation for the model was the possibility of CP-violation in the Higgs sector. It is also possible to have a dark matter candidate in the 2HDM, such a model is called Inert Higgs doublet model [206–209]. A strong motivation for additional doublets is that they are necessary in supersymmetric models. The reasons will be discussed in chapter 3. The 2HDM is a minimal model for Higgs sector flavor violation.

The first constraint for an extension of the SM scalar sector is to correctly reproduce the pattern of electroweak symmetry breaking. For a general scalar sector the ratio of W and Z boson masses squared, usually called the ρ parameter, is at tree-level [210]

$$\rho \equiv \frac{m_W^2}{m_Z^2 \cos^2\theta_W} = \frac{\sum_{T,Y}(4T(T+1)-Y^2)|V_{T,Y}|^2 c_{T,Y}}{\sum_{T,Y} 2Y^2|V_{T,Y}|^2}, \tag{2.19}$$

where T is the isospin, Y the hypercharge and $V_{T,Y}$ the vacuum expectation value of the Higgs field. The factor $c_{T,Y}$ is one for a complex representation and $1/2$ for a real representation.

Experimentally we know that the ρ parameter is close to one. At tree-level $\rho = 1$ for any number of doublets with hypercharge $Y = \pm 1$.

The most general renormalizable and gauge invariant scalar potential for two $Y = 1$ doublets is of the form

$$\begin{aligned} V(\Phi_1,\Phi_2) &= m_{11}^2\Phi_1^\dagger\Phi_1 + m_{22}^2\Phi_2^\dagger\Phi_2 + (m_{12}^2\Phi_1^\dagger\Phi_2 + h.c.) + \frac{\lambda_1}{2}(\Phi_1^\dagger\Phi_1)^2 \\ &+\frac{\lambda_2}{2}(\Phi_2^\dagger\Phi_2)^2 + \lambda_3\Phi_1^\dagger\Phi_1\Phi_2^\dagger\Phi_2 + \lambda_4(\Phi_1^\dagger\Phi_2)(\Phi_2^\dagger\Phi_1) \\ &+\frac{\lambda_5}{2}((\Phi_1^\dagger\Phi_2)^2 + h.c.) + [(\lambda_6(\Phi_1^\dagger\Phi_1) + \lambda_7(\Phi_2^\dagger\Phi_2))(\Phi_1^\dagger\Phi_2) + h.c.] \end{aligned} \tag{2.20}$$

with $\Phi_1 = (\varphi_1^+ \ \varphi_1^0)^T$, $\Phi_2 = (\varphi_2^+ \ \varphi_2^0)^T$. We denote the VEV of φ_i^0 by $v_i/\sqrt{2}$. Since the potential is Hermitian, m_{11}^2, m_{22}^2 and $\lambda_{1,2,3,4}$ are real, other parameters can be complex.

Vacuum stability gives some conditions on the quartic couplings. For instance if $\lambda_6 = \lambda_7 = 0$ the conditions are [206] $\lambda_1 > 0$, $\lambda_2 > 0$, $\lambda_3 + \sqrt{\lambda_1\lambda_2} > 0$, $\lambda_3 + \lambda_4 \pm |\lambda_5| + \sqrt{\lambda_1\lambda_2} > 0$.

Type	Φ_1	Φ_2	$hu\bar{u}$	$hd\bar{d}$	$h\ell^+\ell^-$	hVV
Type I	—	u, d, ℓ	$\frac{\cos\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$\sin(\beta-\alpha)$
Type II	d, ℓ	u	$\frac{\cos\alpha}{\sin\beta}$	$-\frac{\sin\alpha}{\cos\beta}$	$-\frac{\sin\alpha}{\cos\beta}$	$\sin(\beta-\alpha)$
Lepton-specific	ℓ	u, d	$\frac{\cos\alpha}{\sin\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$-\frac{\sin\alpha}{\cos\beta}$	$\sin(\beta-\alpha)$
Flipped	d	u, ℓ	$\frac{\cos\alpha}{\sin\beta}$	$-\frac{\sin\alpha}{\cos\beta}$	$\frac{\cos\alpha}{\sin\beta}$	$\sin(\beta-\alpha)$
Type III	u, d, ℓ	u, d, ℓ				$\sin(\beta-\alpha)$

Table 2.4: Various types of 2HDM. Up-type quarks are denoted by u, down-type quarks by d and charged leptons by ℓ. There are various conventions for the naming of these types in the literature. The two latter ones can be called Type IB and Type IIB or Type X and Type Y. We also tabulated the couplings of the SM-like CP-even scalar to various particles compared to Higgs couplings in the SM. The couplings in Type III are those of Type I added with a flavor changing part.

The vacuum state can be either electroweak conserving ($\langle\Phi_1\rangle = \langle\Phi_2\rangle = 0$), charge conserving ($\langle\varphi_1^+\rangle = \langle\varphi_2^+\rangle = 0$) or charge breaking. There can be at most one charge breaking minimum and at most two charge conserving minima [211–213]. For a given potential charged and neutral minima can never coexist.

Gauge boson masses coincide with the SM as long as the vacuum is charge conserving and $v_1^2 + v_2^2 = v_{\rm SM}^2$. It is common to denote $v_1 = v\cos\beta$ and $v_2 = v\sin\beta$ and to use $\tan\beta = v_2/v_1$ to parametrize the VEVs of the two Higgs fields.

Flavor changing neutral currents (FCNC) are known experimentally to be small. Hence the couplings to fermions should be designed so that there will be no large new sources of FCNC. The most common solution is natural flavor conservation [214], which means that each type of fermions gets its mass from one doublet only. This leads to four alternatives of fermionic couplings with natural flavor conservation, given in Table 2.4. The fermionic couplings determine what we mean by the two doublets in Eq. (2.20) — without them we could freely rotate the doublets and choose whatever basis we want for the Higgs sector. Supersymmetric models are of Type II and the inert Higgs doublet model is of Type I having $\cot\beta = 0$.

The option of coupling fermions to both doublets is known as type-III 2HDM. There are several ways to suppress FCNC in such a setting: discrete symmetries [215], approximate U(1) symmetries [216] or just simply constrain the Yukawa couplings experimentally [217].

The hierarchies of fermion masses may get a partial explanation in some types of 2HDM, since they can be produced by two VEVs of different size. For example if in Type II we had $\tan\beta \simeq m_t/m_b$, the two Yukawa couplings would be of the same size and the Yukawa coupling for τ would not be much smaller.

Also the most general Higgs potential (2.20) can mediate FCNC via the $\lambda_{6,7}$ couplings and the mixing term proportional to m_{12}^2. Typically one avoids Higgs sector FCNC simply by setting $\lambda_6 = \lambda_7 = 0$. This leads to an approximate Z_2 symmetry in the Higgs sector which is broken softly by the dimension-2 mixing term. This soft breaking will not generate FCNC at tree-level.

After electroweak symmetry breaking there are five physical scalars, three neutral ones and a charged pair ($H^{\pm}$). If CP is conserved, two of the neutral scalars are CP-even (h, H) and one CP-odd (A). The 125 GeV scalar is one of the CP-even states. It may be the lighter or the heavier one [218], though in the latter case the viable parameter space is quite limited.

The mass eigenstates are linear combinations of the doublets. In the CP-conserving case the CP-odd and charged eigenstates are given by the VEV mixing angle β but in the CP-even sector the mixing is nontrivial:

$$\begin{aligned} h &= -(\sqrt{2}\Re\varphi_1^0 - v_1)\sin\alpha + (\sqrt{2}\Re\varphi_2^0 - v_2)\cos\alpha, && (2.21)\\ H &= (\sqrt{2}\Re\varphi_1^0 - v_1)\cos\alpha + (\sqrt{2}\Re\varphi_2^0 - v_2)\sin\alpha, && (2.22)\\ A &= -\sqrt{2}\Im\varphi_1^0\sin\beta + \sqrt{2}\Im\varphi_2^0\cos\beta, && (2.23)\\ H^{\pm} &= -\varphi_1^{\pm}\sin\beta + \varphi_2^{\pm}\cos\beta. && (2.24) \end{aligned}$$

The mixing angles α and β are related at tree-level by

$$\tan 2\alpha = \frac{2m_{12}^2 + v^2(\lambda_3 + \lambda_4 + \lambda_5)\sin 2\beta}{m_{12}^2(\tan\beta - \cot\beta) + v^2(\lambda_1\cos^2\beta - \lambda_2\sin^2\beta)} \tag{2.25}$$

assuming $\lambda_6 = \lambda_7 = 0$.

Since the Higgs coupling to vector bosons is proportional to the VEV and there does not seem to be a large deviation from the SM couplings, the 125 GeV scalar is the mass eigenstate with the larger VEV.

The CP-odd Higgs and the charged Higgs have tree-level masses

$$\begin{aligned} m_A^2 &= -\frac{2m_{12}^2}{\sin 2\beta} - \left[\frac{\lambda_6\cos^2\beta + \lambda_7\sin^2\beta}{\sin 2\beta} + \lambda_5\right]v^2, && (2.26)\\ m_{H^\pm}^2 &= m_A^2 + \frac{\lambda_5 - \lambda_4}{2}v^2. && (2.27) \end{aligned}$$

The tree-level CP-even mass matrix is of the form

$$M^2 = \begin{pmatrix} A_1\tan\beta + B_1v^2 & A_3 + B_3v^2 \\ A_3 + B_3v^2 & A_2\cot\beta + B_2v^2 \end{pmatrix}, \tag{2.28}$$

where $A_1 = -m_{12}^2 + (3\lambda_6\cos^2\beta - \lambda_7\sin^2\beta)v^2/2$, $A_2 = -m_{12}^2 + (3\lambda_7\sin^2\beta - \lambda_6\cos^2\beta)v^2/2$, $A_3 = m_{12}^2 + (3\lambda_6\cos^2\beta + 3\lambda_7\sin^2\beta)v^2/2$, $B_1 = \lambda_1\cos^2\beta$, $B_2 = \lambda_2\sin^2\beta$ and $B_3 = \frac{1}{2}(\lambda_3 + \lambda_4 + \lambda_5)\sin 2\beta$.

One sees immediately that the charged Higgs and the CP-odd Higgs are roughly degenerate if they are much heavier than 125 GeV. Also one of the CP-even states is nearly degenerate with them and one will be lighter. This comes from the fact that if $\lambda_{6,7} \simeq 0$ as the smallness of FCNC suggests, the determinant of the "A-parts" will be nearly[6] zero. Treating the "B-terms" as a perturbation there will be one state getting its mass from the B-terms $\mathcal{O}(\lambda v^2)$

[6] All terms are proportional to either λ_6 or λ_7.

and another with a mass mainly from A-terms $\mathcal{O}(A_1 + A_2) \simeq m_A^2$. This case is known as the decoupling limit.

In 2HDM the Higgs couplings are modified compared to the SM. The SM couplings are multiplied by factors given in Table 2.4[7]. The couplings to gauge bosons are always suppressed but couplings to fermions can either be enhanced or suppressed. This will modify the production and decay rates compared to the SM. In Type I and the lepton-specific model the cross section gluon fusion will be modified by a factor $(\cos\alpha/\sin\beta)^2$. In Type II and the flipped model the gluon fusion cross section may be different if the coupling of the SM-like Higgs to bottom quarks is enhanced. In that case also $b\bar{b}h$ production will be enhanced. The electroweak production modes will be suppressed by a factor $\sin^2(\beta - \alpha)$. The branching ratios can also be altered through the modified couplings. The Higgs decay to photons or $Z\gamma$ will get a new contribution from the charged Higgs, but the contribution from scalar loops is small compared to vector loops. The phenomenology of 2HDM has been reviewed in [219, 220].

Since the 125 GeV scalar does not seem to show large deviations from the SM Higgs, the viable parameter space is near the alignment limit. This means that the vacuum expectation value is aligned in the direction of one of the CP-even mass eigenstates so that $\sin(\beta - \alpha) = 1$. In such a case there will be a scalar that has couplings like the SM Higgs. The eigenstates are close to the alignment limit if we are at another limit, the decoupling limit, discussed above. However, alignment is possible without decoupling [221, 222].

In the alignment limit the couplings of the other CP-even state to vector bosons are small since they are proportional to $\cos(\beta - \alpha)$. The CP-odd state does not couple to WW or ZZ. Hence the heavier neutral states are mainly produced via gluon fusion and they decay mainly to fermions of the third generation. On the other hand charged Higgs bosons production rates are smaller since one either needs quarks in the initial state (like $gb \to tH^-$) or production of heavy quarks in the final state (like $gg \to t\bar{b}H^-$). This will make finding the other scalar states difficult.

Models with more than two doublets have a very large number of free parameters unless there are symmetries that forbid a subset of terms from the Lagrangian. One also has more freedom in choosing the fermionic couplings. One may for instance use three doublets to give masses to up-type quarks, down-type quarks and leptons [223]. One can have one or several doublets that couple to fermions and simultaneously inert doublets, which are viable dark matter candidates [224].

Most of the studies have concentrated on the three Higgs doublet model (3HDM). All finite symmetry groups of the Higgs sector of 3HDM have been identified [225]. Examples of recent phenomenological studies of 3HDM are [226–228], where some features and constraints of the Higgs sector have been studied. However the diversity of the 3HDM phenomenology still remains largely unexplored.

[7] The couplings of all scalars to fermions can be found *e.g.* in [219], Tables 2 and 3.

2.4.2 Singlets

Singlet scalars can also be easily added to the Standard Model, since they do not affect the ρ parameter. There are several motivations to study such an extension. First of all one may use the singlet as a dark matter candidate [229,230]. Another motivation is that cosmic inflation needs a scalar field called inflaton, which could be a singlet under the SM gauge group [231]. If there is a hidden sector, the messenger between the hidden sector and the visible sector could also be a singlet scalar particle coupling to the Higgs with renormalizable couplings. Such a scenario is called the Higgs portal [232]. A singlet, when mixing with the doublet Higgs, can also help to stabilize the scalar potential [47] or modify the electroweak phase transition [233–236].

The simplest possible singlet extension of the Standard Model is to simply add a (real or complex) singlet scalar to the SM and impose a Z_2 symmetry. This allows the singlet to be a dark matter candidate. The scalar potential of such a model is

$$V(\Phi, S) = m_\Phi^2 \Phi^\dagger \Phi + m_S^2 S^\dagger S + \lambda_\Phi (\Phi^\dagger \Phi)^2 + \lambda_S (S^\dagger S)^2 + \lambda_{\Phi S} (\Phi^\dagger \Phi)(S^\dagger S). \quad (2.29)$$

If $m_S^2 > 0$ the VEV of the singlet will be zero and its Z_2 symmetry will remain unbroken. Hence the singlet will be stable.

The singlet can also be a viable dark matter candidate in the freeze-in scenario [237,238], where the dark matter component is not in thermal equilibrium with the other particle species. However in that case the portal coupling $\lambda_{\Phi S}$ is so small that it will make the singlet unobservable at collider experiments.

If one does not impose the Z_2 symmetry, one can also include the cubic S^3 term to the Lagrangian. Such a term can help in having a strong first order electroweak phase transition [233].

If we allow a singlet VEV, the two scalars will mix. If the singlet-dominated state is heavier, the mixing will push the mass of the doublet-dominated Higgs down. Hence one will need a larger quartic coupling to have a 125 GeV Higgs than in the SM. If $\lambda_{\Phi S} \gtrsim 0.03$ this effect is large enough to make the scalar potential stable to arbitrarily high energies [47]. One may notice that the singlet VEV is not constrained by the ρ parameter so it may be small or large and hence the singlet-dominated state can be light or heavy.

One may also extend a model with more than one doublet with singlets. Such an extension is rather natural in supersymmetric models, where the addition of a singlet state solves the so called μ problem. We shall discuss such models in section 3.4.1.

The main difference in phenomenology is that mixing with a singlet state will result in an universal suppression on all production rates but branching ratios will be unaltered. If the singlet is lighter than $m_h/2$ it can contribute to the invisible Higgs branching ratio if it is stable. The invisible branching ratio can be large [239]. If the singlet can decay, it will look like a SM Higgs of that mass and could lead to decay chains like $H_2 \to 2H_1 \to 4b, 2b2\tau$ [240].

2.4.3 Triplets

Triplet scalar extensions can be motivated in various ways. SU(2) triplets can be used to give masses to neutrinos, this is commonly known as type-II seesaw [12–16]. Another type of models, where triplets often exist, is that of left-right symmetric models [241, 242], where parity is broken spontaneously. This breaking is often due to right-handed scalar triplets. Left-right symmetric models will be discussed in detail in chapter 5.

Unlike doublets and singlets, triplet scalars, if their neutral components have VEVs, change the ρ parameter. For a triplet with $Y = 0$ Eq. (2.19) gives $\rho \simeq 1 + 4(v_T/v)^2$, where v_T is the VEV of the triplet field. The current constraints on the ρ parameter[8] require the VEVs of left-handed triplets to be below a few GeV's. One may still construct models, where $\rho = 1$ at tree-level even if triplet VEVs are larger by arranging suitable Higgs representations. An example of such a model is the one by Georgi and Machacek [243], where there is a complex $Y = 2$ triplet and a real $Y = 0$ triplet in addition to the usual doublet.

It is possible to write a dimension four term for neutrino masses with $\mathrm{SU}(2)_L$ triplets [12–16]. Gauge invariance requires the hypercharge to be $Y = 2$. The simplest model for this is called the Higgs triplet model (HTM), where one has one doublet and one triplet, usually given in the form

$$\Phi = \begin{pmatrix} \varphi^+ \\ \varphi^0 \end{pmatrix}, \quad \Delta = \begin{pmatrix} \delta^+/\sqrt{2} & \delta^{++} \\ \delta^0 & -\delta^+/\sqrt{2} \end{pmatrix}. \tag{2.30}$$

The scalar potential of the HTM can be written in the form [244, 245]

$$\begin{aligned} V(\Phi, \Delta) = &-m^2(\Phi^\dagger\Phi) + M^2\mathrm{Tr}(\Delta^\dagger\Delta) + \frac{\lambda_1}{2}(\Phi^\dagger\Phi)^2 + \frac{\lambda_2}{2}[\mathrm{Tr}(\Delta^\dagger\Delta)]^2 \\ &+ \lambda_3\mathrm{Det}(\Delta^\dagger\Delta) + \lambda_4(\Phi^\dagger\Phi)\mathrm{Tr}(\Delta^\dagger\Delta) + \lambda_5(\Phi^\dagger\tau_i\Phi)(\Delta^\dagger\tau_i\Delta) \\ &- \frac{M\lambda_6}{\sqrt{2}}(\Phi^T i\sigma_2\Delta^\dagger\Phi + \mathrm{h.c.}). \end{aligned} \tag{2.31}$$

For such a model Eq. (2.19) gives $\rho \simeq 1 - 2(v_\delta/v)^2$, so we need $v_\delta \ll v$. The λ_6 term is linear in Δ so v_δ is always nonzero. Assuming v_δ to be small and hence neglecting quartic terms, one gets $v_\delta = \lambda_6 v^2/4M$. Thus if the triplet states are heavy, the triplet VEV will be small.

The neutrino masses will be generated from the term

$$\mathcal{L} = y_{ij}^{(\nu)} L_i^T i\sigma_2\Delta L_j + \mathrm{h.c.} \tag{2.32}$$

This will give neutrinos a Majorana mass $\mathcal{O}(y^{(\nu)}v_\delta)$. This also favors small v_δ. The same coupling couples also the doubly charged Higgs to a pair of same

[8] In the $\overline{\mathrm{MS}}$ scheme, when the top quark loop contribution is removed, one has $\rho = 1.00037 \pm 0.00023$ [90].

sign charged leptons and the singly charged triplet Higgs to a charged lepton-neutrino pair. The Majorana mass term violates lepton number by two units and it can contribute to leptogenesis [246].

The important feature of triplets is that one may combine a triplet with two doublets to form a dimension three term in the Lagrangian. This can introduce a mixing between triplets and doublets that is larger than that of doublets with larger representations[9]. Such a coupling is also possible in the superpotential of supersymmetric models.

If the triplet states are heavy, the only low-energy feature of HTM will be neutrino masses. If, however, the small neutrino masses are due to small couplings, the triplet states could be within the reach of the LHC. The doubly charged Higgs could change the $h \to \gamma\gamma$ rate considerably due to its charge if the λ_4 coupling is large [247].

The charged Goldstone will be mostly doublet-like, so the physical charged Higgs is mostly triplet-like. This leads to two interesting features. First, from Eq. (2.32), we find that the leptonic decay modes are affected by the neutrino Yukawa couplings, which can alter the branching ratios completely. However, this requires that the doublet part of the physical charged Higgs is smaller than the ratio y_ν/y_τ, where y_ν is a typical element of the matrix $y^{(\nu)}$. For $M = 500$ GeV, $\lambda_2 = \lambda_4 = 0.1$, $\lambda_5 = 0$, this happens when $v_\delta \lesssim 0.01$ GeV. With small v_δ the mixing in the CP-even neutral sector is small and the triplet-like state will have a large branching ratio to neutrinos. If $v_\delta \simeq 1$ GeV, the mixing is so large that the doublet part dominates the fermionic decay widths.

The second feature is that there is a coupling between $H^\pm W^\mp Z$ opening the decay channel $H^+ \to W^+ Z$. In models with only doublets such a coupling is zero [248]. This coupling increases with v_δ and when v_δ is close to the limits from the ρ parameter, could lead to a substantial branching ratio for $H^+ \to W^+ Z$ [249].

2.5 Experimental constraints on extended Higgs sectors

2.5.1 Neutral scalars

Neutral scalars have been searched in various ways. Many of the Standard Model Higgs searches can be reinterpreted as searches for other neutral scalars.

For masses below m_W the scalars must couple very weakly to the Z-boson. The LEP limits on the hZZ-coupling constrain it to be less than 4% of the SM coupling [250,251]. Hence, if a light scalar exists, its dominant component must be a singlet under SU(2) or its VEV must be small. The constraints are a lot milder at masses 80 GeV$< m <$ 110 GeV and above 110 GeV the LEP is no more sensitive.

[9] Such mixing will be suppressed by the small VEV of the larger Higgs representation.

At higher masses the couplings to WW and ZZ are constrained by the LHC searches in the various decay channels of WW and ZZ, most restricting often being $H \to ZZ \to 4\ell$. These constrain the production cross section times branching ratio to be about 10% of the SM value for a range of masses from 145 GeV to 400 GeV (except the ZZ threshold around 180 GeV) and below the SM value up to 1 TeV [252].

If we are close to the alignment limit in the 2HDM, the decays to gauge bosons are rare. Instead the other CP-even state must be searched through its decays to third generation fermions. Since quarks are produced copiously at LHC[10] the decay $H \to \tau^+\tau^-$ will have the lowest background. The sensitivity to the Higgs signal is enhanced by requiring associated production with at least one tagged bottom quark. There are no excesses in the data [253–255].

The decays $H \to \gamma\gamma$ and $H \to Z\gamma$ are possible due to the coupling to charged fermions. In the diphoton data of Run I there are no excesses above two standard deviations at ATLAS between 65 GeV and 600 GeV [256] and at CMS between 150 GeV and 850 GeV [257]. Both experiments reported an excess close to 750 GeV in the first year of Run II, but the 2016 data set does not support the excess [258]. The analysis of the $Z\gamma$ final state has been performed at masses below 160 GeV [194, 195]. The sensitivity is not good enough to give a clear signal of the SM Higgs so the constraints on additional scalars are very weak.

If the other CP-even Higgs is heavier than $2m_h$, the decay mode $H \to hh$ becomes possible. So far the non-resonant double Higgs production has not been observed [259] and the limits for resonant dihiggs production are also weak [259–261]. The problem is that the main decay channels of the Higgs have a large background and the overall production rate is not large enough to allow enough statistics for the rare decay modes. There is a small excess roughly at 750 GeV in searches for resonant dihiggs production [259, 262].

For CP-odd scalars the decay to gauge bosons is not allowed. Hence it is similar to the search of the other CP-even state in the alignment limit. The most restrictive bound comes typically from searches of the decay $A \to \tau^+\tau^-$. In Type II 2HDM the $A\tau\tau$ coupling is proportional to $\tan\beta$ so one can exclude high values of $\tan\beta$ by these searches. The region $\tan\beta > 10$ is excluded up to $m_A = 300$ GeV and $\tan\beta > 50$ even up to $m_A = 1$ TeV [253, 254, 263].

For the CP-odd state the decay mode $A \to Zh$ is possible on-shell if $m_A >$ 216 GeV. This is basically studying VH production and looking for resonant production. No significant excesses have been found [260, 264, 265]. The limits for $\sigma(pp \to A) \times B(A \to Zh) \times B(h \to \tau\tau)$ are below 100 fb for a large range of masses.

The most important low-energy constraint comes from the rare decay $B_s \to \mu\mu$. The branching ratio does not seem to deviate largely from its SM value [63, 64]. In extended Higgs models the CP-odd state can be a mediator in this process and change the branching ratio significantly [266]. At large values of

[10]Since the couplings to gauge bosons are small, VBF or VH tagging does not help to cut the background.

$\tan\beta$ the limits depend on $(\tan\beta/m_A)^4$ in Type II 2HDM and hence exclude high values of $\tan\beta$ if the CP-odd state is not too heavy.

For light neutral scalars the bounds are quite restrictive unless the scalar has a large singlet component. Above 300 GeV the viable parameter space becomes larger. There are no convincing hints of new scalar states.

2.5.2 Singly charged scalars

Charged scalars have been searched mostly in the context of two Higgs doublet models. They can either be produced directly or they may be found from top decays. The dominant decay modes for charged Higgs bosons are $\tau\nu$ and $c\bar{s}$ once they are kinematically allowed. If the charged Higgs boson is heavy enough the decay modes to $t\bar{b}$ and W^+h are also important.

Searches at LEP gave a model-independent limit $m_{H^\pm} \geq 79.3$ GeV [267]. The only assumption was that $B(H^+ \to \tau^+\nu) + B(H^+ \to c\bar{s}) = 1$. Recently a combination of charged Higgs pair production was also released, that excludes $m_{H^\pm} < 80$ GeV in Type II 2HDM and $m_{H^\pm} < 72.5$ GeV in Type I 2HDM [268].

Thereafter charged Higgs bosons have been searched in top quark decays at the Tevatron [269–271] and the LHC [75, 76, 272]. CDF could constrain the branching ratio of $t \to H^+b$ below 5.9% assuming the charged Higgs decays solely to $\tau\nu$. The bounds from the LHC on the branching ratio are around one percent. The LHC can exclude the existence of a charged Higgs in Type II 2HDM in the region 80 GeV$< m_{H^\pm} < 155$ GeV if $B(H^\pm \to \tau^\pm\nu) = 1$ [75, 76]. In Type I only low values of $\tan\beta$ are excluded.

The cross section for direct charged Higgs production is rather small, one picobarn or less even at $\sqrt{s} = 14$ TeV for $m_{H^\pm} > m_t$ [273]. Hence only a small part of the parameter space at large values of $\tan\beta$ (or $\cot\beta$) can be excluded [75, 76].

As discussed in section 2.4.3, in the case of triplet models, the charged Higgs can decay to WZ. Searches for such resonances have been made, though typically not interpreted as searches of the charged Higgs. No significant excess is seen [274, 275].

There are also constraints from low-energy precision measurements. Especially the decay $b \to s\gamma$ is sensitive to the charged Higgs. In Type II 2HDM the contribution from the charged Higgs is always positive whereas it can be positive or negative in Type I 2HDM [276]. The branching ratio of $\bar{B} \to X_s\gamma$ has been computed to NNLO in QCD [277,278], giving $(3.15\pm0.23)\times10^{-4}$. The averaged experimental branching fraction is $(3.49 \pm 0.19) \times 10^{-4}$ [90], slightly more than one standard deviation above the SM prediction. In the 2HDM this can be put to a bound $m_{H^\pm} \gtrsim 350$ GeV, using the results of [279]. One must note that this result assumes that there are no other charged particles than those of the 2HDM.

2.5.3 Doubly charged scalars

Doubly charged scalars have only a few possible decay modes. The first option is same sign leptons, the second one is same sign W-bosons if the doubly charged Higgs is heavier than $2m_W$ and the third one is same sign singly charged Higgs bosons if the doubly charged Higgs is heavier than $2m_{H^\pm}$. Also $W^\pm H^\pm$ is possible if the charged Higgs has a triplet component. Doubly charged Higgses can be produced in pairs via $e^+e^-/q\bar{q} \to \gamma^*, Z^* \to H^{++}H^{--}$ [280]. A single doubly charged Higgs can be produced via the fusion of same sign W bosons or singly charged Higgs bosons [281, 282].

DELPHI found a bound $m_{H^{\pm\pm}} \geq 97.3$ GeV if the doubly charged Higgs decays to same sign taus [283]. L3 considered all dilepton final states and got bounds $m_{H^{\pm\pm}} \geq 95.5$ GeV or better [284]. The Tevatron was able to deduce limits from $m_{H^{\pm\pm}} \geq 112$ GeV up to 136 GeV depending on the final state [285, 286].

The LHC has searched for the doubly charged Higgs in the leptonic channels. ATLAS was able to exclude a right-handed doubly charged Higgs up to 374 GeV and a left-handed doubly charged Higgs up to 468 GeV in the final states containing electrons or muons [287]. CMS studied all options for left-handed doubly charged Higgses. For the ditau channel their bound is 396 GeV and for all other channels 481–712 GeV depending on the channel [288] assuming 100% branching ratio to each of these channels. For right-handed doubly charged Higgs bosons the bounds are slightly weaker due to the different couplings to Z bosons and the absence of vector boson fusion.

The bounds on the $W^\pm W^\pm$ decay mode are weak. In the diboson decay mode only a bound of $m_{H^{\pm\pm}} \geq 60$ GeV has been found [289]. They extrapolate their results so that a dedicated search with full Run I data should be sensitive to masses close to 100 GeV.

2.6 Unitarity constraints

In addition to constraints from direct searches of new scalar states, there is also another class of constraints for the quartic scalar interactions. The argument is based on assuming that weak interactions are weak at all energy scales. This is not trivial since massive gauge bosons have a longitudinal polarization state, whose polarization is $\epsilon^\mu_L = (E, 0, 0, p)/m$. At high energies the amplitude for the 4-point interaction in $V_L V_L \to V_L V_L$ scattering grows like $(E_{cm}/m_V)^4$ for any single Feynman diagram.

The total cross section in the forward direction is limited by the optical theorem. If the tree-level cross section increases at high energies beyond the limit of the optical theorem, higher orders in perturbation theory must cancel a part of the tree-level amplitude so that the optical theorem is satisfied. Once this happens the theory becomes non-perturbative.

In Figure 2.5 we have diagrams for $VV \to VV$ scattering neglecting the Higgs contribution. The sum of the amplitudes of diagrams *(a), (b)* and *(c)*

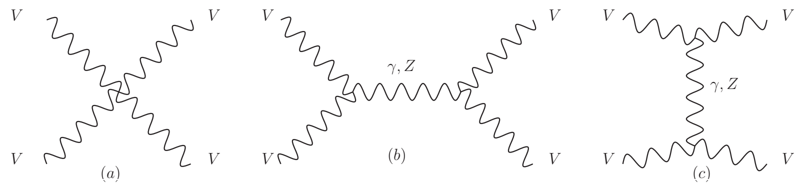

Figure 2.5: Diagrams contributing to $VV \to VV$ ($V = W^{\pm}$, Z) scattering. The amplitude of any single diagram grows like $(E_{cm}/m_V)^4$ but the sum of the diagrams grows more slowly, only proportional to E_{cm}^2. If one includes also the s- and t-channel Higgs exchange diagrams, the amplitude grows only logarithmically.

of Fig. 2.5 grows more slowly than any of the single diagrams due to gauge invariance. However, one needs to add also the Higgs exchange diagrams to get a cross section, which grows only logarithmically at large energies. The Higgs exchange diagrams can be neglected at energies $E_{cm} \ll m_h$ so requiring that the tree-level amplitude does not violate the optical theorem, one gets an upper limit for the Higgs boson mass. This kind of an analysis was first done by Lee, Quigg and Thacker [151, 290]. One writes the tree-level scattering amplitude in terms of partial waves and then sets the bound given by the optical theorem on these.

The bound by Lee, Quigg and Thacker was $m_h \lesssim 1$ TeV. One may also impose a more stringent constraint, as first noted in [291], so that the bound improves by a factor of $\sqrt{2}$ down to $m_h \lesssim 700$ GeV.

After the Higgs discovery we know that there is a state that at least approximately has the couplings of the SM Higgs to gauge bosons. If the couplings are not exactly those of the SM, the amplitude will reach the unitarity bound at a higher energy than if there was no Higgs at all. Such an analysis was performed in [292]. They found that if the Higgs signal strength is at 0.5 (0.8) times the SM value, there will be no violations of unitarity up to 1.7 (2.7) TeV. Similar studies can be found also in [293, 294].

Often in deriving the bounds it is easiest to use the Goldstone boson equivalence theorem, where one may replace the longitudinal vector bosons by the Goldstone bosons of the scalar theory [295] with errors $\mathcal{O}(m_V/E_{cm})$. Hence one essentially will constrain four-point scalar interactions. Such constraints are relevant also in extended Higgs models.

Unitarity constraints have been considered in the singlet extended SM [296–298], the 2HDM [299–302] and also lately in the 3HDM [228]. The unitarity constraints cannot usually be converted to constraints on the full Higgs spectrum, since there will be dimensionful parameters that are not constrained by vacuum conditions, like m_{12}^2 in the 2HDM. We see from equations (2.26)–(2.28) that if $|m_{12}^2| \gg v^2$ the Higgs spectrum can be arbitrarily heavy except for the lightest CP-even state.

Chapter 3

Supersymmetry and Higgs physics

Fundamental scalar masses are not protected by any symmetries in the Standard Model. Hence light fundamental scalars are considered unnatural. In any model with higher mass scales, the Higgs mass has to be fine-tuned and in general this procedure is not stable under radiative corrections.

There are several approaches to this problem. One of them is to assume the Higgs to be composite [303,304]. In this scenario one assumes an asymptotically free gauge symmetry, whose confinement scale is beyond current experimental reach. The Higgs-like state is composed of fermions so there are no elementary scalars that would cause a hierarchy problem. The Higgs-like state can be light if it is the pseudo-Goldstone boson of the theory [305, 306].

Another option that has been considered are extra dimensions. The introduction of compactified extra dimensions brings the Planck scale closer to the electroweak scale, if gravity is the only force that propagates through the extra dimensions. The phenomenologically most interesting alternatives are several flat dimensions [307] or a single extra dimension with a warped metric [308]. Another type of a model is that of universal extra dimensions (UED) [309], where also matter fields propagate through the compactified dimensions.

The third option, which will be studied in more depth in this thesis, is supersymmetry.

3.1 Supersymmetry

Coleman and Mandula studied the possible symmetry groups in particle physics. They proved a theorem [310] which states that a symmetry group that commutes with an analytical S-matrix is always a direct product of the Poincaré group and an internal symmetry group or the S-matrix becomes trivial. Supersymmetry (SUSY) extends conventional quantum field theory by introducing anticommuting symmetry generators and hence the assumptions of the Coleman-Mandula theorem do not apply. Although the first models with supersymmetry were introduced already earlier [311–314], Haag, Łopuszanski and Sohnius [315] were the first ones to generalize the discussion of Coleman and Mandula to the case of generators with a fermionic nature. They found that the only possibility was to have generators belonging to either the $(\frac{1}{2}, 0)$ or the $(0, \frac{1}{2})$ representation of the Poincaré group. In addition the algebra is very restricted.

The algebra of the Poincaré group is

$$
\begin{aligned}
[P^\mu, P^\nu] &= 0, && (3.1)\\
[M^{\mu\nu}, P^\rho] &= i(g^{\nu\rho} P^\mu - g^{\mu\rho} P^\nu), && (3.2)\\
[M^{\mu\nu}, M^{\rho\sigma}] &= i(g^{\mu\sigma} M^{\nu\rho} + g^{\nu\rho} M^{\mu\sigma} - g^{\mu\rho} M^{\nu\sigma} - g^{\nu\sigma} M^{\mu\rho}), && (3.3)
\end{aligned}
$$

where $M^{\mu\nu}$ is an antisymmetric tensor with components $M^{ij} = \epsilon^{ijk} J_k$, J_k being the generator of rotations around the axis k, and $M^{0i} = -K^i$ with K^i being the generator of boosts in the direction i.

Supersymmetry extends this algebra with anticommuting generators $Q_\alpha \in (\frac{1}{2}, 0)$ and $\bar{Q}_{\dot{\alpha}} \in (0, \frac{1}{2})$, which satisfy the following commutation and anticommutation relations:

$$
\begin{aligned}
\{Q_\alpha, Q_\beta\} &= 0 = \{\bar{Q}_{\dot{\alpha}}, \bar{Q}_{\dot{\beta}}\}, && (3.4)\\
[Q_\alpha, P^\mu] &= 0 = [\bar{Q}_{\dot{\alpha}}, P^\mu], && (3.5)\\
\{Q_\alpha, \bar{Q}_{\dot{\alpha}}\} &= 2\sigma^\mu_{\alpha\dot{\alpha}} P_\mu, && (3.6)\\
[M^{\mu\nu}, Q_\alpha] &= -i(\sigma^{\mu\nu})_\alpha{}^\beta Q_\beta, && (3.7)\\
\left[M^{\mu\nu}, \bar{Q}^{\dot{\alpha}}\right] &= -i(\bar{\sigma}^{\mu\nu})^{\dot{\alpha}}{}_{\dot{\beta}} \bar{Q}^{\dot{\beta}}. && (3.8)
\end{aligned}
$$

Here the indices on Weyl spinors can be raised and lowered by $\epsilon_{\alpha\beta} = i\sigma^2 = \epsilon_{\dot{\alpha}\dot{\beta}}$ and $\epsilon^{\alpha\beta} = -i\sigma^2 = \epsilon^{\dot{\alpha}\dot{\beta}}$: $\bar{Q}^{\dot{\alpha}} = \epsilon^{\dot{\alpha}\dot{\beta}} \bar{Q}_{\dot{\beta}}$. We also define $\sigma^{\mu\nu} = \frac{1}{4}(\sigma^\mu \bar{\sigma}^\nu - \sigma^\nu \bar{\sigma}^\mu)$ and $\bar{\sigma}^{\mu\nu} = \frac{1}{4}(\bar{\sigma}^\mu \sigma^\nu - \bar{\sigma}^\nu \sigma^\mu)$.

In principle one could have more than one set of fermionic generators. If one restricts oneself to models with no elementary particles with spins $s > 2$, the maximum number of generators is eight. These models are called extended supersymmetries, which are of significant theoretical interest (*e.g.* [316–323]), but they are not phenomenologically viable, since fermions are either in the adjoint representation of the gauge group or fermions of both chirality are in the same supermultiplet. We shall restrict ourselves to one set of SUSY generators, so called $N = 1$ supersymmetry only.

By using the rules of addition of angular momentum, one finds that applying supersymmetric generators produces a state that differs by half a unit of spin from the original state. Since $[Q_\alpha, P^\mu] = 0$, the states are degenerate in mass. Hence supersymmetry predicts states having the same mass and differing by half a unit of spin. Such particles have not been observed, so supersymmetry must be broken.

The anticommutator (3.6) also implies that the energy of a supersymmetric vacuum, for which $Q_\alpha|\Omega\rangle = 0$ and $\bar{Q}^{\dot{\alpha}}|\Omega\rangle = 0$, is zero.

3.1.1 Superfields

The first supersymmetric models in four dimensions were constructed by Wess and Zumino [324]. They constructed what are now known as the chiral supermultiplet and the vector supermultiplet by more or less brute force and intelligent guessing. One can construct supermultiplets directly by applying SUSY transformations of the form $e^{i(\xi Q+\bar{\xi}\bar{Q})}$ and starting with states of definite spin. The more systematic way to construct supermultiplets uses the concepts of superspace and superfields, introduced by Salam and Strathdee [325–327].

In superspace one has in addition to the usual 4D spacetime coordinates x^μ four dimensions of Grassmannian coordinates θ_α and $\bar{\theta}^{\dot{\alpha}}$, both being two-component anticommuting variables. We define $\theta\theta = \theta^\alpha\theta_\alpha$ and $\bar{\theta}\bar{\theta} = \bar{\theta}_{\dot{\alpha}}\bar{\theta}^{\dot{\alpha}}$.

The most general commuting field in superspace can be written as

$$\begin{aligned}\Phi(x,\theta,\bar{\theta}) = \varphi(x) + \theta\psi(x) + \bar{\theta}\bar{\chi}(x) + \theta\theta M(x) + \bar{\theta}\bar{\theta}N(x) + \theta\sigma^\mu\bar{\theta}V_\mu(x) \\ + \theta\theta\bar{\theta}\bar{\lambda}(x) + \bar{\theta}\bar{\theta}\alpha(x) + \theta\theta\bar{\theta}\bar{\theta}D(x). \end{aligned} \quad (3.9)$$

This contains four complex scalar fields, four complex Majorana spinors (two of both chirality) and a vector field. This representation is reducible.

The most systematic way to find the irreducible representations is to introduce supercovariant derivatives D_α and $\bar{D}_{\dot{\alpha}}$, which anticommute with SUSY generators and commute with the generators of the Poincaré group [327]. Then one imposes constraints with the help of these supercovariant derivatives. These constraints will then be invariant under supersymmetry.

The form of the supercovariant derivatives depends on the choice of SUSY transformations[1] [328]. If one uses a linear representation in where [329]

$$Q_\alpha = i\left[\frac{\partial}{\partial\theta^\alpha} + i\sigma^\mu_{\alpha\dot{\alpha}}\bar{\theta}^{\dot{\alpha}}\partial_\mu\right], \quad \bar{Q}^{\dot{\alpha}} = i\left[\frac{\partial}{\partial\bar{\theta}_{\dot{\alpha}}} + i\bar{\sigma}^{\mu\alpha\dot{\alpha}}\theta_\alpha\partial_\mu\right] \quad (3.10)$$

the supercovariant derivatives are of the form

$$D_\alpha = \frac{\partial}{\partial\theta^\alpha} - i\sigma^\mu_{\alpha\dot{\alpha}}\bar{\theta}^{\dot{\alpha}}\partial_\mu, \quad \bar{D}_{\dot{\alpha}} = \frac{\partial}{\partial\bar{\theta}^{\dot{\alpha}}} - i\theta^\alpha\sigma^\mu_{\alpha\dot{\alpha}}\partial_\mu \quad (3.11)$$

and these satisfy $\{D_\alpha, \bar{D}_{\dot{\alpha}}\} = -2i\sigma^\mu_{\alpha\dot{\alpha}}\partial_\mu$.

A superfield is called left-chiral if it satisfies the equation $\bar{D}_{\dot{\alpha}}\Phi = 0$. Such a field can be given in the form

$$\Phi_L(x,\theta,\bar{\theta}) = \varphi(y) + \theta\psi(y) + \theta\theta F(y), \quad (3.12)$$

where $y^\mu = x^\mu - i\theta\sigma^\mu\bar{\theta}$. A right-chiral superfield is defined by the equation $D_\alpha\Phi = 0$ and it depends on $x^\mu + i\theta\sigma^\mu\bar{\theta}$ and $\bar{\theta}$. The chiral superfield contains a scalar field φ, a chiral spinor field ψ and an auxiliary field F, which can be given in terms of the other fields once the equations of motion are solved.

[1] Since Grassmann variables do not commute, $e^{i(\xi Q+\bar{\xi}\bar{Q})} \neq e^{i\xi Q}e^{i\bar{\xi}\bar{Q}}$ etc.

The auxiliary field $F(y)$ transforms as a total derivative under supersymmetry transformations. It can thus be considered as a candidate Lagrangian for supersymmetric theories with chiral superfields.

The supercovariant derivatives follow the generalized Leibniz rule $D_\alpha(\Phi_1\Phi_2) = (D_\alpha\Phi_1)\Phi_2 \pm \Phi_1(D_\alpha\Phi_2)$, where the sign is plus (minus) when Φ_1 is bosonic (fermionic), so the product of two left-chiral (right-chiral) superfields is also a left-chiral (right-chiral) superfield. The superpotential is a polynomial of superfields with one chirality only. The F-term of a trilinear superpotential contains Yukawa interactions and quartic scalar interactions. If the superpotential is gauge invariant, the F-term and hence the Lagrangian will also be gauge invariant.

The scalar potential for a given superpotential W can be computed from

$$V = \sum_i \left| \frac{\partial W}{\partial \varphi_i} \right|^2, \tag{3.13}$$

where the sum is taken over all of the fields and the superfields are replaced by their scalar components.

The Yukawa interactions are given by

$$\mathcal{L} = -\frac{1}{2} \sum_{i,j} \frac{\partial^2 W}{\partial \varphi_i \partial \varphi_j} \Psi_i \Psi_j, \tag{3.14}$$

where the sum is taken over all superfields and the superfields in W are replaced by their scalar components.

The other superfield needed to construct the supersymmetric extension of the SM is the vector superfield. It can be constructed by two methods. Historically the first one was to impose a reality condition for the general superfield (3.9). The second way is to combine a left-handed and a right-handed superfield, since the $(\frac{1}{2}, \frac{1}{2})$ representation of the Poincaré group corresponds to a vector field.

When the reality condition has been imposed on the general superfield (3.9), it may be put to the form

$$\begin{aligned} V(x,\theta,\bar\theta) &= C(x) + i\theta\chi(x) - i\bar\theta\bar\chi(x) + \frac{i}{2}\theta\theta(M(x)+iN(x)) \\ &\quad - \frac{i}{2}\bar\theta\bar\theta(M(x)-iN(x)) + \theta\sigma^\mu\bar\theta V_\mu + i\theta\theta\bar\theta\left(\bar\lambda(x) + \frac{i}{2}\bar\sigma^\mu\partial_\mu\chi(x)\right) \\ &\quad - i\bar\theta\bar\theta\theta\left(\lambda(x) + \frac{i}{2}\sigma^\mu\partial_\mu\bar\chi(x)\right) + \frac{1}{2}\theta\theta\bar\theta\bar\theta\left(D(x) - \frac{1}{2}\partial^\mu\partial_\mu C(x)\right). \end{aligned} \tag{3.15}$$

When V is written with this rather cumbersome form, $V_{\mu\nu} \equiv \partial_\mu V_\nu - \partial_\nu V_\mu$, λ, $\bar\lambda$ and D transform among themselves under supersymmetry transformations and hence form an irreducible representation containing a vector field and a spinor field.

The D-field is an auxiliary field, which transforms as a total derivative under SUSY transformations. It gives the candidate Lagrangian for vector superfields.

In a more general setting, the D-term contribution comes from the Kähler potential

$$K(\Phi^\dagger, \Phi) = G_{mn}\Phi_m^\dagger \Phi_n + \dots, \tag{3.16}$$

where G_{mn} is called the Kähler metric[2], Φ's are chiral superfields and the dots represent nonrenormalizable terms.

One may also construct a chiral superfield from the vector superfield. This is known as the field strength superfield, which can be defined as

$$W_\alpha = \bar{D}^2 D_\alpha V(x, \theta, \bar{\theta}). \tag{3.17}$$

This is clearly a chiral superfield, since $\bar{D}^3 \equiv 0$. The F-term of $\frac{1}{32}W^\alpha W_\alpha$ gives the kinetic terms for the vector field and the associated spinor field and a D-field contribution to the scalar potential in the form of $\frac{1}{2}D^2$.

Chiral and vector superfields are the only ones needed for the supersymmetric generalization of the SM. One may still construct other types of superfields. A real linear superfield [328, 330, 331] satisfies the constraints

$$D^2\Phi = 0, \quad \bar{D}^2\Phi = 0. \tag{3.18}$$

The real linear superfield consists of a real scalar field, a Majorana spinor and an antisymmetric rank-2 tensor field. None of these are auxiliary fields.

One can also define a complex linear superfield by requiring only one of the conditions (3.18), but the quantization of such a field needs an infinite number of ghost fields [332].

3.1.2 Supersymmetric gauge transformations

The SM is based on gauge invariance. In order to make a supersymmetric generalization of the SM one needs a prescription for supersymmetric gauge transformations.

The supersymmetric generalization of quantum electrodynamics (QED) was developed by Wess and Zumino [333]. They noticed that adding a chiral superfield to the vector superfield induces a gauge transformation for the vector field, leaves the other parts of the irreducible multiplet unchanged and allows to eliminate the other degrees of freedom completely by choosing a suitable gauge, now known as the Wess-Zumino gauge.

The supersymmetric non-Abelian gauge transformations were first derived by Ferrara and Zumino [334]. The vector superfields are transformed like $V(x,\theta,\bar{\theta}) \to V(x,\theta,\bar{\theta}) + i(\Lambda - \Lambda^\dagger)$, where Λ is a chiral superfield characterizing the gauge transformation.

If we assume a chiral superfield to transform under a gauge transformation like $\Phi \to e^{-2igt^a\Lambda_a}\Phi$, we find that the combination

$$\Phi^\dagger e^{2gV}\Phi, \tag{3.19}$$

[2]Often it is assumed that the Kähler metric is flat, *i.e.* G_{mn} is a unit matrix.

where $V = t^a V_a$, is gauge invariant. It is a vector superfield, whose D-term is the candidate Lagrangian. It contains the kinetic terms for the components of the chiral superfield, the gaugino interactions with fermion-sfermion pairs and quartic scalar interactions.

To get the F-term contribution from the generalization of the field strength superfield, one needs to generalize the supercovariant derivatives to be also gauge invariant. One may choose $\nabla_{\dot{\alpha}} = D_{\dot{\alpha}}$ and $\nabla_\alpha = e^{-V} D_\alpha e^V$, where $V \equiv 2gt^a V_a$. These operators satisfy $\{\nabla_\alpha, \nabla_{\dot{\alpha}}\} = 2i\sigma^\mu_{\alpha\dot{\alpha}}\partial_\mu$. The generalized field strength superfield is now

$$W_\alpha = \frac{1}{2g}\bar{\nabla}^2\nabla_\alpha = \frac{1}{2g}\bar{D}^2 e^{-V}(D_\alpha e^V). \tag{3.20}$$

The kinetic terms for the gauge fields and gauginos and the self-interactions of the gauge field can be obtained from $\frac{1}{64}(W^\alpha W_\alpha)_F$. This gives also a contribution to the scalar potential. After eliminating the auxiliary fields, the scalar potential becomes

$$V = \frac{1}{2}D^a D_a = \frac{1}{2}\sum_a g^2|\varphi^i t^a_{ij}\varphi^j|^2, \tag{3.21}$$

where the sum runs over all of the generators of the gauge group.

One can write the F- and D-terms formally as Grassmann integrals, since $\int \mathrm{d}\theta\ \theta = 1$. The Lagrangian can be written as

$$\mathcal{L} = \int \mathrm{d}^2\theta \left(W(\Phi) + \frac{1}{64}W^\alpha W_\alpha\right) + \text{h.c.} + \int \mathrm{d}^2\theta \mathrm{d}^2\bar{\theta}\ \Phi^\dagger e^{2gV}\Phi. \tag{3.22}$$

3.1.3 Nonrenormalization theorems

The most important theoretical feature of supersymmetry from the point of view of model building is the cancellation of radiative corrections in the supersymmetric limit. Formally this can be stated in the form of nonrenormalization theorems [335–338].

One may prove that any perturbative correction can be written as a D-term and hence the superpotential will not receive loop corrections at any order in perturbation theory. D-terms give the kinetic terms of the Lagrangian and they need only wave function renormalization. The important consequences are that quadratic divergences do not appear in any order of perturbation theory and that new terms are not generated to the superpotential by radiative corrections even if they were allowed by symmetries. In conventional quantum field theories one needs to include all terms allowed by symmetries to get all counterterms needed in the renormalization procedure.

Also the loop contribution to the vacuum energy vanishes to all orders in perturbation theory in theories with exact supersymmetry [339, 340].

Superfield	Fermion	Boson	Superpartner name
$\hat{Q} = (\hat{U}\ \hat{D})^T$	$(u_L\ d_L)^T$	$(\tilde{u}_L\ \tilde{d}_L)^T$	Left-handed up- and down-type squarks
$\hat{U}^c$	$\bar{u}_R$	$\tilde{u}_R^*$	Right-handed up-type squarks
$\hat{D}^c$	$\bar{d}_R$	$\tilde{d}_R^*$	Right-handed down-type squarks
$\hat{L} = (\hat{N}\ \hat{E})^T$	$(\nu_L\ \ell_L^-)^T$	$(\tilde{\nu}_L\ \tilde{\ell}_L^-)^T$	Left-handed sleptons
$\hat{E}^c$	ℓ_R^+	$\tilde{\ell}_R^+$	Right-handed charged sleptons

Table 3.1: The superfields corresponding to SM fermions. All superfields are left-chiral superfields. The superscript c denotes charge conjugation. The representations of these fields under the SM gauge group are the same as in Table 2.1.

3.2 The MSSM and its Higgs sector

The supersymmetric model with the minimal field content that contains the Standard Model is called the Minimal Supersymmetric Standard Model (MSSM), first formulated by Fayet[3] [341, 342]. As we discuss later, there are reasons to believe that the MSSM can not be the final theory but it still is a useful benchmark for studies of supersymmetry.

3.2.1 Field content

The MSSM contains all the SM particles. None of these have a full set of identical quantum numbers and a 1/2 unit difference in spins so none of the superpartners have been found. Hence supersymmetrizing the SM gives two new complex scalars per each Dirac fermion and one Majorana fermion per each massless gauge boson. The resulting scalar superpartners of SM fermions are called squarks and sleptons and the superpartners of SM bosons get an 'ino' suffix.

The superpotential needs to be constructed of superfields with one chirality only. Since SM fermions have both left and right chiralities, one needs to introduce right-chiral fermions as left-handed charge conjugated superfields, which have similar transformation properties. The MSSM superfields corresponding to SM fermions are given in Table 3.1. The physical squark and slepton states are mixtures of the left- and right-handed[4] states, although the mixing is very small in all but the stop sector.

However, as seen from Eq. (2.11), in the SM one needs both the fundamental and the charge conjugated representation of the Higgs doublet to generate masses for all fermions. The charge conjugated representation of a single Higgs

[3]Fayet assumed that the neutrino is the Goldstone fermion associated with spontaneous supersymmetry breaking. Nowadays it is known that neutrinos are not massless. Goldstone fermions can be "eaten" by the super-Higgs mechanism so they are not thought to be in the physical spectrum.

[4]Since squarks and sleptons are scalars, the particles are not left- or right-handed. They are often called so since they are superpartners of left- and right-handed fermions.

SM boson		Superpartner		Superpartner	
Gluon	g	Gluino	$\tilde{g}$	Gluino	$\tilde{g}$
W-boson Charged Higgs	$W^\pm$ $H^\pm$	Wino Higgsino	$\tilde{W}^\pm$ $\tilde{H}^\pm$	Chargino	$\tilde{\chi}_1^\pm$,$\tilde{\chi}_2^\pm$
Neutral W Hypercharge boson Higgs bosons	W^0 B^0 $h,\ H,\ A$	Neutral wino Bino Higgsinos	$\tilde{W}^0$ $\tilde{B}^0$ $\tilde{H}_1,\ \tilde{H}_2$	Neutralino	$\tilde{\chi}_1^0,\ \tilde{\chi}_2^0$ $\tilde{\chi}_3^0,\ \tilde{\chi}_4^0$

Table 3.2: The SM bosons and their superpartners. Except for the gluinos the mass eigenstates and the gauge eigenstates do not coincide. The gauge eigenstates are given in the second column and the mass eigenstates in the third column. The photon and the Z-boson are linear combinations of the W^0- and B^0-fields as explained on p. 13.

doublet is of the wrong chirality to be included in the superpotential. The solution is to add a $Y = -1$ Higgs doublet superfield. This also cancels the triangle anomaly that would be otherwise generated by the Higgsinos.

The MSSM superpotential is

$$W = y^{(u)}(Q \cdot H_u)U^c + y^{(d)}(Q \cdot H_d)D^c + y^{(\ell)}(L \cdot H_d)E^c + \mu(H_u \cdot H_d). \quad (3.23)$$

The dot product is $A \cdot B \equiv \epsilon_{ab}A^aB^b$. The Yukawa terms are those of the SM, but there is a term with a dimensionful parameter μ. The μ-term has to be nonzero or there will be a light chargino that would have been detected at LEP. On the other hand, there is a tree-level relation

$$\frac{1}{2}m_Z^2 = -\frac{m_{H_u}^2 \tan^2\beta - m_{H_d}^2}{\tan^2\beta - 1} - |\mu|^2, \quad (3.24)$$

where $m_{H_u}^2$ and $m_{H_d}^2$ are soft supersymmetry breaking parameters. This relation suggests that μ should not be too much above the electroweak scale or otherwise it should cancel almost exactly the contribution from the soft supersymmetry breaking parameters. If there are high mass scales, like that of grand unification, one could naively expect μ to get such a high value. The cancellation in Eq. (3.24) would then seem unlikely. There are models, where the μ-term is related to the supersymmetry breaking scale as was shown by Giudice and Masiero [343]. In such a case the cancellation is not so unnatural.

The gauge sector needs vector superfields for the SM gauge fields. In the fermionic sector the fermionic partners of the electroweak gauge fields get mixed with the fermionic partners of the Higgs fields. The mixtures of gauge eigenstates are called neutralinos and charginos. The SM bosons and their superpartners are given in Table 3.2.

3.2.2 R-parity

The MSSM superpotential (3.23) is not the most general gauge invariant and renormalizable one that could be constructed out of MSSM superfields. One could also add the following terms to the superpotential:

$$W_{\mathrm{RPV}} = \lambda_{ijk}(L_i \cdot L_j)E^c_k + \lambda'_{ijk}(L_i \cdot Q_j)D^c_k + \lambda''_{ijk}U^c_i D^c_j D^c_k + \mu_i(L_i \cdot H_u). \quad (3.25)$$

These terms violate either lepton number or baryon number, both of which are so called accidental symmetries of the SM. There is no evidence of lepton or baryon number violation experimentally so all of the terms in Eq. (3.25) must be small. The most constraining limit comes from the absence of proton decay, which restricts some products of $\lambda\lambda''$- and $\lambda'\lambda''$-terms to be $\mathcal{O}(10^{-26})$ [344,345], though some combinations are constrained only to a level of $\mathcal{O}(10^{-10})$ [346]. The constraints from proton decay are absent if either all baryon number violating or all lepton number violating terms are zero. Constraints from various other processes are reviewed *e.g.* in [347–350].

Even though in supersymmetry one could set these terms to zero by hand and due to nonrenormalization theorems they would not be regenerated radiatively, one can also impose a symmetry that forbids these terms. That is a discrete Z_2 symmetry called R-parity [341,351,352]. The quantum numbers are assigned so that SM particles have $R = +1$ and their superpartners $R = -1$. The R-parity can be defined *e.g.* by

$$R = (-1)^{3(B-L)+2s}, \quad (3.26)$$

where B and L are baryon and lepton number, respectively, and s is spin.

R-parity conservation implies that R-odd particles can only be produced in pairs and in the decay chain of a R-odd particle there will always be an odd number of R-odd particles. Hence the lightest R-odd particle must be stable. If it is neutral, it is the dark matter candidate of supersymmetry. R-parity conserving models where the lighest supersymmetric particle (LSP) is charged or colored are not considered viable.

Relaxing the assumption of R-parity will be discussed in chapter 4.

3.2.3 Soft supersymmetry breaking

Nature has shown us that supersymmetry must be broken. In the MSSM supersymmetry is broken by adding the most general set of soft supersymmetry breaking terms. By soft breaking we mean that the introduction of SUSY breaking does not generate quadratic divergences. One may add mass terms for the scalar components of chiral superfields, Majorana mass terms for the fermionic components of vector superfields and trilinear scalar interactions without introducing quadratic divergences [353].

There are several ways to introduce the soft terms. The agnostic way is to simply write the soft terms as independent parameters. When the soft breaking is given by the Lagrangian

$$
\begin{aligned}
\mathcal{L} = &- \Big[m_{\tilde{Q}}^2 Q^\dagger Q + m_{\tilde{U}}^2 \tilde{u}_R^* \tilde{u}_R + m_{\tilde{D}}^2 \tilde{d}_R^* \tilde{d}_R + m_{\tilde{L}}^2 L^\dagger L + m_{\tilde{E}}^2 \tilde{e}_R^* \tilde{e}_R + m_{H_u}^2 H_u^\dagger H_u \\
&+ m_{H_d}^2 H_d^\dagger H_d \Big] - \frac{1}{2} M_1 \tilde{B} \tilde{B} - \frac{1}{2} M_2 \tilde{W}_i \tilde{W}_i - \frac{1}{2} M_3 \tilde{g}_j \tilde{g}_j + \Big[A^{(u)} (\tilde{Q}_L \cdot H_u) \tilde{u}_R^* \\
&+ A^{(d)} (\tilde{Q}_L \cdot H_d) \tilde{d}_R^* + A^{(\ell)} (\tilde{L}_L \cdot H_d) \tilde{e}_R^* - m_3^2 H_u \cdot H_d + \text{h.c.} \Big] \qquad (3.27)
\end{aligned}
$$

the MSSM contains 105 new physical parameters compared to the SM [354].

There are various ways to reduce the number of free parameters. One is to assume that scalar masses have a common value m_0^2, gaugino masses a common value $M_{1/2}$ and trilinear scalar couplings a common value A_0 (that multiplies the superpotential coupling) at some high scale and then use renormalization group running [355] to obtain their values at the electroweak scale. Such a scenario can be motivated by supergravity [356, 357].

Universal soft parameters at the high scale is not the most general alternative [358]. If the Kähler potential does not have a flat metric, the soft SUSY breaking parameters will carry the pattern of the Kähler metric. Some of the simplest alternatives are those of non-universal Higgs masses, where the soft Higgs masses differ from the other soft scalar masses [359–361]. One could also have nonuniversal masses for the third generation of sfermions [362]. For the two first generations degeneracy is favored due to flavor constraints.

Gaugino masses can also be non-universal [363, 364]. In that case the masses at high scale have non-trivial relations. If one assumes grand unification at some high scale, the ratio of gaugino masses and gauge couplings squared is a constant in the case of universal gaugino masses, leading to the prediction that gluinos are heavier than electroweak gauginos. Non-universal gaugino masses can lead to different realizations of the mass pattern [365].

An intermediate approach is that of the phenomenological MSSM (pMSSM) [366]. With the help of certain simplifying assumptions, the number of free parameters is reduced to 19. These include three gaugino masses, soft scalar masses for the first two generations, different soft masses for the third sfermion generation, two parameters for the Higgs sector ($\tan\beta$ and m_A^2), the μ parameter and trilinear scalar couplings for the third generation. These can be then defined at any convenient scale.

The models with a small number of SUSY breaking parameters are predictive and can be constrained easily. The part of the parameter space where squarks and gluinos are below 1 TeV has been already excluded [367]. This motivates to explore also the alternatives, where the soft parameters are not universal as they allow a larger variety of collider signatures. The exclusion of a large part of the parameter space in the simplest models also is a strong motivation to study models beyond the MSSM.

3.2.4 Origin and mediation of supersymmetry breaking

The soft supersymmetry breaking discussed in the previous section is believed to have a dynamical origin. Supersymmetry can be spontaneously broken by non-zero values of F-terms [368] or D-terms [369]. The breaking cannot originate from the MSSM sector, since there are tree-level sum rules, which state that the supertrace of the mass squared matrix vanishes [370]

$$\mathrm{Str}(M^2) = \sum_i (-1)^{2s_i}(2s_i+1)m_i^2 = 0, \tag{3.28}$$

where s_i is the spin and m_i the mass of particle species i. The validity of this sum rule requires that the gauge groups are anomaly-free, which is true in the MSSM. Since the mass matrices of particles with different charges do not mix, this would mean that some of the sleptons need to be lighter than the tau and they would have been observed already.

Hence supersymmetry is assumed to be broken in a sector that is currently unknown (dubbed hidden sector) and the breaking effects are then mediated to the observable sector.The breaking can be mediated to the observable sector in three known ways: gravity mediated supersymmetry breaking, gauge mediated supersymmetry breaking (GMSB) and anomaly mediated supersymmetry breaking (AMSB). Their basic features are summarized in the following.

Supersymmetry, like internal symmetries, can be promoted to a local symmetry. As one needs gauge fields to compensate the additional terms in conventional theories, one needs a new superfield in local supersymmetry. Since the transformation carries spin, the new field will have spin 3/2. It is the superpartner of the graviton, which has spin 2. Theories with the graviton supermultiplet are called supergravities [371, 372].

Spontaneous supersymmetry breaking will create a Goldstone fermion called Goldstino [369]. In supergravity the supersymmetry breaking will be seen as a massive gravitino. A massive spin 3/2 field has four spin states, a massless one has only two. The two additional spin states come from the Goldstino in the super-Higgs mechanism [373–375] similarly to how the Goldstone boson is eliminated in the conventional Higgs mechanism. The gravitino mass can be used as an order parameter of supersymmetry breaking.

Gravity can mediate supersymmetry breaking to the observable sector. The hidden sector and the mediators are assumed to be singlets under the SM gauge group. In the simplest case of using the Polonyi superpotential[5] for the hidden sector and a flat Kähler metric, the low energy scalar potential will have universal soft scalar masses to the scalar components of chiral MSSM superfields and trilinear scalar couplings [356, 357]. The scalar masses will be $\mathcal{O}(m_{3/2}^2)$ and the trilinear couplings will be $\mathcal{O}(\lambda m_{3/2})$, where λ is the corresponding superpotential coupling. The problem of generating large gaugino masses is a more difficult one [376], but models with gaugino masses $\mathcal{O}(m_{3/2})$ do exist [377].

[5] The Polonyi superpotential is $W = m^2(\Phi + \beta)$, where m and β are constants and Φ a chiral superfield.

Gravity will always break supersymmetry. When we talk about the other mechanisms, we assume that the effect from gravity to the supersymmetry breaking terms is subleading or negligible.

Gauge mediation [378] assumes that the mediating sector is charged under the SM gauge group. The hidden sector then couples to these mediators either via superpotential couplings or via some other gauge interactions under which the MSSM fields are singlets.

The soft gaugino masses are generated at one-loop order and the soft scalar masses and trilinear scalar couplings at two-loop order. However, since scalar masses have a dimension mass squared, they are of the same order as gaugino masses. Trilinear scalar couplings have the dimension of mass, so they are suppressed compared to the other soft terms. Hence they are often neglected completely. However, it is possible to extend the minimal setup for gauge mediation in a way that large A-terms are generated [379].

In GMSB the gravitino is light, since the contribution from gravity is small. Hence often the gravitino will be the LSP. On the other hand the gravitino couples very weakly to matter so the lifetime of the next-to-lightest supersymmetric particle (NLSP) is typically so large that it will not decay within the detector system[6].

The third option for mediating supersymmetry breaking is anomalies. Randall and Sundrum noted that the super-Weyl anomaly can mediate the supersymmetry breaking to the auxiliary fields [382]. Similar results were obtained by Giudice *et al.* [383]. The difference to other models of supersymmetry breaking is that the ratios of supersymmetry breaking parameters are determined by the β-functions, which determine the renormalization group evolution and the anomalous dimensions. The gaugino mass pattern differs from other mechanisms, the wino being the lightest gaugino. For a realistic mass spectrum the gravitino needs to be rather heavy. The contribution from gravity to soft supersymmetric terms can be suppressed if the hidden sector is chosen appropriately.

The problem of pure anomaly mediation is that it predicts negative masses squared for sleptons. The minimal setup to cure this problem is to assume that there is a universal contribution to soft scalar masses from some other SUSY breaking sector [384].

3.2.5 The Higgs sector at tree-level

From the superpotential (3.23) we see that the Higgs sector is that of Type II 2HDM, with one doublet having $Y = +1$ and the other $Y = -1$. In addition, the quartic couplings are determined by supersymmetry. This gives $\lambda_6 = \lambda_7 = 0$ (in the notation of Eq. (2.20)) so Higgs sector tree-level FCNC are absent.

The other quartic couplings are $\lambda_1 = \lambda_2 = \frac{1}{4}(g^2 + g'^2)$, $\lambda_3 = \frac{1}{4}(g^2 - g'^2)$, $\lambda_4 = -\frac{1}{2}g^2$ and $\lambda_5 = 0$. Inserting these to equations (2.26) and (2.27) gives immediately the relation

$$m_{H^\pm}^2 = m_A^2 + m_W^2, \tag{3.29}$$

[6] At high energies and low gravitino masses the gravitino can be effectively replaced by the goldstino state so the decays to gravitinos may also be relevant [380, 381].

which implies that the charged Higgs cannot be light in the MSSM. The CP-odd state has a mass $m_A^2 = 2m_3^2/\sin 2\beta$ so it can be considered as a free parameter.

In the CP-even sector the mass matrix is reduced to

$$\begin{pmatrix} m_3^2 \tan\beta + m_Z^2 \cos^2\beta & m_3^2 - m_Z^2 \sin\beta\cos\beta \\ m_3^2 - m_Z^2 \sin\beta\cos\beta & m_3^2 \cot\beta + m_Z^2 \sin^2\beta \end{pmatrix}. \tag{3.30}$$

The eigenvalues of this matrix are

$$m_{h,H}^2 = \frac{m_A^2 + m_Z^2 \pm \sqrt{(m_A^2 + m_Z^2)^2 - 4m_A^2 m_Z^2 \cos^2 2\beta}}{2}. \tag{3.31}$$

The smaller eigenvalue gets its maximum value when $\cos^2 2\beta = 1$, so we get the bound

$$m_h \leq \min(m_A, m_Z) \leq 91.2 \text{ GeV}. \tag{3.32}$$

The heavier CP-even Higgs mass is bounded from below. The smallest bound comes in the limit $m_A \to 0$ so that $m_H \geq m_Z$. If $m_A \gg m_Z$ we have the decoupling limit, where H, A and $H^\pm$ are roughly degenerate with relative mass differences $\mathcal{O}(m_Z^2/m_A^2)$.

As one may anticipate from Eq. (3.30), also the relationship between the Higgs sector mixing angles is simplified compared to the general 2HDM. The mixing angles are related by

$$\tan 2\alpha = \tan 2\beta \cdot \frac{m_A^2 + m_Z^2}{m_A^2 - m_Z^2}. \tag{3.33}$$

These tree-level relations were first derived in [385, 386].

The bound on the lightest Higgs mass and the observation of a 125 GeV SM-like Higgs is a severe constraint for supersymmetric models. Since the smallest eigenvalue of a matrix is smaller than the smallest eigenvalue of any submatrix, the bound can always be computed from the 2×2 matrix of Eq. (3.30). There are four ways out:

1. The 125 GeV state is not the lightest one. The searches for light scalars favor a singlet-dominated lighter state as discussed in section 2.5.1. If this mixes with the lighter doublet state, the mixing can lift the mass at tree-level.

2. There are additional tree-level contributions from F-terms to the matrix (3.30). This requires superfields that couple to the Higgs doublets in the superpotential and whose scalar components get a non-zero VEV.

3. There are additional tree-level contributions from D-terms to the matrix (3.30). This requires an extended gauge group.

4. There are corrections from loops to the matrix (3.30).

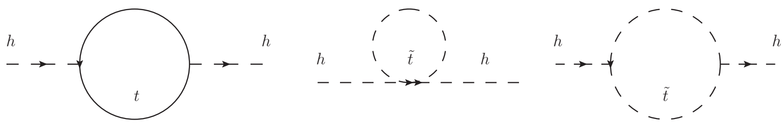

Figure 3.1: The top and stop loops give a large positive correction to the Higgs mass. It is possible to lift the Higgs mass to 125 GeV if the stops are heavier than about 1.5 TeV.

In the MSSM only the last solution is possible[7]. Each of these options can be realized in various nonminimal models, which is a strong motivation for studying them. We shall consider each of the four mechanisms in the later parts of this thesis.

3.2.6 Higgs masses at the loop-level

The MSSM would be excluded if loop corrections could not produce a large positive correction to the lightest Higgs boson mass. The most common way to compute loop corrections to Higgs masses is to use the effective potential approach [390]. The one-loop effective potential may be computed from

$$V^{(1)} = \frac{1}{64\pi^2}\mathrm{Str}\left\{M^4(\varphi_i)\left[\ln\left(\frac{M^2(\varphi_i)}{Q^2}\right) - C_s\right]\right\}, \tag{3.34}$$

where $M(\varphi_i)$ denotes the field-dependent mass, Q is the renormalization scale, the supertrace is defined as the sum over all particles weighted with a factor $(-1)^{2s}(2s+1)$ and C_s is a spin-dependent constant, which gets values $C_0 = 3/2$, $C_{1/2} = 3/2$, $C_1 = 5/6$ in the $\overline{\mathrm{DR}}$ scheme [391].

The dominant correction to the Higgs boson masses comes from the incomplete cancellation of top and stop loops of Fig. 3.1. The one-loop contribution to the Higgs sector masses was first computed by Li and Sher [392], but they assumed $m_t < 60$ GeV and got an upper limit $m_h < 95$ GeV. After the experimental bound on the top quark mass had increased, these results were re-evaluated in the early 1990's and a large positive correction to the lightest Higgs mass was found by Ellis, Ridolfi and Zwirner [393, 394] and others [395, 396].

The tree-level mass matrix for stops is

$$M_{\tilde{t}}^2 = \begin{pmatrix} m_{\tilde{t}_L}^2 + m_t^2 + m_Z^2\cos 2\beta\left(\frac{1}{2} - \frac{2}{3}\sin^2\theta_W\right) & m_t(-A_t + \mu\cot\beta) \\ m_t(-A_t + \mu\cot\beta) & m_{\tilde{t}_R}^2 + m_t^2 + \frac{2}{3}m_Z^2\cos 2\beta\sin^2\theta_W \end{pmatrix}. \tag{3.35}$$

[7] The first solution is practically ruled out because if the 125 GeV state were the heavier one, m_A needs to be small. The charged Higgs would then be light. The constraints from [75, 76] exclude the part of parameter space that was still allowed in [387, 388] and was not in tension with other low energy constraints. The authors of [389] find that it is still possible to fit the 125 GeV state as the heavier one, but there is a clear tension with the $b \to s\gamma$ results due to the light charged Higgs. In the nonsupersymmetric 2HDM the 125 GeV state can still be the heavier one [218].

We shall denote $X_t = A_t - \mu\cot\beta$. The one-loop corrected CP-odd scalar mass is

$$m_A^2 = (m_A^2)_{\rm tree} - \frac{3g^2 m_t^2 A_t \mu}{16\pi^2 m_W^2 \sin 2\beta \sin^2\beta (m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2)}[f(m_{\tilde{t}_1}^2) - f(m_{\tilde{t}_2}^2)], \quad (3.36)$$

where $f(x) = 2x(\ln(x/Q^2) - 1)$, Q is the renormalization scale and $m_{\tilde{t}_{1,2}}$ are the eigenvalues of the stop mass matrix.

The one-loop contribution to (3.30) can be written in the form [394]

$$\begin{aligned}
\delta_{11} &= \frac{3g^2 m_t^4}{16\pi^2 m_W^2 \sin^2\beta}\left(\frac{\mu X_t}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2}\right)^2 g(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2), \\
\delta_{12} &= -\frac{3g^2 m_t^4 \mu X_t}{16\pi^2 m_W^2 (m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2)\sin^2\beta}\left[\ln\left(\frac{m_{\tilde{t}_1}^2}{m_{\tilde{t}_2}^2}\right) - \frac{A_t X_t}{m_{\tilde{t}_1}^2 - m_{\tilde{t}_2}^2} g(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2)\right], \\
\delta_{22} &= \frac{3g^2 m_t^4}{16\pi^2 m_W^2 \sin^2\beta}\left[\ln\left(\frac{m_{\tilde{t}_1}^2 m_{\tilde{t}_2}^2}{m_t^4}\right) - \frac{2A_t X_t}{m_{\tilde{t}_1^2} - m_{\tilde{t}_2}^2}\ln\left(\frac{m_{\tilde{t}_1}^2}{m_{\tilde{t}_2}^2}\right)\right. \\
&\quad \left. + \left(\frac{A_t X_t}{m_{\tilde{t}_1^2} - m_{\tilde{t}_2}^2}\right)^2 g(m_{\tilde{t}_1}^2, m_{\tilde{t}_2}^2)\right], \qquad (3.37)
\end{aligned}$$

where $g(x,y) = 2 - \frac{x+y}{x-y}\ln(x/y)$ and using the one-loop corrected value for m_A^2 in Eq. (3.31).

The one-loop correction is large if the stops are heavy, there is large mixing between them and $\tan\beta$ is large. At large values of $\tan\beta$ also the bottom-sbottom correction can be large. It will decrease the lightest Higgs mass when $\tan\beta \gtrsim m_t/m_b$ [397].

At two loops there is a significant negative contribution from the QCD corrections to the one-loop result [398, 399]. The contribution $\mathcal{O}(\alpha_t^2)$ is positive canceling a part of the QCD correction [400]. These have also been computed in the effective potential approach [401, 402]. The two-loop QCD corrections to bottom-sbottom loops were computed in [397]. Typically the total shift is 3–8 GeV downwards. The dependence on the renormalization scale is a lot milder than in the one-loop approximation. Some benchmarks for the Higgs mass at NNLO computed with HDECAY [188] (which uses FeynHiggs [403]) are plotted in Fig 3.2.

The leading three-loop correction $\mathcal{O}(\alpha_t^2\alpha_s)$ to the lighter CP-even Higgs have also been computed. The three-loop corrections are small, typically below 1 GeV [404, 405].

The one-loop correction to the charged Higgs mas has been computed in [406]. Also the leading two-loop corrections have been computed [407, 408]. The loop corrections to the rest of the Higgs sector are of similar order of magnitude than to the lightest Higgs mass but if the tree-level values are a lot larger, the relative error from loop corrections is rather small.

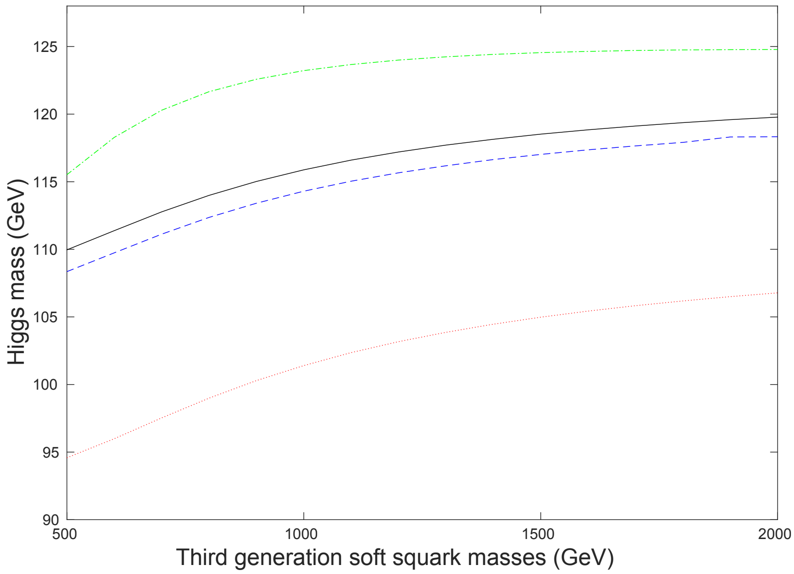

Figure 3.2: The Higgs mass as a function of third generation soft squark masses. We have chosen $m_A = \mu = 500$ GeV, $M_{\tilde{g}} = 2$ TeV and used three values for $\tan\beta$: 3 (red dots), 10 (blue dashes) and 30 (black solid). The third generation soft masses are equal $m_{\tilde{Q}} = m_{\tilde{t}_R} = m_{\tilde{b}_R}$ and trilinear couplings A_b and A_t are set equal to the soft masses. We also show a case where $A_b = A_t = 2m_{\tilde{Q}}$ and $\tan\beta = 30$ (green dashdots). Large mixing in the stop sector increases the Higgs mass.

The Higgs mass is probably the best indication of where superpartners should be. Within the MSSM, the stops should be somewhere between 1.5 TeV (large mixing) and 10 TeV (small mixing) and at least the left-handed sbottom should be at the same scale. Unfortunately only the lower end of this range can be probed at the LHC and it will require a large integrated luminosity at 14 TeV.

3.3 MSSM Higgs phenomenology at the LHC

The Higgs phenomenology was reviewed extensively in chapter 2 so here we shall concentrate only on the features that arise from supersymmetry.

3.3.1 SUSY effects on Higgs production

Supersymmetry modifies the Higgs phenomenology even from the Type II 2HDM, since at least at the loop-level superpartners will contribute.

In gluon fusion one must take into account the contribution from stop and sbottom loops. The left-handed stop coupling to the lighter CP-even Higgs is

$$\lambda_{h\tilde{t}_L\tilde{t}_L^*} = \frac{igm_z}{\cos\theta_W}\left(\frac{1}{2} - \frac{2}{3}\sin^2\theta_W\right)\sin(\alpha+\beta) - \frac{igm_t^2\cos\alpha}{m_W\sin\beta} \tag{3.38}$$

and other couplings between the Higgs and squarks have similar expressions (the complete set of couplings can be found from [409]).

When the LO contribution of squark loops is computed, the part from the couplings proportional to gm_Z cancels, when all squarks are summed over. Only the contribution proportional to gm_q^2/m_W survives [410]. This cancellation makes the LO contribution small if the squarks are heavier than 500 GeV. In the MSSM this holds, since the lightest Higgs mass will not reach 125 GeV without TeV-scale stops. The corrections have also been computed at NLO [411–414] and approximately at NNLO [415–417]. The uncertainty of the SM computation is larger than the contribution from squark loops.

For the heavy CP-even and the CP-odd Higgs the QCD corrections give a similar enhancement to the production cross section as the SM Higgs gets [418]. The contribution from scalar loops can enhance the cross section significantly if the Higgs mass is close to $2m_{\tilde{q}}$ [413], otherwise the contribution is rather small.

The SUSY-QCD corrections to VBF and Higgs-strahlung were computed in [419] and they are at most at a percent level, thus smaller than the uncertainty in the SM result.

In the $t\bar{t}h$ production there will be a modification of the Higgs-top coupling due to SUSY-QCD effects. However, the tree-level Yukawa coupling is large and th contribution is loop-suppressed. For $b\bar{b}h$ production the SUSY-QCD effects are larger relative to the tree-level coupling and they are enhanced at large $\tan\beta$. The enhancement could be large enough to distinguish the signal from the SM [420], but the $b\bar{b}h$ production could be enhanced also in the non-supersymmetric Type-II 2HDM.

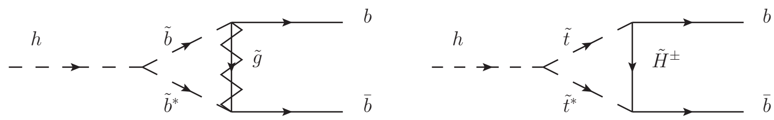

Figure 3.3: The one-loop superpartner corrections to the Higgs-bottom coupling. The contribution from both of these diagrams is enhanced at large $\tan\beta$ and can lead to a substantial deviation from the SM prediction even in the alignment limit.

By analyzing the SM-like Higgs production rates it will be very hard to distinguish the MSSM from a non-supersymmetric Type-II 2HDM. The $b\bar{b}h$ production has the largest potential to provide a deviation from the SM predictions.

3.3.2 SUSY effects on Higgs decays

Supersymmetry will also affect Higgs decays. The most dramatic effect would be if the lightest superpartner was lighter than $m_h/2$ so that it would contribute to the invisible decay width. The LEP and LHC bounds on MSSM superpartner masses[8] allow only the lightest neutralino to be light enough for the decay to be kinematically allowed [90].

The Higgs decay to neutralinos was studied in [422]. With a positive μ the invisible branching ratio can go up to 50% and with a negative μ it may be around 10%. The latest results on the invisible decay width exclude the part with the a branching ratio to neutralinos larger than 24% [202]. With more data this scenario will become more constrained or give an indication of a dark matter candidate.

If the decays to superpartners are not kinematically allowed, the largest effect is the correction from superpartners to the $h \to b\bar{b}$ decay. The corrections arise from the diagrams of Fig. 3.3. These lead to a modified relation between the bottom Yukawa coupling and the bottom mass [423, 424]. We may write $m_b = y_b v \cos\beta(1+\Delta_b)/\sqrt{2}$, where

$$\Delta_b = \left[\frac{2\alpha_s \mu M_3}{3\pi} I(m^2_{\tilde{b}_1}, m^2_{\tilde{b}_2}, M_3^2) + \frac{y_t^2 \mu A_t}{16\pi^2} I(m^2_{\tilde{t}_1}, m^2_{\tilde{t}_2}, \mu^2)\right] \tan\beta, \tag{3.39}$$

and $I(a,b,c) = -[ab\ln(a/b) + bc\ln(b/c) + ca\ln(c/a)]/[(a-b)(b-c)(c-a)]$. This leads to the loop-corrected Higgs-bottom coupling [425]

$$\lambda^{(1)}_{hb\bar{b}} = -\frac{\sqrt{2}m_b}{v}\frac{\sin\alpha}{\cos\beta}\left(1 - \frac{\Delta_b}{1+\Delta_b}\cot\alpha\cot\beta\right). \tag{3.40}$$

[8]The bound on the lightest neutralino $m_{\chi_1} > 46$ GeV assumes universal gaugino masses and can be relaxed if gaugino masses are non-universal and the lightest neutralino is bino-like [421].

The correction can be large at large $\tan\beta$ and it can either enhance or suppress the $h \to b\bar{b}$ decay width [426]. Since $h \to b\bar{b}$ is the primary decay channel, this can affect largely all the signal strengths. The two-loop contribution [427] mainly makes the dependence on the renormalization scale milder. If all of the superpartners are heavy, this contribution will become small.

Also the heavier CP-even and the CP-odd Higgs will get a large correction to their coupling to bottom quarks [426] as will the charged Higgs to the coupling to $t\bar{b}$ [428–430].

The correction of the top Yukawa coupling comes from similar diagrams but the chargino diagram gives a correction proportional to y_b^2 and is hence small compared to the tree-level coupling. In addition there is no enhancement at large $\tan\beta$. The tau Yukawa coupling gets electroweak loop corrections, which are enhanced at large $\tan\beta$, but due to the relatively small gauge couplings these corrections are smaller than in the bottom sector [424]. The different behaviour of $h \to b\bar{b}$ and $h \to \tau^+\tau^-$ can distinguish the MSSM from the Type-II 2HDM, but the lepton-specific and flipped versions of the 2HDM may mimic this behaviour.

The charged Higgs and charged superpartners will enter the loops in the decays $h \to \gamma\gamma$ and $h \to Z\gamma$. Since most of the new particles are scalars and their contribution is small, the effect of the MSSM to the loop-induced decays is somewhat limited. There can be sizable effects if there are light staus and $\tan\beta$ is large [431–433] or there are light charginos [434].

The Higgs decays have a good potential to reveal a deviation from the SM predictions. The dominant decay channel $h \to b\bar{b}$ gets a potentially large correction from superpartners even in the alignment limit. Only in the case where all superpartners are heavy, the corrections are small.

3.4 Extensions of the MSSM

3.4.1 Singlet extensions

The most studied extension of the MSSM is that with an additional singlet field. Whereas one may have other choices for the superpotential [341, 356, 435], the most common one is the purely trilinear version [436]

$$W = Y^{(U)}(Q \cdot H_u)U^C + Y^{(D)}(Q \cdot H_d)D^c + Y^{(L)}(L \cdot H_d)E^c + \lambda(H_u \cdot H_d)S - \frac{\kappa}{3}S^3, \tag{3.41}$$

where S is the singlet superfield and other superfields are those of the MSSM. In addition the following soft terms are added

$$\mathcal{L}_{\text{soft,NMSSM}} = \mathcal{L}_{\text{soft,MSSM}} - m_S^2|S|^2 + [A_\lambda(H_u \cdot H_d)S - A_\kappa S^3/3 + \xi S + \text{h.c.}]. \tag{3.42}$$

This model is called the next-to-minimal supersymmetric standard model (NMSSM). In the NMSSM the VEV of the scalar component of the singlet field generates the μ term of the MSSM. The S^3 term is needed to break a global $U(1)$ symmetry [437] that would otherwise lead to an axion when the scalar components of the superfields acquire VEVs.

The superpotential has a Z_3 symmetry, which will be spontaneously broken by the VEVs of the neutral scalar fields. This could lead to the formation of domain walls in the early universe. This symmetry can be broken by nonrenormalizable terms so that a preferred vacuum exists [438]. This solution is not viable in a straightforward way as pointed out in [439], since there will be a large tadpole term for the singlet. The problem can be solved by introducing an additional symmetry, which is broken only by the soft supersymmetry breaking terms. This will lead to a tadpole term $\xi \sim m_{\mathrm{SUSY}}^3$ [440]. The domain wall problem could also be solved by inflation [441].

The physical Higgs sector consists of three CP-even (h, H_1, H_2) and two CP-odd (A_1, A_2) neutral scalars and a charged pair ($H^{\pm}$). In the NMSSM there are two additional mechanisms that can lift the Higgs mass compared to the MSSM. First, the $\lambda(H_u \cdot H_d)S$-term of the superpotential generates additional terms to the 2×2 mass matrix formed by H_u and H_d. The submatrix becomes

$$M_{H,2\times 2}^2 = \begin{pmatrix} m_Z^2 \sin^2\beta - C_1 \cot\beta & -\frac{1}{2} m_Z^2 \sin 2\beta \left(1 - \frac{4\lambda^2}{g^2+g'^2}\right) + C_1 \\ -\frac{1}{2} m_Z^2 \sin 2\beta \left(1 - \frac{4\lambda^2}{g^2+g'^2}\right) + C_1 & m_Z^2 \cos^2\beta - C_1 \tan\beta \end{pmatrix}, \tag{3.43}$$

where $C_1 = A_\lambda v_s/\sqrt{2} - k\lambda v_s^2/2$. From this we get the smaller eigenvalue

$$m_h^2 = \frac{m_Z^2 + C_2 - \sqrt{(m_Z^2 + C_2)^2 - 4m_Z^2 C_2 \cos^2 2\beta - \frac{m_Z^2 y}{2}(m_Z^2(1-y) + C_2)\sin^2 2\beta}}{2}, \tag{3.44}$$

where $C_2 = -2C_1/\sin 2\beta$, $y = 4\lambda^2/\bar{g}^2$ and $\bar{g}^2 = g^2 + g'^2$. The expression under the square root can be put to the form $[C_2 + m_Z^2(1-2x)]^2 + m_Z^4[4(x - x^2) + y(y-1)\sin^2 2\beta]$, where $x = \cos^2 2\beta + \frac{y}{2}\sin^2 2\beta$. The term proportional to m_Z^4 is non-negative and hence we get the bound [442]

$$m_h^2 \leq m_Z^2 \left(\cos^2 2\beta + \frac{2\lambda^2}{\bar{g}^2}\sin^2 2\beta\right). \tag{3.45}$$

At low $\tan\beta$ and large values of λ the lightest Higgs can be heavier than m_Z at tree-level. This mechanism is known as λSUSY. The remarkable thing is that the bound does not depend on v_s and the Higgs mass can be lifted no matter how heavy the singlet-dominated state is.

To saturate the limit one must be able to decouple the singlet from the doublets. In the NMSSM it is possible to do so by choosing the soft terms so that they make the mixing terms small and/or making κ and v_s so large that the singlet becomes heavy. The limit holds also in the case of several singlets, since one may always choose a linear combination of them such that only one of the singlets has a VEV.

The second option is that the lightest CP-even scalar is a singlet-dominated state that mixes with the lighter doublet state. This mixing will lift the mass of the SM-like Higgs. The production cross section for such a singlet-dominated state is small and hence it can remain unobserved. Such a scenario was considered in [443]. They find that the mixing can lift the Higgs mass by at most

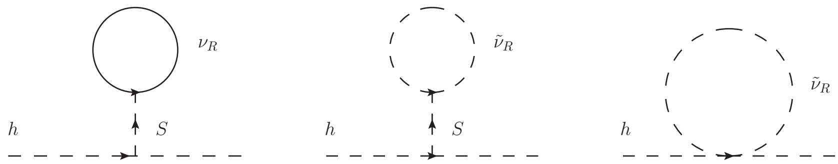

Figure 3.4: The one-loop correction from neutrinos and sneutrinos to the SM-like Higgs mass. The correction is proportional to $\lambda_H \lambda_N$ and can be a few GeV's if the singlet state is not too heavy.

7 GeV. On the other hand the mixing with the singlet suppresses all production modes and hence the overall signal strength should be smaller than in the SM.

The NMSSM allows the third generation squarks to be lighter than in the MSSM. Hence they are easier to discover directly [444–447] and their impact on the Higgs production and decay modes at loop-level can be larger than in the MSSM.

Especially in the case of λSUSY some of the trilinear scalar couplings in the Higgs sector are large and hence $H_{1,2}$ can have a substantial branching fraction to pairs of Higgs bosons [448,449]. In the case where the heavier doublet Higgs is the heaviest one, we may have processes like $pp \to H_2 \to hH_1$ with a production cross section and branching ratio which can lead to observable signatures. The singlet-like H_1 can decay to a pair of singlet-like CP-odd scalars, which again can decay *e.g.* to b-quarks or photons. Hence signatures with six photons, a b-quark pair and four photons etc. are possible. Such final states are very rare in the (MS)SM.

Another chiral superfield that is a singlet under the SM gauge group is the right-handed neutrino. The difference is that the neutrino superfield carries lepton number. The MSSM does not have a mechanism for neutrino mass generation and hence the introduction of right-handed neutrinos is well motivated. One can add the lepton number conserving term $y^{(\nu)}(L \cdot H_u)N^c$ to the superpotential. This generates Dirac masses for the neutrinos. One can also introduce a gauge invariant, but lepton number violating term MN^cN^c, which gives a Majorana mass to the right-handed neutrinos and a small mass to left-handed neutrinos via the seesaw mechanism. In models with R-parity the cubic N^3 term is excluded.

One may also generate the Majorana mass dynamically via a SN^cN^c-term. The superpotential for such a model is [450]

$$W = Y^{(U)}(Q \cdot H_u)U^C + Y^{(D)}(Q \cdot H_d)D^c + Y^{(L)}(L \cdot H_d)E^c + Y^{(\nu)}(L \cdot H_u)N^c + \lambda_H (H_u \cdot H_d)S + \frac{\lambda_N}{2} SN^cN^c + \frac{\lambda_S}{3} S^3. \quad (3.46)$$

The λ_N coupling between the singlets has an impact on the SM-like Higgs. Besides generating the Majorana mass for the right-handed neutrino needed in the seesaw mechanism, it enables a positive correction from the neutrino sector to the SM-like Higgs mass via the loop diagrams of Fig. 3.4 [451]. This correction can be a few GeV's.

3.4.2 Triplet extensions

The MSSM can also be extended by triplet fields. The motivations of non-supersymmetric triplet models apply also to the supersymmetric versions.

In the context of supersymmetry there is one additional reason. The tree-level upper bound for the Higgs mass comes from the 2×2 submatrix of the two doublets. The only renormalizable gauge invariant superpotential couplings which can give a contribution to this submatrix come from singlets or triplets. Hence these are the only alternatives to make the Higgs mass more reachable via F-terms at tree-level.

Most motivated triplet extensions of the MSSM have triplets with $Y = 0$ or $Y = \pm 2$. The simplest triplet extension has one triplet with $Y = 0$ [452]. The Higgs sector consists of

$$H_1 = \begin{pmatrix} H_1^0 \\ H_1^- \end{pmatrix}, \quad H_2 = \begin{pmatrix} H_2^+ \\ H_2^0 \end{pmatrix}, \quad \Sigma = \begin{pmatrix} \xi^0/\sqrt{2} & -\xi_2^+ \\ \xi_1^- & -\xi^{0*}/\sqrt{2} \end{pmatrix} \tag{3.47}$$

and the superpotential is

$$W = W_{\text{MSSM}} + \lambda (H_1^T \cdot \Sigma) H_2 + \mu_\Sigma \text{Tr}(\Sigma^2). \tag{3.48}$$

The physical Higgs sector has three CP-even and two CP-odd neutral scalars and three pairs of singly charged scalars. The VEV of the neutral component of the triplet must be small so that the ρ parameter does not deviate largely from one.

The corrections to the 2×2 submatrix change the bound on the lightest CP-even scalar to (assuming $v_\Sigma \ll v$) [453]

$$m_h^2 \leq m_Z^2 \left(\cos^2 2\beta + \frac{\lambda^2}{\bar{g}^2} \sin^2 2\beta \right). \tag{3.49}$$

Also the one-loop corrections from the triplet sector can increase the SM-like Higgs mass [454] so that a 125 GeV Higgs is possible without corrections from quarks and squarks. Hence the third generation squark masses are not bounded by Higgs physics.

There are several additional charged scalars and charginos in the model. These can change the Higgs decay width to $\gamma\gamma$ and $Z\gamma$ if they are light [455,456]. The deviation can be tens of percents at small values of $\tan\beta$.

The three-point couplings between the scalars are not as large as they can be in the NMSSM, since the triplet VEV is small. Hence the fermionic decay channels dominate for the heavier Higgses. On the other hand for the charged Higgs there is a decay channel $H^+ \to W^+ Z$, which is forbidden in the MSSM due to custodial symmetry [248]. The branching ratio $B(H^+ \to W^+ Z)$ can be up to 50% [249].

In the case of $Y = \pm 2$ triplets, anomaly cancellation requires at least two triplets with opposite hypercharges. We shall denote them

$$\Delta = \begin{pmatrix} \delta^+/\sqrt{2} & \delta^{++} \\ \delta^0 & -\delta^+/\sqrt{2} \end{pmatrix}, \quad \overline{\Delta} = \begin{pmatrix} \bar{\delta}^-/\sqrt{2} & \bar{\delta}^0 \\ \bar{\delta}^{--} & -\bar{\delta}^-/\sqrt{2} \end{pmatrix}. \tag{3.50}$$

There are several options when choosing the superpotential. The most general one is

$$W = W_{\mathrm{MSSM}} + M_\Delta \mathrm{Tr}(\Delta\overline{\Delta}) + Y^{(\nu)}((L^T\cdot\Delta)L + \mathrm{h.c.}) + \lambda_1(H_1^T\cdot\Delta)H_1 + \lambda_2(H_2^T\cdot\overline{\Delta})H_2. \tag{3.51}$$

The tree-level bound on the lightest Higgs mass becomes (assuming $v_\delta, v_{\bar{\delta}} \ll v$) [453]

$$m_h^2 \leq m_Z^2 \left(\cos^2 2\beta + \frac{2\lambda_1^2}{\bar{g}^2}\cos^4\beta + \frac{2\lambda_2^2}{\bar{g}^2}\sin^4\beta\right). \tag{3.52}$$

Hence it is possible to lift the Higgs mass with the F-term contribution from $Y = \pm 2$ triplets at large $\tan\beta$ (or $\cot\beta$), whereas other extensions increase the Higgs mass only when $\tan\beta$ is close to one.

One could expect that the doubly charged Higgses would affect the $h \to \gamma\gamma$ channel. In the supersymmetric case the Higgs self-interactions come from D-terms so the couplings are smaller than in the non-supersymmetric case. In addition, the coupling between the SM-like Higgs and doubly charged Higgses vanishes if the triplets are degenerate [457]. Hence the diphoton rate cannot vary as much as in the non-supersymmetric case [247].

3.5 Searching SUSY at the LHC

We review briefly some of the superpartner searches in R-parity conserving supersymmetry as a benchmark to supersymmetry without R-parity, which will be considered in the next chapter. We shall limit ourselves to the particle content of the MSSM.

There are three approaches from which one may discover new physics or draw limits on it. The simplest approach is effective field theory, where one assumes an effective coupling between particles and then constrains it experimentally. The second approach is that of simplified models, where only a few particles are included. Some of the masses and branching ratios are fixed and some are considered as free parameters. The most ambitious approach is to consider full models, but the number of free parameters is so large that only benchmark points can typically be considered, unless universal parameters are assumed in the soft SUSY breaking sector.

R-parity conservation makes the lightest superpartner stable. Hence all supersymmetric particles decay as cascades ending with (at least) one LSP which is seen as missing momentum. In the SM all events with missing momentum involve neutrinos, *i.e.* either $W \to \ell\nu$ or $Z \to 2\nu$. The former ones can be identified due to the hard lepton and the background from the latter ones can be estimated from similar processes, where $Z \to \ell^+\ell^-$.

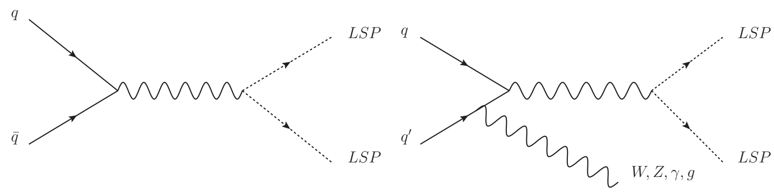

Figure 3.5: The pair production of LSP's can look like nothing happened if it proceeds like the diagram on the left. However, initial state radiation (right) can reveal the LSP production. The mediator can be *e.g.* a Z-boson or a Higgs.

3.5.1 How to find the invisible?

The LSP candidate needs to be neutral and uncolored, so in the MSSM the options are the lightest neutralino, the lightest sneutrino and, if we include supergravity, the gravitino. Sneutrino dark matter is excluded because a light sneutrino would have been seen at LEP and a heavier sneutrino would have been seen in direct detection experiments [458, 459].

The LSP can be searched in cascade decays if the heavier superpartners can be produced and there are large mass differences so that there will be a clear missing transverse momentum signal. Most of the searches find no evidence for such a signal beyond the SM expectations [460–464]. There are a few larger deviations from the SM background [465, 466], but they are not seen in both experiments.

If the supersymmetric spectrum is too heavy so that other superpartners cannot be produced, the LSP pair production will look like nothing happened, *i.e.* $pp \to$ invisible. However, the LSP can be revealed by initial state radiation. The LSP's will be produced either from the collision of a quark and an antiquark or the collision of two gluons. Either of these can radiate a gluon and the quarks may radiate a photon, W or Z as shown in Fig. 3.5. These will then be seen as monojets, monophotons etc. with missing transverse momentum.

Such searches have also been performed. The data are well described by the SM [467–470]. This leads to an upper limit of effective interactions between the LSP and hadronic matter. This can be compared with the other searches. Typically collider searches are more sensitive at the lower end of the mass spectrum, where the recoil is smaller in direct detection experiments.

3.5.2 Neutralinos and charginos

The lightest neutralino is usually assumed to be the LSP. The decay chains of heavier superpartners usually contain neutralinos and charginos so we shall consider their production and decay modes briefly.

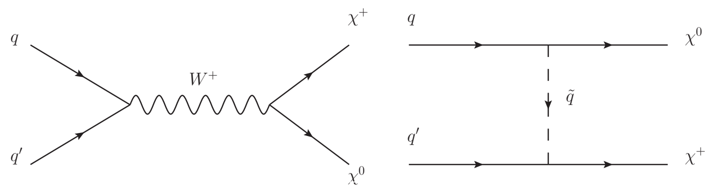

Figure 3.6: Examples of chargino and neutralino production. The t-channel diagrams are subleading if superpartners are a lot heavier than the SM particles. Similar diagrams can be drawn for chargino (mediated by Z, γ, $\tilde{q}$) and neutralino (mediated by Z, $\tilde{q}$) pair production.

The neutralino and chargino mass matrices in the MSSM are

$$M_{\chi^0} = \begin{pmatrix} M_1 & 0 & g'v_d/2 & -g'v_u/2 \\ 0 & M_2 & -gv_d/2 & gv_u/2 \\ g'v_d/2 & -gv_d/2 & 0 & -\mu \\ -g'v_u/2 & gv_u/2 & -\mu & 0 \end{pmatrix} \tag{3.53}$$

using the basis $\{\tilde{B}, \tilde{W}, \tilde{H}_d, \tilde{H}_u\}$ and

$$M_{\chi^\pm} = \begin{pmatrix} M_2 & -gv_u/\sqrt{2} \\ -gv_u/\sqrt{2} & \mu \end{pmatrix} \tag{3.54}$$

using the basis $\tilde{W}^\pm, \tilde{H}^\pm$. The masses are basically free parameters, although $|\mu|$ can be argued not to be too large or there will be fine tuning as discussed on p. 40. The nature of the LSP depends on which of the quantities $|M_1|$, $|M_2|$, $|\mu|$ is the smallest. If $|\mu| \ll |M_1|$, $|M_2|$ the LSP will be higgsino-like, in the opposite case gaugino-like.

Neutralinos and charginos can be produced in pairs via the s-channel or t-channel. A few examples are given in Fig. 3.6. The mediating particle in the t-channel is a superpartner, so with the current bounds on superpartner masses, one may assume that the dominant contribution comes from s-channel production. The production of bino-dominated neutralinos is suppressed since its couplings to SM gauge bosons are small.

If we assume that sfermions are heavier than (at least some of the) neutralinos and charginos, the gaugino-like states decay to gauge bosons and the lighter neutralinos, *i.e.* the decay channels are $\chi_1^\pm \to W^\pm \chi_1^0$, $\chi_2^0 \to Z\chi_1^0$, $\chi_2^\pm \to W^\pm \chi_i^0, Z\chi_1^\pm$ etc. The higgsino-like states can also decay to Higgs bosons and lighter superpartners.

The easiest way to reduce the background is to look at the cases, when at least one of the bosons decays leptonically. The signature is then multileptons, missing transverse momentum and possibly jets. Such searches have been performed [460,471,472] and the results are compatible with SM expectations. The sensitivity of these searches is limited if $m_{\chi_2^0} - m_{\chi_1^0}$ and $m_{\chi_1^\pm} - m_{\chi_1^0}$ are small

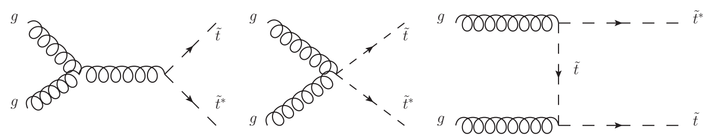

Figure 3.7: Some of the leading order diagrams of stop pair production. The first and third diagrams can also have quarks in the initial state in which case the mediator in the t-channel is a gluino.

so that there will not be much missing transverse momentum from the SUSY decays.

3.5.3 Third generation squarks

The LHC produces mostly strongly interacting particles, so squarks and gluinos, if they are light enough, should be produced copiously. Third generation squarks receive more attention basically because of three reasons. First, we have a rough idea of their mass scale from the SM-like Higgs mass. Second, the lighter stop is usually the lightest squark, because the large Yukawa gives a negative contribution to the SUSY breaking soft mass in renormalization group running and the large mixing in the stop sector pushes the lighter stop mass downwards. Third, the decay products usually include top and bottom quarks, which helps to reduce the background.

The stop mass matrix was given in Eq. (3.35). The mixing in the sbottom sector is not as large, because the mixing term is $m_b(-A_b + \mu \tan\beta)$ and hence the mass eigenstates of sbottoms nearly coincide with the left- and right-handed states.

Squarks can be produced via gluon fusion or quark-antiquark annihilation. The t-channel squark and gluino exchange diagrams will be subleading, since large superpartner masses will suppress the corresponding amplitudes. Some of the production channels are shown in Fig. 3.7.

The main decay channels are $\tilde{t} \to t\tilde{\chi}^0, b\tilde{\chi}^+$ and $\tilde{b} \to b\tilde{\chi}^0, t\tilde{\chi}^-$, where the charginos and neutralinos are mostly higgsino-like due to the large top Yukawa coupling if the decay is kinematically allowed. The neutralinos and charginos will then decay as described in the previous section. To reduce the background, one usually requires that there is at least one hard lepton in the final state.

The final states will then consist of b-jets, one or several leptons, possibly jets and missing transverse momentum. The data are well described by the SM background [473–475] and stop masses up to 800 GeV are ruled out for neutralino masses below 200 GeV except when the mass difference is small. The case of compressed spectrum has been analyzed by looking at spin correlations of $t\bar{t}$ production [476] or using a boost from initial state radiation [477].

Chapter 4

Supersymmetry without R-parity

4.1 R-symmetries, R-charges and R-parity

Perhaps the most natural framework for R-symmetries are extended supersymmetric models. There one can have an internal symmetry among the SUSY generators, under which the SUSY generators are charged. One may show that such a symmetry must be represented by unitary matrices with a dimension not higher than the number of SUSY generators. If we look at one supersymmetry generator, there will be one linear combination of the generators of the symmetry group, which acts nontrivially on that generator. Hence for $N = 1$ supersymmetry the internal symmetry reduces to a U(1)-symmetry known as R-symmetry [351, 478]. The R-charges can be normalized so that the generators Q_α, $\bar{Q}^{\dot{\alpha}}$ have charge ± 1. Also the superspace coordinates θ, $\bar{\theta}$ are charged under the R-symmetry.

The fact that SUSY generators are charged under $\mathrm{U}(1)_\mathrm{R}$ means that different components of a superfield have different R-charges. For the Lagrangian to be R-symmetric, the F-term of the superpotential needs to be R-symmetric and the superpotential itself should carry a R-charge $R(W) = +2$.

R-symmetry is a global U(1) symmetry. All such symmetries are expected to be broken by gravity. The most common assumption is that $\mathrm{U}(1)_\mathrm{R}$ will be broken to a discrete Z_2 subgroup, R-parity [352]. If the quantum numbers are chosen so that for SM particles $R = +1$ and for superpartners $R = -1$ one may discard all renormalizable baryon or lepton number violating operators from the MSSM superpotential. However this does not exclude baryon and lepton number violating operators, which are of higher dimension, like $QQQL$.

Next we shall study supersymmetry without R-parity conservation. The R-symmetries as such are a well motivated alternative. As mentioned, gravity will break the R-symmetry, but if this effect is small, one may still use the R-symmetry as an approximate symmetry. In R-symmetric models one assumes that the effect of gravity on the matter sector is small, so supersymmetry breaking must be mediated by either gauge interactions or anomalies.

The other class of models will be those where R-parity is broken spontaneously. The R-parity violating (RPV) superpotential (3.25) has a large number of free parameters. In spontaneous RPV the operators are generated by the VEVs of sneutrino fields. This makes the model more predictive. Since sneutrino VEVs do not break baryon number conservation, the proton will be stable.

R-parity violation implies that the lightest supersymmetric particle (LSP) will not be stable. Hence the problem of dark matter remains. The only viable dark matter candidate could be the gravitino, since it couples weakly to all particles and hence its lifetime could be greater than the age of the Universe. We shall simply assume that there will be a viable dark matter candidate, possibly from a completely unknown sector.

4.2 Construction of R-symmetric models

4.2.1 Neutralino and chargino masses

R-symmetry forbids Majorana masses for gauginos, the μ-term and trilinear scalar interactions (A-terms). In such a case the electroweakinos would get small masses from EWSB, but the gluinos would be massless. Since massless or light gauginos do not exist, one needs to add new chiral superfields in the adjoint representation of $\mathrm{SU}(3)_C$, $\mathrm{SU}(2)_L$ and $\mathrm{U}(1)_Y$ to allow Dirac mass terms for the gauginos [479]. A strong motivation for R-symmetric models is that they are a natural framework for Dirac gauginos. Models with Dirac gauginos have become more popular since the production cross section of gauginos is suppressed compared to Majorana gauginos [480, 481] and hence the bounds from the LHC are not as severe.

Dirac gauginos need new superfields $\hat{S}$, $\hat{T}$ and $\hat{O}$, which are in the adjoint representation of U(1), SU(2) and SU(3), respectively. The field strength superfield can be combined with the chiral superfields if we use a spurion superfield $M^\alpha = m\theta^\alpha$. We may now write the Lagrangian

$$\mathcal{L} = -\int \mathrm{d}^2\theta\; M^\alpha(c_1 B_\alpha S + c_2 \mathrm{Tr}(W_\alpha T) + c_3 \mathrm{Tr}(G_\alpha O)) + \text{h.c.}, \tag{4.1}$$

where c_i are dimensionless[1] constants.

If there is no μ-term, there will be a light chargino that would already have been detected. This problem can be solved by introducing new superfields R_u and R_d, which carry the quantum numbers of $H_{u,d}$ but have a R-charge $+2$. Then one may add the terms $\mu_u H_u R_d + \mu_d R_u H_d$ to the superpotential, which give masses for the charginos. Both R_u and R_d need to be inert so that they would not break the R-symmetry spontaneously and lead to an R-axion.

[1] Usually the spurion field is a D-component of a field strength superfield and hence of mass dimension two. In such a case the coeffiecients c_i should have dimension -1. We have simply absorbed this to the constant m, which has mass dimension one.

Model	Q	U^c	D^c	L	E^c	H_u	H_d	R_u	R_d	S	T	O
$U(1)_{\mathrm{R_0}}$	1	1	1	1	1	0	0	2	2	0	0	0
$U(1)_{\mathrm{R=B}}$	4/3	2/3	2/3	1	1	0	0	2	2	0	0	0
$U(1)_{\mathrm{R=L}}$	1	1	1	2	0	0	0	2	2	0	0	0
$U(1)_{\mathrm{R=-L}}$	1	1	1	0	2	0	0	2	2	0	0	0

Table 4.1: Alternatives for the R-charges of the minimal R-symmetric model. The first line is the model of [483]. The other models have the R-charges of SM fermions identified with the baryon and ±lepton numbers. Such choices give unconventional baryon and lepton numbers for the squarks and sleptons.

4.2.2 R-charge assignments

The R-charges for the MSSM fields can be given in various ways. A phenomenological constraint is that the scalar fields which get VEVs must have $R(\varphi) = 0$ to avoid physical Goldstone bosons. Various options of choosing the R-charges have been considered in [482]. The generic option is to set the R-charges of the Higgs fields to zero and then set the R-charges of all sfermions to $+1$ so that the MSSM Yukawa couplings are allowed [483]. Such a choice eliminates completely the $d = 5$ terms leading to proton decay, but allows the Weinberg operator $(LH_u)^2$. Other options are to set the R-charges equal to the baryon number or the lepton number. More options can be obtained by taking linear combinations of these. The R-charges of these options are given in table 4.1.

R-symmetries allow large flavor violation in the sfermion sector without violating current bounds from flavor physics [483–486]. This is due to negligible left-right mixing and the absence of soft trilinear scalar couplings, which give the main constraints for supersymmetry in the flavor sector [487].

If the R-charges are identified with baryon or lepton numbers, this leads to sleptons carrying baryon number or squarks carrying lepton number. Hence the sfermions can decay to channels, which conserve B or L but would be considered R-parity violating in the context of MSSM. The important feature of R-symmetries is that since baryon or lepton number is conserved, the constraints from baryon or lepton number violating processes do not produce constraints on the corresponding couplings. This is a clear difference compared to the MSSM extended with RPV. We studied squark decays in the case of the $R = -L$ model in **III** and showed that in some parts of the parameter space the unconventional decay modes may dominate and give a clear signal.

Some features of the $R = B$ model have been discussed in [482, 488]. We shall discuss the $R = -L$ model in more detail in section 4.3.

4.2.3 The Higgs mass in R-symmetric models

There are some subtleties concerning the SM-like Higgs mass in models with Dirac gauginos. If the Dirac masses come from supersoft D-term SUSY breaking like in Eq. (4.1), the gauge sector is basically a vector supermultiplet of $N = 2$ supersymmetry and in this limit the quartic scalar couplings vanish [489] pushing the tree-level Higgs mass down.

Solutions to this problem have been proposed. First, gravity breaks the R-symmetry and will reintroduce the quartic scalar couplings [489, 490]. There will always be a suppression compared to the MSSM since the effect is seen as the mixing elements between the SM-like Higgs and the adjoint neutral scalars S and T^0 [490, 491].

Second, the chiral superfields S and T as singlets and triplets may increase the tree-level value at low $\tan\beta$ as discussed in section 3.4 if the superpotential couplings are large. In such a case they may also provide positive loop corrections to the SM-like Higgs mass [492, 493]. Altogether one may reach the 125 GeV Higgs mass also in models with Dirac gauginos and R-symmetry.

Another approach to this is to assume that the gaugino masses come from R-preserving F-term supersymmetry breaking [494]. This choice does not make the quartic scalar couplings to vanish. In order for the Dirac gaugino masses to be large, the scale of supersymmetry breaking should be rather low, since the gaugino masses are of the order $\langle F\rangle^2/M^3$, whereas soft scalar masses are suppressed by only two powers of M.

4.3 Phenomenology of the $R = -L$ model

4.3.1 Superpotential and sneutrino as a Higgs

We shall have a closer look on the $\mathrm{U}(1)_R$ lepton number model considered in **III**. In addition to the fields given in table 4.1, we added a right-handed neutrino N^c with an R-charge +2.

The idea of identifying the R-charge with lepton number is an old one. In the 1970's it was thought that the neutrino could be the superpartner of the photon [341]. This requires that the SUSY generators do not commute with lepton number and hence carry lepton number. That idea didn't work out since neutrinos are doublets under SU(2) and photons are not. However, identifying the R-charges with lepton number and extending the MSSM minimally does produce a phenomenologically interesting model.

The most general superpotential for the model is

$$\begin{aligned} W = Y^{(U)}(Q\cdot H_u)U^c + Y^{(D)}(Q\cdot H_d)D^c + Y^{(L)}(L\cdot H_d)E^c + f(L\cdot H_u)N^c \\ + \mu_u(H_u\cdot R_d) + \mu_d(R_u\cdot H_d) + \lambda_N(H_u\cdot H_d)N^c + \lambda_S(H_u\cdot R_d)S + \lambda_S'(R_u\cdot H_d)S \\ + \lambda_T H_u T R_d + \lambda_T' R_u T H_d - M_R S N^c + \lambda(L\cdot L)E^c + \lambda'(Q\cdot L)D^c. \end{aligned} \tag{4.2}$$

The superpotential contains a subset of the couplings of Eq. (3.25), which are R-parity and lepton number violating in the MSSM. Here they are R-symmetric and lepton number conserving.

There are several F-term contributions that may increase the SM-like Higgs mass at tree-level. Each of the terms with coefficients $\lambda_{N,S,T}$, $\lambda'_{S,T}$ can lift the Higgs mass if the coupling is large and $\tan\beta$ is small. These couplings can also introduce new loop contributions to the SM-like Higgs mass.

The interesting feature of this model is that the sneutrino has the same quantum numbers as the down-type Higgs. It does not carry lepton number so its VEV will not create Majorana neutrino masses nor can it mediate lepton flavor violating processes. Hence the sneutrino can have a large VEV and mix with the Higgses. One may even make H_d heavy and its VEV so small that $\tilde{\nu}_L$ may take its place [495]. In such a case also the neutrino Yukawa coupling f can lift the Higgs mass via the λSUSY mechanism [492]. In the sneutrino as Higgs case the model can be simplified by assuming the fields R_u and H_d heavy and integrating them out. This still leaves the $\lambda_{S,T}$ couplings to lift the Higgs mass.

If the sneutrino takes the role of the Higgs, the λ and λ' couplings take the role of Yukawa couplings. However, due to the antisymmetry of the first two indices in λ_{ijk}, the charged lepton, whose sneutrino has a VEV, cannot get a mass via this mechanism at tree-level. It can get a mass at loop-level [482], so it is natural to assume that this is the electron. Hence the electron sneutrino is the one with the VEV.

4.3.2 Collider phenomenology

The phenomenology of this model has been studied recently. Electroweak precision measurements were analyzed in [495, 496] and the constraints allow quite a large variation for the unknown parameters. Somewhat heavy gauginos are favored. Neutrino physics has been considered in [492, 497, 499] and a sterile neutrino as a dark matter candidate in [498]. The experimental constraints from neutrino masses and mixings can be satisfied within this model.

The decay signatures and experimental bounds for neutralinos, charginos, squarks and sleptons were discussed in [482, 500]. In this model squarks are leptoquarks, *i.e.* they have a non-zero baryon and lepton number. The mixing even between stops is negligible, since both of the terms of the mixing element in (3.35) are small. The A-terms are forbidden by the R-symmetry and can come only from the R-breaking part. The part proportional to μ comes from the $\mu_u H_u R_d$-term of the superpotential. When differentiating with respect to H_u, there is a term $\mu_u R_d \tilde{t}_L \tilde{t}^*_R$, but since $\langle R_d \rangle = 0$, the contribution to the mixing element vanishes.

The studies have shown that the "R-parity violating" modes can have a substantial branching fraction although the typical modes with SM particles and neutralinos or charginos are often dominant. Sometimes the neutralinos and charginos decay with neutrinos in the final state so that the signature at the LHC is a similar cascade as it would be in the MSSM. Altogether the dominant decay channels depend on the values of the parameters and hence there is a large variety of possible final states.

The fact that the electron sneutrino has the VEV gives a unique pattern for the stop decays. Our benchmark points in **III** show that the decay $\tilde{t}_L \to b e^+$ can be dominant if the Higgsino is heavier than the stop and comparable to $t\chi^0$ and $b\chi^+$ in any case. For the right-handed stop only the conventional decay modes are possible. Currently CMS has not performed a search for this final state

— the bounds are derived from the search of first generation leptoquarks and assuming that the tagger has missed a b-jet [500, 501]. ATLAS has performed a search on RPV stops decaying to $b\mu^+$ and be^+ [502]. ATLAS uses a cut on hadronic transverse momentum together with a Z-veto for the invariant mass of the lepton pair. Our selection of signal events was somewhat different. In addition to the Z-veto we used the a cut on the transverse momentum of the lepton with the highest transverse momentum. The SM background produces very few events with $p_T > 200$ GeV, whereas they are typical in the signal as can be seen from the figures of section VII in **III**.

The stop can decay as a leptoquark also if R-parity is violated [503, 504]. In the RPV case the leptoquark final state is dominant only if the stop is the LSP. Otherwise the constraints from lepton number violating processes constrain the coupling so small that the R-parity conserving decay modes may dominate and the signature is a top or bottom and multileptons coming from the LSP [505, 506]. In the $R = -L$ model the stop need not to be the LSP for the branching fraction to be^+ to be large.

The case where the stop is lighter than the top quark was analyzed in [507]. They find that the decay mode $\tilde{t} \to t^*\nu \to Wb\nu$ can be kinematically distinguished from $\tilde{t} \to Wb\tilde{G}$, where $\tilde{G}$ is the gravitino. In general the stealth stop case is hard to dinstinguish from top pair production, since the kinematics are similar and the production cross section is of the same order as the uncertainty of the SM cross section for top production (see sec. VIII of **III**).

4.4 Models of spontaneous R-parity violation

There are several models, in which non-zero sneutrino VEVs can exist. If the model conserves lepton number, sneutrino VEVs lead to a Goldstone boson, usually called majoron [244, 508]. This gives a strong phenomenological constraint, since the majoron is always accompanied with a light scalar and these will give a contribution to the invisible decays of the Z-boson [509, 510]. Such a contribution for doublets corresponds to a half of a neutrino generation, which is experimentally excluded.

There are three ways out. One is to build the model so that the majoron will be mostly a gauge singlet so that it will not couple to the Z-boson. The second one is to gauge the lepton number by extending the gauge group with $\mathrm{U}(1)_{\mathrm{B-L}}$ so that the majoron will become the longitudinal polarization state of the Z'-boson. The third one is to include explicit lepton number violation so that the majoron will get a mass.

A model representing the first solution is the one proposed by Masiero, Romao, Santos and Valle [511, 512]. It is based on the superpotential[2]

$$W = Y^{(U)}(Q\cdot H_u)U^c + Y^{(D)}(Q\cdot H_d)D^c + Y^{(L)}(L\cdot H_d)E^c + \lambda((H_u\cdot H_d) - \epsilon^2)S$$
$$+ \mu(H_u\cdot H_d) + Y^{(\nu)}(L\cdot H_u)N^c + \kappa S N^c \nu_S + M N^c \nu_S + M' S^2 + \lambda S^3, \quad (4.3)$$

where ν_S are singlets with a lepton number $L = +1$. If $\langle\nu_S\rangle, \langle N^c\rangle \gg \langle\nu_L\rangle$ the majoron will be mostly singlet.

The second solution was introduced by Valle from a superstring-inspired setup [514]. The phenomenological implications at LEP were discussed in [515]. In this case one must extend the particle content with right-handed neutrinos. The Z'-boson needs to get a mass clearly larger than m_Z and this requires a larger VEV than what the left-handed sneutrino can have as shall be discussed later. The superpotential of [515] contains in addition to the MSSM Yukawa couplings the neutrino Yukawa term and the $S(H_u \cdot H_d)$-term.

A model representing the third solution is that proposed by Kitano and Oda [450], which we studied in **I**. The superpotential was given in Eq. (3.46). Spontaneous R-parity violation is generated by assuming large A-terms compared to soft supersymmetry breaking scalar masses. There is an upper bound for the ratio[3] $|A/m| < 3$, above which the scalar field with the smallest Yukawa coupling will acquire a VEV [516]. In the MSSM going beyond the limit would lead to a charge (and possibly color) breaking vacuum. In [450] it was assumed that the neutrino Yukawa couplings would be the smallest ones and then the deepest minimum would be the one, where the sneutrino gets a VEV. Such a small neutrino Yukawa coupling also suppresses the left-handed sneutrino VEV.

The term $\frac{\lambda_N}{2}N^2S$ violates lepton number by two units and hence the would-be-majoron gains a mass. Since the mass comes from lepton number violating terms, it will vanish in the limit $\lambda_N \to 0$. The constraint from invisible Z-boson decays then gives a lower bound for $\lambda_N \gtrsim 0.1$ [517].

There are some common features to all models with spontaneous RPV. First the VEV of the doublet sneutrino is constrained by neutrino masses. If the SUSY breaking scale is $\mathcal{O}(\mathrm{TeV})$ the doublet sneutrino VEVs should be less than 1 MeV [518]. The VEVs of singlet sneutrinos are almost unconstrained. Hence, in many cases one may simply neglect the VEVs of doublet sneutrinos and consider only the VEVs of singlet sneutrinos. The neutrino Yukawa terms then generate the bilinear R-parity violating terms $\mu_i(L_i \cdot H_u)$, where $\mu_i = Y^{(\nu)}_{ij}\langle N^c\rangle_j$. Many phenomenological features of spontaneous R-parity violation are equivalent to bilinear RPV.

The effective bilinear terms can also be converted to trilinear terms with a

[2]The original model of [511] was shown not to have a stable R-parity violating vacuum [513]. The extended model of [512] has additional terms, which help to satisfy experimental constraints and can give an R-parity violating minimum.

[3]Here m denotes the soft supersymmetry breaking scalar mass, all of which are assumed equal and A multiplies the terms of the superpotential in the trilinear terms of the soft supersymmetry breaking Lagrangian.

change of basis. By defining

$$H' = \frac{\mu H_d + \sum \mu_i L_i}{\sqrt{\mu^2 + \sum \mu_i^2}} \tag{4.4}$$

the only bilinear term is the $H_u H'$-term but Yukawa couplings generate R-parity violating terms when written in the new basis. These terms are proportional to the Yukawa couplings, although they depend also on how the new lepton superfields are defined. The sneutrinos will have VEVs in this basis. Another common basis choice is the one with

$$H'' = \frac{v_d H_d + \sum v_i L_i}{\sqrt{v_d^2 + \sum v_i^2}} \tag{4.5}$$

where the new Higgs field H'' is having a VEV but the sneutrinos have $\langle\tilde{\nu}\rangle = 0$. Also in this case trilinear terms proportional to the Yukawa couplings will be generated but there will be bilinear LH_u-terms also. In addition to the bilinear RPV terms in the superpotential there will be corresponding soft supersymmetry breaking terms. Those terms cannot in general be rotated away with the same basis transformation as the superpotential terms.

Since lepton number and R-parity are not conserved, the leptonic sector will mix with superpartners having the same charge and spin. This includes the mixing of neutrinos and neutralinos, charged leptons and charginos and Higgs bosons and sneutrinos. The mixing in the spin-1/2 sector needs to be small to satisfy phenomenological constraints as will be discussed below, but the mixing of the Higgs and the right-handed sneutrino may be substantial as was discussed in **I**.

Spontaneous RPV generates neutrino masses [519, 520]. This constrains the parameters related to RPV. Neutrino masses from bilinear R-parity violation were considered in [521]. Viable mass spectra were achieved with μ_i being a few MeV's. Hence, if the right-handed sneutrino VEVs are at the TeV scale, the neutrino Yukawa couplings should be rather small, $\mathcal{O}(10^{-6})$.

The smallness of neutrino masses also explains why models of spontaneous R-parity violation satisfy the constraints of lepton flavor violation. The smallness of doublet sneutrino VEVs suppresses lepton flavor violation originating from the doublet sector and the smallness of neutrino Yukawas (or singlet sneutrino VEVs) suppresses lepton flavor violation from the singlet sector. There can be rather large mixing in the scalar sector (if $\lambda_H \lambda_N$ is large in the case of superpotential (3.46) or if $\lambda\kappa$ is large in the case of superpotential (4.3)), but that contribution to lepton flavor violation will be suppressed by loop factors and small Yukawa couplings at low energies.

Since R-parity is a discrete symmetry, its spontaneous breaking leads to a potential domain wall problem. In some models the problem is solved trivially. If lepton number is conserved, the degenerate vacua are also connected via the continuous $\mathrm{U}(1)_\mathrm{L}$ symmetry and hence long-range correlations between field variables can exist. Since all of the minima can be continuously transformed to each other, there will be no domain walls.

The case, where lepton number is explicitly broken, is more troublesome as the Goldstone mode does not connect the degenerate vacua. The way to avoid domain walls in the NMSSM with nonrenormalizable terms and symmetries is somewhat questionable, since that would mean that we need to introduce RPV terms in the nonrenormalizable part of the superpotential, whereas the renormalizable part would be R-parity conserving. Technically that is possible, although various phenomenological constraints need to be satisfied. The inflationary solution, discussed in the context of NMSSM in [441], works also in the case of spontaneous R-parity violation.

4.5 Collider imprints of supersymmetry without R-parity

The crucial difference between RPV and R-parity conserving (RPC) supersymmetry is that RPV allows the LSP to decay to SM particles. Usually it is assumed that the heavier superpartners decay promptly to SM particles and one or several LSPs. The lifetime of the LSP leads to three scenarios. If the lifetime is large enough for the particles to escape the detector, the collider signature is the same as in RPC supersymmetry if the LSP is neutral. If the lifetime is $10^{-12} \dots 10^{-8}$ s, the LSP will be seen as displaced vertices. If the lifetime is even shorter, the decays will be prompt.

We shall first have a look at the implications of RPV to collider searches with an example, stop pair production. We assume that both λ, λ' type couplings are nonzero and all superpartners decay promptly. As mentioned on p. 58, the stops are expected to be the lightest squarks. In RPV supersymmetry the LSP is not a dark matter candidate, so it may be charged. Hence the stop can also be the LSP and in such a case it will decay directly to SM particles via $\tilde{t} \to q\ell^+$. If it is not the LSP, it will decay via $\tilde{t} \to t\tilde{\chi}^0, b\tilde{\chi}^+$, where the higgsino-like channel is preferred if it is kinematically allowed. The lighter superpartner then decays to SM particles, and the heavier one either directly or via the lighter state.

If a neutralino is the LSP, it will decay to either $\tilde{\chi}^0 \to \ell^\pm \tilde{\ell}^{*\mp} \to \ell^+\ell'^-\nu$, $\ell qq'$, $\tilde{\chi}^0 \to \nu\tilde{\nu}^* \to \ell^+\ell'^-\nu$, $\nu qq'$ or $\tilde{\chi}^0 \to q\tilde{q} \to qq'\ell$ depending on the masses of the superpartners and the sizes of the RPV couplings. If the chargino is the LSP, it will decay via $\tilde{\chi}^+ \to \tilde{\ell}^+\nu, \ell^+\tilde{\nu}, q\tilde{q}$ and the sleptons decaying either to leptons or jets and squarks to a quark and a lepton. Hence the overall signature of the decaying LSP will be one or more leptons, possibly jets and some missing transverse energy — in the R-parity conserving case this all would simply be missing transverse energy. Some examples of stop decay modes are given in figure 4.1.

Even in the rather simple case of a stop pair, there is a multitude of possible final states. If the stop is the LSP, the signature consists of leptons and quarks, but in all other cases there will be b-jets, charged leptons and possibly decay products of W-bosons, jets from lighter quarks and some missing transverse energy. In the R-parity conserving case the signature would be a b-jet, a decay-

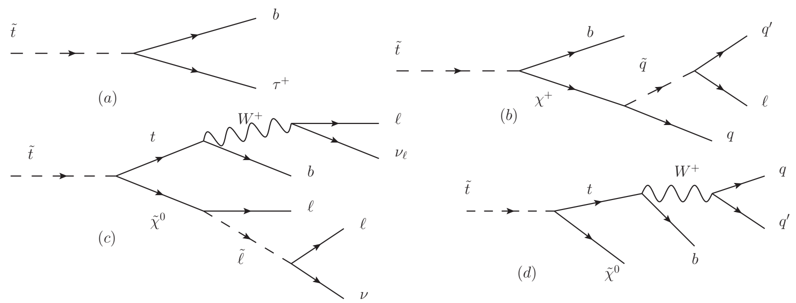

Figure 4.1: Stop decays in various scenarios. Diagram (a) represents a case where the stop is the LSP and it decays directly to SM particles. Diagrams (b) and (c) are examples of cases, where the stop decays to lighter superpartners, which subsequently decay to SM particles. Diagram (d) is an example of stop decays in R-parity conserving supersymmetry. R-parity violation allows a larger variety of final states and produces less missing transverse energy.

ing W-boson and missing transverse energy assuming that there are only a few superpartners lighter than the stop.

In general the collider signatures depend on the LSP and how R-parity is broken. If we have spontaneous R-parity violation, only lepton number violating operators are present and hence the LSP will decay leptonically or semileptonically. Those decay modes are also possible in the case of the $R = -L$ model, since the λ, λ' couplings are allowed, even though they are not lepton number violating. If only the λ'' couplings were allowed, the LSP would decay to hadrons.

We review some of the limits of RPV SUSY searches. Searches for displaced hadronic and leptonic vertices have been performed [522, 523]. No signal was observed and ATLAS gave limits of 0.15 fb for the production cross section of displaced vertices from new physics. Both ATLAS and CMS gave exclusion limits for squark production as a function of $c\tau$, where τ is the mean life time. Squarks of 700 GeV were excluded over a large range of lifetimes leading to displaced vertices.

In the case of prompt decays spontaneous RPV leads to (multi)lepton final states and possibly some jets or missing transverse energy. Both CMS and ATLAS have performed such searches, though they are interpreted in terms of explicit R-parity violation and often assuming a single dominant trilinear RPV coupling. Whereas the RPV terms involving the third generation are usually larger due to the proportionality to Yukawa couplings, there are typically more couplings than just a single one that are relevant. Hence the bounds on superpartner masses are only indicative in the case of spontaneous R-parity violation. The typical signature of spontaneous RPV consists of decay modes coming from both LLE-type and LQD-type operators.

None of the analyses sees any statistically significant signal. Assuming a neutralino LSP, which decays leptonically, ATLAS gets a bound of 1.35 TeV for a gluino NLSP, 750 GeV for a chargino NLSP and 400 GeV for a slepton NLSP [524]. This search analyzes many possible final states assuming only the nature of the LSP and NLSP and hence does not rely too heavily on the single coupling dominance. ATLAS searches also heavy resonances decaying to opposite sign leptons of different flavor [525]. This can be interpreted as bounds on sneutrino masses (which are beyond 1.7 TeV) but this decay mode is dominant only if sneutrino is the LSP.

CMS has analyzed the case of a stop LSP/NLSP (neutralino LSP if stop is the NLSP) and found bounds up to 1 TeV but these depend strongly on the dominant coupling [526]. For the channel $\tilde{t} \to b\tilde{\chi}^+$ and the chargino decaying to leptons the bound is 890 GeV [527].

Leptoquark searches [528,529] can also be interpreted as searches for squarks in RPV supersymmetry, but the bounds are derived assuming a single dominant coupling. The bounds are only slightly weaker in RPV supersymmetry due to other nonzero RPV couplings if the squark is the LSP. If there are lighter superpartners, the squark will mainly decay via RPC decay modes and the leptoquark mode may have a small branching ratio.

In many cases leptonic RPV leads to hard leptons, which are easy to trigger. Especially at the higher end of leptonic p_{T} the SM background is small. Hence the bounds on masses are often higher than in R-parity conserving supersymmetry, typically $\mathcal{O}(1\text{ TeV})$ for strongly interacting particles. The baryon number violating RPV couplings lead to signatures that have a significant background from standard QCD processes, which can hide the signal. The bounds in such a case are not as restrictive [530–532] as in the leptonic decay modes.

4.6 Higgs and spontaneous R-parity violation

Spontaneous R-parity violation alters the phenomenology of the Higgs sector. First, since lepton number is not conserved, sneutrinos have to be counted in the Higgs sector, too. The mixing between $\tilde{\nu}_L$ and the Higgses is suppressed due to the smallness of the sneutrino VEV. Hence the left-handed sneutrinos essentially decouple from the rest of the Higgs sector.

The right-handed sneutrinos may mix largely with the Higgs. In fact we showed in **I** that in the model of Kitano and Oda [450] it is not possible to decouple the right-handed sneutrino from the Higgs sector as the mixing terms grow if we try to make the sneutrino-dominated state heavy. Hence, as we showed, even though there are additional terms to the 2×2 submatrix from the couplings with singlets, the lightest scalar cannot be heavier than m_Z once R-parity is violated spontaneously. The only way the Higgs mass can be lifted beyond m_Z at tree-level is to have a lighter scalar. In such a case the mixing between the lightest scalar and the Higgs can increase the Higgs mass well beyond m_Z. Such a mixing is possible generally in models of spontaneous RPV.

The bound that we found on the lightest scalar mass is not a generic feature of spontaneous RPV. It is due to the fact that the singlet S couples to both $\tilde{N}$ with $H_{u,d}$ in the superpotential and that in the soft supersymmetry breaking Lagrangian the only mixing term between $\tilde{N}$ and H_u is proportional to the very small combination $A_\nu v_{\tilde{\nu}}$. If we had not included the explicitly lepton number violating Majorana mass term (like the model of [515]) in the superpotential in **I**, we would have been able to decouple the sneutrino from the MSSM Higgs sector.

In [533] the authors studied spontaneous RPV in the context of the $\mathrm{U(1)_{B-L}}$ model. Since there only the neutrino Yukawa coupling was added to the MSSM superpotential, there will be no large mixing between the Higgs and the right-handed sneutrino. Hence their predictions for the lightest Higgs mass did not differ much from the MSSM.

Even though spontaneous RPV creates lepton number violation, the SM-like Higgs does not have lepton number violating decays unless the decay $h \to \tilde{\chi}\tilde{\chi}$ is kinematically allowed. The subsequent signals from the lepton number violating decays of the neutralinos[4] have been analyzed in [534]. These decays lead to multilepton final states with missing transverse energy and hence they have a reasonably low background at the LHC.

The heavy Higgs does have a lepton number violating decay mode[5] $H \to h\tilde{N}$ if it is kinematically allowed. The SM-like Higgs can have an additional decay mode $h \to \tilde{N}\tilde{N}$ if it is kinematically allowed. A light right-handed sneutrino does not have too many decay modes available. It may decay to photons via a charged Higgs loop or undergo a three-body decay via νh^*, the latter being suppressed by the small neutrino Yukawa coupling. For the same reason the Higgs-sneutrino mixing does not alter the branching ratios too much. The larger effect may come from the deviation of the $hb\bar{b}$-coupling from its SM value as we discussed in section 3.3 of **I**.

The Higgs may have lepton flavor violating decay modes with RPV, although they are suppressed by either the small left-handed sneutrino VEV or the sneutrino Yukawa. In models with explicit RPV one may get a branching ratio $\mathcal{O}(10^{-5})$ for the decay $h \to \mu\tau$ [535]. So far the sensitivity of the searches for this mode are not at that level of precision.

Altogether the SM-like Higgs does not produce significant constraints for RPV. The absence of any nonstandard Higgs decay modes constrains mostly the case of superpartners lighter than $m_h/2$. Eventually one will start to constrain the universal suppression of all Higgs production and decay modes due to the mixing with sneutrinos.

[4] The authors had a model with explicit RPV. In this case there is no essential difference to spontaneous RPV.

[5] We denote here with $\tilde{N}$ the sneutrino-dominated physical state.

Chapter 5

Left-right symmetric supersymmetry

5.1 Left-right symmetric models

The basic idea of left-right symmetric models [241, 242, 536] is that parity is a symmetry of Nature, which is broken spontaneously. The gauge group of the SM must be extended by introducing right-handed weak interactions. In a left-right symmetric theory one should introduce the right-handed fermions as doublets. If one then generalizes the definition of electric charge by the most evident way by setting

$$Q = T_L^3 + T_R^3 + X, \tag{5.1}$$

where X is the quantum number related to the U(1) symmetry, one finds that $X = (B - L)/2$, where B is baryon number and L is lepton number. Hence the gauge group of left-right symmetric models is $\mathrm{SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)}_{B-L}$.

There are different left-right symmetric models depending on how the gauge symmetry is broken to the SM gauge group. One may also introduce some features that are not explicitly symmetric under parity, *e.g.* assign different values to the gauge couplings of $\mathrm{SU}(2)_L$ and $\mathrm{SU}(2)_R$.

5.1.1 Field content

Left-right symmetry implies that right-handed fermions are in doublets of $\mathrm{SU(2)_R}$. One may notice immediately that a right-handed neutrino is a necessary part of left-right symmetric models. Hence one may introduce Dirac masses for the neutrinos. However, the lepton number violating Majorana mass term for the right-handed neutrino is forbidden by the gauge symmetry.

The fact that left- and right-handed fermions are doublets means that the Higgs that gives masses to the SM fermions must be a bidoublet, *i.e.* be in the $(2, 2)$ representation of $\mathrm{SU(2)_L \times SU(2)_R}$:

$$\Phi = \begin{pmatrix} \varphi_1^0 & \varphi_1^+ \\ \varphi_2^- & \varphi_2^0 \end{pmatrix}. \tag{5.2}$$

The bidoublet has two neutral components, which couple to different components of the doublets and hence it could in principle give masses to up- and down-type quarks and leptons. If one would use only a single Yukawa coupling for all quarks, we should have $m_b/m_t = m_s/m_c = m_d/m_u = \langle\varphi_2^0\rangle/\langle\varphi_1^0\rangle$, which is not satisfied. As we discuss later, such a setup would also provide a large mixing between the left- and right-handed gauge bosons, which is experimentally excluded. To avoid a large left-right mixing in the gauge sector, one of the neutral components of the bidoublet must be inert or the VEV should be tiny. In the following we shall simply assume that only one bidoublet component gets a nonzero VEV.

In the nonsupersymmetric case one may proceed like in the SM and use the charge-conjugated bidoublet $\Phi^c = \tau_2\Phi^*\tau_2$ to get two Yukawa matrices for quarks, which allow to generate masses and CKM-mixings [537–540]. In the supersymmetric version of left-right symmetry one needs two bidoublet superfields, which we denote by Φ and X in order to get the correct pattern of masses and mixings.

The bidoublets will give the gauge bosons of $\mathrm{SU(2)_L}$ and $\mathrm{SU(2)_R}$ a mass but they will be equal. One needs still additional scalars, which are singlets under $\mathrm{SU(2)_L}$ but charged under $\mathrm{SU(2)_R}$ to make the right-handed gauge bosons heavy. The most common choice is to use triplets [14, 541, 542], since they simultaneously allow to generate masses for neutrinos. However, one may also use doublets to break $\mathrm{SU(2)_R}$ and assume neutrino masses to be generated via operators of higher dimension.

In nonsupersymmetric theories it is enough to have just one triplet Δ_R in the $(1,1,3,2)$ representation, which then implies via left-right symmetry that there should also be a triplet Δ_L in the $(1,3,1,2)$ representation. The triplets are usually written in the 2×2 form like

$$\Delta_{L,R} = \begin{pmatrix} \delta^+/\sqrt{2} & \delta^{++} \\ \delta^0 & -\delta^+/\sqrt{2} \end{pmatrix}. \tag{5.3}$$

In supersymmetric models one also needs to add triplets with $B-L=-2$ to cancel anomalies. Hence the Higgs content of left-right symmetric models is very large. In the nonsupersymmetric case even the minimal field content[1] requires four CP-even, two CP-odd, two singly charged and two doubly charged physical scalars. In the supersymmetric version there are eight CP-even, six CP-odd, six singly charged and four doubly charged states.

The superpotential for a left-right symmetric model with triplets is

$$\begin{aligned} W = {} & Y_Q^1 Q_L^T \Phi Q_R + Y_Q^2 Q_L^T X Q_R + Y_L^1 L_L^T \Phi L_R + Y_L^2 L_L^T X L_R + Y_L^3 L_L^T \Delta_{1L} L_L \\ & + Y_L^4 L_R^T \Delta_{2R} L_R + \mu_L \mathrm{Tr}(\Delta_{1L}\Delta_{2L}) + \mu_R \mathrm{Tr}(\Delta_{1R}\Delta_{2R}) + \mu \mathrm{Tr}(\Phi\tau_2 X\tau_2) \\ & + \mu_\Phi \mathrm{Tr}(\Phi\tau_2\Phi\tau_2) + \mu_X \mathrm{Tr}(X\tau_2 X\tau_2), \end{aligned} \tag{5.4}$$

[1]If $\mathrm{SU(2)_R}$ triplets are chosen to break the symmetry.

where $Y^i_{Q,L}$ are 3×3 Yukawa matrices, τ_2 is the second Pauli matrix and Δ_{1L} and Δ_{2R} are the triplets with $B - L = +2$. One can also introduce a singlet superfield S and generate the couplings between the Higgs multiplets dynamically by replacing $\mu_i \to \lambda_i S$.

The field content of left-right symmetric supersymmetry (LRSUSY) has one clear advantage compared to the MSSM: Since there are no right-handed singlet superfields and the Higgses are bidoublet fields, none of the R-parity violating terms of (3.25) are possible. You need not to impose R-parity, it is dictated by the symmetries and the particle content. This is due to $B - L$ being a part of the gauge symmetry, which prevents the interactions leading to fast proton decay [543, 544].

5.1.2 Bounds on the gauge sector

One may immediately notice that left-right symmetric models predict the existence of new gauge bosons, denoted by $W^{\pm}_R$ and Z_R[2]. They have decay modes to dijets and dileptons and searches for such resonances have been made. The leptonic decay channels of the W_R differ from its left-handed counterpart, unless neutrinos are Dirac fermions. If there is a Majorana mass term — as is the case in LRSUSY with triplets — the right-handed neutrinos will decay.

The decay mode will be a charged lepton and an off-shell W_R, which then subsequently decays. In the case of the lightest right-handed neutrino the decay will be to a pair of jets and that is the dominant decay channel for the other generations, too. Hence the signal will be two charged leptons and two jets. Due to the Majorana nature of the neutrino, the charged lepton coming from its decay may be of either sign. Hence one should see a similar signature on both same sign and opposite sign leptons [545]. Such a signature would be a signal of both the extension of the gauge group and the Majorana character of neutrinos.

CMS made a dedicated search for right-handed muon neutrinos and W_R bosons [546]. There were no significant excesses. Assuming $g_R = g_L$ they were able to exclude W_R masses up to 2.5 TeV depending on the right-handed neutrino mass. ATLAS has made a search in the dijet mode and got a bound of 2.6 TeV [547], whereas CMS got a bound of 2.7 TeV [548]. These bounds apply if there are no allowed decay modes to non-SM particles. If the decays to superpartners are kinematically allowed, the bounds become slightly weaker.

Searches for the Z_R have also been made [547, 549, 550]. The bounds are typically around 3 TeV. However, if the W_R is beyond 2.5 TeV, the Z_R will be beyond 4 TeV in left-right symmetric models. Hence the bounds on W_R are more restrictive. The mass bound on W_R requires the right-handed triplet VEV to be in the multi-TeV range. This will make many — but not all — triplet states heavy both in the scalar and the fermionic sector.

[2] All three neutral physical gauge bosons are mixtures of the pure gauge eigenstates B^0, W^0_L and W^0_R. Hence the heavy boson is not only right-handed unlike the right-handed W. The largest component is W^0_R, which motivates the notation used here.

In addition to the bound on masses, there is a bound on mixing between the W_L and W_R. The mixing term between W_L and W_R in LRSUSY is $\frac{1}{2}g^2(\langle\varphi_1^0\varphi_2^0\rangle + \langle\chi_1^0\chi_2^0\rangle)$, where φ_i^0 and χ_i^0 are the neutral components of Φ and X, respectively. The mixing angle can be constrained from *e.g.* muon decays or kaon physics [551] and the bounds are $\mathcal{O}(0.03)$ when light right-handed neutrinos are not assumed. Hence only one component of a bidoublet can have a large VEV.

5.1.3 Vacuum stability

The LRSUSY model where $\mathrm{SU(2)_R}$ is broken by triplets has a problem with its vacuum [552, 553]. In the supersymmetric limit at tree-level there is a set of degenerate charge-breaking vacua, which are connected via a doubly charged Goldstone mode. When supersymmetry is broken, a preferred vacuum is picked and this will be charge breaking and the doubly charged Higgs mass matrix will have a negative eigenvalue. There have been various approaches to solve this problem.

The first solution was the one suggested by Kuchimanchi and Mohapatra. Spontaneous R-parity violation can stabilize the vacuum [552]. Right-handed sneutrinos can indeed have large VEVs as was discussed in chapter 4 and that may indeed provide a stable vacuum. This makes the LSP unstable and hence one has problems with finding a dark matter candidate but otherwise this setup leads to a viable model.

Mohapatra and Rasin then showed that nonrenormalizable terms in the superpotential may be enough to guarantee a charge conserving vacuum [554,555]. The most important operator will be $\frac{\lambda}{M}[\mathrm{Tr}(\Delta_{1R}\tau\Delta_{2R})]^2$. By choosing the sign of λ one may make the energy of the charge conserving minimum lower than the charge breaking one. A larger set of nonrenormalizable terms was analyzed in [556] with similar conclusions.

Aulakh *et al.* [557] proposed the use of yet more triplets $\Omega_L \in (1, 1, 3, 0)$ and $\Omega_R \in (1, 3, 1, 0)$, which can also help to achieve a stable vacuum. In such a setup spontaneous R-parity violation would result in a charge breaking vacuum [558]. Hence if the vacuum is neutral, R-parity is conserved. There are a few phenomenological studies on this model [559, 560].

Babu and Mohapatra noted that the vacuum problem can also be cured by radiative corrections [561]. The one-loop corrections from the right-handed leptons and sleptons had an effect on the right-handed triplet masses. The radiative corrections could make the charge conserving vacuum stable and lift the doubly charged Higgs mass up to a few hundred GeV's. If the left-handed fields had VEVs, their radiatively corrected masses would be $\mathcal{O}(v_L)$ and since left-handed triplet VEVs are severely constrained by the ρ parameter, they would be too light [562]. However, there is no phenomenological problem with the left-handed triplets being inert.

If the corrections from the lepton-slepton sector are dominant, this entangles the right-handed triplet VEV to the soft supersymmetry breaking slepton masses. Namely, if the sleptons were significantly lighter than v_R, the one-loop correction would turn negative [563]. The full one-loop corrections were analyzed in [564]. The authors noted that also the gauge sector and the Higgs sector can contribute significantly to the doubly charged Higgs mass and thus this opens more options to achieve a stable and neutral R-parity conserving vacuum.

We used the radiative corrections to lift the doubly charged Higgs mass in **II** and **IV**. In **II** only the corrections from the right-handed lepton sector were taken into account. This correction is the dominant one, if sleptons are heavy and the Yukawa coupling between the leptons and the triplet is large. In **IV** we wanted to study a sneutrino LSP and hence the main correction needed to come from the gauge and Higgs sectors of the model. That was indeed possible if the corresponding couplings were large and the gauginos somewhat heavy.

The doubly charged Higgs masses we were able to obtain were at most slightly above 500 GeV. The bounds from same sign dilepton searches [287, 288] exclude a doubly charged Higgs at these masses unless it decays dominantly to same sign taus. In that channel the current mass bound is 396 GeV for left-handed doubly charged Higgses [288]. For right-handed doubly charged Higgses the bound is lower due to the reduced coupling to Z-bosons. CMS does not quote a bound for right-handed doubly charged Higgses, but one can estimate that the bound is 300 GeV or a bit more.

The right-handed neutrinos get their masses through the same Yukawa couplings that govern the doubly charged Higgs decay and hence also the other couplings must be nonzero, but an order of magnitude smaller so that the decays to these modes are suppressed.

5.1.4 Neutralinos and charginos

The model has also a rich neutralino and chargino spectrum. In addition to three neutral gauginos, the bidoublets and triplets give eight more neutralinos and if a singlet is included, the total number of neutralinos is 12. The full mass matrix of neutralinos is given in [565].

The spectrum depends on further assumptions. If we don't include singlet self-interactions in the superpotential as in **II** there will be a light singlet-dominated neutralino. The spectrum depends mostly on the soft supersymmetry breaking terms and the (effective) μ-term. At tree-level two bidoublet Higgsinos form a Dirac fermion with a mass $|\mu_{\text{eff}}|$ and the left-handed triplet neutralinos decouple if the triplet scalars are inert. Otherwise all of the neutralinos mix and hence the analysis is complicated.

If the gauginos are lighter than the higgsinos and we have universal gaugino masses, the bino-dominated state will be the lightest neutralino. Since the gauge group is $\mathrm{U}(1)_{B-L}$ and not $\mathrm{U}(1)_Y$ as in the MSSM, the bino will have somewhat different properties. For instance, its coupling to the SM-like Higgs is small, since the bidoublets are not charged under $\mathrm{U}(1)_{B-L}$.

If the bino-like state is the LSP, there will often be leptons in the decay chains of superpartners and this gives the best chance to see a signal at hadron colliders. In [565] some benchmarks were considered and in the case of light sleptons, the signals in dilepton and multilepton channels would be stronger than in the MSSM. Some of the benchmarks considered are now excluded.

There are six charginos and two doubly charged charginos. In addition to the MSSM particle content there is a new charged gaugino, a new charged Higgsino from the bidoublets and two charginos from the triplets. The doubly charged charginos come from the triplets. The phenomenology of doubly charged charginos was considered in [566, 567]. The multilepton $+\not{E}_T$ signature has a small background at hadron colliders and allows to find or exclude doubly charged Higgsinos up to $\mathcal{O}(500)$ GeV's.

5.2 The SM-like Higgs in LRSUSY

5.2.1 The Higgs mass

The number of neutral scalars in LRSUSY is large. There are two MSSM-like bidoublet Higgses and two inert bidoublets and the $\mathrm{SU}(2)_\mathrm{R}$ breaking fields provide also new neutral scalars. The model we studied in **II** has nine CP-even scalars and seven physical CP-odd scalars. The size of the mass matrices makes the analysis challenging.

Like any multi-Higgs model, the tree-level bound on the lightest Higgs mass can be derived by using only the 2×2 submatrix. Due to the bidoublet nature of the Higgs, the expression $g^2 + g'^2$ of the SM will be replaced by $g_L^2 + g_R^2$ in LRSUSY. If $g_L = g_R$, the tree-level bound will be $\sqrt{2}m_W$ [568], the more general expression

$$m_h^2 \leq m_W^2 \left(1 + \frac{g_R^2}{g_L^2}\right) \cos^2 2\beta \tag{5.5}$$

was derived in [569], though there is also an argument that the bound is too optimistic [570]. In any case it is clear that the upper bound at tree-level is higher than in the MSSM. One may also note that when $\tan\beta$ is close to one, the tree-level mass goes to zero. Hence the 125 GeV Higgs mass implies a lower bound for $\tan\beta$, although the bound depends on the superpartner masses.

Although there are triplets in the model, they do not help in lifting the tree-level bound, since one cannot write a gauge invariant superpotential coupling between bidoublets and triplets, which would contribute to the 2×2 submatrix. If we have a singlet, the triplets can make a contribution to the Higgs mass at the loop level since the F-terms in the scalar potential introduce a coupling between the bidoublets and triplets[3] [562]. The loop contribution is large, when the triplet VEVs are large, but this also increases the tree-level mixing between the triplets and the bidoublets. The mixing will push the lightest scalar mass down and hence the main extra contribution to the Higgs mass comes from the extended gauge structure of the model.

[3]At tree-level the effect of these terms cancel when the smaller eigenvalue of the 2×2 submatrix is considered.

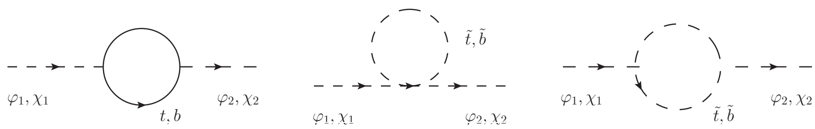

Figure 5.1: Diagrams which lead to the mixing of inert and active bidoublets. The vertices are proportional to y_t and y_b, $y_t y_b$, A_t and A_b, respectively. The mixing increases at large $\tan\beta$, when $y_b = \sqrt{2}m_b/v\cos\beta$ becomes large.

5.2.2 The Higgs couplings

The 125 GeV Higgs couplings were analyzed in **II**. The most significant difference compared to the SM is the possibility of a loop-induced mixing in the Higgs sector, which mixes the inert bidoublets with the SM-like state. The inert bidoublets couple to down-type quarks via the up-type Yukawa couplings and vice versa, so especially the coupling to bottom quarks may get a part that is proportional to the top Yukawa. This can either enhance or suppress the bottom decay width. Since the bottom has the largest branching ratio, any change in it will lead to a large deviation in all other signal strengths.

The diagrams that contribute to the mixing of the inert and active bidoublets are shown in figure 5.1. The largest term mixing the inert and active[4] bidoublets in the mass matrix will be proportional to $y_t y_b m_t^2 \ln(m_{\tilde{t}}^2/m_t^2)$. The mixing is significant, when $\tan\beta$ is large, since then $y_b = \sqrt{2}m_b/v\cos\beta$ increases. The neutrino masses are generated via the type-II seesaw at TeV-scale so the corresponding Yukawa couplings are small. Hence the tau coupling will not be affected.

The model has a relatively light doubly charged Higgs in addition to the other new charged states. One could expect that these could affect the loop-induced decays. The couplings to the SM-like Higgs are however suppressed. This is due to two effects. First, the mass eigenstates are near to the symmetric and antisymmetric combinations of the gauge eigenstates. For such combinations the contribution to the h-H^{++}-H^{--} coupling from the gauge couplings vanishes. Second, the contribution from F-terms is proportional to $\sin 2\beta$, which is rather small at the values of $\tan\beta$, which allow a 125 GeV Higgs. This is somewhat similar to the supersymmetric version of the Higgs triplet model, where the doubly charged Higgs contribution to $h \to \gamma\gamma$ is also suppressed compared to the nonsupersymmetric case [457].

At loop-level the large number of new charged superpartners can also cause deviations from the SM Higgs coupling predictions. The MSSM-like charged Higgs is not too heavy and the singly and doubly charged charginos can also be rather light. Typically these additional particles generate a 10% spread in the effective couplings between the Higgs and $\gamma\gamma$ or $Z\gamma$.

The latest limits on the doubly charged Higgs mass [288] exclude about 80%

[4] We call a scalar field active, when it has a nonzero VEV.

of the data points of **II**. This and the Higgs data, which shows no large deviation from the SM predictions, favor moderate values of $\tan\beta$ and somewhat heavy superpartners.

5.3 Dark matter in LRSUSY

5.3.1 Dark matter — constraints and searches

There are good reasons to believe that the Universe has a large amount of matter in the form of particles of a species so far unknown to us. The evidence for dark matter (DM) comes from the CMB power spectrum [21, 22], rotational motion in galaxies [23, 24] and the fact that after a galactic collision in the so called Bullet cluster, gravitational lensing shows that the centers of mass for the galaxies are not where they would be based on the luminous matter [571]. The option of astrophysical objects as dark matter is disfavored, because searches for non-luminous massive objects via gravitational lensing give a strong upper bound on the density of such objects [572, 573]. The option of primordial black holes as dark matter is not ruled out as long as they are massive enough. In the following we shall discuss dark matter that is made of particles.

Dark matter can be formed via two mechanisms. If dark matter interacts so weakly that it is not in thermal equilibrium with the rest of the Universe, it can be produced via the so called freeze-in mechanism [574]. In many models of particle physics, the interactions between the dark sector and the SM are so strong that thermal equilibrium is achieved.

In such a case the number density of dark matter particles is governed by the Boltzmann equation

$$\frac{dn}{dt} = -3Hn - \langle\sigma v\rangle(n^2 - n_{eq}^2), \tag{5.6}$$

where n is the number density of dark matter particles, H is the Hubble constant, $\langle\sigma v\rangle$ is the thermally averaged annihilation cross section times the relative velocity and n_{eq} is the number density in thermal equilibrium.

In a cooling Universe the number density drops exponentially until the annihilation of dark matter particles is not fast enough to compete with the expansion of the Universe. In such a case the dark matter freezes out and it dilutes only due to the expansion of the Universe. The relic density can be computed by solving the Boltzmann equation [575]. If there are other superpartners that are close in mass to the LSP, coannihilations may alter the relic density [576].

In many models the constraint from the relic density, measured by Planck to be $\Omega_{CDM}h^2 = 0.1198 \pm 0.0015$ [577], is the one that constrains the parameter space most, since often one gets a too large relic density. A relic density that is below the Planck limit can be acceptable if dark matter has several components.

Dark matter can be searched with colliders. The methods were described in section 3.5.1, where the searches for the LSP were reviewed. Such methods apply to dark matter candidates in other models, too.

The direct detection experiments are based on dark matter collisions with ordinary matter and detecting the recoil. The experiments are most sensitive in the intermediate mass range, 10 GeV$< m_{LSP} <$ 100 GeV, since at lower masses the recoil becomes small and at larger masses the smaller density of dark matter particles reduces the frequency of collisions. Currently the strongest constraints come from the measurements of LUX collaboration [578], where the limit on the spin-independent DM-nucleon cross section is 2.2×10^{-46} cm^2 for a 50 GeV dark matter particle.

Indirect detection experiments try to find signals of decaying dark matter from the space. Dark matter annihilation to SM particles will in the end lead to electron-positron pairs, photon pairs or nucleon-antinucleon pairs. Typically one tries to find either photons, positrons or antiprotons. However, there is always a background coming from astrophysical sources. There have been some claims of a signal at various energies [579–581], but none of the claimed signals has a large statistical significance. Recently the AMS-02 experiment claimed to see an excess of positrons and antiprotons at high energies that would be compatible with a TeV-scale dark matter candidate [582].

5.3.2 Right-handed sneutrino as a dark matter candidate

The MSSM and the NMSSM have a limited number of possible dark matter candidates. The neutralino is by far the most studied dark matter candidate (for MSSM, see *e.g.* [583–589], for NMSSM *e.g.* [590–592]). Also the gravitino is a viable candidate in theories with local supersymmetry. In the case of a gravitino LSP, the next-to-lightest supersymmetric particle or the lightest messenger typically has a long lifetime and the relic density is determined by its production and decays to the gravitino [593–595].

The third type of neutral superpartners in the MSSM, the left-handed sneutrino, is outruled. If the left-handed sneutrino were light it would have given a contribution to the invisible Z-boson decay width. The heavier sneutrinos are ruled out by direct detection experiments in the range where they would give the observed relic density [458, 459].

LRSUSY has naturally another dark matter candidate, the right-handed sneutrino. In models with the (MS)SM gauge structure right-handed (s)neutrinos are gauge singlets. In such a case it could be possible that the right-handed sneutrino only couples to other particles via the neutrino Yukawa coupling, which is often tiny. In such a case the annihilation cross section would be too small and the corresponding relic density too large.

One may enhance the annihilation cross section in various ways. One way is resonant annihilation. In such a case one sets the mass of the LSP close to one half of *e.g.* the Higgs mass. The annihilation cross section will get a sharp maximum at some LSP mass. There will typically be a small interval of masses, which produces the correct relic density but even a small deviation in mass could change the relic density by an order of magnitude.

In the superpotential (3.46) the singlet S induces a coupling between the sneutrino and the Higgs in the F-term scalar potential. Such a coupling need not to be small, which can then increase the annihilation cross section [596–599]. On the other hand the coupling is unknown. The annihilation cross section can also be enhanced by introducing mixing between the left- and right-handed sneutrinos [600–603], which has also been studied in LRSUSY [604]. The left-handed part of the sneutrino can annihilate via the Z-boson interactions and the mixing can make the annihilation cross section right. However, such a large mixing is not typical at least if we have universal A-terms, *i.e.* $A_\nu \propto y_\nu$.

In LRSUSY there is a D-term coupling between the right-handed sneutrino and the Higgs. If we assume the Higgs to be close to the alignment limit the coupling depends only on the value of the right-handed gauge coupling. This coupling determines the annihilation cross section via a s-channel Higgs, which is the dominant annihilation channel. Since the only free parameter in the alignment limit is the sneutrino mass, we may deduce a range of masses for the LSP sneutrino from the observed relic density.

The gauge coupling is so strong that the s-channel Higgs annihilation produces the right relic density without resonant annihilation. Hence the LSP mass does not need to be tuned like in the case of resonant annihilation. In **IV** we found viable relic densities with sneutrino masses between $250 \ldots 290$ GeV.

On the other hand sneutrino pair production via the reversed annihilation diagram is negligible, since the s-channel Higgs would be off-shell and the propagator suppresses the production cross section to the attobarn level. Resonant production via heavier Higgs bosons is suppressed, since the three-point coupling is proportional to the VEV of the mass eigenstate and close to the alignment limit the VEVs of other bidoublet states are nearly zero. On the other hand the triplets have masses of several TeVs.

The best option to find supersymmetry in the case of sneutrino dark matter is to produce a right-handed W-boson and search for its decays to sleptons. For TeV-scale sleptons there will be enough phase space for the branching fraction to be a few percent. Often the final state will have two rather hard leptons from the initial decay products and missing transverse momentum. If the W_R is just above the experimental bound, there will be a clear signal with 100 fb^{-1} of data. If W_R is heavier, the production cross section is smaller, but it is partially compensated by a larger branching fraction to sleptons due to the larger phase space and harder leptons, which help in distinguishing the signal from the background.

In **IV** we studied four benchmark points. The benchmarks with $m_{W_R} \simeq$ 2.7 TeV gave a large signal with both sneutrino and neutralino dark matter in the dilepton $+\not{E}_T$ channel. The benchmarks with $m_{W_R} \simeq 3.5$ TeV also show a clear deviation from the background with 100 fb^{-1}, but a discovery would need $\mathcal{O}(500)$ fb^{-1}.

It seems possible to distinguish the LSP candidates by comparing the decay channels with two and three leptons. This is due to the fact that with a neutralino LSP the decay goes dominantly as $W_R \to \tilde{\ell}\tilde{\nu}_R \to \ell\tilde{\chi}^0 N_R \tilde{\chi}^0 \to \ell\ell jj + \not{E}_T$. With a sneutrino LSP the decay chains are often more complicated, which can lead to additional leptons. The number of three-lepton events compared to dilepton events is then higher in the sneutrino LSP case than in the neutralino

LSP case, but the overall number of trilepton events will be rather small with 100 fb^{-1} so we will need high luminosities to be able to distinguish the nature of the LSP.

Chapter 6

Summary and outlook

The first run of the LHC has confirmed that the Standard Model is a very good effective theory at present collider energies. The discovery of a scalar particle with couplings roughly similar to the Standard Model Higgs has also given us some knowledge of the mechanism that breaks the electroweak symmetry. However, the Standard Model is, for reasons discussed in section 1.1, incomplete and there are several well motivated directions to go beyond the Standard Model.

Supersymmetry remains as a well motivated framework beyond the Standard Model. The 125 GeV Higgs poses a challenge for supersymmetry in its minimal form, as it requires heavy superpartners, which makes the model somewhat fine-tuned. In nonminimal supersymmetric models there are additional contributions both at tree-level and loops to the SM-like Higgs mass so that the fine-tuning will not be as severe. The additional particles in nonminimal models may also give experimental signatures not present in the minimal supersymmetric Standard Model.

In this thesis we have studied some nonminimal supersymmetric models and their phenomenological implications.

R-parity conservation is not theoretically necessary, although it has certain phenomenological advantages. Models without R-parity lead to different collider signatures and especially the canonical signature of missing transverse momentum is softened. We studied the phenomenology of an R-symmetric model and a model with spontaneous R-parity violation. Both of these have the advantage that the proton is stable, since there are no baryon number violating operators.

We derived a new tree-level bound for the lightest CP-even scalar in the model with spontaneous R-parity violation and showed that the mixing with a light right-handed sneutrino can increase the mass of the SM-like Higgs up to the observed value without the need for heavy supersymmetric partners. The mixing effect needed to get sub-TeV stops is still experimentally viable.

In the R-symmetric model we identified the R-charges with the negative of the lepton number. This means that squarks carry lepton number. Also in the R-symmetric model it was possible to have sub-TeV stops and a 125 GeV Higgs. R-symmetry forbids the terms that mix left- and right-handed stops, while large mixing in the stop sector is typical in other supersymmetric models. One of the stops has a unique decay mode $\tilde{t} \to be^+$, whereas the other one has MSSM-like decay modes. The discovery potential in the be^+ mode is well beyond 1 TeV stop masses.

In left-right symmetric models parity is broken spontaneously. Left-right symmetry extends the gauge sector of the SM and has plenty of Higgs bosons, both neutral, singly and doubly charged ones. The neutral vacuum of left-right supersymmetry is unstable at tree-level. If the instability is cured by radiative corrections, a rather light doubly charged Higgs boson is a feature of the model.

We noticed that even though there is a light doubly charged Higgs, its effect to the loop-induced Higgs decays to $\gamma\gamma$ and $Z\gamma$ is rather limited. This is due to the fact that the coupling of the physical doubly charged Higgs to the Higgs is suppressed. This is similar to the supersymmetric Higgs triplet model discussed on page 55, where the effect of the doubly charged Higgses is a lot smaller than it can be in the nonsupersymmetric case.

The largest new effect in the Higgs couplings is a part proportional to the top Yukawa in the Higgs coupling to bottom quarks. This is induced by a loop-level mixing of the bidoublet Higgses. Since the Higgs branching ratio to bottom quarks is large, this can enhance or suppress all other decay modes significantly.

In left-right supersymmetry the right-handed sneutrino is a part of a doublet. That makes a difference to other models, where it has been studied as a dark matter candidate. The annihilation of sneutrinos in the early Universe proceeds mainly via an s-channel Higgs and the corresponding coupling comes from D-terms, *i.e.* is given by the gauge couplings. We found that the sneutrino can be a viable dark matter candidate without resonant annihilation or sizable left-right mixing in the sneutrino sector. If we assume the equality of left- and right-handed gauge couplings, we could predict a range of masses for the sneutrino from the observed relic density.

In a scenario, where sleptons and electroweak gauginos are the lightest superpartners and all strongly interacting superpartners are heavy, it will be challenging to find any signs of supersymmetry at the LHC. In left-right supersymmetry the right-handed gauge sector opens a new possibility as the production cross section of the gauge bosons is rather large and they may decay also to superpartners. We showed that if the W_R is just slightly heavier than the current experimental bounds, we should be able to see a signal of supersymmetry in the sleptonic decays of W_R. The analysis of the supersymmetric spectrum and the nature of the dark matter candidate will need even higher integrated luminosities.

The Higgs discovery showed us that the idea of spontaneously broken symmetries is realized in Nature. It solved, at least partially, the problems of electroweak symmetry breaking and unitarity, but we are still left with several important and interesting problems, which require new physics. The experiments so far have shown that minimal models with small fine-tuning are not the solution to these problems.

Many of us wish that the LHC would find hints of new physics at the energy frontier in the form of new particles. The LHC will hopefully have many years of operation with the full design energy, which would allow to explore the terrain to a few TeVs for particles with strong interactions and close to a TeV for electroweakly interacting particles.

If that will not lead to a discovery, the hope has to be put to the precision frontier. Low energy experiments looking for neutrinoless double beta decay and charged lepton flavor violation, testing the anomalous magnetic moment of the muon, direct detection experiments of dark matter or searches for rare decays could give us hints of new physics. The interpretation of such discoveries will be model dependent but it could lead us to the right track.

The certain thing is that we will get more precise data on the 125 GeV Higgs couplings. As discussed in the previous chapters, if we find deviations from the Standard Model, we may be able to favor some models over other ones. Even if everything looks like the Standard Model, we will be able to constrain the models with better Higgs data.

The optimism of finding new physics quickly after the start of the LHC is gone, but the hard work to explore the routes to new physics continues both on the phenomenological side and the experimental side. Many of these paths will eventually be dead ends, possibly also the ones that we studied, but exploring the unknown and finding new ideas that lead us forward have given joy and excitement to those who participate in these journeys.

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